



XIII Advanced Research Workshop on High Energy Spin Physics

(DSPIN-09)

Proceedings

**XIII Advanced Research Workshop
on High Energy Spin Physics
(DSPIN-09)**

Dubna, September 1–5, 2009

Proceedings

Edited by *A.V. Efremov* and *S.V. Goloskokov*

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The collection includes contributions presented at the XIII Advanced Research Workshop on High Energy Spin Physics (DSPIN-09), (Dubna, September 1–5, 2009), on different theoretical, experimental and technical aspects of this branch of physics. Dedicated to the memory of Yan Paweł Nassalski.

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Contents

Welcome address	
<i>A. Sissakian</i>	9
Jan Paweł Nassalski (1944-2009)	
<i>A. Sandacz</i>	10
Hard collisions of spinning protons: history & future	
<i>A.D. Krisch</i>	12

Theory of spin physics 23

Microscopic Stern-Gerlach effect and Thomas spin precession as an origin of the SSA	
<i>V.V. Abramov</i>	25
Towards study of light scalar mesons in polarization phenomena	
<i>N.N. Achasov</i>	29
Recursive fragmentation model with quark spin. Application to quark polarimetry	
<i>X. Artru</i>	33
Constituent quark rest energy and wave function across the light cone	
<i>M.V. Bondarenco</i>	41
Towards a model independent determination of fragmentation functions	
<i>E. Christova</i>	45
Transversity GPDs from $\gamma N \rightarrow \pi \rho_T N'$ with a large $\pi \rho_T$ invariant mass	
<i>M. El Beiyad</i>	49
Infrared properties of the spin structure function g_1	
<i>B.I. Ermolaev</i>	53
Asymmetries at the ILC energies and B-L gauge models	
<i>E.C.F.S. Fortes</i>	56
On the $q\bar{q}$ -glueball-mixed 0^{++} -meson states in a simple model approach	
<i>S.B. Gerasimov</i>	60
Polarization at Photon Collider. Using for physical studies	
<i>I.F. Ginzburg</i>	64
Cross sections and spin asymmetries in vector meson leptonproduction	
<i>S.V. Goloskokov</i>	69
Two-photon exchange in elastic electron-proton scattering: QCD factorization approach	
<i>N. Kivel</i>	73
$O(\alpha_s)$ spin effects in $e^+e^- \rightarrow q\bar{q}(g)$	
<i>J.G. Körner</i>	77
Transversity in exclusive meson electroproduction	
<i>P. Kroll</i>	82
Spin correlations of the electron and positron in the two-photon process $\gamma\gamma \rightarrow e^+e^-$	
<i>V.V. Lyuboshitz</i>	91
Another new trace formula for the computation of the fermionic line	
<i>M. Mekhfi</i>	95

Analytic perturbation theory and proton spin structure function g_1^p	
<i>R.S. Pasechnik</i>	98
Transverse momentum dependent parton distributions and azimuthal asymmetries in light-cone quark models	
<i>B. Pasquini</i>	102
Colour modification of factorisation in single-spin asymmetries	
<i>P.G. Ratcliffe</i>	108
Impact parameter representation and x -dependence of the transversity spin structure of the nucleons	
<i>O.V. Selyugin</i>	118
Polarized parton distributions from NLO QCD analysis of world DIS and SIDIS data	
<i>O. Shevchenko</i>	122
Polarized PDFs: remarks on methods of QCD analysis of the data	
<i>A.V. Sidorov</i>	126
Potential for a new muon $g - 2$ experiment	
<i>A.J. Silenko</i>	131
Evolution equations for truncated Mellin moments of the parton densities	
<i>D. Strózik-Kotlorz</i>	135
New developments in the quantum statistical approach of the parton distributions	
<i>J. Soffer</i>	139
Polarized hadron structure in the valon model and the nucleon axial coupling constants: a_3 and a_8	
<i>F. Taghavi Shahri</i>	143
Axial anomalies, nucleon spin structure and heavy ions collisions	
<i>O.V. Teryaev</i>	147
Orbital momentum effects due to a liquid nature of transient state	
<i>S.M. Troshin</i>	151
Identification of extra neutral gauge bosons at the ILC with polarized beams	
<i>A.V. Tsytrinov</i>	155
Quark intrinsic motion and the link between TMDs and PDFs in covariant approach	
<i>P. Zavada</i>	159

Experimental results 165

Search for narrow pion-proton states in s-channel at EPECUR: experiment status	
<i>I.G. Alekseev</i>	167
Measurement of tensor polarization of deuteron beam passing through matter	
<i>L.S. Azhgirey</i>	171
Prospects of measuring ZZ and WZ polarization with ATLAS	
<i>G. Bella</i>	175
Latest Results on Deeply Virtual Compton Scattering at HERMES	
<i>A. Borissov</i>	179
Polarization transfer mechanism as a possible source of the polarized antiprotons	
<i>M.A. Chetvertkov</i>	185

Probing the hadron structure with polarised Drell-Yan reactions in COMPASS	
<i>O. Denisov</i>	189
Polarization of valence, non-strange and strange quarks in the nucleon determined by COMPASS	
<i>R. Gazda</i>	197
New Monte-Carlo generator of polarized Drell-Yan processes	
<i>O. Ivanov</i>	205
High energy spin physics with the PHENIX detector at RHIC	
<i>D. Kowall</i>	209
Measurements of the A_{yy} , A_{xx} , A_{xz} and A_y analyzing powers in the $^{12}\text{C}(\vec{d}, P)^{13}\text{C}^*$ reaction at the energy $T_d=270$ MeV	
<i>A. S. Kiselev</i>	217
Overview of recent HERMES results	
<i>V.A. Korotkov</i>	221
NLO QCD predictions for gluon polarization from open-charm asymmetries measured at COMPASS	
<i>K. Kurek</i>	229
Study of light nuclei spin structure from $p(d, p)d$, $^3\text{He}(d, p)^4\text{He}$ and $d(d, p)^3\text{H}$ reactions.	
<i>A.K. Kurilkin</i>	235
Exclusive electroproduction of ρ^0 , ϕ and ω mesons at HERMES	
<i>S.I. Manayenkov</i>	239
Semi-inclusive DIS and transverse momentum dependent distribution studies at CLAS	
<i>M. Mirazita</i>	244
The completion of single-spin asymmetry measurements at the PROZA setup	
<i>V.V. Mochalov</i>	250
The Generalized Parton Distribution experimental program at Jefferson Lab	
<i>C. Muñoz Camacho</i>	256
Lambda physics at HERMES	
<i>Yu. Naryshkin</i>	262
Spin physics at NICA	
<i>A. Nagaytsev</i>	266
Measurements of transverse spin effects in the forward region with the STAR detector	
<i>L.V. Nogach</i>	270
The first stage of polarization program SPASCHARM at the accelerator U-70 of IHEP	
<i>S.B. Nurushev</i>	274
Polarization measurements in photoproduction with CEBAF large acceptance spectrometer	
<i>E. Pasyuk</i>	282
Investigation of dp -elastic scattering and dp -breakup at ITS at Nuclotron	
<i>S.M. Piyadin</i>	288
Spin physics with CLAS	
<i>Y. Prok</i>	292

Measurements of G_{Ep}/G_{Mp} to high Q^2 and search for 2γ contribution in elastic ep at Jefferson Lab <i>V. Punjabi and Ch. Perdrisat</i>	299
Longitudinal spin transfer to Λ and $\bar{\Lambda}$ in polarized proton-proton collisions at $\sqrt{s} = 200$ GeV <i>Qinghua Xu</i>	313
Longitudinal spin transfer of the Λ and $\bar{\Lambda}$ hyperons in DIS at COMPASS <i>V.L. Rapatskiy</i>	319
The GPD program at COMPASS <i>A. Sandacz</i>	323
Azimuthal asymmetries in production of charged hadrons by high energy muons on polarized deuterium targets <i>I.A. Savin</i>	331
Transverse spin and momentum effects in the COMPASS experiment <i>G. Sbrizzai</i>	338
Delta-Sigma experiment – The results obtained. $\Delta\sigma_T(np)$ measurements planned at the Nuclotron-M <i>V.I. Sharov</i>	347
COMPASS results on gluon polarisation from high p_T hadron pairs <i>L. Silva</i>	351
Status of the PAX experiment <i>E. Steffens</i>	359
Asymmetry measurements in the elastic pion-proton scattering in the resonance energy range <i>D.N. Svirida</i>	365
Charge asymmetry and symmetry properties <i>E. Tomasi-Gustafsson</i>	369
First measurement of the interference fragmentation function in $e^+ e^-$ at BELLE <i>A. Vossen</i>	377

Technics and new developments 387

Proton beam polarization measurements at RHIC <i>A. Bazilevsky</i>	389
Development of high energy deuteron polarimeter based on dp -elastic scattering at the extracted beam of Nuclotron-M <i>Yu.V. Gurchin</i>	397
Spin-isotopic analysis at superlow temperatures <i>Yu.F. Kiselev</i>	401
Transparent spin resonance crossing in accelerators <i>A.M. Kondratenko</i>	405
Deuteron beam polarimetry at the internal target station at Nuclotron-M at GeV energies <i>P.K. Kurilkin</i>	411
Spin-manipulating polarized deuterons and protons <i>M.A. Leonova</i>	415

A study of polarized metastable helium-3 atomic beam production	
<i>Yu.A. Plis</i>	419
Proton polarimeter at 200 MeV energy	
<i>M.F. Runtso</i>	423
Spin-orbit potentials in neutron-rich helium isotopes	
<i>S. Sakaguchi</i>	427
Spin control by RF fields at accelerators and storage rings	
<i>Yu.M. Shatunov</i>	431
Related problems	437
Regularization of source of the Kerr-Newman electron by Higgs field	
<i>A. Burinskii</i>	439
About spin particle solution in Born-Infeld nonlinear electrodynamics	
<i>A.A. Chernitskii</i>	443
Spindynamics	
<i>I.B. Pestov</i>	447
Remark on spin precession formulae	
<i>L. Ryder</i>	451
Dynamics of spin in nonstatic spacetimes	
<i>A.J. Silenko</i>	455
DSPIN-09 workshop summary	
<i>J. Soffer</i>	459
List of participants of DSPIN-09	466

WELCOME ADDRESS

Alexei Sissakian
Director, JINR

Dear Colleagues,

To the great sorrow of all us, the co-chair of this workshop, Jan Nassalski, suddenly passed away this August. I would like to ask to honor his memory by a minute of silence... Jan will be remembered by many of his friends and colleagues in Dubna, The special talk about him will be delivered by Andrzej Sandacz in a few minutes. Jan played a very important role as a Chairman of the PAC in Particle Physics and last years he paid a lot of attention to the expertise of the project NICA, the main domestic particle experiment, which will also be discussed at this conference. I will say a few introductory words about it.

JINR has a long lasting traditions of experimental and theoretical studies of spin phenomena. Back in 1981, the first workshop in high energy spin physics was organized in Dubna by Lev Iosifovich Lapidus. These meetings became regular due to the initiative of Anatoly Vasilevich Efremov, their chairman for many years. Last meetings were co-chaired by him together with Jan Nassalski.

The current workshop is taking place in the important period of development of particle physics at JINR related to NICA accelerator complex based on the existing Nuclotron facility.

The main physical goal of NICA is a study of dense QCD matter. The investigation of its important properties, like critical point and mixed phase, do not require very large energies and are suitable just for NICA energy range. The NICA program is complementary to the program of FAIR@GSI.

Another important direction of experiments at NICA is represented by Spin Physics which will be achieved by exploration of polarized deuteron and proton beams and construction of the second detector purposed on Spin Physics. The main goal is the study of Single Spin Asymmetries in Drell-Yan processes and investigation of new distribution functions. At the same time, we plan also study spin asymmetries in the production of heavy quarkonia, spin-related signals in heavy ion collisions and, in particular asymmetries in elastic scattering (Krisch effect)

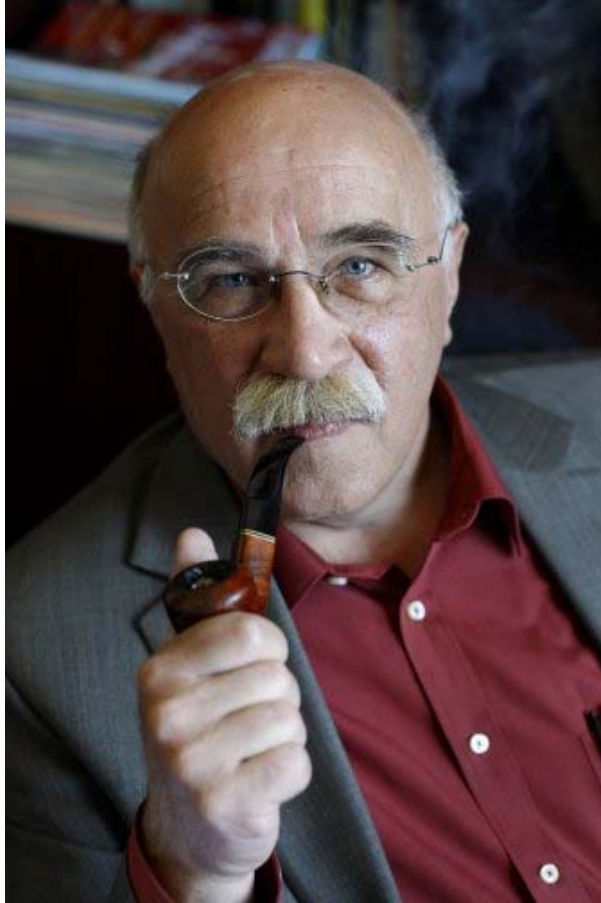
It is a pleasure for me to congratulate Alan Krisch, who is with us today, with his 70th Birthday, and to wish him a good health and many years of successful scientific activity. Alan is known as a brilliant experimentalist, key figure in the development of high energy spin physics, and also as a person who always supported Russian science and scientists, in particular, in their most difficult years. We hope that his fruitful collaboration with our good friends and colleagues in Protvino, and now also with NICA physicists, will last for many years.

Welcome to Dubna and JINR! The Workshop is opened!

JAN PAWEŁ NASSALSKI
1944 - 2009

A. Sandacz

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With deep sadness we learned of unexpected passing away of professor Jan Nassalski on August 5, 2009. We lost the prominent physicist, active organizer of research activities, dedicated teacher and promoter of science, and co-organizer of a number of Dubna-Spin conferences.

For more than 25 years Jan was the leader of Warsaw groups involved in CERN experiments with muon beams. More recently he also headed the Sołtan Institute for Nuclear Studies as Scientific Director. He was a Polish delegate to the CERN Council, the chairman of the JINR Council Advisory Committee on the High Energy Physics program and a member of the Scientific Council at DESY. In Poland he was also the coordinator of activities of High Energy Commission at National Atomic Energy Agency.

Jan Nassalski graduated in 1966 from Department of Mathematics and Physics of Warsaw University. He started his career at the Warsaw Technical University and in 1971

he joined the group of prof. Przemysław Iwo Zieliński at the Institute of Nuclear Research (now the Sołtan Institute for Nuclear Studies). During early years, in the 60's and 70's, he collaborated with physicists from University College London, Dubna, Rutherford Laboratory and Fermilab.

Since 1981 Jan participated in a series of experiments at CERN with muon beams (EMC, NMC, SMC and COMPASS) using deep inelastic scattering to study the internal structure of the nucleons. In the early 80's he set-up a group of Polish physicists involved in the muon experiments. Jan participated in the ground-breaking discovery of the EMC that the quark spins contribute little to the nucleon spin (coined as 'the nucleon spin crisis'). He was a key contributor to the precise studies of the structure functions by the NMC, resulting among others in observation of symmetry breaking in unpolarised parton distributions for non-strange sea quarks, seen through the violation of the Gottfried sum-rule. Since the 90's he focused on high precision experiments of the spin structure of the nucleon. He was co-author of the first experimental test (SMC) of the Bjorken sum rule and of the results on the gluon polarisation in the nucleon (COMPASS) measured independently in the production of open charm and in the high- p_t hadron pair production.

Jan Nassalski's group from the Sołtan Institute for Nuclear Studies also successfully contributed to the NA48 experiment at CERN dedicated to studies of the CP violation in decays of K^0 mesons.

Jan's scientific output consists of about 200 publications in international scientific journals and of numerous conference presentations. In 2005 professor Nassalski was awarded in Poland with Golden Order of Merit.

He was a devoted guide of younger colleagues in their scientific growth. In Poland he was tireless in outreach activities popularizing science and was publishing widely in Polish media. He gave interviews to journalists and helped them to reach interesting places and information. He participated in conferences for school students and teachers and supported the organization of such events. The educational exhibition on the LHC, circulating in Poland, was created with his essential contribution. He played a vital role in making the CERN's high school teacher program a great success in Poland. He was particularly proud of this work, and justifiably so.

Those people who met him could agree that he was a man of great integrity, rigorous in thinking and clear communicator. Although softly spoken, he knew how to carry an argument. Polite and with a natural kindness and sense of humour he was a true gentleman. Physics was his great passion, but not the only one. He was fond of various outdoor activities in the country, at the forests or in the mountains.

We will remember Him as as an excellent physicist and a good friend.

HARD COLLISIONS OF SPINNING PROTONS: HISTORY & FUTURE

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Abstract

There will be a review of the history of polarized proton beams, and a discussion of the unexpected and still unexplained large transverse spin effects found in several high energy proton-proton spin experiments at the ZGS, AGS, Fermilab and RHIC. Next, there will be a discussion of possible future experiments on the violent elastic collisions of polarized protons at IHEP-Protvino's 70 GeV U-70 accelerator in Russia and the new high intensity 50 GeV J-PARC at Tokai in Japan.

I will first discuss the violent elastic collisions of unpolarized protons. Fig. 1 shows the cross section for proton-proton elastic scattering plotted against a scaled P_t^2 variable that was proposed in 1963 [1] and 1967 [2] following Serber's [3] optical model. This plot is from updates by Peter Hansen and me [4, 5]. Notice that at small P_t^2 the cross-section drops off with a slope of about 10 (GeV/c)^{-2} . Fourier transforming this slope gives the size and shape of the proton-proton interaction in the diffraction peak; it is a Gaussian with a radius of about 1 Fermi. At medium P_t^2 there is a component with a slope of about 3 (GeV/c)^{-2} ; however, this component disappears rapidly with increasing energy. At lab energies of a few TeV it has totally disappeared; then, there is a sharp destructive interference between the small- P_t^2 diffraction peak and the large- P_t^2 hard-scattering component. Since the diffraction peak is mostly *diffractive*, its amplitude must be mostly imaginary, as has been experimentally verified. Hence, the sharp destructive interference implies that the large- P_t^2 component is also mostly imaginary; thus, it is probably mostly *diffractive*. This large- P_t^2 component is probably the elastic *diffractive* scattering due to the *direct* interactions of the proton's constituents; its slope of about 1.5 (GeV/c)^{-2} implies that these *direct* interactions occur within a Gaussian-shaped region of radius about 0.3 Fermi.

Since the medium- P_t^2 component disappears at high energy, it is probably the *direct elastic scattering* of the two protons. This view is supported by the experimental fact that proton-proton elastic scattering is the only exclusive process that still can be precisely measured at TeV energies. To understand this, note that *direct* elastic scattering and all other exclusive processes must compete with each other for the total p-p cross-section, which is less than 100 millibarns. At TeV energies, there are certainly more than 10^5 exclusive channels in this competition; thus, each channel has an average cross section of less than 1 microbarn. Moreover, since the medium- P_t^2 elastic component does not interfere strongly with either the large- P_t^2 or small- P_t^2 components, its amplitude is probably real. Also note that the large- P_t^2 component intersects the cross section axis at about 10^{-5} below the small- P_t^2 *diffractive* component.

An earlier version of Fig. 1 got me started in the spin business. In 1966, we measured p-p elastic scattering at the ZGS at exactly 90°_{cm} from 5 to 12 GeV [6]; the sharp slope-change, shown by the stars, was apparently the first direct evidence for constituents in the

proton. Dividing these 90°_{cm} p-p elastic cross sections by 4 (due to the protons' particle identity) made all then existing proton-proton elastic data, above a few GeV, fit on a single curve [2]. During a 1968 visit to Ann Arbor, Prof. Serber informed me that, by dividing the 90°_{cm} points by 4, I had made an assumption about the ratio of the spin singlet and triplet p-p elastic scattering amplitudes. I was astounded and said that I knew nothing about spin and certainly had not measured the spin of either proton. He said with a smile that both statements might be true; nevertheless, my nice fit required this assumption. Prof. Serber, as usual, spoke quietly; however, as a student, I had learned that he was almost always right. Thus, I looked for data on proton-proton elastic scattering data, in the singlet or triplet spin states, above a few GeV. I found that none existed and decided

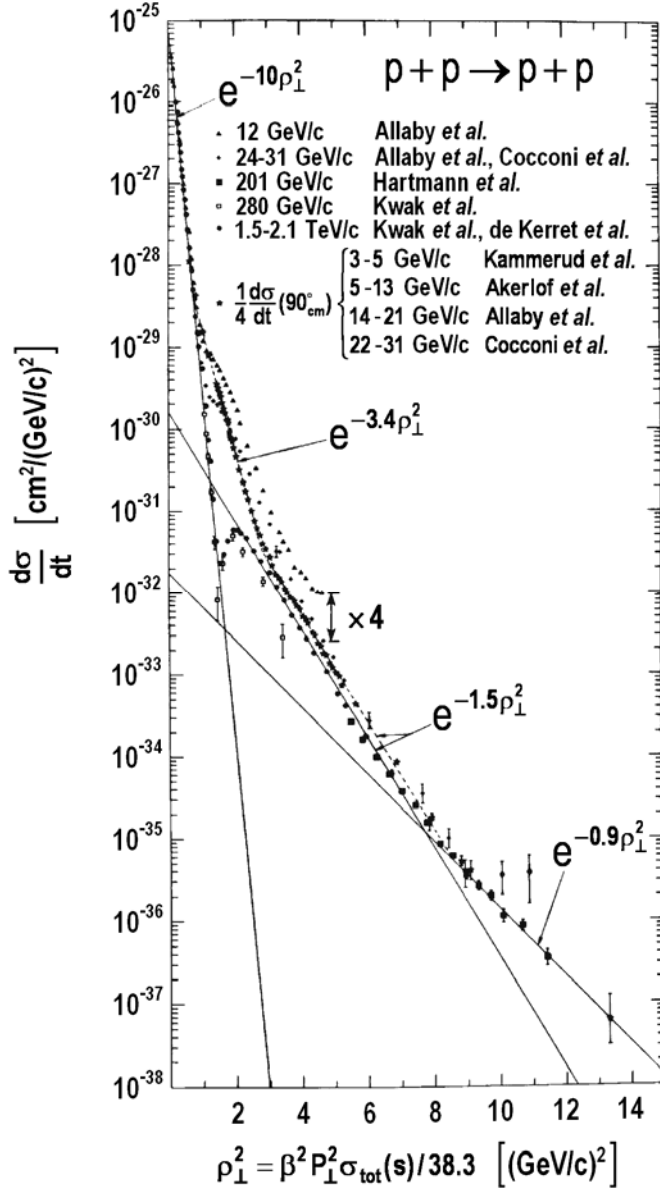


Figure 1: Proton-proton elastic cross-sections plotted vs. the scaled P_t^2 variable [4, 5]. The 12 GeV/c Allaby *et al.* data were not corrected for 90°_{cm} particle identity effects.

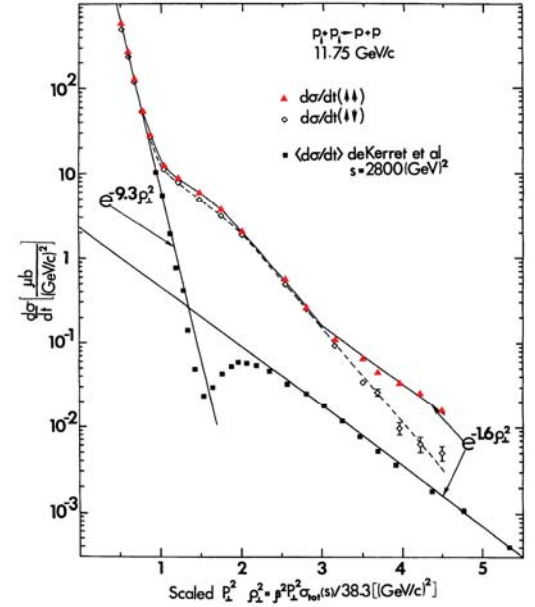


Figure 2: The proton-proton elastic cross section near 12 GeV in pure initial spin states plotted vs. the scaled P_t^2 -variable [11].

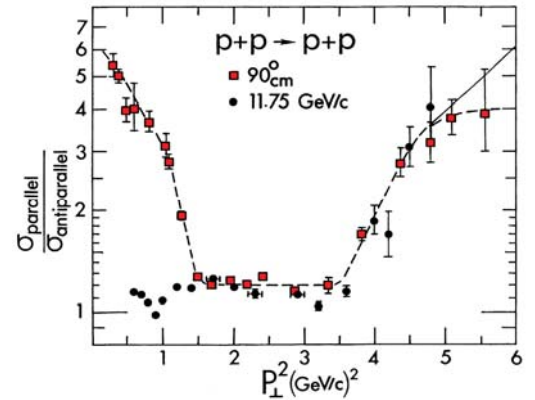


Figure 3: The measured spins-parallel/spins-antiparallel elastic cross-section ratio ($\sigma_{\uparrow\uparrow}/\sigma_{\uparrow\downarrow}$) plotted vs. P_t^2 [12].

to try to polarize the protons in the ZGS.

At the 1969 New York APS Meeting, I learned that EG&G was the representative for a new polarized proton ion source made by ANAC in New Zealand. I discussed this with my long-time colleague, Larry Ratner, and then with Bruce Cork, Argonne's Associate Director, and Robert Duffield, Argonne's Director. They apparently decided it was a good idea; Duffield soon hired me as a consultant to Argonne at \$100 per month. In 1973, after a lot of hard work by many people, the ZGS accelerated the world's first high energy polarized proton beam [7].

The ZGS needed some hardware to overcome both intrinsic and imperfection depolarizing resonances. Fortunately, both types of resonances were fairly weak at the 12 GeV ZGS, which was the highest energy weak focusing accelerator ever built. All higher energy accelerators wisely use strong focusing [8], which makes the depolarizing resonances much stronger. If we had first tried to accelerate polarized protons at a strong focusing accelerator, such as the AGS, we probably would have failed and abandoned the polarized proton beam business. Fortunately, it worked at the weak focusing ZGS [7, 9], and experiments [10] soon showed that the p-p total cross-section had significant spin dependence; this surprised many people, including me.

Figure 2 shows our perhaps most important result [11] from the ZGS polarized proton beam. The 12 GeV proton-proton elastic cross section in pure initial spin states is plotted against the scaled P_t^2 -variable; in the diffraction peak the spin-parallel and spin-antiparallel cross-sections are essentially equal to each other and to the unpolarized data from the CERN ISR at $s = 2800 \text{ GeV}^2$; thus, in small-angle *diffractive* scattering, the protons in different spin states (and at different energies) all have about the same cross-section. The medium- P_t^2 component, which still exists near 12 GeV, has only a small spin dependence; again note that it has totally disappeared at 2800 GeV². However, the behavior of the large- P_t^2 hard-scattering component was a great surprise. When the protons' spins are parallel, they seem to have exactly the same behavior as the much higher energy unpolarized ISR data; however, when their spins are antiparallel their cross-section drops with the medium- P_t^2 component's steeper slope. When this data first appeared in 1977 and 1978, people were totally astounded; most had thought that spin effects would disappear at high energies. In the years following, many theoretical papers tried to explain this unexpected behavior; none were fully successful. In particular, the theory that is now called QCD, has been unable to deal with this data; Glashow once called this experiment "...the thorn in the side of QCD". In his summary talk at Blois 2005, Stan Brodsky called this result "...one of the unsolved mysteries of Hadronic Physics".

I learned something important from questions during two 1978 seminars about this result. Two distinguished physicists, Prof. Weisskopf at CERN and then Prof. Bethe at Copenhagen a week later, asked the same question apparently independently. Each said that our big spin effect at large- P_t^2 was quite interesting; but at 12 GeV, the spins-parallel/spins-antiparallel ratio was only large near 90°_{cm} , where particle identity was important for p-p scattering. They asked: How could one be sure that our large spin effect was due to hard-scattering at large- P_t^2 , rather than particle identity near 90°_{cm} ? One would be foolish to ignore the comments of two such distinguished theorists; moreover, they were related to Prof. Serber's comment 10 years earlier.

It seemed that their question could not be answered theoretically. Thus, we tried to answer it experimentally with a second ZGS experiment, which varied P_t^2 by holding the

p-p scattering angle fixed at exactly 90°_{cm} , while varying the energy of the proton beam. This 90°_{cm} p-p elastic fixed-angle data [12] is plotted against P_t^2 in Fig. 3, along with the fixed-energy data [11] of Fig. 2. There are large differences at small P_t^2 , where the 90°_{cm} data are at very low energy; however, above P_t^2 of about 1.5 (GeV/c)², the two sets of data fall right on top of each other. The point at $P_t^2 = 2.5$ (GeV/c)², where the ratio is near 1, is just as much at 90°_{cm} , as the 5 (GeV/c)² point, where the ratio is 4. This data apparently convinced Profs. Bethe and Weisskopf that the large spin effect was not due to 90°_{cm} particle identity and it was a large- P_t^2 hard-scattering effect.

Figure 4 shows A_{nn} for p-p elastic scattering plotted against the lab momentum, P_{Lab} [5]; it includes the ZGS data from Fig. 3 plus some lower energy data obtained from Willy Haeberli, who is an expert on low energy p-p spin experiments. At the lowest momentum (near $T = 10$ MeV) A_{nn} is very close to -1 ; thus, two protons with parallel spins can never scatter at 90°_{cm} . Next A_{nn} goes rapidly to $+1$; then protons with antiparallel spins can never scatter at 90°_{cm} . Then at medium energy, there are some oscillations that were once thought to be due to dibaryon resonances, but may be due to the onset of N^* resonance production. In the ZGS region, A_{nn} first drops rapidly; it is next small and constant over a large range; it then rises rapidly to 0.6. These huge and sharp oscillations of A_{nn} seem impressive.

I now turn to funding. In 1972 the AEC had agreed to shut down the ZGS in 1975 to get funding for PEP at SLAC. When the unique ZGS polarized beam started operating in 1973, the wisdom of this decision was questioned; AEC then set up a committee which extended ZGS operations through 1977. A second committee in 1976 extended operations of the ZGS polarized beam until 1979 [13]. Henry Bohm, the President of AUA, which operated Argonne, asked ERDA, which had replaced AEC, to set up a third committee to again extend ZGS running. However, OMB objected, so there was no third committee; nevertheless, his efforts had some benefit. When James Kane, of ERDA, responded negatively to Dr. Bohm, his justification was that it might now be possible to accelerate polarized protons in a strong focusing accelerator such as the AGS; moreover, Dr. Kane officially copied me on his letter.

We had started interacting with Ernest Courant and others at Brookhaven about polarizing the AGS, first at a 1977 Workshop in Ann Arbor [14]. Then at a 1978 Polarized AGS Workshop at Brookhaven [15], Brookhaven's Associate Director, Ronald Rau, asked me for a copy of Dr. Kane's letter. He used it to convince William Wallenmeyer, the long-time Director of High Energy Physics at AEC, ERDA and DoE, to provide about \$8 Million to Brookhaven, and about \$2 Million split between Michigan, Argonne, Rice and Yale, for the challenging project of accelerating polarized protons in the strong-focusing AGS, and later in the 400 GeV ISABELLE collider. ISABELLE was canceled in

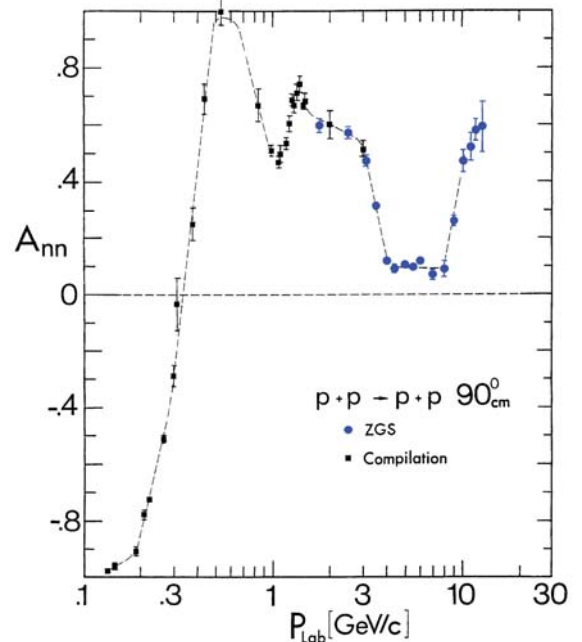


Figure 4: $A_{nn} \equiv (\sigma_{\uparrow\uparrow} - \sigma_{\uparrow\downarrow}) / (\sigma_{\uparrow\uparrow} + \sigma_{\uparrow\downarrow})$ is plotted against P_{Lab} . [5]

1983, but was later reborn as RHIC, which is now colliding 250 GeV polarized protons.

It was far more difficult to accelerate polarized protons in the strong focusing AGS than in the weak focusing ZGS. The strong focusing principal, invented by Courant, Livingston and Snyder [8], made possible all modern large circular accelerators by using alternating quadrupole magnetic fields to strongly focus the beam and thus keep it small. Unfortunately, these strong quadrupole fields were very good at depolarizing protons. To accelerate polarized protons to 22 GeV at the AGS, one had to overcome 45 strong depolarizing resonances. This required: some very challenging hardware; significantly upgrading the AGS controls; and spending lots of time individually overcoming the 45 depolarizing resonances. Michigan built 12 ferrite quadrupole magnets to allow the AGS to overcome its 6 intrinsic resonances by rapidly jumping its vertical betatron tune through each resonance. Brookhaven was building their 12 power supplies; but each power supply had to provide 1500 Amps at 15,000 Volts (about 22 MW) during each quadrupole's 1.6 μ sec rise-time. Overcoming the many imperfection depolarizing resonances (occurring every 520 MeV) required programming the AGS's 96 small correction dipole magnets to form a horizontal B-field wave of 4 oscillations at the instant when the proton energy passed through the $G\gamma = 4$ imperfection resonance. Then, about 20 msec later in the AGS cycle, when $G\gamma$ was 5, the 96 magnets had a horizontal B-field wave with 5 oscillations, etc. ($G = 1.79285$ is the proton's anomalous magnetic moment, while $\gamma = E/m$).

After all this hardware was installed, an even larger problem was tuning the AGS. In 1988, when we accelerated polarized protons to 22 GeV, we needed 7 weeks of exclusive use of the AGS; this was difficult and expensive. Once a week, Nicholas Samios, Brookhaven's Director, would visit the AGS Control Room to politely ask how long the tuning would continue and to note that it was costing \$1 Million a week. Moreover, it was soon clear that, except for Larry Ratner (then at Brookhaven) and me, no one could tune through these 45 resonances; thus, for some weeks, Larry and I worked 12-hour shifts 7-days each week. After 5 weeks Larry collapsed. While I was younger than Larry, I thought it unwise to try to work 24-hour shifts every day. Thus, I asked our Postdoc, Thomas Roser, who until then had worked mostly on polarized targets and scattering experiments, if he wanted to learn accelerator physics in a hands-on way for 12 hours every day. Apparently, he learned well, and now leads Brookhaven's Collider-Accelerator Division.

One benefit from this difficult 7-week period [16, 7] was learning that our method of individually overcoming each resonance, which had worked so well at the ZGS [9, 7], might work at the AGS, but would not be practical at higher energy accelerators. This lesson helped to launch our Siberian snake programs at IUCF [17, 7] and then SSC [18, 19].

In the 1980's, a new proton collider, the SSC, was being planned; it was to have two 20 TeV proton rings each about 80 km in circumference. Owen Chamberlain and Ernest Courant encouraged me to form a collaboration to insure that polarized protons would be possible in the SSC. We first organized a 1985 Workshop in Ann Arbor, with Kent Terwilliger. This Workshop [18] concluded that it should be possible to accelerate and maintain the polarization of 20 TeV protons in the SSC, but only if the new Siberian snake concept of Derbenev and Kondratenko [20] really worked; otherwise, it would be totally impractical. Recall that it took 49 days to correct the 45 depolarizing resonances at the AGS, about one per day. Each 20 TeV SSC ring would have about 36,000 depolarizing resonances to correct. Moreover, these higher energy resonances would be much stronger and harder to correct; but even at one per day, this would require about 100 years of

tuning for each ring. The Workshop also concluded that one must prove experimentally that the *too-good-to-be-true* Siberian snakes really worked; otherwise, there would be no approval to install the 26 Siberian snakes needed in each SSC ring.

Indiana's IUCF was then building a new 500 MeV synchrotron Cooler Ring [7]. Some of us Workshop participants then collaborated with Robert Pollock and others at IUCF to build and test the world's first Siberian snake in the Cooler Ring. We brought experience with synchrotrons and high energy polarized beams, while the IUCF people brought experience with low energy polarized beams and the CE-01 detector, which was our polarimeter. In 1989, we demonstrated that a Siberian snake could easily overcome a strong imperfection depolarizing resonance [17, 7]. For 13 years we continued these experiments and learned many things about spin-manipulating polarized beams. After the IUCF Cooler Ring shut down in 2002, this polarized proton and deuteron beam experiment program was continued at the 3 GeV COSY in Juelich, Germany from 2002-2009 [7].

In 1990 we formed the SPIN Collaboration and submitted to the SSC: Expression of Interest EOI-001 [19]. It proposed to accelerate and store polarized protons at 20 TeV; and to study spin effects in 20 TeV p-p collisions. It was submitted a week before the deadline, which made it SSC EOI-001. Thus, we made the first presentation to the SSC's PAC before a huge audience that included many newspaper reporters and TV cameras. Perhaps partly due to this publicity, we were soon *partly* approved by SSC Director Roy Schwitters. By *partly* I mean that he decided to add 26 empty spaces for Siberian snakes in each SSC Ring; each space was about 20 m long, which added about 0.5 km to each Ring. Unfortunately, the SSC was canceled around 1993, before it was finished, but after \$2.5 Billion was spent. Nevertheless, our detailed studies of the behavior and spin-manipulation of polarized protons at IUCF and COSY helped in developing polarized beams around the world: Brookhaven now has 250 GeV polarized protons in each RHIC ring [21, 7]; perhaps someday CERN's 7 TeV LHC might have polarized protons.

We eventually accelerated polarized protons to 22 GeV in the AGS [16, 7] and obtained some A_{nn} data [22, 23, 7] well above the ZGS energy of 12 GeV; but we never had enough polarized-beam *data-time* to get precise A_{nn} data at high- P_t^2 . But, during tune-up runs for the A_{nn} experiment, we used the unpolarized AGS proton beam to test our polarized proton target and double-arm magnetic spectrometer by measuring A_n in 28 GeV proton-proton elastic scattering; this data resulted in an interesting surprise. Despite QCD's inability to explain the big A_{nn} from the ZGS, our QCD friends had made a firm prediction that the one-spin A_n must go to 0; they also said this prediction would become more firm at higher energies and in more violent collisions. But above $P_t^2 = 3 \text{ (GeV/c)}^2$, A_n instead began to deviate from 0 and was quite large at $P_t^2 = 6 \text{ (GeV/c)}^2$. This led to more controversy [23]; some QCD supporters said that our A_n data must be wrong.

Experimenters take such accusations seriously. Thus, we started preparing an experiment that could study A_n at high- P_t^2 with better precision. Our spectrometer worked well, but we could only use about 0.1% of the AGS beam intensity, because a higher intensity beam would heat our Polarized Proton Target (PPT) and depolarize it. Thus, we started building a new PPT [24, 7] that could operate with 20 times more beam intensity; this required ^4He evaporation cooling at 1 K, which has much more cooling power than our earlier ^3He evaporation PPT at 0.5 K. However, to maintain a target polarization near 50% at 1 K required increasing the B-field from 2.5 to 5 Tesla. Thus, we ordered a 5 T superconducting magnet from Oxford Instruments, with a B-field uniformity of a

few 10^{-5} over the PPT's 3 cm diameter volume. We also obtained a Varian 20 W at 140 GHz Extended Interaction Oscillator; it was apparently the highest power 140 GHz microwave source available. As the PPT assembly started in 1989, we worried that, if the PPT model was wrong, the polarization might be only 10%; but we were very lucky; it was 96% [24, 7].

Moreover, the target polarization averaged 85% for a 3-month-long run with high-intensity AGS beam in early 1990. As shown in Fig. 5, this let us precisely measure A_n at even larger P_t^2 . When these precise new data were published [25], some theorists seemed quite unhappy; they still believed the QCD prediction that A_n must go to 0, but they now refused to state at what P_t^2 or energy this prediction would become valid. They also now said that QCD might not work for elastic scattering, which they now considered less fundamental than inelastic scattering, where they said QCD should work. Thus, one result of our experiments was to make both elastic scattering experiments and spin experiments unpopular in some circles.

Experimenters had also started doing inclusive polarization experiments at Fermilab. Figure 6 shows the 400 GeV inclusive hyperon polarization experiments from the 1970s and 1980s, led by Pondrum, Devlin, Heller and Bunce [26]; it clearly shows a small polarization at small P_t and a larger polarization at larger P_t . Moreover, their data is consistent with 12 GeV data from the KEK PS and with 2000 GeV data from the CERN ISR. These data do not support the QCD prediction that inelastic spin effects disappear at high energy or high P_t^2 .

Another group at Fermilab, led by Yokosawa, developed a *secondary* polarized proton beam using the polarized protons from polarized hyperon decay. The beam's intensity

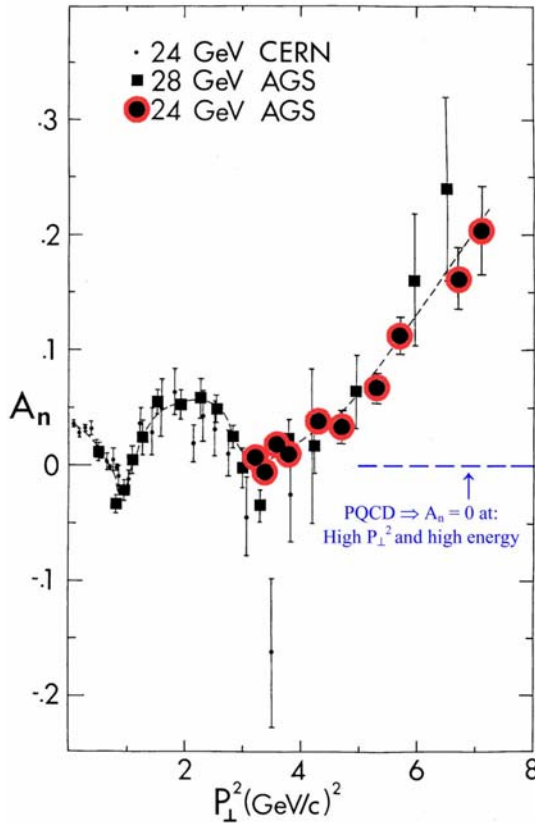
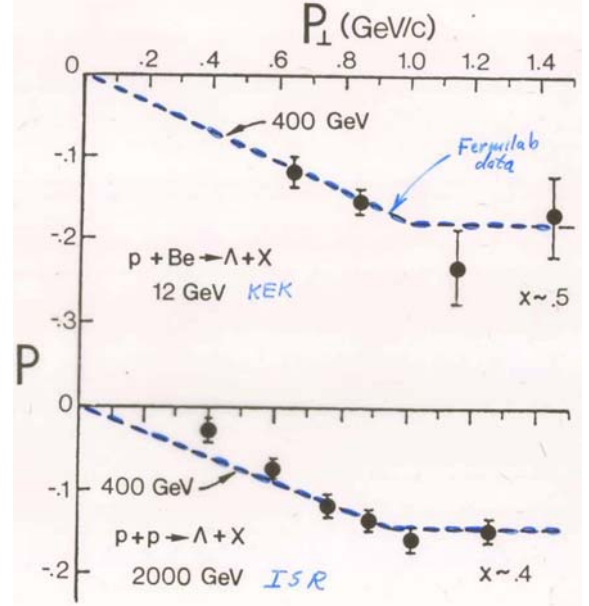


Figure 5: (Left) $A_n \equiv (\sigma_{\uparrow} - \sigma_{\downarrow})/(\sigma_{\uparrow} + \sigma_{\downarrow})$ is plotted vs. P_t^2 for p-p elastic scattering [25].

Figure 6: (Bottom) The inclusive Λ polarization is plotted against P_t [26].



was only about 10^5 per second, but its polarization was about 50% and its energy was about 200 GeV. They obtained some nice A_n data on inclusive π meson production [27], which are shown in Fig. 7. The A_n values for π^+ and π^- mesons are both large but with opposite signs, while A_n for the π^0 data is 50% smaller and is positive. These 200 GeV data provide little support for QCD.

We tried to measure spin effects in very high energy p-p scattering at UNK, which IHEP-Protvino started building around 1986 [7]; IHEP and Michigan signed the NEPTUN-A Agreement in 1989. Michigan's main contribution was a 12 Tesla at 0.16 K Ultra-cold Spin-polarized Jet [7]. UNK's circumference would be 21 km with 3 rings: a 400 GeV warm ring and two 3 TeV superconducting rings; its injector was IHEP's existing 70 GeV accelerator, U-70 [7]. By 1998 the UNK tunnel and about 80% of its 2200 warm magnets were finished, and 70 GeV protons were transferred into its tunnel with 99% efficiency. However, progress became slower each year due to financial problems; in 1998 Russia's MINATOM placed UNK on long-term standby [7].

IHEP Director, A.A. Logunov, had earlier suggested moving our experiment to IHEP's existing 70 GeV U-70 accelerator. By March 2002 the resulting SPIN@U-70 Experiment on 70 GeV p-p elastic scattering at high P_t^2 was fully installed [7], except for our detectors and Polarized Proton Target (PPT) [24, 7]. However, just before our 4 tons of detectors, electronics and computers were to be shipped, the US Government suspended the US-Russian Peaceful Use of Atomic Energy Agreement started by President Eisenhower in 1953. Nevertheless, DoE asked us to send the shipment, since under the terms of the PUoAE Agreement, the experiment should have been done exactly as planned; DoE faxed us a copy of the Agreement. Thus, we sent the shipment. It arrived at Moscow airport on March 11, 2002, where it was impounded for 8 months before it was returned to Michigan.

Despite this problem, we remain friends with our IHEP colleagues and there have been four SPIN@U-70 test runs using Russian detectors and an unpolarized target; we participated in the November 2001 and April 2002 runs. We hope that the US-Russian PUoAE Agreement will soon be restarted so that the SPIN@U-70 experiment can continue and measure A_n at P_t^2 near 12 (GeV/c) $^{-2}$. However, if this international problem continues, we may try to do a similar experiment at Japan's new high-intensity 50 GeV proton accelerator, J-PARC [7] that is now starting to run. If J-PARC could accelerate polarized protons to 50 GeV, then one could study the large and mysterious elastic spin effects in both A_n and A_{nn} for the first time in decades.

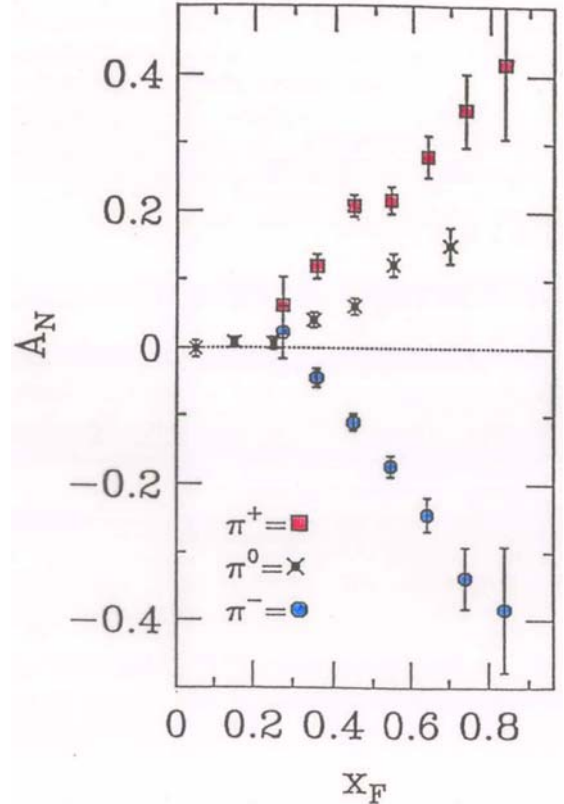


Figure 7: A_n for inclusive π -meson production is plotted vs. X_F [27].

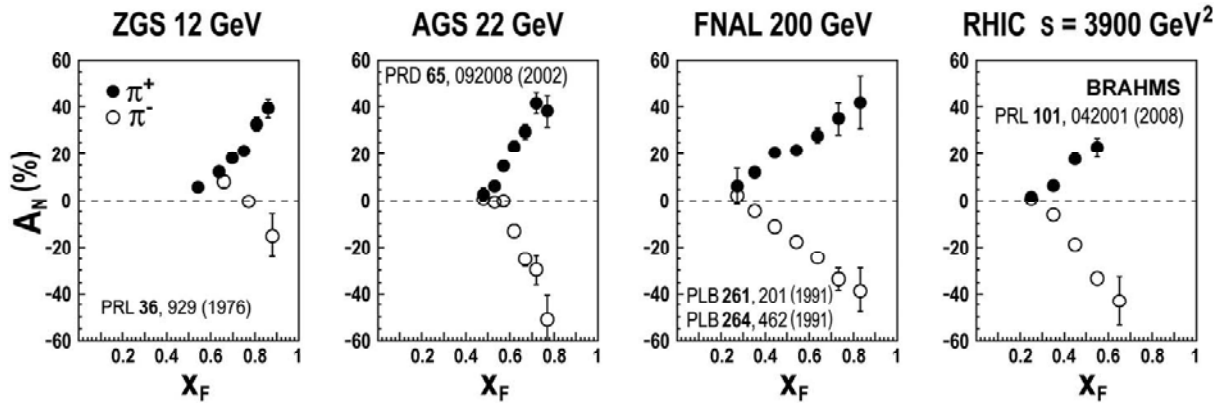


Figure 8: Inclusive pion asymmetry in proton-proton collisions [27].

To summarize, for the past 30 years QCD-based calculations have continued to disagree with the ZGS 2-spin and AGS 1-spin elastic data, and the ZGS, AGS, Fermilab and now RHIC [28] inclusive data. To be specific:

- * These large spin effects do not go to zero at high-energy or high- P_t , as was predicted.

- * No QCD-based model can yet explain simultaneously all these large spin effects.

There is a BASIC PRINCIPLE OF SCIENCE:

- * If a theory disagrees with reproducible experimental data, then it must be modified. Precise spin experiments could provide experimental guidance for the required modification of the theory of Strong Interactions. New experiments at higher energy and higher P_t on the proton-proton elastic cross-section's: $d\sigma/dt$, A_{nn} and A_n could provide further guidance for these modifications, just as the RHIC inclusive A_n experiment [28] confirmed the earlier Fermilab experiments [27]. Elastic scattering is especially important because:

- * It is the only exclusive process large enough to be measured at TeV energy. This is probably because proton-proton elastic scattering is dominated by the diffraction due to the millions of inelastic channels that compete for the total cross-section of only about 100 millibarns at TeV energies. Many people may have forgotten this simple but essential geometrical approach [1], which I learned from Prof. Serber's optical model in 1963 [3]; perhaps it should now be learned or relearned by others.

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THEORY OF SPIN PHYSICS

MICROSCOPIC STERN-GERLACH EFFECT AND THOMAS SPIN PRECESSION AS AN ORIGIN OF THE SSA

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Abstract

The single-spin asymmetry and hadron polarization data are analyzed in the framework of a phenomenological effective-color-field model. Global analysis of the single-spin effects in hadron production is performed for h+h, h+A, A+A and lepton+N interactions. The model explains the dependence of the data on x_F , p_T , collision energy $\sqrt{s}$ and atomic weights A_1 and A_2 of colliding nuclei. The predictions are given for not yet explored kinematical regions.

In this report we discuss a semi-classical mechanism for the single-spin phenomena in inclusive reaction $A + B \rightarrow C + X$. The major assumptions of the model are listed below.

1) An effective color field (ECF) is a superposition of the QCD string fields, created by spectator quarks and antiquarks after the initial color exchange.

2) The constituent quark Q of the detected hadron C interacts with the nonuniform chromomagnetic field via its chromomagnetic moment $\mu_Q^a = sg_Q^a g_S / 2M_Q$ and with the chromoelectric field via its color charge g_S .

3) The microscopic Stern-Gerlach effect in chromomagnetic field $\mathbf{B}^a$ and Thomas spin precession in chromoelectric field $\mathbf{E}^a$ lead to the large observed SSA. The ECF is considered as an external with respect to the quark Q of the observed hadron.

4) The quark spin precession in the ECF (chromomagnetic and chromoelectric) is an additional phenomenon, which leads to the specific SSA dependence (oscillation) as a function of kinematical variables (x_F , p_T or scaling variable $x_A = (x_R + x_F)/2$).

The longitudinal field $\mathbf{E}^a$ and the circular field $\mathbf{B}^a$ of the ECF are written as

$$E_Z^{(3)} = -2\alpha_s \nu_A / \rho^2 \exp(-r^2/\rho^2), \quad B_\varphi^{(2)} = -2\alpha_s \nu_A r / \rho^3 \exp(-r^2/\rho^2), \quad (1)$$

where r is the distance from the string axis, ν_A is the number of quarks at the end of the string, $\rho = 1.25R_c$, R_c - confinement radius, $\alpha_s \approx 1$ - running coupling constant [1, 2].

The Stern-Gerlach type forces act on a quark moving inside the ECF (flux tube):

$$f_x = \mu_x^a \partial B_x^a / \partial x + \mu_y^a \partial B_y^a / \partial x, \quad f_y = \mu_x^a \partial B_x^a / \partial y + \mu_y^a \partial B_y^a / \partial y. \quad (2)$$

The quark Q of the observed hadron C , which gets p_T kick of Stern-Gerlach forces and undergo a spin precession in the ECF is called a “probe” and it “measures” the fields $\mathbf{B}^a$ and $\mathbf{E}^a$. The ECF is created by spectator quarks and antiquarks and obeys quark counting rules. Spectators are all quarks which are not constituents of the observed hadron C . In the case of $p^\uparrow + p \rightarrow \pi^+ + X$ reaction (see Fig. 1) polarized probe u quark from π^+ feels the field, created by the spectator quarks with weight $\lambda = -|\Psi_{qq'}(0)|^2 / |\Psi_{q\bar{q}}(0)|^2 \approx -1/8$, by antiquarks with weight 1, and by target quarks with weight $-\tau\lambda$, respectively. Spectator

quarks from the target B have an additional negative factor $-\tau = -0.0562 \pm 0.003$, since these quarks are moving in opposite direction in cm reference frame. The value of color factor $\lambda = -0.1321 \pm 0.0012$, obtained in a global fit of 68 inclusive reactions, is close to the expected one, which is a strong argument in favor of the ECF model [2].

Another important phenomenon is quark spin precession in the ECF. We assume that the spin precession is described by the Bargman-Michel-Telegdi eqs. (3)-(4) [2]:

$$d\xi/dt = a[\xi \mathbf{B}^a] + d[\xi[\mathbf{E}^a \mathbf{v}]], \quad (3)$$

$$a = g_S(g_Q^a - 2 + 2M_Q/E_Q)/2M_Q, \quad d = g_S[g_Q^a - 2E_Q/(E_Q + M_Q)]/2M_Q. \quad (4)$$

The precession frequency depends on the color charge g_S , the quark mass M_Q and its energy E_Q , and on the color g_Q^a -factor. The value $\Delta\mu^a = (g_Q^a - 2)/2$ is called a color anomalous magnetic moment and it is large and negative in the instanton model. Spontaneous chiral symmetry breaking leads to an additional dynamical mass $\Delta M_Q(q)$ and $\Delta\mu^a(q)$, both depend on momentum transfer q [4]. Kochelev predicted $\Delta\mu^a(0) = -0.2$ [5] and Diakonov predicted $\Delta\mu^a(0) = -0.744$ [4]. The global data analysis results are closer to the Diakonov's predictions.

Due to the microscopic Stern-Gerlach effect the probe quark Q gets an additional spin-dependent transverse momentum δp_x , which causes an azimuthal asymmetry or observed hadron polarization:

$$\delta p_x = \frac{g_Q^a \xi_y^0}{2\rho(g^a - 2 + 2M_Q/E_Q)} \left[\frac{1 - \cos(\phi_A)}{\phi_A} + \epsilon \phi_A \right], \quad (5)$$

where $\phi_A = \omega_A x_A$ is a quark spin precession angle in the fragmentation region of the beam particle A , and a "frequency" of A_N oscillation as a function of x_A is

$$\omega_A = \frac{g_S \alpha_s \nu_A S_0 (g_Q^a - 2 + 2M_Q/E_Q)}{M_Q \rho^2 c}. \quad (6)$$

The length S_0 of the ECF is 0.6 ± 0.2 fm. The constituent quark masses M_Q and $\Delta\mu^a$ values are given in [2]. The parameter $\epsilon = -0.00419 \pm 0.00022$ is small due to subtraction of the Thomas precession term from $\epsilon = 1/2$ for chromomagnetic contribution to the δp_x .

Let us consider in more detail the Thomas precession effect in the ECF. Due to the Thomas precession an additional term $U = \mathbf{s} \cdot \boldsymbol{\omega}_T$ appears in the effective Hamiltonian. The Thomas frequency $\omega_T \approx [\mathbf{F}\mathbf{v}]/M_Q$ depends on the force $\mathbf{F}$, the quark velocity $\mathbf{v}$ and its mass M_Q . The quark polarization due to the Thomas precession, $\delta \mathbf{P}_N = -\omega_T/\Delta E$, is directed opposite to the frequency vector ω_T direction since ΔE is positive [6].

The sign and magnitude of the force $\mathbf{F} = g_S \mathbf{E}^a$ is determined by the quark-counting rules for the ECF. For example, $F_Z \approx -2g_S \alpha_S [1 + \lambda - 3\tau\lambda]/\rho^2 < 0$ for the $pp \rightarrow \Lambda^\dagger + X$ reaction at $\sqrt{s} < 70$ GeV. In case of the reactions $pp \rightarrow \Xi^{0\dagger} + X$ and $p^\dagger + p \rightarrow \pi^+ + X$ the factor in square brackets is $[2 + 2\lambda - 3\tau\lambda] < 0$ and $[3\lambda - 3\tau\lambda] > 0$, respectively.

Recombination potential is more attractive for negative $U = \mathbf{s} \cdot \boldsymbol{\omega}_T$. As a result, for the $pp \rightarrow \Lambda^\dagger + X$ reaction the additional Thomas precession contribution to P_N is positive and opposite in sign to the dominating chromomagnetic term, which gives $P_N < 0$, and to the

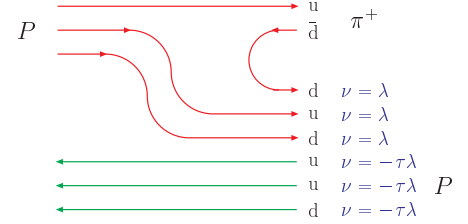


Figure 1: The quark flow diagram for the reaction $p^\dagger + p \rightarrow \pi^+ + X$. The weight ν of each spectator quark contribution to the ECF is indicated.

DeGrand-Miettinen model predicted $P_N < 0$ [6]. The additional transverse momentum δp_x is related to the analyzing power or polarization by the relation $A_N = -D\delta p_x$, where $D = 5.68 \pm 0.13 \text{ GeV}^{-1}$ is an effective slope of the invariant cross section. In the Ryskin model $\delta p_x \approx 0.1 \text{ GeV}/c$ is a constant [7]. In the ECF model we have dynamical origin of A_N or P_N dependence on the kinematical variables (p_T , x_A , $x_B = (x_R - x_F)/2$, x_F) and on the number of (anti)quarks in hadrons A , B and C , and also on quark g_Q^a -factor and its mass M_Q . This dependence is due to the microscopic Stern-Gerlach effect and quark spin precession in the ECF [2].

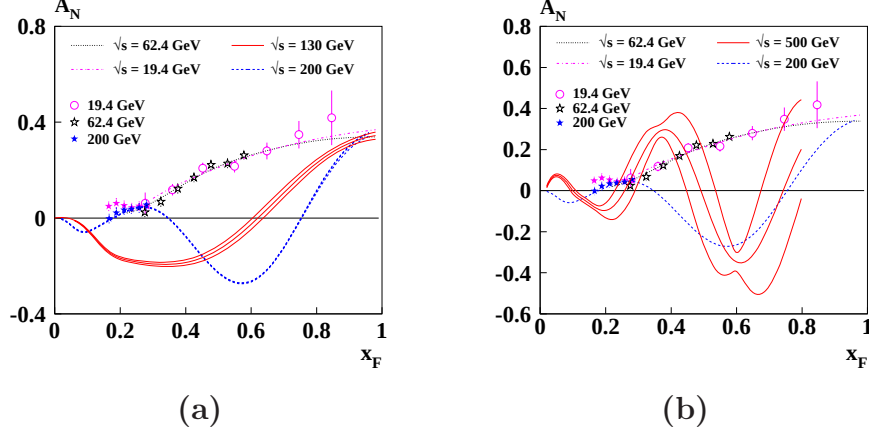


Figure 2: The dependence $A_N(x_F)$ for $p^\dagger + p \rightarrow \pi^+ + X$ reaction. Predictions are for (a) $\sqrt{s} = 130 \text{ GeV}$ and (b) $\sqrt{s} = 500 \text{ GeV}$, respectively. The cm production angle is 4.1° .

The ECF increases dramatically at energy $\sqrt{s} > 70 \text{ GeV}$ or in collisions of nuclei. In the case of nuclei collisions the effective number of quarks in a projectile nuclei, which contributes to the ECF, is equal to its number in a tube with transverse radius limited by the confinement. The new quark contribution to the ECF depends on kinematical variables:

$$f_N = n_q \exp(-W/\sqrt{s})(1 - x_N)^n, \quad x_N = [(p_T/p_N)^2 + x_F^2]^{1/2}, \quad (7)$$

where $W \approx 238(A_1 A_2)^{-1/6} \text{ GeV}$, $n \approx 0.91(A_1 A_2)^{1/6}$, $p_N \approx 28 \text{ GeV}/c$ and $n_q \approx 4.52$ [2].

Due to the dependence of the ECF on $\sqrt{s}$ and atomic weights A_1 , A_2 of colliding particles we expect very unusual behaviour of A_N and P_N as a function of kinematical variables. The data and model predictions of A_N for the π^+ production in pp-collisions are shown in Fig. 2a as a function of x_F . The data are from the E704 [8] and BRAHMS [9] experiments for cm energies 19, 62 and 200 GeV. The model predictions describe the data. The dashed curve is for 200 GeV and the solid curve is for 130 GeV. We expect a negative A_N for 200 GeV and x_F around 0.6 and also for 130 GeV and x_F in the range from 0.1 to 0.6. The negative A_N values are due to the u -quark spin precession in a strong ECF. Even more unusual oscillating behaviour of A_N is expected for 500 GeV (see Fig. 2b).

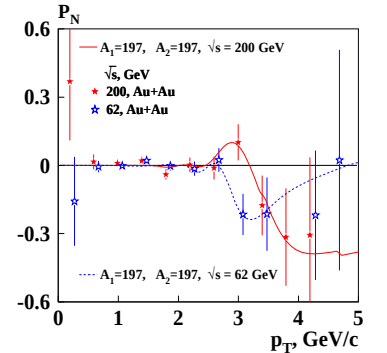


Figure 3: The dependence $P_N(p_T)$ of Λ -hyperon polarization in Au+Au collisions.

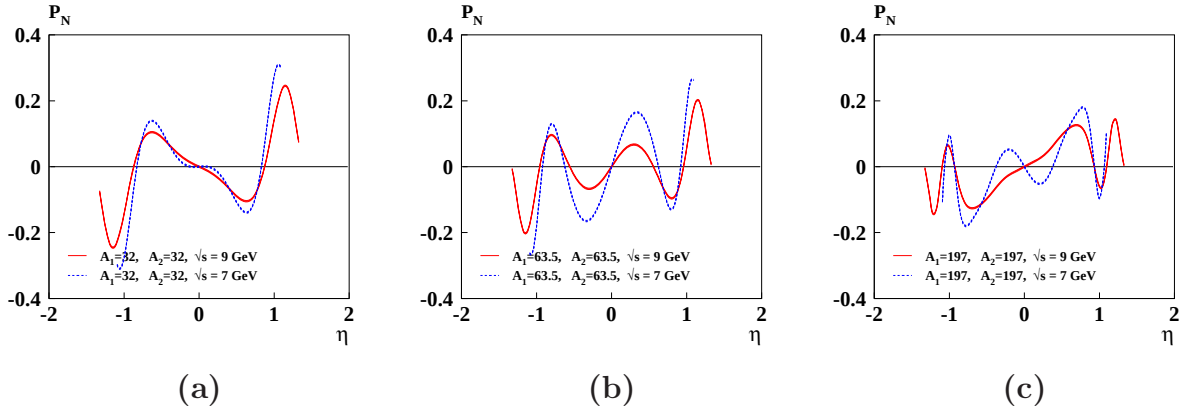


Figure 4: The transverse polarization of Λ vs $\eta = -\ln(\tan\theta_{CM}/2)$ at $p_T = 2.35$ GeV/c in (a) S+S, (b) Cu+Cu and (c) Au+Au collisions. Solid line corresponds to $\sqrt{s} = 9$ GeV and dashed line to $\sqrt{s} = 7$ GeV, respectively.

In Fig. 3 the global polarization P_N of Λ hyperon in Au+Au collisions is shown as a function of p_T . The data are from STAR experiment at 62 and 200 GeV [10]. The ECF model predictions reproduce the data. Oscillating behaviour of P_N is due to the s-quark spin precession in very strong color fields. Similar oscillating behaviour of P_N , as a function of pseudorapidity, is expected for low energies, 7 and 9 GeV in cm. The predictions are shown in Fig. 4 for S+S, Cu+Cu, and Au+Au collisions, respectively.

Conclusion: a semi-classical mechanism is proposed for single-spin phenomena. The effective color field of QCD strings, created by spectator quarks and antiquarks is described by the quark-counting rules. The microscopic Stern-Gerlach effect in the chromomagnetic field and the Thomas spin precession in the chromoelectric field lead to the SSA. The energy and atomic weight dependence of the effective color fields, combined with the quark spin precession phenomenon, lead to the oscillating behaviour of A_N and P_N . The model predictions for different reactions and energies can be checked at the existing accelerators. The work was partially supported by RFBR grant 09-02-00198.

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TOWARDS STUDY OF LIGHT SCALAR MESONS IN POLARIZATION PHENOMENA

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Abstract

After a short review of the production mechanisms of the light scalars which reveal their nature and indicate their quark structure, we suggest to study the mixing of the isovector $a_0^0(980)$ with the isoscalar $f_0(980)$ in spin effects.

1 Introduction.

The scalar channels in the region up to 1 GeV became a stumbling block of QCD. The point is that both perturbation theory and sum rules do not work in these channels because there are not solitary resonances in this region.

As experiment suggests, in chiral limit confinement forms colourless observable hadronic fields and spontaneous breaking of chiral symmetry with massless pseudoscalar fields. There are two possible scenarios for QCD realization at low energy: 1. $U_L(3) \times U_R(3)$ non-linear σ model, 2. $U_L(3) \times U_R(3)$ linear σ model. The experimental nonet of the light scalar mesons suggests $U_L(3) \times U_R(3)$ linear σ model.

2 $SU_L(2) \times SU_R(2)$ Linear σ Model, Chiral Shielding in $\pi\pi \rightarrow \pi\pi$ [1]

Hunting the light σ and κ mesons had begun in the sixties. But the fact that both $\pi\pi$ and πK scattering phase shifts do not pass over 90° at putative resonance masses prevented to prove their existence in a conclusive way.

Situation changes when we showed that in the linear σ model there is a negative background phase which hides the σ meson [1]. It has been made clear that shielding wide lightest scalar mesons in chiral dynamics is very natural. This idea was picked up and triggered new wave of theoretical and experimental searches for the σ and κ mesons.

Our approximation is as follows (see Fig. 1): $T_0^{0(tree)} = \frac{m_\pi^2 - m_\sigma^2}{32\pi f_\pi^2} \left[5 - 3 \frac{m_\pi^2 - m_\sigma^2}{m_\sigma^2 - s} - 2 \frac{m_\pi^2 - m_\sigma^2}{s - 4m_\pi^2} \right] \times \ln \left(1 + \frac{s - 4m_\pi^2}{m_\sigma^2} \right)$, $T_0^0 = \frac{T_0^{0(tree)}}{1 - i\rho_{\pi\pi} T_0^{0(tree)}} = \frac{e^{2i\delta_{bg} + \delta_{res}} - 1}{2i\rho_{\pi\pi}} = \frac{e^{2i\delta_0^0} - 1}{2i\rho_{\pi\pi}} = T_{bg} + e^{2i\delta_{bg}} T_{res}$, $T_{res} = \frac{\sqrt{s} \Gamma_{res}(s) / \rho_{\pi\pi}}{M_{res}^2 - s + \text{Re} \Pi_{res}(M_{res}^2) - \Pi_{res}(s)} = \frac{e^{2i\delta_{res}} - 1}{2i\rho_{\pi\pi}}$, $T_{bg} = \frac{e^{2i\delta_{bg}} - 1}{2i\rho_{\pi\pi}} = \frac{\lambda(s)}{1 - i\rho_{\pi\pi} \lambda(s)}$, $\lambda(s) = \frac{m_\pi^2 - m_\sigma^2}{32\pi f_\pi^2} \left[5 - 2 \times \frac{m_\pi^2 - m_\sigma^2}{s - 4m_\pi^2} \ln \left(1 + \frac{s - 4m_\pi^2}{m_\sigma^2} \right) \right]$, $\text{Re} \Pi_{res}(s) = -\frac{g_{res}^2(s)}{16\pi} \lambda(s) \rho_{\pi\pi}^2$, $\text{Im} \Pi_{res}(s) = \sqrt{s} \Gamma_{res}(s) = \frac{g_{res}^2(s) \rho_{\pi\pi}}{16\pi}$, $g_{res}(s) = \frac{g_{\sigma\pi\pi}}{[1 - i\rho_{\pi\pi} \lambda(s)]}$, $M_{res}^2 = m_\sigma^2 - \text{Re} \Pi_{res}(M_{res}^2)$, $\rho_{\pi\pi} = \sqrt{1 - \frac{4m_\pi^2}{s}}$, $g_{\sigma\pi\pi} = \sqrt{\frac{3}{2}} g_{\sigma\pi^+\pi^-}$, $= \sqrt{\frac{3}{2}} \frac{m_\pi^2 - m_\sigma^2}{f_\pi}$, $T_0^{2(tree)} = \frac{m_\pi^2 - m_\sigma^2}{16\pi f_\pi^2} \left[1 - \frac{m_\pi^2 - m_\sigma^2}{s - 4m_\pi^2} \ln \left(1 + \frac{s - 4m_\pi^2}{m_\sigma^2} \right) \right]$, $T_0^2 = \frac{T_0^{2(tree)}}{1 - i\rho_{\pi\pi} T_0^{2(tree)}} = \frac{e^{2i\delta_0^2} - 1}{2i\rho_{\pi\pi}}$.

The results in our approximation are: $M_{res} = 0.43 \text{ GeV}$, $\Gamma_{res}(M_{res}^2) = 0.67 \text{ GeV}$, $m_\sigma = 0.93 \text{ GeV}$, $\Gamma_{res}^{renorm}(M_{res}^2) = \frac{\Gamma_{res}(M_{res}^2)}{1 + d \text{Re} \Pi_{res}(s) / ds|_{s=M_{res}^2}} = 0.53 \text{ GeV}$, $g_{res}(M_{res}^2) / g_{\sigma\pi\pi} = 0.33$, $a_0^0 = 0.18 m_\pi^{-1}$, $a_0^2 = -0.04 m_\pi^{-1}$, the Adler zeros $(s_A)_0^0 = 0.45 m_\pi^2$ and $(s_A)_0^2 = 2.02 m_\pi^2$. The chiral shielding of the $\sigma(600)$ meson in $\pi\pi \rightarrow \pi\pi$ is illustrated in Fig. 2 with the help of the $\pi\pi$ phase shifts δ_{res} , δ_{bg} , δ_0^0 (a), and with the help of the corresponding cross sections (b).

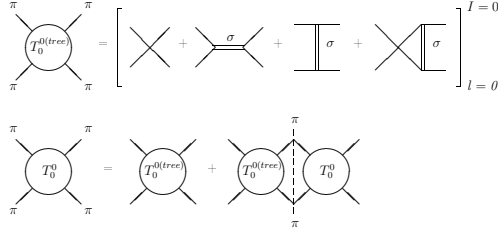


Figure 1: The graphical representation of the S wave $I = 0$ $\pi\pi$ scattering amplitude T_0^0 .

3 The σ Propagator [1]

$1/D_\sigma(s) = 1/[M_{res}^2 - s + \text{Re}\Pi_{res}(M_{res}^2) - \Pi_{res}(s)]$. The σ meson self-energy $\Pi_{res}(s)$ is caused by the intermediate $\pi\pi$ states, that is, by the four-quark intermediate states. This contribution shifts the Breit-Wigner (BW) mass greatly $m_\sigma - M_{res} \approx 0.50$ GeV. So, half the BW mass is determined by the four-quark contribution at least. The imaginary part dominates the propagator modulus in the region $0.3 \text{ GeV} < \sqrt{s} < 0.6 \text{ GeV}$. So, the σ field is described by its four-quark component at least in this energy (virtuality) region.

4 Four-quark Model

The nontrivial nature of the well-established light scalar resonances $f_0(980)$ and $a_0(980)$ is no longer denied practically anybody. As for the nonet as a whole, even a cursory look at PDG Review gives an idea of the four-quark structure of the light scalar meson nonet, $\sigma(600)$, $\kappa(700 - 900)$, $f_0(980)$, and $a_0(980)$, inverted in comparison with the classical P wave $q\bar{q}$ tensor meson nonet $f_2(1270)$, $a_2(1320)$, $K_2^*(1420)$, $\phi'_2(1525)$. Really, while the scalar nonet cannot be treated as the P wave $q\bar{q}$ nonet in the naive quark model, it can be easy understood as the $q^2\bar{q}^2$ nonet, where σ has no strange quarks, κ has the s quark, f_0 and a_0 have the $s\bar{s}$ pair. Similar states were found by Jaffe in 1977 in the MIT bag [2].

5 Radiative Decays of the ϕ Meson and the K^+K^- Loop Model [3]

Ten years later we showed that $\phi \rightarrow \gamma a_0 \rightarrow \gamma\pi\eta$ and $\phi \rightarrow \gamma f_0 \rightarrow \gamma\pi\pi$ can shed light on the problem of the $a_0(980)$ and $f_0(980)$ mesons. Now these decays are studied not only theoretically but also experimentally. When basing the experimental investigations, we suggested one-loop model $\phi \rightarrow K^+K^- \rightarrow \gamma a_0/f_0$, see Fig. 3. This model is used in the data treatment and is ratified by experiment, see Fig. 4. Gauge invariance gives the conclusive arguments in favor of the K^+K^- loop transition as the principal mechanism of the $a_0(980)$ and $f_0(980)$ meson production in the ϕ radiative decays.

6 The K^+K^- Loop Mechanism is Four-Quark Transition [3]

In truth this means that the $a_0(980)$ and the $f_0(980)$ are seen in the ϕ meson radiative decays owing to the K^+K^- intermediate state. So, the mechanism of the $a_0(980)$ and $f_0(980)$ production in the ϕ meson radiative decays is established at a physical level of proof. We are dealing with the four-quark transition. A radiative four-quark transition between two $q\bar{q}$ states requires creation and annihilation of an additional $q\bar{q}$ pair, i.e., such a transition is forbidden by the OZI rule, while a radiative four-quark transition between $q\bar{q}$ and $q^2\bar{q}^2$ states requires only creation of an additional $q\bar{q}$ pair, i.e., such a transition is allowed by the OZI rule. The large N_C expansion supports this conclusion.

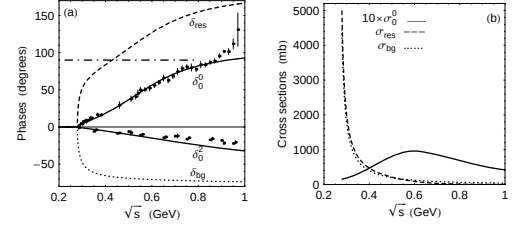


Figure 2: The σ model. Our approximation. $\delta_0^0 = \delta_{res} + \delta_{bg}$. $(\sigma_0^0, \sigma_{res}, \sigma_{bg}) = \frac{32\pi}{s} (|T_0^0|^2, |T_{res}|^2, |T_{bg}|^2)$.

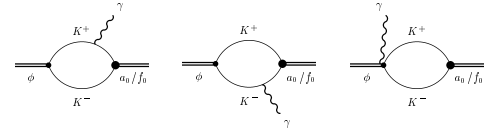


Figure 3: The K^+K^- loop model.

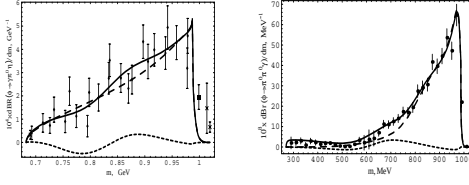


Figure 4: The left (right) plot shows the fit to the KLOE data for the $\pi^0\eta$ ($\pi^0\pi^0$) mass spectrum in the $\phi \rightarrow \gamma\pi^0\eta$ ($\phi \rightarrow \gamma\pi^0\pi^0$) decay caused by the $a_0(980)$ ($\sigma(600)+f_0(980)$) production through the K^+K^- loop mechanism.

7 Scalar Nature and Production

Mechanisms in $\gamma\gamma$ collisions [4]

Twenty seven years ago we predicted the suppression of $a_0(980) \rightarrow \gamma\gamma$ and $f_0(980) \rightarrow \gamma\gamma$ in the $q^2\bar{q}^2$ MIT model, $\Gamma_{a_0 \rightarrow \gamma\gamma} \sim \Gamma_{f_0 \rightarrow \gamma\gamma} \sim 0.27$ keV. Experiment supported this prediction.

Recently the experimental investigations have made great qualitative advance. The Belle Collaboration published data on $\gamma\gamma \rightarrow \pi^+\pi^-$, $\gamma\gamma \rightarrow \pi^0\pi^0$, and $\gamma\gamma \rightarrow \pi^0\eta$, whose statistics are huge [5], see Fig. 5. They not only proved the theoretical expectations based on the four-quark nature of the light scalar mesons, but also have allowed to elucidate the principal mechanisms of these processes. Specifically, the direct coupling constants of the $\sigma(600)$, $f_0(980)$, and $a_0(980)$ resonances with the system are small with the result that their decays into $\gamma\gamma$ are the four-quark transitions caused by the rescatterings $\sigma(600) \rightarrow \pi^+\pi^- \rightarrow \gamma\gamma$, $f_0(980) \rightarrow K^+K^- \rightarrow \gamma\gamma$ and $a_0(980) \rightarrow K^+K^- \rightarrow \gamma\gamma$ in contrast to the $\gamma\gamma$ decays of the classic P wave tensor $q\bar{q}$ mesons $a_2(1320)$, $f_2(1270)$ and $f_2'(1525)$, which are caused by the direct two-quark transitions $q\bar{q} \rightarrow \gamma\gamma$ in the main. As a result the practically model-independent prediction of the $q\bar{q}$ model $g_{f_2\gamma\gamma}^2 : g_{a_2\gamma\gamma}^2 = 25 : 9$ agrees with experiment rather well. The two-photon light scalar widths averaged over resonance mass distributions $\langle \Gamma_{f_0 \rightarrow \gamma\gamma} \rangle_{\pi\pi} \approx 0.19$ keV, $\langle \Gamma_{a_0 \rightarrow \gamma\gamma} \rangle_{\pi\eta} \approx 0.3$ keV and $\langle \Gamma_{\sigma \rightarrow \gamma\gamma} \rangle_{\pi\pi} \approx 0.45$ keV. As to the ideal $q\bar{q}$ model prediction $g_{f_0\gamma\gamma}^2 : g_{a_0\gamma\gamma}^2 = 25 : 9$, it is excluded by experiment.

8 Summary of the Above [1, 3, 4]

(i) The mass spectrum of the light scalars, $\sigma(600)$, $\kappa(800)$, $f_0(980)$, $a_0(980)$, gives an idea of their $q^2\bar{q}^2$ structure. (ii) Both intensity and mechanism of the $a_0(980)/f_0(980)$ production in the $\phi(1020)$ radiative decays, the $q^2\bar{q}^2$ transitions $\phi \rightarrow K^+K^- \rightarrow \gamma[a_0(980)/f_0(980)]$, indicate their $q^2\bar{q}^2$ nature. (iii) Both intensity and mechanism of the scalar meson decays into $\gamma\gamma$, the $q^2\bar{q}^2$ transitions $\sigma(600) \rightarrow \pi^+\pi^- \rightarrow \gamma\gamma$ and $[f_0(980)/a_0(980)] \rightarrow K^+K^- \rightarrow \gamma\gamma$, indicate their $q^2\bar{q}^2$ nature also.

9 The $a_0^0(980) - f_0(980)$ Mixing in Polarization Phenomena [6]

The $a_0^0(980) - f_0(980)$ mixing as a threshold phenomenon was discovered theoretically in 1979 in our work [6]. Now it is timely to study this phenomenon experimentally.¹

¹In Ref. [7] the search program of the $a_0^0(980) - f_0(980)$ mixing at the C/τ factory has been proposed. Recently the VES Collaboration published the data on the first effect of the $a_0^0(980) - f_0(980)$ mixing, $f_1(1420) \rightarrow \pi^0 a_0^0(980) \rightarrow \pi^0 f_0(980) \rightarrow 3\pi$ [8], in agreement with our calculation 1981 [6].

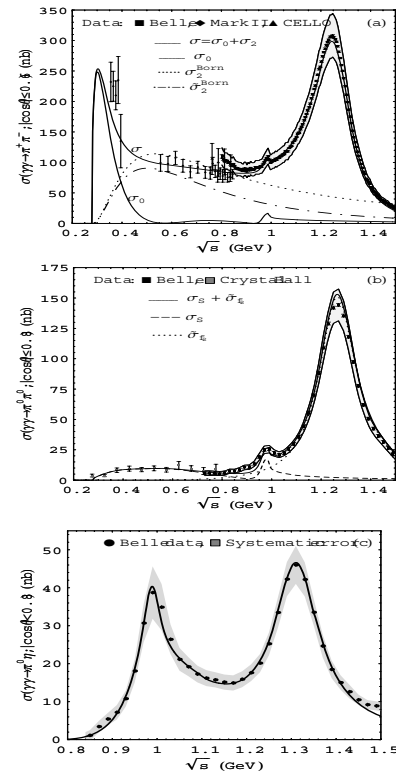


Figure 5: Descriptions of the Belle data on $\gamma\gamma \rightarrow \pi^+\pi^-$ (a), $\gamma\gamma \rightarrow \pi^0\pi^0$ (b), and $\gamma\gamma \rightarrow \pi^0\eta$ (c).

The main contribution originates from the $a_0^0(980) \rightarrow (K^+K^- + K^0\bar{K}^0) \rightarrow f_0(980)$ transition, see Fig. 6. Between the $K\bar{K}$ thresholds

$$|\Pi_{a_0f_0}(m)| \approx \frac{|g_{a_0K^+K^-} - g_{f_0K^+K^-}|}{16\pi} \sqrt{\frac{2(m_{K^0} - m_{K^+})}{m_{K^0}}} \approx 0.127 \frac{|g_{a_0K^+K^-} - g_{f_0K^+K^-}|}{16\pi} \simeq 0.03 \text{ GeV}^2.$$

It dominates for two reasons. i) It has the $\sqrt{m_d - m_u} \sim \sqrt{\alpha}$ order. ii) The strong coupling of the $a_0^0(980)$ and $f_0(980)$ to the $K\bar{K}$ channels, $|g_{a_0K^+K^-} - g_{f_0K^+K^-}|/4\pi \simeq 1 \text{ GeV}^2$.

We noted in 2004 [6] that the phase jump, see Fig. 4(b), suggests the idea to study the $a_0^0(980) - f_0(980)$ mixing in polarization phenomena. If a process amplitude with a suitable spin configuration is dominated by the $a_0^0(980) - f_0(980)$ mixing then a spin asymmetry of a cross section jumps near the $K\bar{K}$ thresholds. An example is the reaction $\pi^- p \uparrow \rightarrow (a_0^0(980) + f_0(980)) n \rightarrow a_0^0(980)n \rightarrow \eta\pi^0 n$, see Fig. 7. Performing the polarized target experiments on the reaction $\pi^- p \rightarrow \eta\pi^0 n$ at high energy could unambiguously and very easily establish the existence of the $a_0(980) - f_0(980)$ mixing phenomenon through the presence of a strong (~ 1) jump in the normalized azimuthal spin asymmetry of the S wave $\eta\pi^0$ production cross section near the $K\bar{K}$ thresholds. In turn it could give an exclusive information on the $a_0(980)$ and $f_0(980)$ coupling constants with the $K\bar{K}$ channels, $|g_{a_0K^+K^-} - g_{f_0K^+K^-}|/4\pi$.

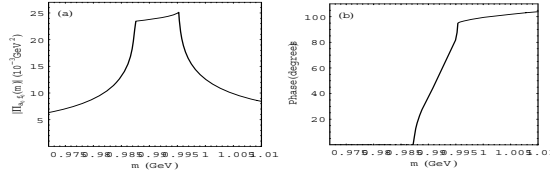


Figure 6: The “resonancelike” behavior of the modulus (a) and phase (b) of the $a_0^0(980) - f_0(980)$ mixing amplitude $\Pi_{a_0f_0}(m)$.

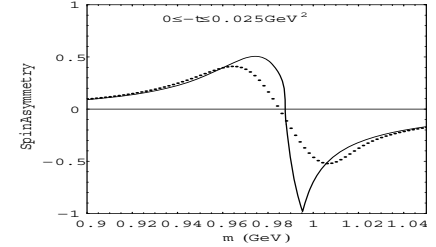


Figure 7: The spin asymmetry in $\pi^- p \uparrow \rightarrow (a_0^0(980) + f_0(980)) n \rightarrow \eta\pi^0 n$.

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RECURSIVE FRAGMENTATION MODEL WITH QUARK SPIN. APPLICATION TO QUARK POLARIMETRY

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Abstract

An elementary recursive model accounting for the quark spin in the fragmentation of a quark into mesons is presented. The quark spin degree of freedom is represented by a two-components spinor. Spin one meson can be included. The model produces Collins effect and jet handedness. The influence of the initial quark polarization decays exponentially with the rank of the meson, at different rates for longitudinal and transverse polarizations.

1 Introduction

Present Monte-Carlo event generators of quark and gluon jets do not include the parton spin degree of freedom, therefore do not generate the Collins [1] and jet handedness [2] effects. These are azimuthal asymmetries bearing on one, two or three hadrons, which can serve as *quark polarimeters*. However the asymmetries may strongly depend, in magnitude and sign, on the quark and hadron flavors and on the transverse momenta $\mathbf{p}_T$ and scaled longitudinal momenta z of these hadrons. Therefore a good knowledge of this dependence is needed for parton polarimetry. Due to the large number of kinematical variables, a hadronisation model which takes spin into account is urgently needed as a guide.

The semi-classical Lund $^3P_0$ mechanism [3], grafted on the string model, can generate a Collins effect [4], but not jet-handedness. Here we propose a fully quantum model of spinning quark fragmentation, based on the multiperipheral model. It reproduces the results of the $^3P_0$ mechanism and also contains the jet-handedness effect.

2 Some recalls about quark fragmentation

Figure 1 describes the creation of a quark " q_0 " and an antiquark " $\bar{q}_{-1}$ " in e^+e^- annihilation or $W^\pm$ decay, followed by the hadronisation,

$$q_0 + \bar{q}_{-1} \rightarrow h_1 + h_2 \dots + h_N . \quad (1)$$

Looking from right to left, one sees it as the recursive process (see [5] and ref. 4 of [6]),

$$\begin{array}{ll} q_0 \equiv q_0 \rightarrow h_1 + q_1 & \text{4-momenta : } k_0 = p_1 + k_1 , \\ q_1 \rightarrow h_2 + q_2 & k_1 = p_2 + k_2 , \\ \dots & \dots \\ q_{N-1} \rightarrow h_N + q_N & k_{N-1} = p_N + k_N . \end{array} \quad (2)$$

$q_N \equiv q_{-1}$ is a "quark propagating backward in time" and $k_N \equiv -k(\bar{q}_{-1})$.

Kinematical notations :

$\mathbf{k}_0 = \mathbf{k}(q_0)$ and $\mathbf{k}(\bar{q}_{-1})$ are in the $+\hat{\mathbf{z}}$ and $-\hat{\mathbf{z}}$ directions respectively. For a quark, $\mathbf{t}_n \equiv \mathbf{k}_{nT}$. For a 4-vector, $a^\pm = a^0 \pm a^z$ and $\mathbf{a}_T = (a^x, a^y)$. We denote by a tilde the *dual transverse vector* $\tilde{\mathbf{a}}_T \equiv \hat{\mathbf{z}} \times \mathbf{a}_T = (-a^y, a^x)$.

In Monte-Carlo simulations, the k_n are generated according to the *splitting distribution*

$$dW(q_{n-1} \rightarrow h_n + q_n) = f_n(\zeta_n, \mathbf{t}_{n-1}^2, \mathbf{t}_n^2, \mathbf{p}_{nT}^2) d\zeta_n d^2\mathbf{t}_n, \quad \zeta_n \equiv p_n^+/k_{n-1}^+.$$

In particular the *symmetric Lund* splitting function [3],

$$f_n \propto \zeta_n^{a_{n-1}-a_n-1} (1 - \zeta_n) \exp[-b(m_n^2 + \mathbf{p}_{nT}^2)/\zeta_n], \quad (3)$$

inspired by the string model, fulfills the requirement of *forward-backward equivalence*.

One can also consider [6] the upper part of Fig.1 as a **multiperipheral** [7] diagram with the Feynman amplitude

$$\begin{aligned} \mathcal{M}_{q_0 + \bar{q}_{-1} \rightarrow h_1 \dots + h_N} &= \bar{v}(k_{-1}, \mathbf{S}_{-1}) \Gamma_{q_N, h_N, q_{N-1}}(k_N, k_{N-1}) \Delta_{q_{N-1}}(k_{N-1}) \dots \\ &\dots \Delta_{q_2}(k_2) \Gamma_{q_2, h_2, q_1}(k_2, k_1) \Delta_{q_1}(k_1) \Gamma_{q_1, h_1, q_0}(k_1, k_0) u(k_0, \mathbf{S}_0). \end{aligned} \quad (4)$$

$\mathbf{S}_0$ and $\mathbf{S}_{-1}$ are the polarization vectors of the initial quark and antiquark. $\mathbf{S}^2 = 1$, $S_z =$ helicity, $\mathbf{S}_T =$ transversity. Γ and Δ are vertex functions and propagators which depend on the quark momenta and flavors. Note that Fig.1 is a loop diagram : k_0 is an integration variable, therefore the "jet axis" is not really defined. Furthermore, in Z_0 or γ^* decay, the spins q_0 and $\bar{q}_{-1}$ are entangled so that one cannot define $\mathbf{S}_0$ and $\mathbf{S}_{-1}$ separately.

Collins and jet-handedness effects. Let us first assume that the *jet axis* (quark direction) is well determined :

- the *Collins effect* [1], in $\vec{q} \rightarrow h + X$, is an asymmetry in $\sin[\varphi(\mathbf{S}) - \varphi(h)]$ for a transversely polarized quark. The fragmentation function reads

$$F(z, \mathbf{p}_T; \mathbf{S}_T) = F_0(z, \mathbf{p}_T^2) (1 + A_T \mathbf{S}_T \cdot \tilde{\mathbf{p}}_T / |\mathbf{p}_T|) \quad (\tilde{\mathbf{p}}_T \equiv \hat{\mathbf{z}} \times \mathbf{p}_T). \quad (5)$$

$A_T = A_T(z, \mathbf{p}_T^2) \in [-1, +1]$ is the Collins analyzing power.

- *jet handedness* [2], in $\vec{q} \rightarrow h + h' + X$, is an asymmetry in $\sin[\varphi(h) - \varphi(h')]$ proportional to the quark helicity. The 2-particle longitudinally polarized fragmentation function is

$$F(z, \mathbf{p}_T, z', \mathbf{p}'_T; S_z) = F_0(z, \mathbf{p}_T^2, z', \mathbf{p}'_T{}^2, \mathbf{p}_T \cdot \mathbf{p}'_T) \left(1 + A_L S_z \frac{\tilde{\mathbf{p}}_T \cdot \mathbf{p}'_T}{|\tilde{\mathbf{p}}_T \cdot \mathbf{p}'_T|} \right). \quad (6)$$

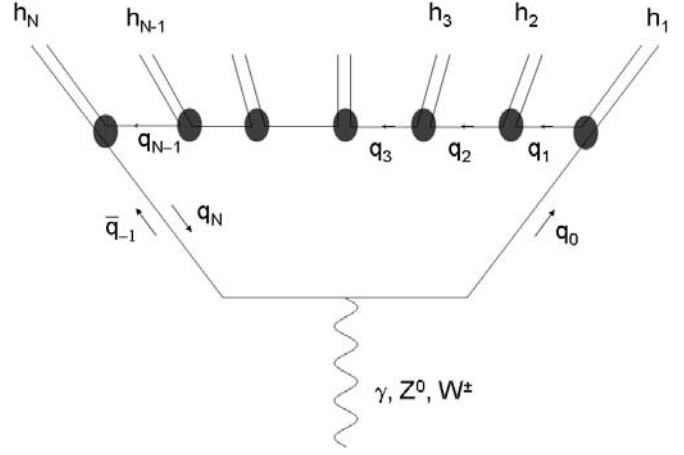


Figure 1: Electroweak boson $\rightarrow q\bar{q} \rightarrow$ mesons.

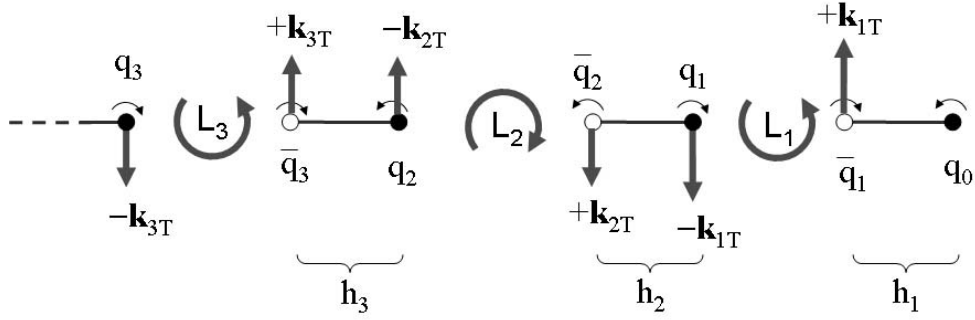


Figure 2: String decaying into pseudoscalar mesons.

$A_L = A_L(z, \mathbf{p}_T^2, z', \mathbf{p}_T'^2, \mathbf{p}_T \cdot \mathbf{p}_T') \in [-1, +1]$ is the handedness analyzing power. $\tilde{\mathbf{p}}_{1T} \cdot \mathbf{p}_T'$ is the same as $\hat{\mathbf{z}} \cdot (\mathbf{p}_T \times \mathbf{p}_T')$.

If the jet axis is not well determined, an additional fast hadron, h' or h'' is needed. The z axis is taken along $\mathbf{P} = (\mathbf{p} + \mathbf{p}')$ (Collins) or $\mathbf{P} = (\mathbf{p} + \mathbf{p}' + \mathbf{p}'')$ (handedness). In this way, we define the *2-particle relative Collins effect* (also called *interference fragmentation*) and the *three-particle jet handedness*, which corresponds to the original definition of [2].

The Lund $^3P_0$ mechanism [3]. Figure 2 depicts the decay of the initial massive string accompanied with the creation of a $q\bar{q}$ pairs. Forgetting transverse oscillations of the initial string, the transverse hadron momenta come from the internal orbital motions of the pairs. After a tunnel effect the q and $\bar{q}$ of a pair become on-shell and their relative position $\mathbf{r} \equiv \mathbf{r}(q) - \mathbf{r}(\bar{q})$ is along $-\hat{\mathbf{z}}$. The pair is assumed to be in the $^3P_0$ state, which has the vacuum quantum number. The relative momentum $\mathbf{k} \equiv \mathbf{k}(q) = -\mathbf{k}(\bar{q})$ and the orbital angular momentum $\mathbf{L} = \mathbf{r} \times \mathbf{k}$ are such that $\hat{\mathbf{z}} \cdot [\mathbf{k}_T \times \mathbf{L}] < 0$. In the $^3P_0$ state $\langle \mathbf{s}_q \rangle = \langle \mathbf{s}_{\bar{q}} \rangle = -\langle \mathbf{L}/2 \rangle$. As a result, the transverse spins of q and $\bar{q}$ are correlated to their transverse momenta :

$$\langle \hat{\mathbf{z}} \cdot [\mathbf{k}_T(\mathbf{q}) \times \mathbf{s}_q] \rangle > 0, \quad \langle \hat{\mathbf{z}} \cdot [\mathbf{k}_T(\bar{\mathbf{q}}) \times \mathbf{s}_{\bar{q}}] \rangle < 0. \quad (7)$$

The correlation can be transmitted to a baryon. Then $\langle \hat{\mathbf{z}} \cdot [\mathbf{p}_T \times \mathbf{s}_B] \rangle$ has the sign of $\langle \mathbf{s}_q \cdot \mathbf{s}_B \rangle$. This can explain transverse spin asymmetries in hyperon production [3].

Application to the Collins effect [4]. In Fig. 2, q_0 is polarized along the direction $+\hat{\mathbf{y}}$ toward the reader and h_1 is a pseudoscalar meson, for which $\langle \mathbf{s}(q_0) \rangle = -\langle \mathbf{s}(\bar{q}_1) \rangle$. Then q_1 and $\bar{q}_1$ are polarised along $-\hat{\mathbf{y}}$ and, according to (7), $\mathbf{k}_T(\bar{q}_1) = \mathbf{p}_T(h_1)$ is in the $+\hat{\mathbf{x}}$ direction. This provides a model for the Collins effect. Fig. 2 also indicates that, for a sequence of pseudoscalar mesons, the Collins asymmetries are of alternate sides. Besides, q_{n-1} and $\bar{q}_n$ go on the same side, which enhances the asymmetry. It may explain why π^- from u -quarks have a strong Collins analyzing power. Note that this effect also enhances $\langle \mathbf{p}_T^2 \rangle$, independently of the q_0 polarization.

3 A simplified multiperipheral quark model

In Eq.(4), let us replace Dirac spinors by Pauli spinors. A minimal model, restricted to the direct emission of pseudoscalar mesons, is built with the following prescriptions :

- 1) replace $u(k_0, \mathbf{S}_0)$ and $\bar{v}(k_{-1}, \mathbf{S}_{-1}) \equiv -\bar{u}(k_{\bar{q}-1}, -\mathbf{S}_{\bar{q}-1}) \gamma_5$ by the Pauli spinors $\chi(\mathbf{S}_0)$ and $-\chi^\dagger(-\mathbf{S}_{-1}) \sigma_z$,
- 2) assume no momentum dependence of Γ ,
- 3) replace γ_5 by σ_z ,
- 4) replace the usual pole $(k^2 - m_q^2)^{-1}$ of $\Delta_q(k)$ by the (k_L, k_T) separable form

$$D_q(k) = g_q(k^+ k^-) \exp(-B \mathbf{t}^2/2) , \quad (8)$$

- 5) replace the usual numerator $m_q + \gamma \cdot k$ by $\mu_q(k^+ k^-, \mathbf{t}^2) + i \boldsymbol{\sigma} \cdot \tilde{\mathbf{t}}$.

These prescriptions respect the invariance under the following transformations :

- rotation about the z -axis,
- Lorentz transformations along the z -axis (longitudinal boost),
- mirror reflection about any plane containing the z -axis (parity),
- forward-backward equivalence.

The jet axis being fixed, full Lorentz invariance is not required, whence the separate dependence of D_q and μ_q in $k^+ k^-$ and $\mathbf{t}^2$. In item 5), $\mu + i \boldsymbol{\sigma} \cdot \tilde{\mathbf{t}}$ is reminiscent of the meson-baryon scattering amplitude $f(s, t) + i g(s, t) \boldsymbol{\sigma} \cdot (\mathbf{p} \times \mathbf{p}')$. Single-spin effects are obtained for $\Im \mu \neq 0$. The choice of putting the spin dependence in the propagators rather than in the vertices is inspired by the $^3P_0$ mechanism : in both models the polarization germinates in the quark line between two hadrons.

For a fast investigation of the model, we make the further approximations :

- *Neglect the influence of the antiquark flavor and polarization in the quark fragmentation region.* This is allowed at large invariant $q_0 + \bar{q}_{-1}$ mass.
- *Discard the interference diagrams.* For a given final state, the *rank ordering* of hadrons in the multiperipheral diagram is not unique and differently ordered diagrams can interfere. This interference (and the resulting Bose-Einstein correlations) will be neglected.
- *Disentangle $k^\pm$ and $\mathbf{k}_T$.* We will assume that $\mu_q(k^+ k^-, \mathbf{t}^2)$ is constant or a function of $\mathbf{t}^2$ only. Thus we have no more "dynamical" correlation between longitudinal and transverse momenta. However there remains a "kinematical" correlation coming from the mass shell constraint

$$(k_{n-1} - k_n)^2 \equiv (k_{n-1}^+ - k_n^+)(k_{n-1}^- - k_n^-) - (\mathbf{t}_{n-1} - \mathbf{t}_n)^2 = m^2(h_n) . \quad (9)$$

In the following we will ignore the $(\mathbf{t}_{n-1} - \mathbf{t}_n)^2$ term. This approximation is drastic for pion emission because $\langle \mathbf{t}^2 \rangle > m_\pi^2$. We only use it here for a qualitative investigation of the spin effects allowed by the multiperipheral model. Thanks to it, the $\mathbf{t}$'s become fully decoupled from the $k_n^\pm$ and kinematically decorrelated between themselves. They remain correlated only via the quark spin.

p_T -distributions in the quark fragmentation region. With the above approximations we can treat the process (2), at least in p_T -space, like a cascade decay of unstable particles, which has no constraint coming from the future. The joint p_T distribution of the n first mesons is proportional to

$$I(\mathbf{p}_{1T}, \mathbf{p}_{2T}, \dots, \mathbf{p}_{nT}) = \exp(-B\mathbf{t}_1^2 - B\mathbf{t}_2^2 \dots - B\mathbf{t}_n^2) \text{Tr} \left\{ \mathbf{M}_{12\dots n} \frac{\mathbf{1} + \mathbf{S}_0 \cdot \boldsymbol{\sigma}}{2} \mathbf{M}_{12\dots n}^\dagger \right\}, \quad (10)$$

with

$$\mathbf{M}_{12\dots n} = \mathbf{M}_n \cdots \mathbf{M}_2 \mathbf{M}_1, \quad \mathbf{M}_r = (\mu_r + i\boldsymbol{\sigma} \cdot \tilde{\mathbf{t}}_r) \sigma_z. \quad (11)$$

3.1 Applications to azimuthal asymmetries

In this section we will calculate azimuthal asymmetries for particles of definite ranks. For comparison with experiments, one should mix the contributions of different rank assignments. For simplicity we take a unique and constant μ for all quark flavors.

First-rank Collins effect. Applying (10-11) for $n = 1$ gives

$$I(\mathbf{p}_{1T}) = \exp(-B\mathbf{t}_1^2) (|\mu|^2 + \mathbf{t}_1^2 - 2\Im(\mu) \tilde{\mathbf{t}}_1 \cdot \mathbf{S}), \quad (12)$$

with $\mathbf{t}_1 = -\mathbf{p}_{1T}$. For complex μ one has a Collins asymmetry (cf Eq.5) with

$$A_T = 2 \frac{\Im(\mu) |\mathbf{p}_{1T}|}{|\mu|^2 + \mathbf{p}_{1T}^2} \in [-1, +1]. \quad (13)$$

If $\Im(\mu) > 0$ it has the same sign as predicted by the $^3P_0$ mechanism.

Joint p_T spectrum of h_1 and h_2 . Applying (10-11) for $n = 2$ one obtains

$$\begin{aligned} I(\mathbf{p}_{1T}, \mathbf{p}_{2T}) &= \exp(-B\mathbf{t}_1^2 - B\mathbf{t}_2^2) \{ (|\mu|^2 + \mathbf{t}_1^2)(|\mu|^2 + \mathbf{t}_2^2) - 4\mathbf{t}_1 \cdot \mathbf{t}_2 \Im(\mu) \\ &+ 2\Im(\mu) \mathbf{S} \cdot \tilde{\mathbf{t}}_1 (2\mathbf{t}_1 \cdot \mathbf{t}_2 - |\mu|^2 - \mathbf{t}_2^2) \\ &+ 2\Im(\mu) \mathbf{S} \cdot \tilde{\mathbf{t}}_2 (|\mu|^2 - \mathbf{t}_1^2) \\ &- 2\Im(\mu^2) \mathbf{S} \cdot (\mathbf{t}_1 \times \mathbf{t}_2) \}, \end{aligned} \quad (14)$$

with $\mathbf{t}_1 = -\mathbf{p}_{1T}$, $\mathbf{t}_2 = -(\mathbf{p}_{1T} + \mathbf{p}_{2T})$ and $\mathbf{S} \cdot (\mathbf{t}_1 \times \mathbf{t}_2) = S_z \tilde{\mathbf{p}}_{1T} \cdot \mathbf{p}_{2T}$.

The last line contains jet handedness (cf Eq.6), of analyzing power

$$A_L = \frac{-2\Im(\mu^2) |\mathbf{p}_{1T} \times \mathbf{p}_{2T}|}{(|\mu|^2 + \mathbf{t}_1^2)(|\mu|^2 + \mathbf{t}_2^2) - 4\mathbf{t}_1 \cdot \mathbf{t}_2 \Im(\mu)} \in [-1, +1]. \quad (15)$$

The second line contains the Collins asymmetry of h_1 . Both 2nd and 3rd lines contribute to the h_2 one after integration over $\mathbf{t}_1$, and to the *relative 2-particle Collins asymmetry*, which bears on

$$\mathbf{r}_{12} = \frac{z_2 \mathbf{p}_{1T} - z_1 \mathbf{p}_{2T}}{z_1 + z_2} = \left[\frac{z_1}{z_1 + z_2} \right] \mathbf{t}_2 - \mathbf{t}_1. \quad (16)$$

Note that Collins and jet-handedness asymmetries are not maximum for the same value of $\arg(\mu)$. This is related to the *positivity* [8] constraint

$$A_L^2(\mathbf{p}_{1T}, \mathbf{p}_{2T}) + A_T^2(\mathbf{p}_{1T}, \mathbf{p}_{2T}) \leq 1. \quad (17)$$

3.2 Evolution of the polarization of the cascading quark

Let us first assume that $\mathbf{t}_1, \mathbf{t}_2, \dots, \mathbf{t}_n$ are fixed and consider the spin density matrix $\rho_n = (\mathbf{1} + \mathbf{S}_n \cdot \boldsymbol{\sigma})/2$ of q_n at the $(n+1)^{\text{th}}$ vertex :

$$\rho_n = R_n / \text{Tr}\{R_n\}, \quad R_n = \mathbf{M}_{12\dots n} \frac{\mathbf{1} + \mathbf{S}_0 \cdot \boldsymbol{\sigma}}{2} \mathbf{M}_{12\dots n}^\dagger. \quad (18)$$

If ρ_0 is a pure state ($\det \rho_0 = 0$), then ρ_n is also a pure state ; no information is lost.

Let us now integrate over $\mathbf{t}_1, \mathbf{t}_2, \dots, \mathbf{t}_n$ (equivalently over $\mathbf{p}_{1T}, \dots, \mathbf{p}_{nT}$). It leads to a loss of information. The spin density matrix of q_n becomes

$$\bar{\rho}_n = \bar{R}_n / \text{Tr}\{\bar{R}_n\}, \quad \bar{R}_n = \int d^2\mathbf{t}_1 \dots \int d^2\mathbf{t}_n \mathbf{M}_{12\dots n} \frac{\mathbf{1} + \mathbf{S}_0 \cdot \boldsymbol{\sigma}}{2} \mathbf{M}_{12\dots n}^\dagger. \quad (19)$$

R_n and $\bar{R}_n$ obey the recursion relations

$$R_n = \mathbf{M}_n R_{n-1} \mathbf{M}_n^\dagger, \quad \bar{R}_n = \int d^2\mathbf{t}_n \mathbf{M}_n \bar{R}_{n-1} \mathbf{M}_n^\dagger. \quad (20)$$

At fixed $\mathbf{t}$'s, the left equation gives (setting $\mu = \mu' + i\mu''$) :

$$\mathbf{S}_n = \frac{1}{C} \left\{ 2\mu'' \tilde{\mathbf{t}}_n + \mathcal{R}[\hat{\mathbf{z}}, \varphi_n] \begin{pmatrix} \mathbf{t}^2 - |\mu|^2 & 0 & -2|\mathbf{t}| \mu' \\ 0 & -|\mu|^2 - \mathbf{t}^2 & 0 \\ 2|\mathbf{t}| \mu' & 0 & |\mu|^2 - \mathbf{t}^2 \end{pmatrix} \mathcal{R}[\hat{\mathbf{z}}, -\varphi_n] \mathbf{S}_{n-1} \right\}, \quad (21)$$

with $C = \text{Tr}\{R_n\} = |\mu|^2 + \mathbf{t}_n^2 - 2\mu'' \tilde{\mathbf{t}}_n \cdot \mathbf{S}_{n-1}$. The rotation $\mathcal{R}[\hat{\mathbf{z}}, \varphi_n]$ about $\hat{\mathbf{z}}$ brings $\hat{\mathbf{x}}$ along $\mathbf{t}_n$. Iterating (12), where we replace $\{\mathbf{S}, \mathbf{t}_1\}$ by $\{\mathbf{S}_{n-1}, \mathbf{t}_n\}$, and (21), we generate the successive transverse momenta with the Monte-Carlo method. From (21) we learn :

- if $\Im \mu \neq 0$, the inhomogeneous term in $\mu'' \tilde{\mathbf{t}}_n$ is a source (or sink) of transverse polarisation : one can have $\mathbf{S}_{nT} \neq 0$ even with $\mathbf{S}_{n-1} = 0$.
- helicity is partly converted into transversity along $\mathbf{t}_n$ and vice-versa.

The last fact explains the mechanism of jet handedness in this model : first, the helicity S_{z0} is partly converted into $\mathbf{S}_{1T}$ parallel to $\mathbf{p}_{1T}$, then $\mathbf{S}_{1T}$ produces a Collins asymmetry for h_2 in the plane perpendicular to $\mathbf{p}_{1T}$.

Let us now consider the $\mathbf{t}$ -integrated density matrix. The right equation in (20) gives

$$S_{n,z} = D_{LL} S_{n-1,z}, \quad \mathbf{S}_{n,T} = D_{TT} \mathbf{S}_{n-1,T}; \quad D_{LL}, D_{TT} \in [-1, +1], \quad (22)$$

$$\left(\frac{D_{LL}}{D_{TT}} \right) = \int d^2\mathbf{t} \exp(-B\mathbf{t}^2) \begin{pmatrix} |\mu|^2 - \mathbf{t}^2 \\ -|\mu|^2 \end{pmatrix} \bigg/ \int d^2\mathbf{t} \exp(-B\mathbf{t}^2) (|\mu|^2 + \mathbf{t}^2). \quad (23)$$

Analytical values : $D_{LL} = (\xi - 1)/(\xi + 1)$ and $D_{TT} = -\xi/(\xi + 1)$ with $\xi = B|\mu|^2$. The geometrical decays of $|S_{n,z}|$ and $|\mathbf{S}_{n,T}|$ along the quark chain occur at different speeds. They are similar to the decays of charge and strangeness correlations. D_{LL} and D_{TT} saturate a Soffer-type [8] positivity condition

$$2|D_{TT}| \leq 1 + D_{LL}. \quad (24)$$

Indeed, $2D_{TT} = -1 - D_{LL}$. This is due to the zero spin of h_n (compare with text after Eq.(4.87) of [8]). The negative value of D_{TT} leads to Collins asymmetries of alternate signs, in accordance with the $^3P_0$ mechanism. It comes from the σ_z vertex for pseudoscalar mesons. For *scalar* mesons we replace σ_z by $\mathbf{1}$. In this case D_{TT} is positive, q_{n-1} and $\bar{q}_n$ tend towards opposite sides and the Collins effect is small, except for h_1 . This is also the prediction of the $^3P_0$ mechanism.

4 Inclusion of spin-1 mesons

For a $J^{PC} = 1^{--}$ vector meson and the associated self-conjugate multiplet, the "minimal" emission vertex written with Pauli matrices is

$$\Gamma = G_L V_z^* \mathbf{1} + G_T \boldsymbol{\sigma} \cdot \mathbf{V}_T^* \sigma_z, \quad (25)$$

where $\mathbf{V}$ is the vector amplitude of the meson normalized to $\mathbf{V} \cdot \mathbf{V}^* = 1$. It is obtained from the relativistic 4-vector V^μ first by a longitudinal boost which brings the hadron at $p_z = 0$, then a transverse boost which brings the hadron at rest.

For a $J^{PC} = 1^{++}$ axial meson of amplitude $\mathbf{A}$, the "minimal" emission vertex is

$$\Gamma = \tilde{G}_T \boldsymbol{\sigma} \cdot \mathbf{A}_T^*. \quad (26)$$

It differs by a σ_z matrix from the second term of (25). A term of the form $\tilde{G}_L A_z^* \sigma_z$ with constant $\tilde{G}_L$ is not allowed by the forward-backward equivalence.

Let us treat the case where the 1st-rank particle is a ρ^+ meson and fix the momenta $p(\pi^+)$ and $p(\pi^0)$ of the decay pions. Then $V^\mu \propto p(\pi^+) - p(\pi^0)$, which is real, corresponding to a linear polarisation. Replacing the σ_z coupling of (11) by (25) we obtain

$$\begin{aligned} I(\mathbf{p}_T, \mathbf{V}) = & \exp(-B\mathbf{t}^2) |G_T|^2 \times \\ & \{ (|\alpha|^2 V_z^2 + \mathbf{V}_T^2) (|\mu|^2 + \mathbf{t}^2) - 4\mathbf{V}_T \cdot \mathbf{t} V_z \Im(\alpha) \Im(\mu) \\ & + 2 \Im(\mu) |\alpha|^2 V_z^2 \mathbf{S} \cdot \tilde{\mathbf{t}} \\ & + 2 \Im(\alpha) (|\mu|^2 + \mathbf{t}^2) V_z \mathbf{V}_T \cdot \tilde{\mathbf{S}} \\ & + 2 \Im(\mu) (\mathbf{V}_T \cdot \tilde{\mathbf{t}} \mathbf{V}_T \cdot \mathbf{S} + \mathbf{V}_T \cdot \mathbf{t} \mathbf{V}_T \cdot \tilde{\mathbf{S}}) \\ & + 4 \Re(\alpha) \Im(\mu) S_z V_z \mathbf{V}_T \cdot \tilde{\mathbf{t}} \} , \end{aligned} \quad (27)$$

with $\mathbf{t} \equiv \mathbf{t}_1 = -\mathbf{p}_T(\rho^+)$ and $\alpha \equiv G_L/G_T$. Let us comment this formula :

- The 2nd line is for unpolarized quark. It gives some tensor polarization.
- The 3rd line is a Collins effect for the ρ^+ as a whole, opposite to the pion one (compare with (12) and only for longitudinal linear polarization, in accordance with the $^3P_0$ mechanism. For $\langle V_z^2 \rangle = 1/3$ (unpolarized ρ^+) and $\alpha = 1$ one recovers the Czyzewski prediction [9] $A_T(\text{leading } \rho)/A_T(\text{leading } \pi) = -1/3$.
- The 4th line gives an oblique polarization in the plane perpendicular to $\mathbf{S}_T$ corresponding to $\hat{h}_1$ or h_{1LT} of [10, 11]. After ρ^+ decay, it becomes a relative $\pi^+ - \pi^0$ Collins effect.
- The 5th line is a new type of asymmetry, in $\sin[2\varphi(\mathbf{V}) - \varphi(\mathbf{t}) - \varphi(\mathbf{S})]$.
- The last line also gives an oblique polarization, but in the plane perpendicular to $\mathbf{p}_T(\rho^+)$. After ρ^+ decay, it becomes jet-handedness. Indeed, ignoring an effect of transverse boost, we have $\mathbf{V}_T \propto \mathbf{p}_T(\pi^+) - \mathbf{p}_T(\pi^0)$, therefore $\mathbf{V}_T \cdot \tilde{\mathbf{t}} \propto -\mathbf{p}_T(\pi^+) \times \mathbf{p}_T(\pi^0)$.

5 Conclusion

For the direct fragmentation of a transversely polarized quark into pseudo-scalar mesons, the model we have presented has essentially one free complex parameter μ and reproduces the results of the semi-classical $^3P_0$ mechanism : large asymmetry for the 2nd-rank meson, Collins asymmetries of alternate sides for the subsequent mesons. In addition, it possesses

a jet-handedness asymmetry, generated in two steps : partial transformation of helicity into transversity, then Collins effect.

We have also considered the inclusion of spin-1 mesons. When longitudinally polarized, a leading vector meson has a Collins asymmetry opposite to that of a pseudoscalar, as also expected from the $^3P_0$ mechanism. The decay pions of a ρ meson exhibit a relative Collins effects as well as jet-handedness. These effects are associated to oblique linear polarizations of the ρ meson. The fact that two pions coming from a ρ show the same spin effects as two successive "direct" pions is reminiscent of duality.

Even for unpolarized initial quarks, the spin degree of freedom of the cascading quark has to be considered. It enhances the $\langle \mathbf{p}_T^2 \rangle$ of the pseudoscalar mesons compared to scalar and longitudinal vector mesons.

A next task for building a realistic Monte-Carlo generator with quark spin is to take into account the $(\mathbf{t}_{n-1} - \mathbf{t}_n)^2$ term in (9). One must also be aware that there exist other mechanisms of spin asymmetries in jets. For example the Collins effect can be generated by the interference between direct emission and the emission via a resonance [12].

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CONSTITUENT QUARK REST ENERGY AND WAVE FUNCTION ACROSS THE LIGHT CONE

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Abstract

It is shown that for a constituent quark in the intra-nucleon self-consistent field the spin-orbit interaction lowers the quark rest energy to values ~ 100 MeV, which agrees with the DIS momentum sum rule. The possibility of violation of the spectral condition for the light-cone momentum component of a bound quark is discussed.

Given the information from the DIS momentum sum rule [1] that half of the nucleon mass is carried by gluons, one may lower the expectations for the constituent quark mass from $M_N/3$ to $\lesssim M_N/6$. This may also provide grounds for formation of a gluonic mean-field in the nucleon, permitting single-quark description of the wave-function, and (by analogy with the situation in atomic mean fields), making LS-interaction dominant over the SS.¹ Here, adopting the mentioned assumptions for the nucleon wave function, we will show that the constituent quark mass $\sim 50 \div 100$ MeV, in fact, results from the empirical values for the nucleon mean charge radius and the magnetic moment. We will also address the issue of valence quark wave function continuity and sizeable value at zero light-front momentum, which owes to rather low a value of the constituent quark mass, and discuss phenomenological implications thereof.

Wave function ansatz. For a spherically symmetric ground state, by parity reasons, the upper Dirac component has orbital momentum $l = 0$, whereas the lower one has $l = 1$ and flipped spin. This is accounted for by writing the single-quark wave function

$$\psi_q(\mathbf{r}, t) = e^{-i\kappa^0 t} \begin{pmatrix} \varphi(r)w \\ -i\boldsymbol{\sigma} \cdot \frac{\mathbf{r}}{r} \chi(r)w \end{pmatrix} \quad (\kappa^0 > 0) \quad (1)$$

with w being an $\mathbf{r}$ -independent Pauli spinor.

The relation between φ and χ to a first rough approximation is determined by the quark rest energy. For instance, for a Dirac particle of mass m moving in a static potential $V(r)$ (being either a 4th component of vector, or a scalar – cf. [2]), in the ground state

$$\chi(r) = \frac{1}{m + E - V(r)} \varphi'(r). \quad (2)$$

Thereat, a reasonable model may result from replacing the denominator of (2) by its average $\langle \kappa^0 \rangle$ (which may somewhat differ from κ^0 in (1)):

$$\chi(r) = \frac{1}{2\langle \kappa^0 \rangle} \varphi'(r).$$

¹Neglecting the SS-interaction and employing the single-quark model for the nucleon wave function, one must be ready that the accuracy is not better than $\frac{M_\Delta - M_N}{M_N} \sim 30 \div 40\%$.

For bag models [1], the proportionality between χ and φ' holds exactly, but the particular shape of bag model wave functions, especially χ , may be not phenomenologically reliable due to the influence of the bag sharp boundary. Instead, we prefer to take φ a smooth function for all r from 0 to ∞ :

$$\psi_q(\boldsymbol{\kappa}) = \left(1 + \frac{1}{2\langle\kappa^0\rangle}\gamma^0\boldsymbol{\gamma}\cdot\boldsymbol{\kappa}\right)\varphi(\kappa)\begin{pmatrix} w \\ 0 \end{pmatrix} \equiv \varphi(\kappa)\begin{pmatrix} w \\ \frac{\boldsymbol{\sigma}\cdot\boldsymbol{\kappa}}{2\langle\kappa^0\rangle}w \end{pmatrix}, \quad (3)$$

$$\varphi(\kappa) \propto e^{-a^2\kappa^2/2}. \quad (4)$$

Thereby, the model involves only 2 parameters: $\langle\kappa^0\rangle$ and the Gaussian radius a . It may describe the valence quark component in the nucleon, and is suitable for calculation of matrix elements of vector (conserved) currents, whereas axial vector currents require the account of sea quarks, which is beyond our scope here.

Proton mean charge radius and magnetic moment.

With recoilless spectators, our model can describe only formfactors at $Q \ll M_N$. Equating Q/M_N to our accuracy $30 \div 40\%$, one finds $Q < 0.4$ GeV. In this region the decrease of the form-factors is approximately linear [1], whereby they are characterized basically by the proton's magnetic moment and the mean charge (or magnetic) radius.²

The form-factors are defined as overlaps of ψ_q with ψ'_q , of the same ground state, but with a shifted momentum and a different polarization w' :

$$\int \frac{d^3\kappa}{(2\pi)^3} \bar{\psi}'_q\left(\boldsymbol{\kappa} + \frac{\mathbf{Q}}{2}\right) \gamma^0 \psi_q\left(\boldsymbol{\kappa} - \frac{\mathbf{Q}}{2}\right) = F_1(Q^2)w'^+w, \quad (F_1(0) = 1) \quad (5)$$

$$\int \frac{d^3\kappa}{(2\pi)^3} \bar{\psi}'_q\left(\boldsymbol{\kappa} + \frac{\mathbf{Q}}{2}\right) \boldsymbol{\gamma} \psi_q\left(\boldsymbol{\kappa} - \frac{\mathbf{Q}}{2}\right) = F_2(Q^2)w'^+ \frac{i[\mathbf{Q} \times \boldsymbol{\sigma}]}{2\langle\kappa^0\rangle} w \quad (6)$$

(the symmetric shift of momenta preserves the parity of the integrand with respect to the origin, simplifying the calculations). Since we neglect the recoil, F_1 , F_2 may be associated equally well with Dirac or with Sachs form-factors. That will be non-contradictory provided in the exact identities relating those types of form-factors

$$F_e = F_1 - \frac{Q^2}{4M_N^2}F_2, \quad F_m = F_2 + F_1 \quad (7)$$

the last terms are small. For the proton, $F_1(Q^2)/F_2(Q^2) \approx F_1(0)/F_2(0) = 1/2.79$, and within our accuracy $30 \div 40\%$ this contribution is indeed negligible. Also, the Foldy term for the proton's mean square charge radius $\frac{3F_2(0)}{2M_N^2} = (0.4\text{fm})^2$ is small compared to

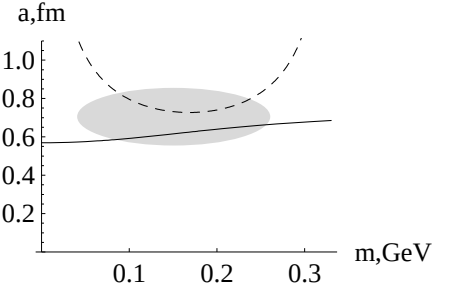


Figure 1: Curves defined by Eq. (8) (solid), Eq. (9) (dashed) and the region of their consistency.

²By the isospin symmetry, neutron's magnetic moment $\mu_n = -\frac{2}{3}\mu_p$, and its magnetic radius is about the same as proton's [1].

$\langle r_{\text{ch}}^2 \rangle \approx 1 \text{ fm}^2$. Now with Eqs. (5, 3), let us actually evaluate the mean square charge radius:

$$\begin{aligned}
\langle r_{\text{ch}}^2 \rangle &= -\frac{6}{F_1(0)} \frac{dF_1(\mathbf{Q}^2)}{d\mathbf{Q}^2} \Big|_{\mathbf{Q}^2 \rightarrow 0} \\
&= -\frac{6}{F_1(0)} \frac{\partial}{\partial \mathbf{Q}^2} \int \frac{d^3 \kappa}{(2\pi)^3} \left[1 + \frac{\kappa^2 - (\mathbf{Q}/2)^2}{4 \langle \kappa^0 \rangle^2} \right] \varphi \left(\kappa + \frac{\mathbf{Q}}{2} \right) \varphi \left(\kappa - \frac{\mathbf{Q}}{2} \right) \Big|_{\mathbf{Q} \rightarrow 0} \\
&= a^2 \left(\frac{3}{2} + \frac{1}{1 + \frac{8 \langle \kappa^0 \rangle^2 a^2}{3}} \right). \tag{8}
\end{aligned}$$

Another constraint on the model parameters comes from the proton magnetic moment, which in our additive quark model with charges $e_u = \frac{2}{3}e_p$, $e_d = -\frac{1}{3}e_p$ appears to equal

$$\frac{\mu_p}{e_p} = \frac{\mu_q}{e_q} = \frac{1}{2 \langle \kappa^0 \rangle} F_2(0) = \frac{1}{2 \langle \kappa^0 \rangle} \int \frac{d^3 \kappa}{(2\pi)^3} \varphi^2(\kappa) = \frac{1}{2 \langle \kappa^0 \rangle} \frac{1}{1 + \frac{3}{8 \langle \kappa^0 \rangle^2 a^2}}. \tag{9}$$

Together, constraints (8-9) may determine both free parameters of the wave function. With the experimental values $\mu_p/e_p \approx \frac{2.79}{2M_N}$, $\langle r_{\text{ch}}^2 \rangle \approx (0.9 \text{ fm})^2$ [1], Eqs. (8, 9) are, strictly speaking, mutually inconsistent (see Fig. 1). But within our accuracy $30 \div 40\%$ they are sufficiently consistent, yielding $a = 0.55 \div 0.85 \text{ fm}$ and $\langle \kappa^0 \rangle = 50 \div 250 \text{ MeV}$. Similar values for a were found by other authors [3].³ As for $\langle \kappa^0 \rangle$, it is about twice smaller than typical quark mass in non-relativistic constituent models [1].

Valence quark distribution function. The constructed model should also be able to describe valence quark distribution at Bjorken $x < 0.4$, i. e., including the region of the quark distribution maximum. The definition of a quark parton distribution (cf., e.g., [4]) leads to

$$\begin{aligned}
q_v(x) &= \int \frac{d^3 \kappa}{(2\pi)^3} \delta \left(x - \frac{\kappa^0 + \kappa^3}{M_N} \right) \bar{\psi}_q(\kappa) (\gamma^0 + \gamma^3) \psi_q(\kappa) \\
&= \frac{M_N}{2\pi} \int \frac{d^2 \kappa_{\perp}}{(2\pi)^2} \left[\left(1 - \frac{\kappa^0}{2 \langle \kappa^0 \rangle} + \frac{x M_N}{2 \langle \kappa^0 \rangle} \right)^2 + \frac{\kappa_{\perp}^2}{4 \langle \kappa^0 \rangle^2} \right] \varphi^2(\kappa) \Big|_{\kappa^3 = x M_N - \kappa^0} \\
&= \frac{M_N a}{\sqrt{\pi} (4 \langle \kappa^0 \rangle^2 a^2 + \frac{3}{2})} \left[(x M_N - \kappa^0 + 2 \langle \kappa^0 \rangle)^2 a^2 + 1 \right] e^{-(x M_N - \kappa^0)^2 a^2}. \tag{10}
\end{aligned}$$

$$\left(\int_{-\infty}^{\infty} q_v(x) dx = 1 \right)$$

Behavior of function (10) at different $\langle \kappa^0 \rangle$ is shown in Fig. 2. Values $\langle \kappa^0 \rangle > 150 \text{ MeV}$ look rather unrealistic. At $\langle \kappa^0 \rangle \simeq 50 \text{ MeV}$ our $x q_v(x)$ holds within the declared $30 \div 40\%$ accuracy with the MRST, CTEQ pdf fits. It seems interesting that the constituent quark “mass” $\langle \kappa^0 \rangle = 50 \div 100 \text{ MeV}$ gets already commensurable with the pion mass scale $m_{\pi} \approx 140 \text{ MeV}$.

³However, if instead of (4) one took a model $\varphi \propto \left(1 + \frac{a^2 \kappa^2}{2\alpha} \right)^{-\alpha}$ with $\alpha < 4$, the value for a might be larger.

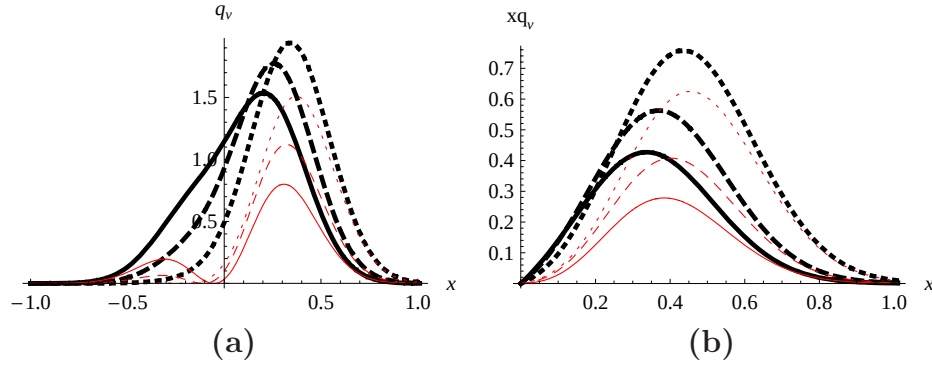


Figure 2: **(a)** Valence quark distribution for $a = 0.7$ fm and $m = 50$ MeV (solid line), $m = 100$ MeV (dashed), $m = 200$ MeV (dotted). In thin red the transverse spin asymmetry $\Delta_T q(x)$ is shown. **(b)** Same as (a) but for the valence quark distribution weighted with x . The momentum sums are, correspondingly, $3 \int_0^1 dx x q_v(x) = 0.5, 0.7, 1.0$.

The features of Fig. 2a are $\int_0^1 dx q_v(x) < 1$ and high value of $q_v(0)$.⁴ That is expectable: since the quark rest energy is smaller than its typical momentum $\sim (0.7 \text{ fm})^{-1} \simeq 300$ MeV, the wave function must reach well across the light cone. Note that $q_v(x)$ at negative values of x should *not* be associated with antiquarks [4], but still they may describe quarks, only ones inaccessible to leading twist DIS. In fact, for strongly bound relativistic states the positivity condition does not apply to light-cone components of momenta. Even considering the DIS handbag diagram in terms of 2-particle intermediate states with Bethe-Salpeter wave functions in nucleon vertices, the κ^- -integration will give a non-zero result for κ^+ beyond the interval $[0, P^+]$, because due the BS vertex functions $G(\kappa)$ depending on κ^+ through κ^3 one can not completely withdraw the integration contour to ∞ .

In conclusion, let us point out that the continuity of a valence quark wave function around $x = 0$ permits its use for perturbative description of hadron-hadron reactions with single quark exchange. Thereat, non-zero $q_v(0)$ opens a specific possibility of transverse polarization generation in reaction $np \rightarrow pn$ [6].

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⁴ $q_v(0)$ is finite also in bag models [5], whereas in light-cone models often $q_v(x) \xrightarrow{x \rightarrow 0} 0$. The latter property seems phenomenologically unlikely, even assuming further Regge enhancements.

TOWARDS A MODEL INDEPENDENT DETERMINATION OF FRAGMENTATION FUNCTIONS

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Abstract

We show that the difference cross sections in unpolarized SIDIS $e+N \rightarrow e+h+X$ and pp hadron production $p+p \rightarrow h+X$ determine, uniquely and in any order in QCD, the two FFs: $D_u^{h-\bar{h}}$ and $D_d^{h-\bar{h}}$, $h = \pi^\pm, K^\pm$. If both $K^\pm$ and K_s^0 are measured, then $e^+e^- \rightarrow K+X$, $e+N \rightarrow e+K+X$ and $p+p \rightarrow K+X$ yield independent measurements of $(D_u - D_d)^{K^++K^-}$. In a combined fit to $K^\pm$ and K_s^0 production data from e^+e^- collisions, $(D_u - D_d)^{K^++K^-}$ is obtained and compared to conventional parametrizations.

1 Introduction

Now, that the new generation of high energy scattering experiments with a final hadron h detected are taking place, it has become clear that in order to obtain the correct information about quark-lepton interactions, not only knowledge of the parton distribution functions (PDFs) are important, but a good knowledge of the fragmentation functions (FFs) D_i^h , that determine the transition of partons i into hadrons h , are equally important. The PDFs and the FFs are the two basic ingredients that have to be correctly extracted from experiment. Here we discuss the FFs.

The most direct way to determine the FFs is the total cross section for one-particle inclusive production in e^+e^- annihilation:

$$e^+e^- \rightarrow h+X, \quad h = \pi^\pm, K^\pm, p/\bar{p}... \quad (1)$$

However, these processes can determine, in principle, only the combinations

$$D_{u+\bar{u}}^h, \quad D_{d+\bar{d}}^h, \quad D_{s+\bar{s}}^h, \quad D_g^{h+\bar{h}} \quad (D_{q+\bar{q}}^h = D_q^{h+\bar{h}} \equiv D_q^h + D_q^{\bar{h}}), \quad (2)$$

i.e. they cannot distinguish the quark and anti-quark FFs. In order to achieve separate *phenomenological* determination of D_q^h and $D_q^{\bar{h}}$, the one-hadron semi-inclusive processes play an essential role:

$$l+N \rightarrow l+h+X \quad \text{and} \quad p+p \rightarrow h+X. \quad (3)$$

The factorization theorem implies that the FFs are universal, i.e. in (1) and (3) the FFs are the same. However, in (3) the hadron structure enters and when analyzing the data, usually different theoretical assumptions have to be made.

Schematically the cross sections for (1) and (3) can be written in the form:

$$d\sigma_{e^+e^-}^h(z) \simeq \sum_c \hat{\sigma}_{e^+e^-}^c \otimes D_c^{h+\bar{h}} \quad (4)$$

$$d\sigma_N^h(x, z) \simeq \sum_{q,c} f_q(x) \otimes \hat{\sigma}_{lq}^c \otimes D_c^h \quad (5)$$

$$d\sigma_{pp}^h(x, z) \simeq \sum_{a,b,c} f_a \otimes f_b \otimes \hat{\sigma}_{ab}^c \otimes D_c^h \quad (6)$$

where $\hat{\sigma}^c$ are the corresponding, perturbatively QCD calculable, parton-level cross sections for producing a parton c , f_a are the unpolarized PDFs. At present several sets of FFs exist and there is a significant disagreement among some of the FFs.

In this talk we present a model independent approach, developed in [1], which suggests that if instead of $d\sigma_N^h$ and $d\sigma_{pp}^h$ one works with the difference cross sections for producing hadrons and producing their antiparticles, i.e. with data on $d\sigma_N^{h-\bar{h}} \equiv d\sigma_N^h - d\sigma_N^{\bar{h}}$ or $d\sigma_{pp}^{h-\bar{h}} \equiv d\sigma_{pp}^h - d\sigma_{pp}^{\bar{h}}$, one obtains information about the non-singlet (NS) combinations $D_q^{h-\bar{h}}$. This is the complementary to $D_q^{h+\bar{h}}$ quantity, measured in (1), that would allow to determine D_q and $D_{\bar{q}}$ without assumptions.

Further this method is applied to charged and neutral kaon e^+e^- -production data to determine directly the non-singlet $(D_u - D_d)^{K^+K^-}$ and compare to conventional global fit analysis.

2 Difference cross sections with $\pi^\pm$ and $K^\pm$

From C-invariance of strong interactions it follows:

$$D_g^{h^+-h^-} = 0, \quad D_q^{h^+-h^-} = -D_{\bar{q}}^{h^+-h^-} \quad (7)$$

which, applied to (3), eliminates $D_g^{h^+-h^-}$ and $D_{\bar{q}}^{h^+-h^-}$ in the difference cross sections:

$$d\sigma_N^{h^+-h^-} = d\sigma_N^{h^+} - d\sigma_N^{h^-} \quad \text{and} \quad d\sigma_{pp}^{h^+-h^-} = d\sigma_{pp}^{h^+} - d\sigma_{pp}^{h^-} \quad (8)$$

This implies that, in any order of QCD, $d\sigma_N^{h^+-h^-}$ and $d\sigma_{pp}^{h^+-h^-}$ are expressed solely in terms of the NS combinations of the FFs. In NLO we have:

$$d\sigma_p^{h^+-h^-}(x, z, Q^2) = \frac{1}{9} [4u_V \otimes D_u + d_V \otimes D_d + s_V \otimes D_s]^{h^+-h^-} \otimes (1 + \frac{\alpha_s}{2\pi} C_{qq}) \quad (9)$$

$$d\sigma_d^{h^+-h^-}(x, z, Q^2) = \frac{1}{9} [(u_V + d_V) \otimes (4D_u + D_d) + 2s_V \otimes D_s]^{h^+-h^-} \otimes (1 + \frac{\alpha_s}{2\pi} C_{qq}) \quad (10)$$

$$\begin{aligned} E^h \frac{d\sigma_{pp}^{h^+-h^-}}{d^3P^h} &= \frac{1}{\pi} \int dx_a \int dx_b \int \frac{dz}{z} \times \\ &\times \sum_{q=u,d,s} [L_q(x_b, t, u) q_V(x_a) + L_q(x_a, u, t) q_V(x_b)] D_q^{h^+-h^-}(z) \end{aligned} \quad (11)$$

where u_V and d_V are the valence quarks PDFs, $s_V = s - \bar{s}$ and L_q , given explicitly in [1], are functions of the *known* quark densities $q + \bar{q}$ and partonic cross sections.

Common for the difference cross sections (9) - (11) is that they all have the same structure: 1) only the non-singlets $D_q^{h-\bar{h}}$ enter and 2) they enter multiplied by $q_V = q - \bar{q}$, i.e. in the combination $q_V D_q^{h-\bar{h}}$. This implies that the contributions of D_u^h and D_d^h are enhanced by the large valence quark densities, while D_s^h is suppressed by the small factor $(s - \bar{s})$. Recently a strong bound on $(s - \bar{s})$ was obtained from neutrino experiments – $|s - \bar{s}| \leq 0.025$ [2], which implies that the contribution from D_s^h can be safely neglected. Thus, the ep , ed and pp semi-inclusive difference cross sections provide three independent measurements for the two unknown FFs $D_u^{h^+-h^-}$ and $D_d^{h^+-h^-}$. Note that the SIDIS cross sections involve only u_V and d_V , which are the best known parton densities, with 2%-3% accuracy at $x \lesssim 0.7$.

Further information can be obtained specifying the final hadrons.

1) If $h = \pi^\pm$ the difference cross sections will determine, without any assumptions, $D_u^{\pi^+-\pi^-}$ and $D_d^{\pi^+-\pi^-}$ which would allow to test the usually made assumption

$$D_u^{\pi^+-\pi^-} = -D_d^{\pi^+-\pi^-}. \quad (12)$$

In [3] it was suggested, for the first time, that this relation might be violated up to 10 %.

2) If $h = K^\pm$ the difference cross sections will determine, without any assumptions, $D_u^{K^+-K^-}$ and $D_d^{K^+-K^-}$ which would allow to test the usually made assumption

$$D_d^{K^+-K^-} = 0. \quad (13)$$

One can formulate the above results also like this: If relation $D_u^{\pi^+-\pi^-} = -D_d^{\pi^+-\pi^-}$ (or $D_d^{K^+-K^-} = 0$) holds, then Eqs. (9), (10) and (11) for $h = \pi^\pm$ (or $h = K^\pm$) are expressed solely in terms of $D_u^{\pi^+-\pi^-}$ (or $D_u^{K^+-K^-}$) and thus look particularly simple.

3 Difference cross sections with $K^\pm$ and K_s^0

If in addition to the charged $K^\pm$ also neutral kaons $K_s^0 = (K^0 + \bar{K}^0)/\sqrt{2}$ are measured, no new FFs are introduced into the cross-sections. We show that the combination

$$\sigma^K \equiv \sigma^{K^+} + \sigma^{K^-} - 2\sigma^{K_s^0} \quad (14)$$

in the considered three types of semi-inclusive processes (1) and (3):

$$e^+ + e^- \rightarrow K + X, \quad (15)$$

$$e + N \rightarrow e + K + X, \quad (16)$$

$$p + p \rightarrow K + X, \quad (17)$$

$K = K^\pm, K_s^0$, always measures only the NS combination $(D_u - D_d)^{K^++K^-}$. This result relies only on SU(2) invariance for the kaons, that relates $D_q^{K_s^0}$ to $D_q^{K^++K^-}$ and does not involve *any* assumptions about PDFs or FFs, it holds in any order in QCD. As $(D_u - D_d)^{K^++K^-}$, obtained in this way, is model independent it would be interesting to compare it to the existing parametrizations from e^+e^- data, obtained using various assumptions.

4 $(D_u - D_d)^{K^++K^-}$ from e^+e^- kaon production

Further we focus on the most precisely measured and theoretically calculated process (15). In NLO we have [1]:

$$d\sigma_{e^+e^-}^{\mathcal{K}}(z, Q^2) = \frac{8\pi\alpha_{em}^2}{s} (\hat{e}_u^2 - \hat{e}_d^2) \left(1 + \frac{\alpha_s}{2\pi} C_q \otimes\right) (D_u - D_d)^{K^++K^-}(z, Q^2). \quad (18)$$

where $\hat{e}_q^2(s)$ are the quark electro-weak charges. Using (18) we determine [6] $(D_u - D_d)^{K^++K^-}$ from the available data on $K^\pm$ and K_s^0 production in $e^+e^- \rightarrow (\gamma, Z) \rightarrow K + X$, $K = K^\pm, K_s^0$ and compare it to those obtained in global fit analysis.

Our analysis has several advantages: it allows for the first time to extract $(D_u - D_d)^{K^++K^-}$ without any assumptions about the FFs (commonly used in global fit analysis) and without any correlations to other FFs (and especially to $D_g^{K^\pm}$), it allows to use data at much lower values of z than in global fit analysis (being a NS the combination $\sigma_{e^+e^-}^{\mathcal{K}}$ does not contain the unresummed soft gluon logarithms), etc.

In the analysis we include the data on $K^\pm$ and K_s^0 production from TASSO, HRS, MARKII, TPC, TOPAZ, ALEPH, DELPHI, OPAL, SLD and CELLO collaborations, in the energy range intervals $\sqrt{s} = 12 - 14.8, 21.5 - 22, 29 - 35, 42.6 - 44, 58, 91.2$ and $183 - 186$ GeV.

In Fig.1 we present $(D_u - D_d)^{K^++K^-}$ in NLO obtained in our approach and from the global fits of the DSS [3], HKNS [4] and AKK08 [5] sets. As seen from the Figure, at $z \gtrsim 0.4$ there is an agreement among the different FFs in shape, but our FF is in general larger. The difference becomes more pronounced at $z \lesssim 0.4$. There could be different reasons for this. Most probably it is due either to inclusion of the small z -data in our fit or to the different assumptions in the global fit parametrizations.

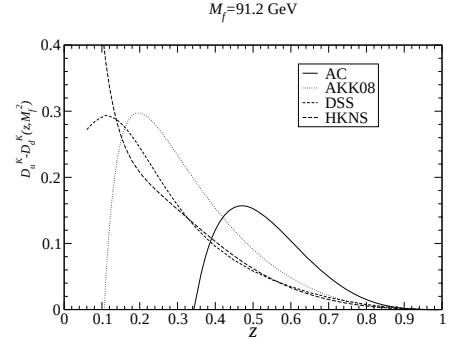


Figure 1: $D_{u-d}^{K^++K^-}$ as obtained in this paper (labeled AC) and as obtained from the parametrizations of DSS, HKNS and AKK08.

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TRANSVERSITY GPDs FROM $\gamma N \rightarrow \pi \rho_T N'$ WITH A LARGE $\pi \rho_T$ INVARIANT MASS

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Abstract

The chiral-odd transversity generalized parton distributions of the nucleon can be accessed experimentally through the exclusive photoproduction process $\gamma + N \rightarrow \pi + \rho_T + N'$, with or without beam and target polarization, provided the vector meson is produced in a transversely polarized state. The kinematical domain of factorization is defined through a large invariant mass of the meson pair and a small transverse momentum of the final nucleon. We calculate perturbatively the scattering amplitude at leading order in α_s and build a simple model for the dominant transversity GPD $H_T(x, \xi, t)$ based on the concept of double distribution. Counting rate estimates are in progress.

Transversity quark distributions in the nucleon remain among the most unknown leading twist hadronic observables. This is mostly due to their chiral odd character which enforces their decoupling in most hard amplitudes. After the pioneering studies [1], much work [2] has been devoted to the exploration of many channels but experimental difficulties have challenged the most promising ones.

On the other hand, tremendous progress has been recently witnessed on the QCD description of hard exclusive processes, in terms of generalized parton distributions (GPDs) describing the 3-dimensional content of hadrons. Access to the chiral-odd transversity GPDs [3], noted H_T , E_T , $\tilde{H}_T$, $\tilde{E}_T$, has however turned out to be even more challenging [4] than the usual transversity distributions: one photon or one meson electroproduction leading twist amplitudes are insensitive to transversity GPDs. The strategy which we follow here, as initiated in Ref. [5], is to study the leading twist contribution to processes where more mesons are present in the final state; the hard scale which allows to probe the short distance structure of the nucleon is $s = M_{\pi\rho}^2 \sim |t'|$ in the fixed angle regime. In the example developed previously [5], the process under study was the high energy photo (or electro) diffractive production of two vector mesons, the hard probe being the virtual "Pomeron" exchange (and the hard scale was the virtuality of this pomeron), in analogy with the virtual photon exchange occurring in the deep electroproduction of a meson. A similar strategy has also been advocated recently in Ref. [6] to enlarge the number of processes which could be used to extract information on chiral-even GPDs.

The process we study here

$$\gamma + N \rightarrow \pi + \rho_T + N', \quad (1)$$

is a priori sensitive to chiral-odd GPDs because of the chiral-odd character of the leading twist distribution amplitude (DA) of the transversely polarized ρ meson. The estimated rate depends of course much on the magnitude of the chiral-odd GPDs. Not much is known about them, but model calculations have been developed in [5, 7–9]; moreover, a few moments have been computed on the lattice [10].

To factorize the amplitude of this process we use the now classical proof of the factorization of exclusive scattering at fixed angle and large energy [11]. The amplitude for the process $\gamma + \pi \rightarrow \pi + \rho$ is written as the convolution of mesonic DAs and a hard scattering subprocess amplitude $\gamma + (q + \bar{q}) \rightarrow (q + \bar{q}) + (q + \bar{q})$ with the meson states replaced by collinear quark-antiquark pairs. This is described in Fig. 1a. The absence of any pinch singularities (which is the weak point of the proof for the generic case $A + B \rightarrow C + D$) has been proven in Ref. [12] for the case of interest here. We then extract from the factorization procedure of the deeply virtual Compton scattering amplitude near the forward region the right to replace in Fig. 1a the lower left meson DA by a $N \rightarrow N'$ GPD, and thus get Fig. 1b. We introduce ξ as the skewness parameter which can be written in terms of the meson pair squared invariant mass $M_{\pi\rho}^2$ as

$$\xi = \frac{\tau}{2 - \tau} \quad , \quad \tau = \frac{M_{\pi\rho}^2}{S_{\gamma N} - M^2} . \quad (2)$$

Indeed the same collinear factorization property bases the validity of the leading twist approximation which either replaces the meson wave function by its DA or the $N \rightarrow N'$ transition to its GPDs. A slight difference is that light cone fractions $(z, 1 - z)$ leaving the DA are positive, while the corresponding fractions $(x + \xi, \xi - x)$ may be positive or negative in the case of the GPD. The calculation will show that this difference does not ruin the factorization property, at least at the Born order we are studying here.

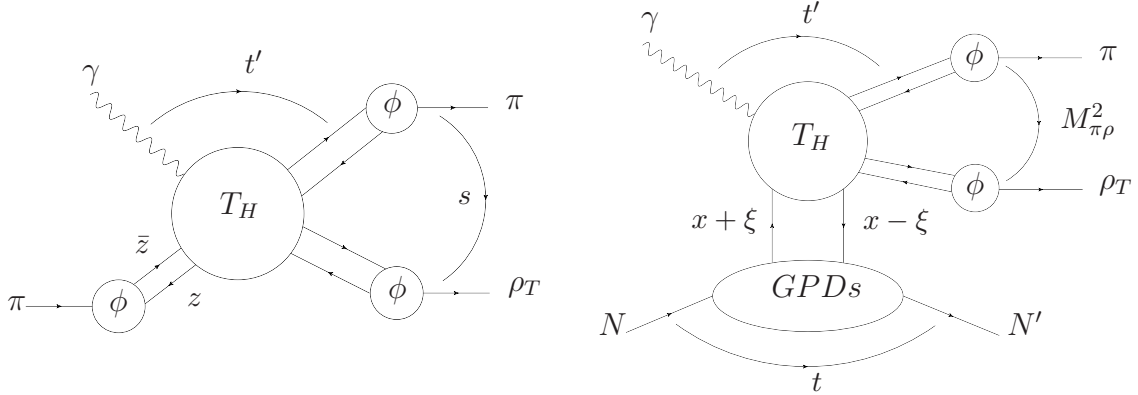


Figure 1: a) (left) Factorization of the amplitude for the process $\gamma + \pi \rightarrow \pi + \rho$ at large s and fixed angle (i.e. fixed ratio t'/s); b) (right) replacing one DA by a GPD leads to the factorization of the amplitude for $\gamma + N \rightarrow \pi + \rho + N'$ at large $M_{\pi\rho}^2$.

In order for the factorization of a partonic amplitude to be valid, and the leading twist calculation to be sufficient, one should avoid the dangerous kinematical regions where a small momentum transfer is exchanged in the upper blob, namely small $t' = (p_\pi - p_\gamma)^2$ or small $u' = (p_\rho - p_\gamma)^2$, and the regions where strong interactions between two hadrons in the final state are non-perturbative, namely where one of the invariant squared masses $(p_\pi + p_{N'})^2, (p_\rho + p_{N'})^2, (p_\pi + p_\rho)^2$ is in the resonance region.

The scattering amplitude of the process (1) is written in the factorized form :

$$\mathcal{A}(t, M_{\pi\rho}^2, p_T) = \int_{-1}^1 dx \int_0^1 dv \int_0^1 dz T^q(x, v, z) H_T^q(x, \xi, t) \Phi_\pi(z) \Phi_\perp(v), \quad (3)$$

where T^q is the hard part of the amplitude and the transversity GPD of a parton q in the nucleon target which dominates at small momentum transfer is defined by [3]

$$\langle N'(p_2), \lambda' | \bar{q} \left(-\frac{y}{2} \right) \sigma^{+j} \gamma^5 q \left(\frac{y}{2} \right) | N(p_1), \lambda \rangle = \bar{u}(p', \lambda') \sigma^{+j} \gamma^5 u(p, \lambda) \int_{-1}^1 dx e^{-\frac{i}{2} x (p_1^+ + p_2^+) y^-} H_T^q,$$

where λ and λ' are the light-cone helicities of the nucleon N and N' . The chiral-odd DA for the transversely polarized meson vector ρ_T , is defined, in leading twist 2, by the matrix element [14]

$$\langle 0 | \bar{u}(0) \sigma^{\mu\nu} u(x) | \rho_T^0(p, \epsilon_\pm) \rangle = \frac{i}{\sqrt{2}} (\epsilon_\pm^\mu(p) p^\nu - \epsilon_\pm^\nu(p) p^\mu) f_\rho^\perp \int_0^1 du e^{-i u p \cdot x} \phi_\perp(u), Dv$$

where $\epsilon_\pm^\mu(p_\rho)$ is the ρ -meson transverse polarization and with $f_\rho^\perp = 160$ MeV.

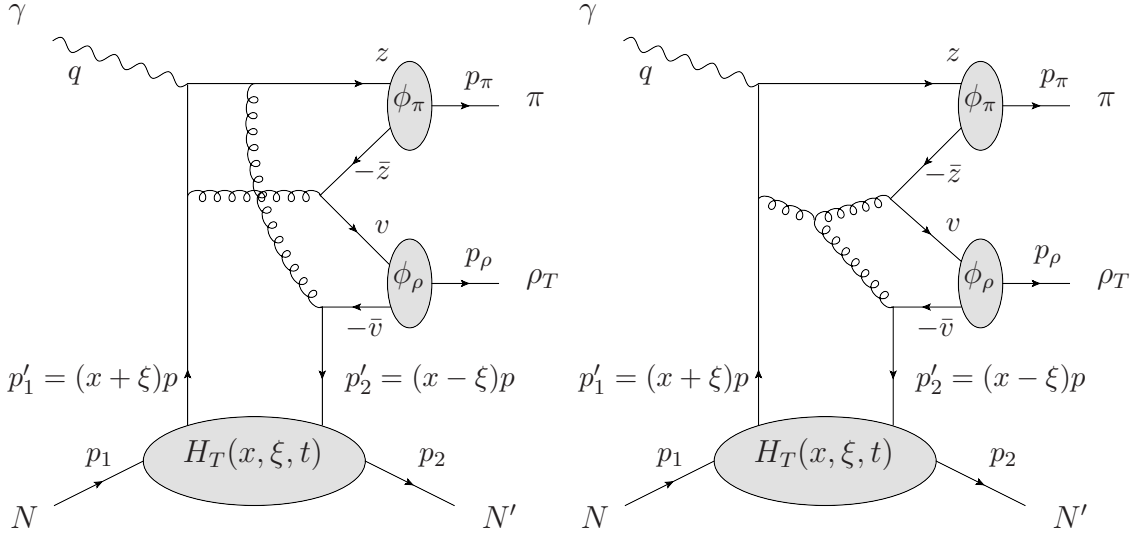


Figure 2: Two representative diagrams without (left) and with (right) three gluon coupling.

Two classes of Feynman diagrams (see Fig.2), without and with a 3-gluon vertex, describe this process. In both cases, an interesting symmetry allows to deduce the contribution of some diagrams from other ones, reducing our task to the calculation of half the 62 diagrams involved in the process. The scattering amplitude gets both a real and an imaginary parts. Integrations over v and z have been done analytically whereas numerical methods are used for the integration over x . Various observables can be calculated with this amplitude. We stress that even the unpolarized differential cross-section $\frac{d\sigma}{dt du' dM_{\pi\rho}^2}$ is sensitive to the transversity GPD. Rate estimates are under way, based on a double distribution model for H_T . Preliminary results allow us to provisionally conclude that the photoproduction of a transversely polarized vector meson on a nucleon target is a

good way to reach information on the generalized chiral-odd quark content of the proton. For instance, if the JLab 12 GeV upgrade gets the nominal luminosity ($\mathcal{L} \sim 10^{35} \text{ cm}^2 \cdot \text{s}^{-1}$), we expect a sufficient number of events per year to extract the dominant transversity GPD.

We are grateful to Igor Anikin, Markus Diehl, Samuel Friot and Jean Philippe Lansberg for useful discussions. This work is partly supported by the French-Polish scientific agreement Polonium 7294/R08/R09, the ECO-NET program, contract 18853PJ, the ANR-06-JCJC-0084, the Polish Grant N202 249235 and the DFG (KI-623/4).

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Abstract

We present analysis of the infrared dependence of the perturbative contributions to the spin structure function g_1 .

1 Introduction

Factorization [1] of the long and short -distance strong interactions is introduced to guarantee applicability of Perturbative QCD for calculating the DIS structure functions. According to this concept, the perturbative calculations involve dealing with virtual partons (quarks and gluons) of the virtualities $\gtrsim \mu^2$ while contributions of the partons with small virtualities $\lesssim \mu^2$ are collected into the initial parton densities. Therefore, parameter μ plays the role of a border between the perturbative and non-perturbative QCD. Being artificial parameter, μ should vanish when both perturbative and non-perturbative QCD contributions are taken into account, i.e. in convolutions of the perturbative expressions for g_1 with initial parton distributions. For example, the non-singlet component of g_1 can be written as

$$g_1^{NS} = g_1^{NS \text{ } pert} \otimes \delta q. \quad (1)$$

Naturally, the perturbative part, $g_1^{NS \text{ } pert}$ depends on μ . This dependence is different in the Standard Approach and beyond it.

2 The μ -dependence of g_1 in Standard Approach

Standard Approach describes g_1 at large x . It is based on the DGLAP evolution equations [2]. It is believed that the μ dependence in this approach appears in the first-loop integration over $k_\perp$ while calculations in higher loops do not depend on μ because of the well-known DGLAP- ordering

$$\mu^2 < k_{1\perp}^2 < k_{2\perp}^2 < \dots < Q^2. \quad (2)$$

However, it was shown [3] there is an additional μ -dependence in this approach coming from the parametrization of α_s . Indeed, α_s in the DGLAP -equations should be replaced by α_s^{eff} given by the following expression:

$$\begin{aligned}
\alpha_s^{eff} &= \alpha_s(\mu^2) + \frac{1}{\pi b} \left[\arctan \left(\frac{\pi}{\ln(k_{\perp}^2/\beta\Lambda^2)} \right) - \arctan \left(\frac{\pi}{\ln(\mu^2/\Lambda^2)} \right) \right] \\
&= \alpha_s(\mu^2) + \frac{1}{\pi b} \left[\arctan \left(\pi b \alpha_s(k_{\perp}^2/\beta) \right) - \arctan \left(\pi b \alpha_s(\mu^2) \right) \right]
\end{aligned} \tag{3}$$

where β is the fraction of the longitudinal moment. When

$$\mu^2 \gg \Lambda^2 e^\pi \approx 23\Lambda^2, \tag{4}$$

α_s^{eff} can be approximated by the much simpler expression:

$$\alpha_s^{eff} \approx \alpha_s(k_{\perp}^2/\beta). \tag{5}$$

If additionally x is large, $\alpha_s^{eff} \approx \alpha_s(k_{\perp}^2)$. Therefore, the QCD coupling in the Standard Approach is always μ -dependent, though implicitly, through Eq. (4).

3 The μ -dependence of g_1 at small x

The μ -dependence of g_1 at small x is even more involved because accounting for contributions $\sim \ln^n(1/x)$ requires to change the DGLAP-ordering by the new one:

$$\mu^2 < k_{1\perp}^2/\beta_1 < k_{2\perp}^2/\beta_2 < \dots < 2pq \tag{6}$$

where β_j are the fractions of the longitudinal parton momenta. In this case the QCD coupling is explicitly μ -dependent as shown in Eq. (3) and besides the ordering (6) demonstrates that, in contrast to the DGLAP, every order of the perturbative QCD brings new μ -dependent terms. Indeed, the first-loop integration at small x looks quite similarly to DGLAP:

$$\int_{\mu^2}^{2pq} \frac{dk_{1\perp}^2}{k_{1\perp}^2} \tag{7}$$

while the second-loop contribution contains the DGLAP-like term

$$\int_{k_{1\perp}^2}^{2pq} \frac{dk_{2\perp}^2}{k_{2\perp}^2} \int_{k_{2\perp}^2}^{\beta_1} \frac{d\beta_2}{\beta_2} \tag{8}$$

and the new μ -dependent term

$$\int_{\mu^2}^{k_{1\perp}^2} \frac{dk_{2\perp}^2}{k_{2\perp}^2} \int_{k_{2\perp}^2}^{\beta_1 k_{2\perp}^2/k_{1\perp}^2} \frac{d\beta_2}{\beta_2}. \tag{9}$$

4 Conclusion

We have demonstrated explicitly that there are two sources for the infrared dependence of the perturbative contributions to the spin structure function g_1 : (a) Explicit or implicit μ dependence of the QCD coupling; (b) Explicit μ -dependence of the integrations over transverse momenta of virtual partons in the first loop and higher loops.

5 Acknowledgement

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ASYMMETRIES AT THE ILC ENERGIES AND B-L GAUGE MODELS

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Abstract

The use of polarized beams of electrons/positrons is one of the main characteristics that can make the proposed ILC (International Linear Collider) to be able to bring results for physics beyond the Standard Model and for unraveling the structure of the underlying physics. In this work we study extra neutral gauge bosons $Z_{B-L} (\equiv Z')$ deriving from two extensions of the standard model with an extra gauge $U(1)_{B-L}$ factor. Among several Z' proposed experiments, we study the capability of e^+e^- colliders to give responses to the Z' existence using both, electron and positron polarized beams. We emphasize $e^+e^- \rightarrow \mu^+\mu^-$ scattering through the demonstration of graphs related to them. We analyzed left-right and polarization asymmetries for the two models.

Most of the well motivated extensions of the standard model (SM) are those which contain, at least, one extra neutral vector boson usually denoted by Z' . For instance, left-right models [1], and grand unified theories based on $SO(10)$ [2] or E_6 [3]. For recent review of extra neutral vector bosons see Ref. [4]. If Z and Z' denote the symmetry eigenstates, then Z_1 and Z_2 are the mass eigenstates.

In particular we shall consider two extensions of the electroweak standard model in which there is an extra $U(1)$ local symmetry and their charges are $B - L$. Both models are based on the gauge symmetry and its breakdown:

$$SU(2)_L \otimes U(1)_\xi \otimes U(1)_{B-L} \rightarrow SU(2)_L \otimes U(1)_Y \rightarrow U(1)_{em}. \quad (1)$$

In the first model $\xi \equiv Y'$, and it is chosen to obtain the hypercharge Y of the standard model, given by $Y = Y' + (B - L)$. Thus, in this case, the charge operator is given by

$$\frac{Q}{e} = I_3 + \frac{1}{2} [Y' + (B - L)]. \quad (2)$$

We shall call this model “flipped” model because the electric charge is partly in $U(1)_{B-L}$ [5]. There are at least two models of this sort depending on the lepton number (L) attributed to the right-handed neutrinos, see Ref. [5]. We consider here the model with 3 right-handed neutrinos carrying $L = 1$.

The other electroweak model has $\xi \equiv Y$ and, the charge operator is as in the SM, i.e., given by [6],

$$\frac{Q}{e} = I_3 + \frac{1}{2} Y. \quad (3)$$

We shall call this model the “secluded” model because the electric charge has no component in $U(1)_z$, this happens even if $z \neq B - L$. It has also three right-handed neutrinos with $L = 1$.

Both models have the same scalar sector: one doublet, H , with $Y = +1$; and one complex singlet, ϕ , with $Y = 0$. However, in general, depending on the z value, the doublet in the secluded model has to carry $U(1)_z$ charge. Only in this case there is mixing between Z and Z' in the mass square matrix. The difference between (2) and (3) distinguishes both models. Only when $z_H = 0$ the factor $U(1)_z$ corresponds to $U(1)_{B-L}$ and this is the case to be considered here. For the secluded model see also Ref. [7]. In most of the paper we compare both models assuming that the breaking of the $B - L$ symmetry occurs at an energy above the TeV scale.

These models have two massive neutral vector bosons that we will denote Z_1 and Z_2 and we will parameterize their neutral currents as follows:

$$\mathcal{L}^{NC} = -\frac{g}{2c_W} \sum_i \bar{\psi}_i \gamma^\mu [(g_V^i - g_A^i \gamma_5) Z_{1\mu} + (f_V^i - f_A^i \gamma_5) Z_{2\mu}] \psi_i, \quad (4)$$

with $Z_1 \approx Z$ and $Z_2 \approx Z'$.

At ILC, the phenomenological constraints on extra neutral gauge bosons can be investigated with the use of polarization. These bosons can be detected by measuring some observables, as Z' decay partial widths, and Z' mass and several kinds of asymmetries and some cross sections as well.

In the recent past, the precision of electroweak measurements was achieved at electron-positron colliders SLC and LEP. Although the SLC data statistics were smaller than the LEP one, the presence of longitudinal polarization allowed complementary and competitive measurements of Z couplings. The ILC is been planed to have both electron and positron beams polarized and we shall see that the polarization is a useful tool to distinguish different kinds of models. Here we present and analyze some polarization asymmetries for the two models discussed above. We study the decay channel $e^+e^- \rightarrow \mu^+\mu^-$. Bhabha and Møller scattering can also be used as additional observables in e^+e^- collisions but these kind of processes are sensitive only to the Z' couplings to electron. Møller scattering has an advantage of profiting from two highly polarized electron beams.

So we work with left-right asymmetry, polarization asymmetry and a mix asymmetry that combines the forward-backward and left-right asymmetries. The definition of these asymmetries are as follows: $A_{FB} = 3N_{FB}/4\sigma_T$, $A_{LR} = 3N_{LR}/4\sigma_T$, and $A_{LR,FB} = 3N_{LR,FB}/4\sigma_T$, where $N_{FB} = \sigma_{LL} - \sigma_{LR} + \sigma_{RR} - \sigma_{RL}$, $N_{LR} = \sigma_{LL} + \sigma_{LR} - \sigma_{RR} - \sigma_{RL}$, and $N_{LR,FB} = \sigma_{LL} - \sigma_{RR} + \sigma_{RL} - \sigma_{LR}$, and $\sigma_T = \sigma_{LL} + \sigma_{LR} + \sigma_{RR} + \sigma_{RL}$. We are not considering the effects of transverse polarization in this situation because they suffer from experimental difficulties.

For the case of $e^+e^- \rightarrow \mu^+\mu^-$ at high energies, the polarization asymmetry is just the left-right asymmetry with the opposite sign. The left-right and polarization asymmetries are very sensitive to the weak mixing angle and the combined left-right and forward-backward asymmetry is a very important statistical tests. At SLC, this kind of asymmetry had given an statistical precision equivalent to the measurements of unpolarized forward-backward asymmetry at LEP [8]. In Fig. 1(a) we show the cross section at the Z' peak as a function of center of mass energy, s ; and in Fig. 1(b), we show the forward-backward asymmetry also in terms of s . Notice that the cross section is more sensitive to the presence of Z' at the peak, but the asymmetry can be relevant even out of the pole.

In the studied models, we set a Z' with a mass of 1 TeV. The decay width for the flipped and secluded model are respectively 17.89 and 19.54 GeV. We consider ideal situations, that is, the values of polarization for electron and positron are respectively $P_{e^-} = -1$ and $P_{e^+} = +1$. As we can see from figures 2, these asymmetries show that the flipped and secluded model can be distinguished successfully at ILC energies. With the choices of the VEVs, u and v the VEVs of ϕ and H , respectively, and also g_{B-L} which implies the same Z' mass in both models, the asymmetries are larger at the peak in the flipped model than in the secluded model. This is because the Z' in the latter model has identical vectorial couplings and the axial couplings vanish when $z_H = 0$.

In fact, the secluded model considered as a $B-L$ gauge model the $g_{V,A}^i$ couplings in Eq. (4) are exactly the same of the standard model, and $f_A^i = 0$ and the f_V^i 's are given

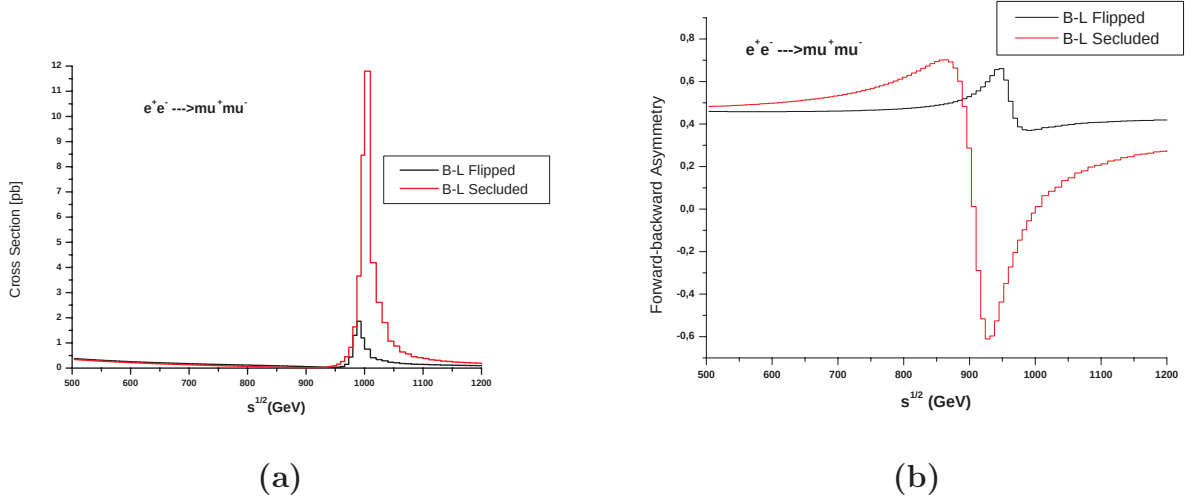


Figure 1: (a) Cross section for $e^+e^- \rightarrow \mu^+\mu^-$ in function of s . (b) Forward-backward asymmetry for the same process.

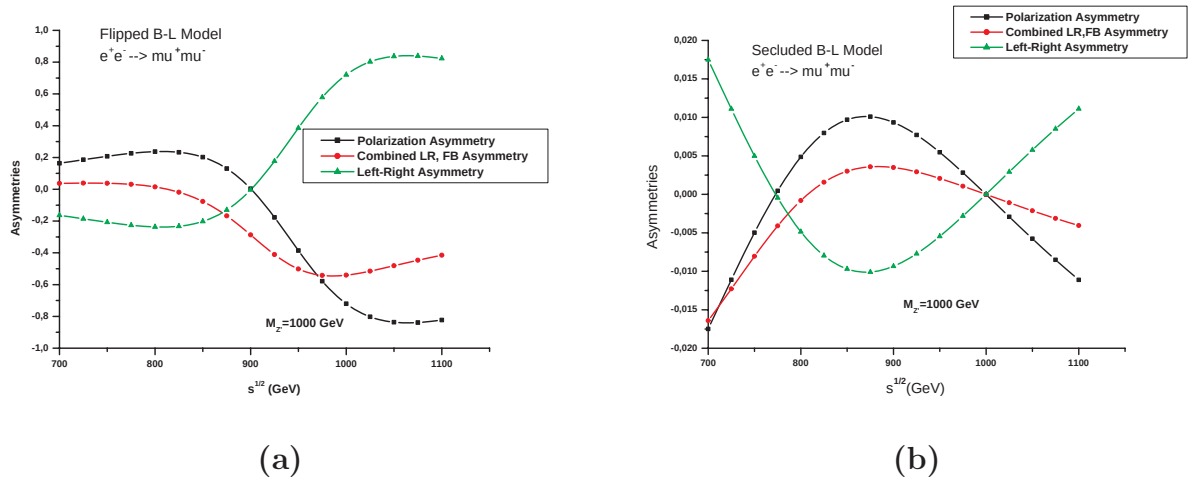


Figure 2: (a) Polarization asymmetries in the flipped model. (b) The same as in (a) for the secluded model.

by $f_V^\nu = f_V^l = f_V^u = f_V^d = -\frac{g_z}{g_Y} s_W$, where $g_z \equiv g_{B-L}$ and g_Y is the standard model $U(1)_Y$ charge. We see that Z' couples universally with all fermions. For the flipped model these couplings are more complicated and we shall show them elsewhere [10]. We recall that the difference among models with additional $U(1)$'s groups not inspired in unified theories and those we are considering here is that the neutral current parameters in the latter case must satisfy some relations [9] that do not exist in the former.

Similar analysis has been done for the Z' bosons of the 3-3-1 models [11] and we shall show elsewhere [10].

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ON THE $q\bar{q}$ -GLUEBALL-MIXED 0^{++} -MESON STATES IN A SIMPLE MODEL APPROACH

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Abstract

The next to lowest mass scalar multiplet is treated as the $q\bar{q}$, P -wave, non-radial-excited nonet, mixed with the $J^{PC} = 0^{++}$ -glueball. The Gell-Mann-Okubo and Schwinger-type mass-formulas are used to derive and discuss the quark-gluon configuration structure of the obtained meson states.

Contents.

1. Underlying reasons
2. Mass-relations - from octets via nonets toward decuplets.
3. The isoscalar "sub-multiplet" in scalar sector.
4. Concluding remarks.

1. We consider the mass region $1.3 \div 1.8$ GeV occupied by the $J^{PC} = 0^{++}$ mesons. In a selected group of the positive parity mesons there are three isoscalar mesons with reasonably close masses which, in the presence of the nearly lying iso-triplet and iso-doublet ones, suggest an overpopulated nonet where a possible glueball is hidden within structures of the three isoscalar state vectors.

In a standard way, we define the 3×3 mass-matrix $\hat{V}(i)$ as acting on the basis vectors N, S, G to transform them into one of three vectors of the physical scalar meson states, $f_0(i)$,

$$\begin{pmatrix} f_0(i) \end{pmatrix} = \hat{V}(i) \cdot \begin{pmatrix} N \\ S \\ G \end{pmatrix} \quad (1)$$

where

$$\begin{aligned} N(8 \cup 1|_{\text{mix}}; \mathbf{I} = \mathbf{Y} = 0) &= \frac{1}{\sqrt{3}} |8; I = Y = 0 \rangle + \sqrt{\frac{2}{3}} |1; I = Y = 0 \rangle = \frac{1}{\sqrt{2}} (u\bar{u} + d\bar{d}), \\ S(8 \cup 1|_{\text{mix}'; I=Y=0}) &= \frac{1}{\sqrt{3}} |1; I = Y = 0 \rangle - \sqrt{\frac{2}{3}} |8; I = Y = 0 \rangle = (s\bar{s}). \end{aligned} \quad (2)$$

and G is the glueball. We consider the mass-matrices $\hat{V}(i)$ taking into account explicitly the different appearance of the two types of gluon effects in mixing states of the differing flavor. In a certain sense, we follow the way proposed in old works by Isgur [1] to connect the strong "non-ideality" of the $SU(3)$ -singlet-octet mixing angle in the lowest pseudoscalar and scalar meson nonet with the overwhelmingly strong, as compared to the respective term in the vector or tensor meson nonets, annihilation term in the mass-matrix inducing the nondiagonal $q\bar{q} \leftrightarrow s\bar{s}$ transitions.

2. We remind, that the celebrated Gell-Mann–Okubo (GMO) [2] formula, written for the vector meson octet

$$3m^2(\mathbf{8}; I = Y = 0) = 4m^2(\mathbf{8}; I = (1/2), Y = \pm 1) - m^2(\mathbf{8}; I = 1, Y = 0)$$

follows as the mass sum rule after exclusion of parameters M^2 and μ^2 introduced into the general mass term of the phenomenological meson Lagrangian

$$M^2 \cdot \text{Tr}(V_8 V_8) - \mu^2 \cdot \text{Tr}(V_8 V_8 \lambda_8).$$

Okubo [3] proposed replacing $V_8 \rightarrow V_9$ in the GMO mass operator and dropping the term proportional to $\text{Tr}(V_9)$. The well-known "ideal mixing" mass relations

$$m^2(\mathbf{8}; I = 1, Y = 0) = m^2(\mathbf{8} \cup \mathbf{1}|_{\text{mix}}; I = Y = 0)$$

$$, \quad 2m^2(\mathbf{8}; I = (1/2), Y = \pm 1) - m^2(\mathbf{8}; I = 1, Y = 0) = m^2(\mathbf{8} \cup \mathbf{1}|_{\text{mix}}; I = Y = 0)$$

are fulfilled for the vector and reasonably well for the tensor nonet, but poorly for the pseudoscalar one. We indicate also the hierarchy of meson masses following from the effective Lagrangian GMO

$$m^2(S) > m^2(q\bar{s}) > m^2(N).$$

The idea to relate the apparently specific situation for the pseudoscalar meson sector with an additional strong annihilation mechanism transforming the quark field combinations into each other was put forward phenomenologically by Isgur [1] and interpreted now as mediated by short-range fluctuations in the quark-gluon vacuum [4]. We follow this idea in a further generalized form via introducing the "bare" scalar glueball mass and nondiagonal glueball-quarkonium transition-mass into the spin-zero meson mass-matrices. Hence, in the $N = \frac{1}{\sqrt{2}}(u\bar{u} + d\bar{d}), S = s\bar{s}$ basis our symmetric mass-matrix acquires the following form:

$$\hat{M}^2 = \begin{pmatrix} M_N^2 + A_Q & \sqrt{2}A_G & \sqrt{2}rA_Q \\ \sqrt{2}A_G & M_G^2 & rA_G \\ \sqrt{2}rA_Q & rA_G & M_S^2 + r^2A_Q \end{pmatrix} \quad (3)$$

The mass-matrix (3) with additional factors r introduced in different models of the $SU(3)$ -symmetry violations can be taken as the ingredient of the generalized Schwinger-type [5] mass formulas for the hadron multiplets. After reducing it to the diagonal form we should get the matrix of the eigenvalues $\hat{M}_{ph}^2$:

$$\hat{M}_{ph}^2 = \begin{pmatrix} M_{f_0}^2(1) & 0 & 0 \\ 0 & M_{f_0}^2(2) & 0 \\ 0 & 0 & M_{f_0}^2(3) \end{pmatrix}$$

and the matrix $\hat{V}(i)$ of the eigenvectors in a chosen basis.

3. We start treating the mass relations with the higher-mass scalar 0^{++} -sector:

$$M_{a_0} = 1474 \pm 19, M_{K^*_0} = 1425 \pm 50, M_{f_0}(1) = 1370 \pm 50, M_{f_0}(2) = 1505 \pm 6, M_{f_0}(3) = 1724 \pm 7$$

where all values are in MeV .

We define the "bare" mass values M_N and M_S devoid of the strong annihilation contributions via

$$M_N = M_{a_0}, M_S^2 = 2M_{K^*_0}^2 - M_{a_0}^2 \quad (4)$$

We note that the second relation is alike to the S -wave vector quarkonia, but we would like to stress the opposite mass hierarchy sequence

$$M^2(S) < M^2(q\bar{s}) < M^2(N)$$

which follows if one changes the sign of the parameter μ^2 in the GMO mass operator. The changed sign of the term including μ^2 is a phenomenological reflection of an additional flavor dependence in the spin-orbit spin-tensor terms of the $q\bar{q}$, P-wave mass operator and of the result of the $q\bar{q}$ -nonet-states mixing with the lower-lying scalar S-wave nonet-states having the structure of the diquark-anti-diquark ($q^2, \bar{q}^2$)-bound states [6, 7].

The intermediate gluon-matter originated terms A_Q and A_G and the "bare" gluon mass M_G are unknown variables which have to be found by solution of the system of three equations representing the equalities of three invariants of the diagonalisation process: the trace, the determinant and the sum of main minors of the matrices under consideration. The small correction factor $r = \sqrt{M_S/M_N}$ is introduced as reflecting the relativistic normalization of the meson state vectors with different masses. Successively excluding unknown variables A_Q and A_G in favor of M_G , we solve numerically the last equation by varying the remaining unknown M_G under constraint $A_G^2 \geq 0$. There are 2 solutions for M_G and $A_G = 0$ and none for $A_G^2 > 0$. We have chosen as physically acceptable the value of the decoupled (because of $A_G = 0$!) physical glueball mass

$$\begin{aligned} M_G^{phen} &\simeq 1730 \text{ MeV}, \\ M_{f_0(3)}^{exp} &= 1724 \pm 7 \text{ MeV} \end{aligned} \quad (5)$$

The state vectors of $f_0(1506)$ and $f_0(1370)$ are obtained by the diagonalisation of the rest 2×2 matrix:

$$\begin{aligned} |f_0(1506)\rangle &= \pm 0.496|S\rangle + 0.868|N\rangle, \\ |f_0(1370)\rangle &= \mp 0.496|N\rangle + 0.868|S\rangle \end{aligned} \quad (6)$$

The choice of signs remains to be made on the physical grounds. The sensitive check of Eqs.(5) and (6) provides the radiative and hadronic decays of J/Ψ -resonance

$$\frac{\tilde{\Gamma}(J/\Psi \rightarrow \gamma + f_0(1506))}{\tilde{\Gamma}(J/\Psi \rightarrow \gamma + f_0(1370))} = \frac{(.868\sqrt{2} \pm .496)^2}{(\mp .496\sqrt{2} + .868)^2} \approx 107(.22) \quad (7)$$

$$\frac{\tilde{\Gamma}(J/\Psi \rightarrow \omega(780) + f_0(1710))}{\tilde{\Gamma}(J/\Psi \rightarrow \phi(1020) + f_0(1710))} = \frac{1}{\tan^2(\theta_V)} \simeq 2.58. \quad (8)$$

We define the "reduced" width by $\tilde{\Gamma}(A \rightarrow B + C) = (\Gamma(A \rightarrow B + C))/|\vec{k}(A \rightarrow B + C)|^3$, $\vec{k}(A \rightarrow B + C)$ being the 3-momentum of final particles, B and C , in the rest frame

of the resonance A , $\theta_V \simeq 39^\circ$. The larger (smaller) value in the ratio (7) refers to the upper(lower) sign in Eq.(6). The experimental value of ratio (8) $(1.19 \pm .34)$ [9] is twice smaller and signals of an additional suppression in the virtual transition ratio $|\langle gluons \rightarrow \phi \rangle|^2 / |\langle gluons \rightarrow \omega \rangle|^2 \sim (m_\omega/m_\phi)^2$.

4. (•) Large amount of strange quarks in $f_0(1370)$ is supported by the experimental observation that the decay $J/\Psi \rightarrow \phi\pi\pi$ cannot be fitted without excitation of the $f_0(1370)$ -resonance unlike the reaction with ϕ replaced by ω [8].

(•) Dominance of the pion(s) decay mode of $f_0(1506)$ [9] is in agreement with a large amount of the non-strange $q\bar{q}$ -quarks in the state-vector of this resonance.

(•) The latest data of BES-II Collaboration seem to signal on a hierarchy of the radiative decay modes [8]:

$$BR(J/\Psi \rightarrow \gamma f_0(1370)) < BR(J/\Psi \rightarrow \gamma f_0(1506)) < BR(J/\Psi \rightarrow \gamma f_0(1710))$$

that is in accord with the upper sign in Eq.(7) and with the large amount of "gluon-matter" in $f_0(1710)$. The color-averaged, one-gluon-exchange (OGE)-"potential" operating in transitions $c\bar{c} \rightarrow gg \rightarrow gg|_{bound}$ is $3/(4/3)$ times larger than in $c\bar{c} \rightarrow gg \rightarrow q\bar{q}|_{bound}$ -transitions thus giving a hint for enlarged value of the $BR(J/\Psi \rightarrow \gamma + f_0(1710))$ for the gluon-matter-composed $f_0(1710)$ -resonance.

(•) We note that the glueball dominance of the $f_0(1710)$ was inferred in a number of earlier works [10–13], though with a different N - and S -composition of the $f_0(1370)$ and $f_0(1506)$ mesons. In particular, our (6) with upper signs shows that $f_0(1370)(f_0(1506))$ is mainly $SU(3)$ -octet(singlet)-state just opposite to the assignment of ref. [13].

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POLARIZATION AT PHOTON COLLIDER. USING FOR PHYSICAL STUDIES

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Abstract

Photon Colliders will have photon beams with high and easily variable polarization (mainly circular). It offer opportunity for new experimental studies in the problems of hadron physics and QCD in a new energy and transfer regions, in the Standard Model physics and in the physics beyond Standard Model.

• **Photon Collider.** (PLC) is specific option of Linear Collider (LC). The focused laser flash meet the electron bunch of LC in the conversion point C. Here a laser photon scatters on high-energy electron taking from it a large portion of energy. Scattered photons travel along the direction of the initial electron, they are focused in the interaction point. Here they collide with opposite electron ($e\gamma$ collider) or photon ($\gamma\gamma$ collider) [1–3]. The PLC will start after 10 years of work of LHC and few years of work of e^+e^- ILC.

With two different laser photon energies two types of PLC are possible, the first one is most suitable for ILC-1 $E_e = 0.25$ TeV, the second one can be realized only for high energy PLC with $E_e = 0.5 \div 1.5$ TeV (ILC-2, CLIC) [2]. The typical expected parameters of PLC in these two variants are enumerated below. In the cases when parameters for second variant differ from those for the first variant they are shown in curly brackets.

The total additional cost is estimated as $\sim 10\%$ from that of LC [3].

The conversion coefficient $e \rightarrow \gamma$ is $0.7 \{0.15\}$.

Characteristic photon energy $E_{\gamma max} \approx 0.8 \{0.95\} E_e$.

For high energy peak ($E_{\gamma 1,2} > 0.7 E_{\gamma max}$), separated well from low energy part of spectrum

Luminosity $\mathcal{L}_{\gamma\gamma} \approx 0.35 \{0.01 \div 0.05\} \mathcal{L}_{ee}$ and $\mathcal{L}_{e\gamma} \approx 0.25 \{0.2\} \mathcal{L}_{ee}$.

with $\int \mathcal{L}_{\gamma\gamma} dt, \int \mathcal{L}_{e\gamma} dt \approx 200 \div 150 \{50 \div 100\} \text{ fb}^{-1}/\text{year}$.

Mean energy spread $\langle \Delta E_\gamma \rangle \approx 0.07 \{0.03\} E_{\gamma max}$.

Mean photon helicity $\langle \lambda_\gamma \rangle \approx 0.95$, with easily variable sign.

• **QCD and hadron physics.** Photon structure function is unique object of QCD, calculable at large enough Q^2 without additional phenomenological parameters [4]. It can be measured here in $e\gamma$ mode with high accuracy, since photon target with its energy and polarization here are practically known. In this case the range of accessible virtualities of photon is limited from **below** by the details of detector set-up.

The region of electron transverse momenta above 50 GeV ($M_Z/2$) can be studied well, providing opportunity to study effect of **Z-boson exchange and $\gamma^* - Z$ interference**.

The manipulation with beam polarizations will be important instrument here.

The other studies like those at HERA are possible here.

• **Right-handed EW currents.** The cross section $e\gamma \rightarrow \nu W \propto (1 - 2\lambda_e)$, it is switched on or off with variation of electron helicity λ_e . It gives very precise test of absence of right handed currents in the interaction of W with the matter far from the nuclear physics point $q_W^2 \approx 0$.

• **Higgs physics.** Higgs mechanism of EWSB can be realized either by minimal Higgs sector with one observable neutral scalar Higgs boson (SM) or by non-minimal Higgs sector with larger number of observable scalars. In this section for definiteness we consider SM and specific non-minimal Higgs sector – Two Higgs Doublet Model (2HDM). The latter is the simplest extension of Higgs sector of SM. It contains 2 complex Higgs doublet fields ϕ_1 and ϕ_2 with v.e.v.'s $v \cos \beta$ and $v \sin \beta$. The physical sector contains charged scalars $H^\pm$ and three neutral scalars h_i , generally having no definite CP parity. In the CP conserving case these three h_i become two scalars h , H and a pseudoscalar A . For definiteness, we assume the Model II for the Yukawa coupling in 2HDM (the same is realized in MSSM).

A. CP violation in Higgs sector. In many extensions of Higgs model (e.g. in 2HDM) observable neutral Higgs bosons h_i have generally **no** definite CP-parity and effectively

$$\mathcal{L}_{\gamma\gamma H} = G_\gamma^{SM} \left[g_\gamma H F^{\mu\nu} F_{\mu\nu} + i\tilde{g}_\gamma H F^{\mu\nu} \tilde{F}_{\mu\nu} \right]; \quad g_\gamma \sim \tilde{g}_\gamma \sim 1. \quad (1)$$

Here $F^{\mu\nu}$ and $\tilde{F}^{\mu\nu} = \varepsilon^{\mu\nu\alpha\beta} F_{\alpha\beta}/2$ are the standard field strength for the electromagnetic field. The relative effective couplings g and $\tilde{g}$ are described with standard triangle diagram $H\gamma\gamma$, they are expressed with known equations via masses of charged fermions and W , and mixing parameters (parameters of 2HDM potential). They are generally complex ($b\bar{b}$ –loop).

Total production cross section varies strong with variation of circular λ_i and linear ℓ_i polarizations of photons and the angle ψ between linear polarization vectors ℓ_i [5]:

$$\begin{aligned} \sigma(\gamma\gamma \rightarrow H) = \sigma_{np}^{SM} \times [& (|g_\gamma|^2 + |\tilde{g}_\gamma|^2)(1 + \lambda_1\lambda_2) + \\ & + (|g_\gamma|^2 - |\tilde{g}_\gamma|^2)\ell_1\ell_2 \cos 2\psi + 2\text{Re}(g_\gamma^*\tilde{g}_\gamma)(\lambda_1 + \lambda_2) + 2\text{Im}(g_\gamma^*\tilde{g}_\gamma)\ell_1\ell_2 \sin 2\psi] . \end{aligned} \quad (2)$$

In particular, violation of CP symmetry in the Higgs sector leads to difference in the $\gamma\gamma \rightarrow H$ production cross sections in the collision of photons with identical total helicity ($\lambda_1 - \lambda_2 = 0$) but with opposite helicities of separate photons (T_-):

$$T_- = \frac{\sigma(\lambda_i) - \sigma(-\lambda_i)}{\sigma_{np}^{SM}} \propto (\lambda_1 + \lambda_2) \text{Re}(g_\gamma \tilde{g}_\gamma^*). \quad (3)$$

Standard calculation of vertexes in the 2HDM at different parameters of model gives typical dependencies, presented in Fig. 1. It is seen that effects are strong and well measured.

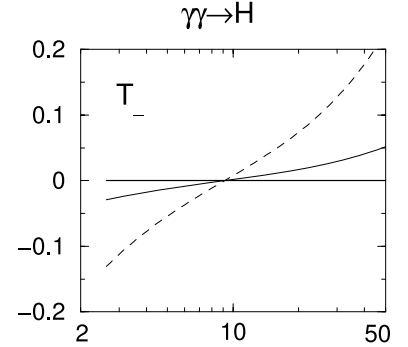


Figure 1: Effect of CP violation in 2HDM, $\lambda_1 = \lambda_2 = \pm 1$.

B. Observation of strong interaction in Higgs sector in $e\gamma \rightarrow eWW$ process at not too high energy. At high values of Higgs boson self-coupling constant, the Higgs mechanism of Electroweak Symmetry Breaking in Standard Model (SM) can be realized without actual Higgs boson but with strong interaction in Higgs sector (SIHS) which will manifest itself as a strong interaction of longitudinal components of W and Z

bosons. It is expected that this interaction will be similar to the interaction of π -mesons at $\sqrt{s} \lesssim 1.5$ GeV and will be seen in the form of $W_L W_L$, $W_L Z_L$ and $Z_L Z_L$ resonances at $1.5 \div 2$ TeV. Main efforts to discover this opportunity are directed towards the observation of such resonant states. It is a difficult task for the LHC due to high background and it cannot be realized at the energies reachable at the ILC in its initial stages.

This strong interaction can be observed in the study of *the charge asymmetry* (specific form of polarization effect) of produced $W^\pm$ in the process $e^- \gamma \rightarrow e^- W^+ W^-$. To explain the set up of the problem we discuss this process in SM [8].

The diagrams for the process are naturally subdivided for three groups. Here subprocesses of main interest are shown in boxes, sign $\otimes$ represents next stage of process.

a) Diagrams $\gamma e^- \rightarrow e^- \gamma^*(Z^*) \otimes \boxed{\gamma \gamma^*(\gamma Z^*) \rightarrow W^+ W^-}$ contain subprocesses $\gamma \gamma^* \rightarrow W^+ W^-$ and $\gamma Z^* \rightarrow W^+ W^-$, modified by the strong interaction in the Higgs sector (*two-gauge contribution*).

b) Diagrams $e^- \gamma \rightarrow e^- \gamma^*(Z^*) \otimes \boxed{\gamma^*(Z^*) \rightarrow W^+ W^-}$ contain subprocesses $\gamma^* \rightarrow W^+ W^-$ and $Z^* \rightarrow W^+ W^-$, modified by the strong interaction in the Higgs sector (*one-gauge contribution*).

c) Diagrams $\gamma \oplus \boxed{e^- \rightarrow W^- W^+ e^-}$ are prepared by connecting the photon line to each charged particle line to the diagram shown inside the box. Strong interaction does not modify this contribution. The latter contributions are switched off at suitable electron polarization.

The subprocess $\gamma \gamma^* \rightarrow W^+ W^-$ (from contribution a)) produces C-even system $W^+ W^-$, the subprocess $\gamma^* \rightarrow W^+ W^-$ (from contribution b)) produces C-odd system $W^+ W^-$. The interference of similar contributions for the production of pions is responsible for large enough charge asymmetry, very sensitive to the phase difference of S (D) and P waves in $\pi\pi$ scattering. This very phenomenon also takes place in the discussed case of W 's. However, for the production of $W^\pm$ subprocesses with the replacement of $\gamma^* \rightarrow Z^*$ are also essential. Therefore, the final states of each type have no definite C -parity. Hence, charge asymmetry appears both due to interference between contributions of types a) and b) and due to interference of γ^* and Z^* contributions each within their own types.

B1. Asymmetries in SM. To observe the main features of the effect of charge asymmetry and its potential for the study of strong interaction in the Higgs sector, we calculated some quantities describing charge asymmetry for $e^- \gamma$ collision at $\sqrt{s} = 500$ GeV with polarized photons. We used CompHEP and CalcHEP packages for simulation.

We denote by $p^\pm$ momenta of $W^\pm$, by p_e – momentum of the scattered electron and $w = \frac{\sqrt{(p^+ + p^-)^2}}{2M_W}$, $v_1 = \frac{\langle (p^+ - p^-) p_e \rangle}{\langle (p^+ + p^-) p_e \rangle}$. We present below dependence of charge asymmetric quantity v_1 on w . The w -dependencies for the other charge asymmetric quantities have similar qualitative features [8].

We applied *the cut in transverse momentum of the scattered electron*,

$$p_\perp^e \geq p_{\perp 0} \quad \text{with} \quad \begin{cases} a) p_{\perp 0} = 10 \text{ GeV}, \\ b) p_{\perp 0} = 30 \text{ GeV}. \end{cases} \quad (4)$$

Observation of the scattered electron allows to check kinematics completely.

B2. Influence of polarization. Fig. 2 (left and central plots) represents distribution in variable v_1 on photon polarization and cut in $p_\perp^e$. We did not study the dependence on

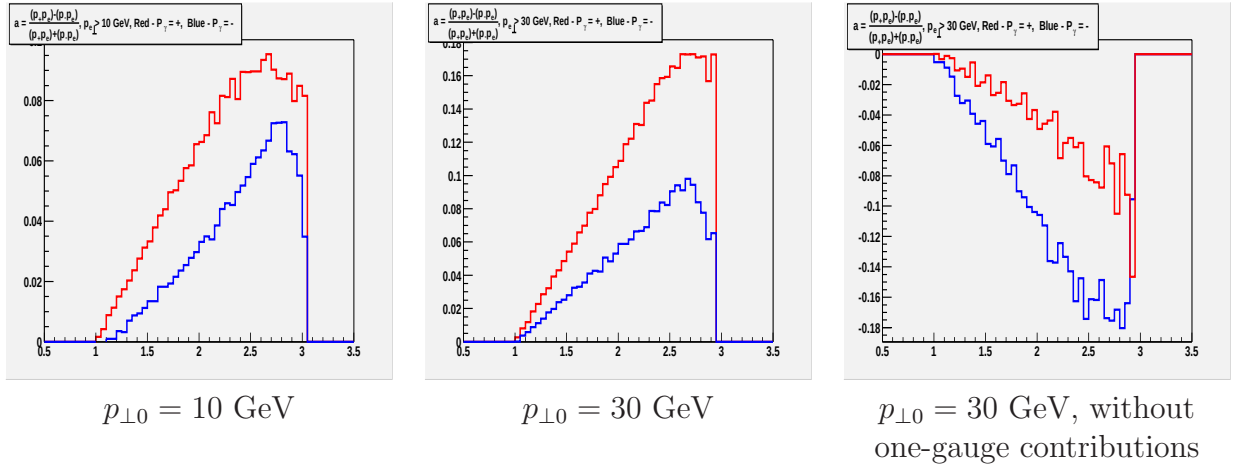


Figure 2: Distribution of v_1 in dependence on w . The upper curves are for right-hand polarized photons, the lower curves are for left-hand polarized photons

electron polarization. This dependence is expected to be weak in SM where main contribution to cross section is given by diagrams of type a) with virtual photons having the lowest possible energy. These photons "forget" the polarization of the incident electron. The strong interaction contribution becomes essential at highest effective masses of WW system with high energy of virtual photon or Z , the helicity of which reproduces almost completely the helicity of incident electron [6]. The study of this dependence will be a necessary part of studies beyond SM.

Significance of different contributions. To understand the extent of the effect of interest, we compared the entire distribution in variable v_1 with that without one-gauge contribution at $p_{\perp 0} = 30$ GeV (right plot in Fig. 2). Strong interaction in the Higgs sector modifies both one-gauge and two-gauge contributions. The study of charge asymmetry caused by their interference will be a source of information on this strong interaction. One can see that one-gauge contribution is so essential that neglecting on it even changes the sign of charge asymmetry (compared to that for the entire process).

Therefore, the charge asymmetry is very sensitive to the interference of two-gauge and one-gauge contributions which is modified under the strong interaction in the Higgs sector. The measurement of this asymmetry will be a source of data on the phase difference of different partial waves of $W_L W_L$ scattering.

• **New particles.** New charged particles will be discovered at LHC and in e^+e^- mode of LC. We expect their decay for final states with invisible particles (like LSP in MSSM). In the measurement of mass, decay modes and spin of these new particles the $\gamma\gamma$ production provides essential advantages compared to e^+e^- collisions.

The cross section of the *pair production* $\gamma\gamma \rightarrow P^+P^-$ ($P = S$ – scalar, $P = F$ – fermion, $P = W$ – gauge boson) not far from the threshold is given by QED with reasonable accuracy (Fig. 3). It is seen that these cross sections decreases slowly with energy growth. Therefore, they can be studied relatively far from the threshold where the decay products are almost non-overlapping.

Near the threshold

$$\sigma(\gamma\gamma \rightarrow P^+P^-) = \sigma^{np}(\gamma\gamma \rightarrow P^+P^-)(1 + \lambda_1\lambda_2 \pm \ell_1\ell_2 \cos 2\phi)$$

with + sign for $P = S$ and – sign for $P = F$ (here ℓ_i are vectors of linear polarization

of photons). This polarization dependence provides the opportunity to determine spin of produced particle P in the experiments with longitudinally polarized photons.

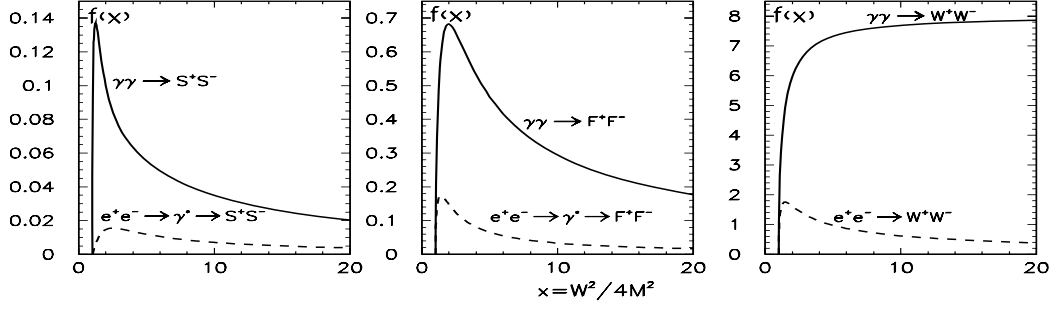


Figure 3: $\frac{\sigma(\gamma\gamma \rightarrow P^+P^-)}{\pi\alpha^2/M_P^2}$, nonpolarized photons, and $\frac{\sigma(e^+e^- \rightarrow \gamma^* \rightarrow P^+P^-)}{\pi\alpha^2/M_P^2}$

The polarization of produced fermion or vector P depends on the initial photon helicity. At the P decay this polarization is transformed into the momentum distribution of decay products. E.g., for the SM processes like $\gamma\gamma \rightarrow \mu^+\mu^- + \text{neutrals}$ (obtained from muon decay modes of $\gamma\gamma \rightarrow WW$, $\gamma\gamma \rightarrow \tau^+\tau^-$, etc.) muons should exhibit charge asymmetry linked to the polarization of initial photons – see [7]. These studies can help to understand the *nature of candidates for Dark Matter particles*.

The possible CP violation in the $P\gamma$ interaction can be seen as a variation of cross section with changing the sign of both photon helicities (like in fig. 1).

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CROSS SECTIONS AND SPIN ASYMMETRIES IN VECTOR MESON LEPTOPRODUCTION

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Abstract

Light vector meson leptonproduction is analyzed on the basis of the generalized parton distributions. Our results on the cross section and spin effects are in good agreement with experiment at HERA, COMPASS and HERMES energies. Predictions for A_{UT} asymmetry for various reactions are presented.

In this report, investigation of vector meson leptonproduction is based on the handbag approach where the leading twist amplitude at high Q^2 factorizes into hard meson electroproduction off partons and the Generalized Parton Distributions (GPDs) [1]. The higher twist (TT) amplitude which is essential in the description of spin effects exhibits the infrared singularities, which signals the breakdown of factorization [2]. These problems can be solved in our model [3] where subprocesses are calculated within the modified perturbative approach in which quark transverse degrees of freedom accompanied by Sudakov suppressions are considered. The quark transverse momentum regularizes the end-point singularities in the TT amplitudes so that it can be calculated.

In the model, the amplitude of the vector meson production off the proton with positive helicity reads as a convolution of the partonic subprocess $\mathcal{H}^V$ and GPDs H^i ($\tilde{H}^i$)

$$\mathcal{M}_{\mu'+,\mu+}^{Vi} = \frac{e}{2} \sum_a e_a C_a^V \sum_\lambda \int_{xi}^1 d\bar{x} \mathcal{H}_{\mu'\lambda,\mu\lambda}^{Vi} H_i(\bar{x}, \xi, t), \quad (1)$$

where i denotes the gluon and quark contribution, sum over a includes quarks flavor a and C_a^V are the corresponding flavor factors [3]; μ (μ') is the helicity of the photon (meson), and $\bar{x}$ is the momentum fraction of the parton with helicity λ . The skewness ξ is related to Bjorken- x by $\xi \simeq x/2$. In the region of small $x \leq 0.01$ gluons give the dominant contribution. At larger $x \sim 0.2$ the quark contribution plays an important role [3].

To estimate GPDs, we use the double distribution representation [4]

$$H_i(\bar{x}, \xi, t) = \int_{-1}^1 d\beta \int_{-1+|\beta|}^{1-|\beta|} d\alpha \delta(\beta + \xi \alpha - \bar{x}) f_i(\beta, \alpha, t). \quad (2)$$

The GPDs are related with PDFs through the double distribution function

$$f_i(\beta, \alpha, t) = h_i(\beta, t) \frac{\Gamma(2n_i + 2)}{2^{2n_i+1} \Gamma^2(n_i + 1)} \frac{[(1 - |\beta|)^2 - \alpha^2]^{n_i}}{(1 - |\beta|)^{2n_i+1}}. \quad (3)$$

The powers $n_i = 1, 2$ (i= gluon, sea, valence contributions) and the functions $h_i(\beta, t)$ are proportional to parton distributions [3].

To calculate GPDs, we use the CTEQ6 fits of PDFs for gluon, valence quarks and sea [5]. Note that the $u(d)$ sea and strange sea are not flavor symmetric. In agreement with CTEQ6 PDFs we suppose that $H_{sea}^u = H_{sea}^d = \kappa_s H_{sea}^s$, with

$$\kappa_s = 1 + 0.68/(1 + 0.52 \ln(Q^2/Q_0^2)) \quad (4)$$

The parton subprocess $\mathcal{H}^V$ contains a hard part which is calculated perturbatively and the $k_\perp$ - dependent wave function. It contains the leading and higher twist terms describing the longitudinally and transversally polarized vector mesons, respectively. The quark transverse momenta are considered in hard propagators decrease the LL amplitude and the cross section becomes in agreement with data. For the TT amplitude these terms regularize the end point singularities.

We consider the gluon, sea and quark GPDs contribution to the amplitude. This permits us to analyse vector meson production from low x to moderate values of x (~ 0.2) typical for HERMES and COMPASS. The obtained results [3] are in reasonable agreement with experiments at HERA [6, 7], HERMES [8], COMPASS [9] energies for electroproduced ρ and ϕ mesons.

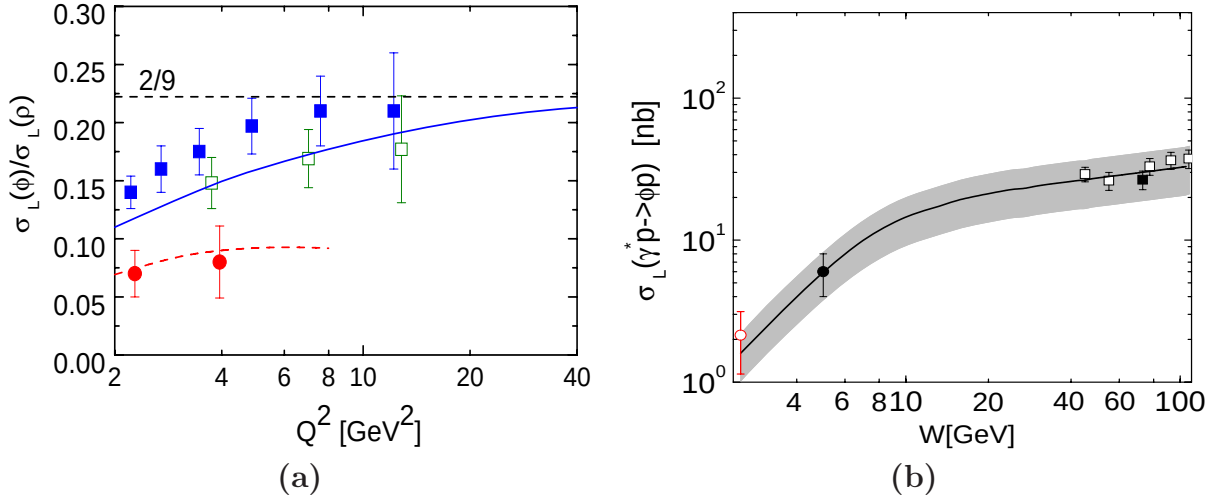


Figure 1: (a) The ratio of cross sections σ_ϕ/σ_ρ at HERA energies- full line and HERMES-dashed line. Data are from H1 -solid, ZEUS -open squares, HERMES solid circles. (b) The longitudinal cross section for ϕ at $Q^2 = 3.8 \text{ GeV}^2$. Data: HERMES, ZEUS, H1, open circle- CLAS data point.

In Fig 1a, we show the strong violation of the σ_ϕ/σ_ρ ratio from $2/9$ value at HERA energies and low Q^2 , which is caused by the flavor symmetry breaking (4) between $\bar{u}$ and $\bar{s}$. The valence quark contribution to σ_ρ decreases this ratio at HERMES energies. It was found that the valence quarks substantially contribute only at HERMES energies. At lower energies this contribution becomes small and the cross section decreases with energy. This is in contradiction with CLAS results which involve essential increasing of σ_ρ for $W < 5 \text{ GeV}$. On the other hand, we found good description of ϕ production at CLAS [10] Fig 1b. This means that we have problem only with the valence quark contribution at low energies.

The results for the $R = \sigma_L/\sigma_T$ ratio are shown in Fig 2a for HERA, COMPASS and HERMES energies. We found that our model describe fine W and Q^2 dependencies of R .

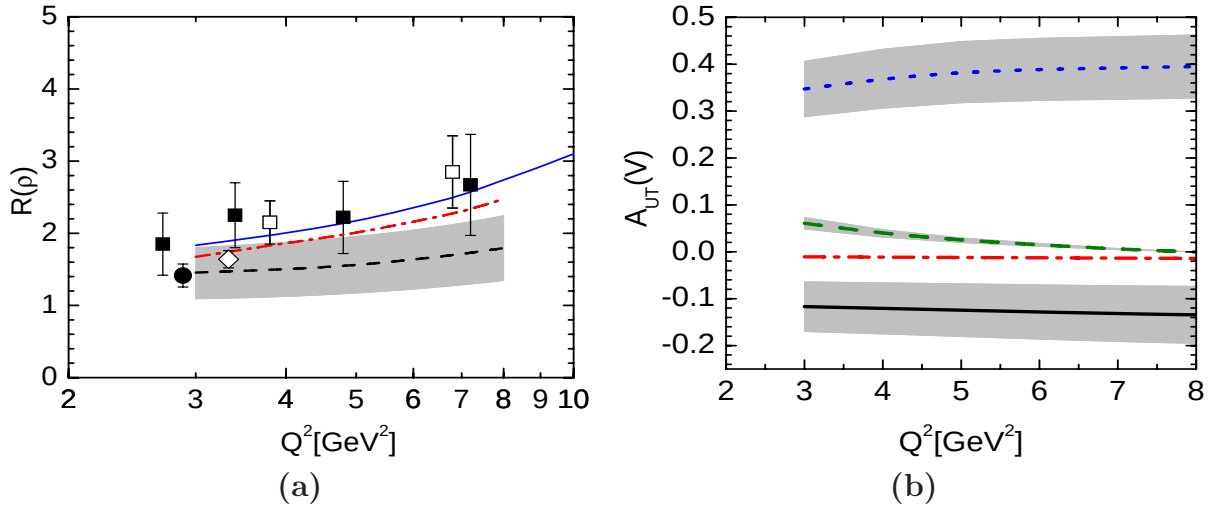


Figure 2: **(a)** The ratio of longitudinal and transverse cross sections for ρ production at low Q^2 . Full line- HERA, dashed-dotted -COMPASS and dashed- HERMES. **(b)** Predicted A_{UT} asymmetry at COMPASS for various mesons. Dotted-dashed line ρ^0 ; full line ω ; dotted line ρ^+ and dashed line K^{*0} .

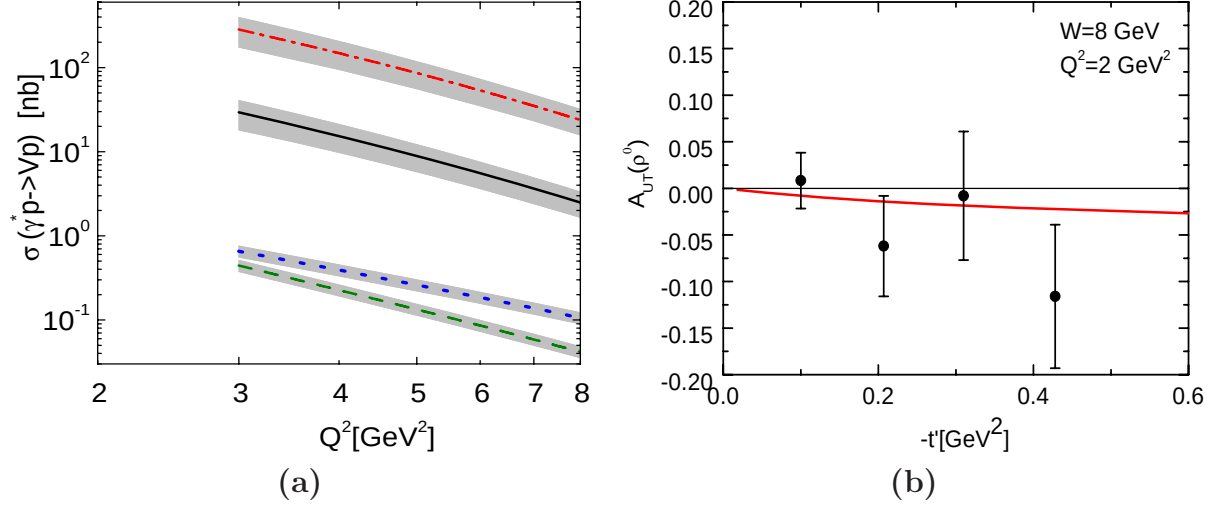


Figure 3: **(a)** The integrated cross sections of vector meson production at $W = 10 \text{ GeV}$. Lines are the same as in Fig 2b. **(b)** Predictions for A_{UT} asymmetry $W = 8 \text{ GeV}$. Preliminary COMPASS data at this energy are shown [11].

The analysis of the target A_{UT} asymmetry for electroproduction of various vector mesons was carried out in our approach too [12]. This asymmetry is sensitive to an interference between H and E GPDs. We constructed the GPD E from double distributions and constrain it by the Pauli form factors of the nucleon, positivity bounds and sum rules. The GPDs H were taken from our analysis of the electroproduction cross section. Predictions for the A_{UT} asymmetry at $W = 10 \text{ GeV}$ are given for ω , ϕ , ρ^+ , K^{*0} mesons [12] in Fig 2b. It can be seen that we predicted not small negative asymmetry for ω and large positive asymmetry for ρ^+ production. In these reactions the valence u and d quark

GPDs contribute to the production amplitude in combination $\sim E^u - E^d$ and do not compensate each other (E^u and E^d GPDs has different signs). The opposite case is for the ρ^0 production where one have the $\sim E^u + E^d$ contribution to the amplitude and valence quarks compensate each other essentially. As a result A_{UT} asymmetry for ρ^0 is predicted to be quite small. Unfortunately, it is much more difficult to analyse experimentally A_{UT} asymmetry for ω and ρ^+ production with respect to ρ^0 because the cross section for the first reactions is much smaller compared to ρ^0 , Fig. 3a. Our prediction for A_{UT} asymmetry of ρ^0 production at COMPASS reproduces well the preliminary experimental data Fig. 3b.

Thus, we can conclude that the vector meson electroproduction at small x is a good tool to probe the GPDs. In different energy ranges, information about quark and gluon GPDs can be extracted from the cross section and spin observables of the vector meson electroproduction.

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TWO-PHOTON EXCHANGE IN ELASTIC ELECTRON-PROTON SCATTERING: QCD FACTORIZATION APPROACH

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Abstract

We estimate the two-photon exchange contribution to elastic electron-proton scattering at large momentum transfer Q^2 . It is shown that the leading two-photon exchange amplitude behaves as $1/Q^4$ relative to the one-photon amplitude, and can be expressed in a model independent way in terms of the leading twist nucleon distribution amplitudes. Using several models for the nucleon distribution amplitudes, we provide estimates for existing data and for ongoing experiments.

Elastic electron-nucleon scattering in the one-photon (1γ) exchange approximation is a time-honored tool for accessing information on the structure of the nucleon. Precision measurements of the proton electric to magnetic form factor ratio at larger Q^2 using polarization experiments [1–3] have revealed significant discrepancies in recent years with unpolarized experiments using the Rosenbluth technique [4]. As no experimental flaw in either technique has been found, two-photon (2γ) exchange processes are the most likely culprit to explain this difference. Their study has received a lot of attention lately, see [5] for a recent review (and references therein), and [6] for a recent global analysis of elastic electron-proton (ep) scattering including 2γ corrections. In this work we calculate the leading in Q^2 behavior of the elastic ep scattering amplitude with hard 2γ exchange.

To describe the elastic ep scattering, $l(k) + N(p) \rightarrow l(k') + N(p')$, we adopt the definitions : $P = (p + p')/2$, $K = (k + k')/2$, $q = k - k' = p' - p$, and choose $Q^2 = -q^2$, $s/Q^2 = \zeta$ and $\nu = K \cdot P$ as the independent kinematical invariants. Neglecting the electron mass, it was shown in [7] that the T -matrix for elastic ep scattering can be expressed through 3 independent Lorentz structures as :

$$T = \frac{e^2}{Q^2} \bar{u}(k') \gamma_\mu u(k) \bar{u}(p') \left(\tilde{G}_M \gamma^\mu - \tilde{F}_2 \frac{P^\mu}{M} + \tilde{F}_3 \frac{\gamma \cdot K P^\mu}{M^2} \right) u(p),$$

where e is the proton charge and M is the proton mass. In Eq. (1), $\tilde{G}_M$, $\tilde{F}_2$, $\tilde{F}_3$ are complex functions of ν and Q^2 . To separate the 1γ and 2γ exchange contributions, it is furthermore useful to introduce the decompositions : $\tilde{G}_M = G_M + \delta\tilde{G}_M$, and $\tilde{F}_2 = F_2 + \delta\tilde{F}_2$, where $G_M(F_2)$ are the proton magnetic (Pauli) form factors (FFs) respectively, defined from the matrix element of the electromagnetic current, with $G_M(0) = \mu_p = 2.79$ the proton magnetic moment. The amplitudes $\tilde{F}_3$, $\delta\tilde{G}_M$ and $\delta\tilde{F}_2$ originate from processes involving at least 2γ exchange and are of order e^2 (relative to the factor e^2 in Eq. (1)).

In the hard regime, where $Q^2, s = (k + p)^2 \gg M^2$ and $s/Q^2 = \zeta$ is fixed, contribution to the 2γ exchange correction to the elastic ep amplitude is given by a convolution integral of the proton distribution amplitudes (DAs) with the hard coefficient function as shown in Fig. 1.

Let us now consider the proton matrix element which appears in the graph of Fig. 1. Following the notation from [8], the proton matrix element is described at leading twist level by three nucleon DAs as :

$$\begin{aligned}
& 4 \langle 0 | \varepsilon^{ijk} u_{\alpha}^i(a_1 \lambda n) u_{\beta}^j(a_2 \lambda n) d_{\sigma}^k(a_3 \lambda n) | p \rangle \\
& = V p^+ \left[\left(\frac{1}{2} \bar{n} \cdot \gamma \right) C \right]_{\alpha\beta} [\gamma_5 N^+]_{\sigma} \\
& + A p^+ \left[\left(\frac{1}{2} \bar{n} \cdot \gamma \right) \gamma_5 C \right]_{\alpha\beta} [N^+]_{\sigma} \\
& + T p^+ \left[\frac{1}{2} i \sigma_{\perp \bar{n}} C \right]_{\alpha\beta} [\gamma^{\perp} \gamma_5 N^+]_{\sigma}, \quad (1)
\end{aligned}$$

with light-cone momentum $p^+ = Q$, where C is charge conjugation matrix : $C^{-1} \gamma_{\mu} C = -\gamma_{\mu}^T$, and where $X = \{A, V, T\}$ stand for the nucleon DAs which are defined by the light-cone matrix element :

$$X(a_i, \lambda p^+) = \int d[x_i] e^{-i\lambda p^+ (\sum x_i a_i)} X(x_i), \text{ with } d[x_i] \equiv dx_1 dx_2 dx_3 \delta(1 - \sum x_i).$$

In the large Q^2 limit, the pQCD calculation of Fig. 1 involves 24 diagrams, and leads to hard 2γ corrections to $\delta\tilde{G}_M$, and $\nu/M^2 \tilde{F}_3$, which are found as [9]:

$$\delta\tilde{G}_M = -\frac{\alpha_{em}\alpha_S(\mu^2)}{Q^4} \left(\frac{4\pi}{3!} \right)^2 (2\zeta - 1) \int d[y_i] d[x_i] \frac{4x_2 y_2}{D} \{\dots\}, \quad (2)$$

$$\frac{\nu}{M^2} \tilde{F}_3 = -\frac{\alpha_{em}\alpha_S(\mu^2)}{Q^4} \left(\frac{4\pi}{3!} \right)^2 (2\zeta - 1) \int d[y_i] d[x_i] \frac{2(x_2 \bar{y}_2 + \bar{x}_2 y_2)}{D} \{\dots\}, \quad (3)$$

where

$$\begin{aligned}
\{\dots\} = & 2Q_u Q_d [V'V + A'A] (1, 3, 2) + Q_u^2 [(V' + A')(V + A) + 4T'T] (3, 2, 1) \\
& + Q_u Q_d [(V' + A')(V + A) + 4T'T] (1, 2, 3), \quad (4)
\end{aligned}$$

with quark charges $Q_u = +2/3$, $Q_d = -1/3$, $\alpha_{em} = e^2/(4\pi)$, $\alpha_s(\mu^2)$ is the strong coupling constant evaluated at scale μ^2 , and the denominator factor D is defined as :

$$D \equiv (y_1 y_2 \bar{y}_2) (x_1 x_2 \bar{x}_2) [x_2 \bar{\zeta} + y_2 \zeta - x_2 y_2 + i\varepsilon] [x_2 \zeta + y_2 \bar{\zeta} - x_2 y_2 + i\varepsilon]. \quad (5)$$

The unprimed (primed) quantities in Eqs. (2, 3) refer to the DAs in the initial (final) proton respectively. Eqs. (2, 3) are the central result of the present work. One notices that at large Q^2 , the leading behavior for $\delta\tilde{G}_M$ and $\nu/M^2 \tilde{F}_3$ goes as $1/Q^4$. In contrast, the invariant $\delta\tilde{F}_2$ is suppressed in this limit and behaves as $1/Q^6$. The similar calculation (up to normalization factor) was also reported in [10].

To evaluate the convolution integrals in Eqs. (2, 3), we need to insert a model for the nucleon twist-3 DAs, V , A , and T . The asymptotic behavior of the DAs and their first conformal moments were given in [8] as :

$$\begin{aligned}
V(x_i) & \simeq 120 x_1 x_2 x_3 f_N [1 + r_+(1 - 3x_3)], \quad A(x_i) \simeq 120 x_1 x_2 x_3 f_N r_-(x_2 - x_1), \\
T(x_i) & \simeq 120 x_1 x_2 x_3 f_N \left[1 + \frac{1}{2} (r_- - r_+) (1 - 3x_3) \right], \quad (6)
\end{aligned}$$

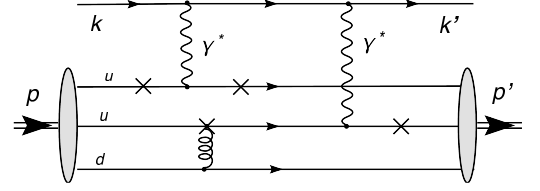


Figure 1: Typical graph for the elastic ep scattering with two hard photon exchanges. The crosses indicate the other possibilities to attach the gluon. The third quark is conventionally chosen as the d -quark. There are other diagrams where the one photon is connected with u - and d - quarks. We do not show these graphs for simplicity.

and depend on three parameters : f_N , r_- and r_+ . In this work, we will provide calculations using two models for the DAs that were discussed in the literature. The corresponding parameters (at $\mu = 1$ GeV) read

$$\begin{aligned} \text{COZ [11]: } & f_N = 5.0 \pm 0.5, r_- = 4.0 \pm 1.5, r_+ = 1.1 \pm 0.3, \\ \text{BLW [12]: } & f_N = 5.0 \pm 0.5, r_- = 1.37, r_+ = 0.35. \end{aligned} \quad (7)$$

We next calculate the effect of hard 2γ exchange, given through Eqs. (2, 3), on the elastic ep scattering observables. The general formulas for the observables including the 2γ corrections $\delta\tilde{G}_M$, $\delta\tilde{F}_2$, and $\nu/M^2\tilde{F}_3$ were derived in [7], to which we refer for the corresponding expressions. We assume that perturbative expansion at moderate Q^2 region is already applicable and fix the renormalization scale to be $\mu^2 = 0.6Q^2$ for each value of Q^2 shown below.

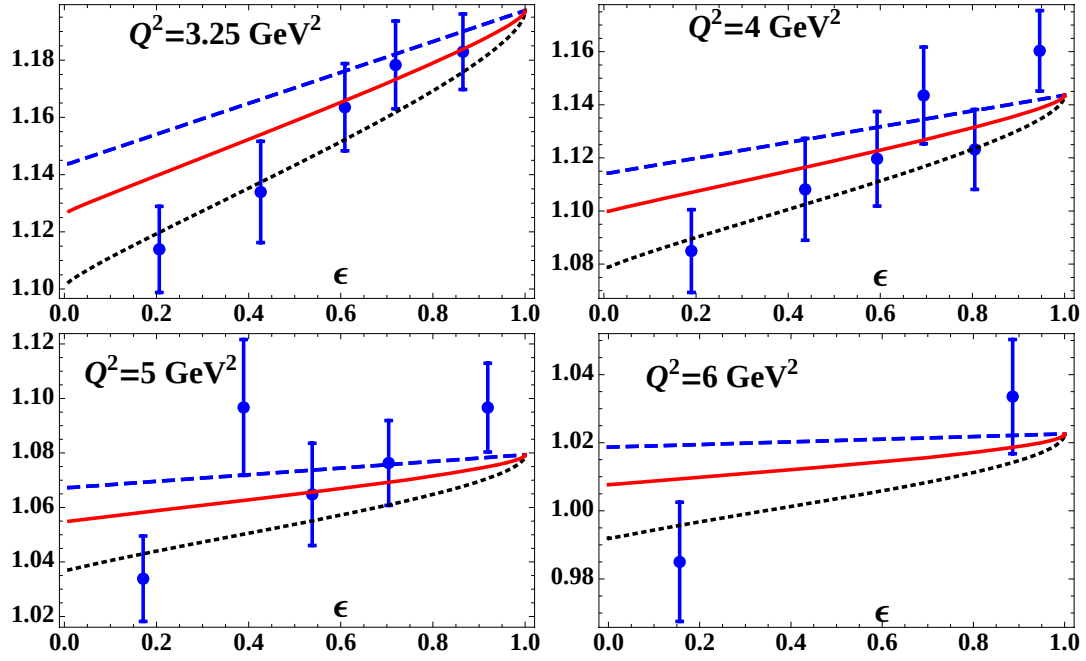


Figure 2: Rosenbluth plots for elastic ep scattering: σ_R divided by $\mu_p^2/(1 + Q^2/0.71)^4$. Dashed (blue) curves: 1γ exchange, using the G_{Ep}/G_{Mp} ratio from polarization data [1–3] and empirical parametrization for G_{Mp} from [13]. Solid red (dotted black) curves show the effect including hard 2γ exchange calculated with the BLW (COZ) model for the proton DAs. The data are from Ref. [4].

In Fig. 2, we calculate the reduced cross section σ_R as a function of the photon polarization parameter ε and different values of Q^2 . In the 1γ exchange, $\sigma_R = G_M(Q^2) + \varepsilon/\tau G_E(Q^2)$, with $\tau = Q^2/(4M^2)$, and the Rosenbluth plot is linear in ε , indicated by the dashed straight lines in Fig. 2. The effect including the hard 2γ exchange is shown for both the COZ and BLW models of the proton DAs. One sees that including the 2γ exchange changes the slope of the Rosenbluth plot, and that sizeable non-linearities only occur for ε close to 1. The inclusion of the hard 2γ exchange is able to well describe the Q^2 dependence of the unpolarized data, when using the polarization data [1–3] for the proton FF ratio G_{Ep}/G_{Mp} as input. Quantitatively, the COZ model for the nucleon DA leads to a correction about twice as large as when using the BLW model. We like to note

here that in contrast to the pQCD treatment of the proton FFs, which requires two hard gluon exchanges, the 2γ correction to elastic ep scattering only requires one hard gluon exchange. One therefore expects the pQCD calculation to set in for Q^2 values in the few GeV^2 range, which, probably, is confirmed by the results shown in Fig. 2. More detailed discussion of various observables can be found in [9].

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$O(\alpha_s)$ SPIN EFFECTS IN $e^+e^- \rightarrow q\bar{q}(g)$

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Abstract

We discuss $O(\alpha_s)$ spin effects in $e^+e^- \rightarrow q\bar{q}(g)$ for the polarization of single quarks and for spin-spin correlations of the final state quarks. Particular attention is paid to residual mass effects in the limit $m_q \rightarrow 0$ which are described in terms of universal helicity flip and helicity non-flip contributions.

1 Introduction

In a series of papers we have investigated $O(\alpha_s)$ final state spin phenomena in $e^+e^- \rightarrow q\bar{q}(g)$. In [1–3] we have provided analytical results for the polarization of single quarks and in [4, 5] for longitudinal spin-spin correlations of the final state quarks including their dependence on beam-event correlations¹. By carefully taking the $m_q \rightarrow 0$ limits of our analytical $O(\alpha_s)$ results we have found that, at $O(\alpha_s)$, the single-spin polarization $P^{(\ell_1)}$ and the spin-spin polarization correlation $P^{(\ell_1\ell_2)}$ do not agree with their $m_q = 0$ counterparts. In fact, averaging over beam-event correlation effects, at $O(\alpha_s)$ one has

$$P_{m_q \rightarrow 0}^{(\ell_1)} = P_{m_q=0}^{(\ell_1)} \left(1 - \left[\frac{2}{3} \frac{\alpha_s}{\pi} \right] \right), \quad P_{m_q \rightarrow 0}^{(\ell_1\ell_2)} = P_{m_q=0}^{(\ell_1\ell_2)} \left(1 - \left[\frac{4}{3} \frac{\alpha_s}{\pi} \right] \right) \quad (1)$$

where the residual $m_q \rightarrow 0$ contributions are encased in square brackets. In Eq.(1) $P^{(\ell_1)}$ denotes the longitudinal single-spin polarization of the final state quark and $P^{(\ell_1\ell_2)}$ denotes the longitudinal spin-spin polarization correlation of the quark and antiquark. From Eq.(1) it is apparent that at $O(\alpha_s)$ $QCD(m_q = 0) \neq QCD(m_q \rightarrow 0)$ in polarization phenomena. We mention that the longitudinal polarization components are the only polarization components that survive in the high energy (or $m_q = 0$) limit. Here and in the following the residual mass contributions will be called anomalous contributions for the reason that the anomalous helicity flip contribution enters as absorptive input in the dispersive derivation of the value of the axial anomaly [7]. We shall see further on that $P_{m_q=0}^{(\ell_1)} = g_{14}/g_{11}$ where g_{14} and g_{11} are electroweak coupling coefficients and $P_{m_q=0}^{(\ell_1\ell_2)} = -1$ independent of the electroweak coupling coefficients. It is important to keep in mind that there are no residual mass effects at $O(\alpha_s)$ in the unpolarized rate.

By an explicit calculation we have checked that the difference of the two results originates from the near-forward region (see [8]) which is very suggestive of an explanation in

¹Numerical $O(\alpha_s)$ results on final state quark polarization effects in $e^+e^- \rightarrow q\bar{q}(g)$ can be found in [6].

terms of universal near-forward contributions. That there is a universal helicity flip contribution in the splitting process $q_{\pm} \rightarrow q_{\mp} + g$ has been noted some time ago in the context of QED [9] (see also [10, 11]). We have found that there also exists a universal helicity non-flip contribution $q_{\pm} \rightarrow q_{\pm} + g$ which is perhaps not so well known in the literature. In fact, the anomalous contribution to the single-spin polarization $P^{(\ell_1)}$ results to 100% from the universal helicity non-flip contribution whereas the anomalous contributions to the spin-spin correlation $P^{(\ell_1\ell_2)}$ is 50% helicity flip and 50% helicity non-flip.

2 Definition of polarization observables

In the limit $m_q \rightarrow 0$ the relevant spin degrees of freedom are the longitudinal polarization components $s_1^\ell = 2\lambda_q$ and $s_2^\ell = 2\lambda_{\bar{q}}$ of the quark and the antiquark. One defines an unpolarized structure function H and single-spin and spin-spin polarized structure functions $H^{(\ell_1)}$, $H^{(\ell_2)}$ and $H^{(\ell_1\ell_2)}$, resp., according to

$$H(s_1^\ell s_2^\ell) = \frac{1}{4} (H + H^{(\ell_1)} s_1^\ell + H^{(\ell_2)} s_2^\ell + H^{(\ell_1\ell_2)} s_1^\ell s_2^\ell). \quad (2)$$

Eq.(2) can be inverted to give

$$\begin{aligned} H &= [H(\uparrow\uparrow)] + H(\uparrow\downarrow) + H(\downarrow\uparrow) + [H(\downarrow\downarrow)], \\ H^{(\ell_1)} &= [H(\uparrow\uparrow)] + H(\uparrow\downarrow) - H(\downarrow\uparrow) - [H(\downarrow\downarrow)], \\ H^{(\ell_2)} &= [H(\uparrow\uparrow)] - H(\uparrow\downarrow) + H(\downarrow\uparrow) - [H(\downarrow\downarrow)], \\ H^{(\ell_1\ell_2)} &= [H(\uparrow\uparrow)] - H(\uparrow\downarrow) - H(\downarrow\uparrow) + [H(\downarrow\downarrow)]. \end{aligned} \quad (3)$$

We have indicated in (3) that the spin configurations $H(\uparrow\uparrow)$ and $H(\downarrow\downarrow)$ can only be populated by anomalous contributions in the $m_q \rightarrow 0$ limit. The normalized single-spin polarization $P^{(\ell_1)}$ and the spin-spin correlation $P^{(\ell_1\ell_2)}$ are then given by

$$P^{(\ell_1)} = \frac{g_{14}}{g_{11}} \frac{H^{(\ell_1)}}{H} \quad P^{(\ell_1\ell_2)} = \frac{H^{(\ell_1\ell_2)}}{H}, \quad (4)$$

where g_{14} and g_{11} are q^2 -dependent electroweak coupling coefficients (see e.g. [3]). For the ratio of electroweak coupling coefficients one finds $g_{14}/g_{11} = -0.67$ and -0.94 for up-type and down-type quarks, respectively, on the Z_0 resonance, and $g_{14}/g_{11} = -0.086$ and -0.248 for $q^2 \rightarrow \infty$.

Considering the fact that one has $H(\uparrow\uparrow) = H(\downarrow\downarrow) = 0$ for $m_q = 0$, $H^{pc}(\uparrow\downarrow) = H^{pc}(\downarrow\uparrow)$ for the p.c. structure functions ($H, H^{(\ell_1\ell_2)}$), and $H^{pv}(\uparrow\downarrow) = -H^{pv}(\downarrow\uparrow)$ for the p.v. structure function $H^{(\ell_1)}$ with $H^{pv}(\uparrow\downarrow) = H^{pc}(\uparrow\downarrow)$, it is not difficult to see from Eq.(3) that $H^{(\ell_1)}/H = +1$ and $H^{(\ell_1\ell_2)}/H = -1$ for $m_{q=0}$ to all orders in perturbation theory.

3 Polarization results for $m_q \rightarrow 0$ and $m_q = 0$

The results of taking the $m_q \rightarrow 0$ limit of our analytical finite mass results in [1–5] can be concisely written as

$$\begin{aligned} H^{pc}(s_1^\ell, s_2^\ell) &= \frac{1}{4} \left(H^{pc} + H^{pc(\ell_1 \ell_2)} s_1^\ell s_2^\ell \right) = N_c q^2 \left((1 - s_1^\ell s_2^\ell) \left(1 + \frac{\alpha_s}{\pi} \right) + \left[\frac{4}{3} \times \frac{\alpha_s}{\pi} s_1^\ell s_2^\ell \right] \right), \\ H^{pv}(s_1^\ell, s_2^\ell) &= \frac{1}{4} \left(H^{pv(\ell_1)} s_1^\ell + H^{pv(\ell_2)} s_2^\ell \right) = N_c q^2 (s_1^\ell - s_2^\ell) \left(1 + \frac{\alpha_s}{\pi} - \left[\frac{2}{3} \times \frac{\alpha_s}{\pi} \right] \right). \end{aligned} \quad (5)$$

We have again indicated the parity nature of the structure functions ((pc): parity conserving; (pv): parity violating). The $m_q = 0$ results are obtained from Eq.(5) by dropping the anomalous square bracket contributions (see [3, 5]).

It is not difficult to see that one has $H^{pc} = -H^{pc(\ell_1 \ell_2)} = H^{pv(\ell_1)} = -H^{pv(\ell_2)}$ for $m_q = 0$ by commuting γ_5 through the relevant $m_q = 0$ diagrams. As Eq.(5) shows these relations no longer hold true for the anomalous contributions showing again that the anomalous contributions originate from residual mass effects which obstruct the simple γ_5 -commutation structure of the $m_q = 0$ contributions.

4 Near-forward gluon emission

In Table 1 we list the helicity amplitudes $h_{\lambda_{q_1} \lambda_{q_2} \lambda_g}$ for the splitting process $q(p_1) \rightarrow q(p_2) + g(p_3)$ and the $\cos \theta$ -dependence of the helicity amplitudes in the near-forward region. We also list the difference of the initial helicity and the final helicities $\Delta\lambda = \lambda_{q_1} - \lambda_{q_2} - \lambda_g$.

$h_{\lambda_{q_1} \lambda_{q_2} \lambda_g}$	$\Delta\lambda$	$\cos \theta$ dependence
$h_{1/2 \ 1/2 \ +1}$	-1	$\sim \sqrt{(1 - \cos \theta)}$
$h_{1/2 \ 1/2 \ -1}$	+1	$\sim \sqrt{(1 - \cos \theta)}$
$h_{1/2 \ -1/2 \ +1}$	0	$\sim m_q/E$
$h_{1/2 \ -1/2 \ -1}$	+2	0

Table 1. Helicity amplitudes and their near-forward behaviour.

Column 2 shows helicity balance $\Delta\lambda = \lambda_{q_1} - \lambda_{q_2} - \lambda_g$.

In the forward direction $\cos \theta = 1$ the only surviving helicity amplitude is $h_{1/2 \ -1/2 \ +1}$ for which the helicities satisfy the collinear angular momentum conservation rule $\Delta\lambda = 0$.

Squaring the helicity flip amplitude $h_{1/2 \ -1/2 \ +1}$ and adding in the propagator denominator factor $2p_2 p_3 = 2E^2 x(1-x)(1 - \sqrt{1 - m^2/E^2})$ one obtains ($x = E_3/E$; E is the energy of the incoming quark and $s = 4E^2$)

$$\frac{d\sigma_{[hf]}}{dx d\cos\theta} = \sigma_{Born}(s) C_F \frac{\alpha_s}{4\pi} x \frac{m_q^2}{E^2} \frac{1}{\left(1 - \cos \theta \sqrt{1 - m_q^2/E^2} \right)^2} \quad (6)$$

Note that the helicity flip contribution $h_{1/2 \ -1/2 \ +1}$ is not seen in a $m_q = 0$ calculation. The helicity flip splitting function $d\sigma_{[hf]}/d\cos\theta$ is strongly peaked in the forward direction.

Using the small (m_q^2/E^2) -approximation one finds that $\sigma_{[hf]}$ has fallen to 50% of its forward peak value at $\cos\theta = 1 - (\sqrt{2} - 1)m_q^2/(2E^2)$.

Integrating Eq.(6) over $\cos\theta$, we obtain

$$\frac{d\sigma_{[hf]}}{dx} = \sigma_{Born}(s) C_F \frac{\alpha_s}{2\pi} x \equiv \sigma_{Born}(s) D_{[hf]}(x) \quad (7)$$

where $D_{[hf]} = C_F(\alpha_s/2\pi)x$ is called the helicity flip splitting function [11]. The integrated helicity flip contribution survives in the $m_q \rightarrow 0$ limit since the integral of the last factor in Eq.(6) is proportional to $1/m_q^2$ which cancels the overall m_q^2 factor in Eq.(6). It is for this reason that the survival of the helicity flip contribution is sometimes called an m/m -effect.

For the helicity non-flip contribution one finds

$$\frac{d\sigma_{nf}}{dx d\cos\theta} = \sigma_{Born}(s) C_F \frac{\alpha_s}{2\pi} \frac{1 + (1-x)^2}{x} \frac{(1 - \cos\theta)}{\left(1 - \cos\theta \sqrt{1 - m_q^2/E^2}\right)^2} \quad (8)$$

The helicity non-flip splitting function $d\sigma_{nf}/d\cos\theta$ vanishes in the forward direction but is strongly peaked in the near-forward direction. Using again the small (m_q^2/E^2) approximation σ_{nf} can be seen to peak at $\cos\theta = 1 - m_q^2/(2E^2)$.

Integrating Eq.(8) one obtains

$$\frac{d\sigma_{nf}}{dx} = \sigma_{Born}(s) C_F \frac{\alpha_s}{\pi} \frac{1 + (1-x)^2}{x} \ln \frac{E}{m_q} \equiv \sigma_{Born}(s) D_{nf}(x) \quad (9)$$

where we have retained only the leading-log contribution. $D_{nf}(x)$ can be seen to be the usual non-flip splitting function.

We now turn to the anomalous helicity non-flip contribution. Let us rewrite the non-flip contribution in the form

$$\sigma_{nf} = \underbrace{\sigma_{nf} + \sigma_{[hf]}}_{\sigma_{total}} - \sigma_{[hf]} \quad (10)$$

By explicit calculation we have seen that the total unpolarized rate σ_{total} has no anomalous contribution. The conclusion is that there is an anomalous contribution also to the non-flip transition with the strength $-D_{[hf]} = -C_F(\alpha_s/2\pi)x$. We conjecture that the same pattern holds true for other processes, i.e. that there are no anomalous contributions to unpolarized rates but that there are anomalous contributions to both helicity flip and non-flip contributions in polarized rates.

Taking into account the anomalous helicity flip $\sigma_{[hf]}$ and non-flip $\sigma_{[nf]}$ contributions the pattern of the anomalous helicity contributions to the various spin configurations can then be obtained in terms of the Born term contributions and the universal flip and non-flip contributions $\pm \int_0^1 D_{[hf]}(x)dx = \pm C_F \alpha_s/(4\pi) = \pm \alpha_s/(3\pi)$. Using a rather suggestive notation for the anomalous contributions one has

$$\begin{aligned} [\uparrow\uparrow] &= (\uparrow\downarrow_{[hf]}) + (\downarrow_{[hf]}\uparrow) = \left((\uparrow\downarrow)_{Born} + (\downarrow\uparrow)_{Born} \right) \left[C_F \frac{\alpha_s}{4\pi} \right] \\ [\uparrow\downarrow] &= (\uparrow\downarrow_{[nf]}) + (\uparrow_{[nf]}\downarrow) = -2(\uparrow\downarrow)_{Born} \left[C_F \frac{\alpha_s}{4\pi} \right], \\ [\downarrow\uparrow] &= (\downarrow\uparrow_{[nf]}) + (\downarrow_{[nf]}\uparrow) = -2(\downarrow\uparrow)_{Born} \left[C_F \frac{\alpha_s}{4\pi} \right], \\ [\downarrow\downarrow] &= (\downarrow\uparrow_{[hf]}) + (\uparrow_{[hf]}\downarrow) = \left((\downarrow\uparrow)_{Born} + (\uparrow\downarrow)_{Born} \right) \left[C_F \frac{\alpha_s}{4\pi} \right] \end{aligned} \quad (11)$$

Since one has $(\uparrow\downarrow)_{Born}^{pc} = (\downarrow\uparrow)_{Born}^{pc} = 2N_c q^2$ and $(\uparrow\downarrow)_{Born}^{pv} = -(\downarrow\uparrow)_{Born}^{pv} = 2N_c q^2$ one obtains the anomalous contributions in Eq.(5) using the factorizing relations (11). In particular one sees that anomalous contribution to the single-spin polarization $P^{(\ell_1)}$ results to 100% from the universal helicity non-flip contribution whereas the anomalous contributions to the spin-spin correlation function $P^{(\ell_1\ell_2)}$ is 50% helicity flip and 50% helicity non-flip. We conclude that the anomalous non-flip contribution are unavoidable in the $O(\alpha_s)$ description of $m_q \rightarrow 0$ spin phenomena in $e^+e^- \rightarrow q\bar{q}(g)$. The normal contributions in Eq.(5) require an explicit calculation. The result $H = 4N_c q^2(1 + \alpha_s/\pi)$ is, of course, well known since many years.

In this talk we did not discuss physics aspects of the anomalous helicity flip contribution. We refer the interested reader to the papers [10–13] which contain a discussion of various aspects of the physics of the anomalous helicity flip contributions in QED and QCD.

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Abstract

In this talk various spin effects in hard exclusive electroproduction of mesons are briefly reviewed and the data discussed in the light of recent theoretical calculations within the frame work of the handbag approach. For π^+ electroproduction it is shown that there is a strong contribution from $\gamma_T^* \rightarrow \pi$ transitions which can be modeled by the transversity GPD H_T accompanied by the twist-3 meson wave function.

1 Introduction

Electroproduction of mesons allows for the measurement of many spin effects. For instance, using a longitudinally or transversely polarized target and/or a longitudinally polarized beam various spin asymmetries can be measured. The investigation of spin-dependent observables allows for a deep insight in the underlying dynamics. Here, in this article, it will be reported upon some spin effects and their dynamical interpretation in the frame work of the so-called handbag approach which offers a partonic description of meson electroproduction provided the virtuality of the exchanged photon, Q^2 , is sufficiently large. The theoretical basis of the handbag approach is the factorization of the process amplitude into a hard partonic subprocess and in soft hadronic matrix elements, the so-called generalized parton distributions (GPDs), as well as wave functions for the produced mesons, see Fig. 1. In collinear approximation factorization has been shown to hold rigorously for hard exclusive meson electroproduction [1, 2]. It has also been shown that the transitions from a longitudinally polarized photon to a likewise polarized vector meson or a pseudoscalar one, $\gamma_L^* \rightarrow V_L(P)$, dominates at large Q^2 . Other photon-meson transitions are suppressed by inverse powers of the hard scale.

Here, in this article a variant of the handbag approach is utilized for the interpretation of the data in which the subprocess amplitudes are calculated within the modified perturbative approach [3], and the GPDs are constructed from reggeized double distributions [4, 5]. In the modified perturbative approach the quark transverse momenta are retained in the subprocess and Sudakov suppressions are taken into account. The partons are still emitted and re-absorbed by the proton collinearly. For the meson wave functions Gaussians in the variable $k_\perp^2/(\tau(1-\tau))$ are assumed with transverse size parameters fitted to experiment [6]. The variable τ denotes the fraction of the meson's momentum the quark entering the meson, carries. In a series of papers [7] it has been shown that with the proposed handbag approach the data on the cross sections and spin density matrix elements (SDMEs) for ρ^0 and ϕ production are well fitted in the kinematical range of $Q^2 \gtrsim 3 \text{ GeV}^2$, $W \gtrsim 5 \text{ GeV}$ (i.e. for small values of skewness $\xi \simeq x_{Bj}/2 \lesssim 0.1$) and

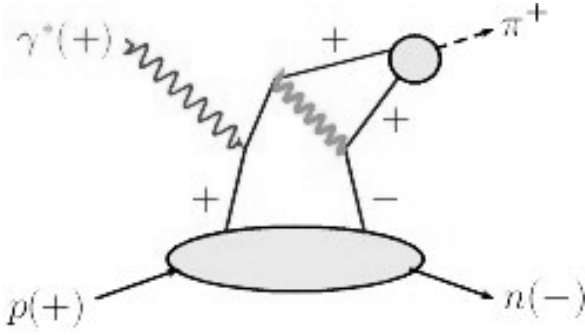


Figure 1: A typical lowest order Feynman graph for meson electroproduction. The signs indicate helicity labels for the contribution from transversity GPDs to the amplitude $\mathcal{M}_{0-,++}$, see text.

voted to a discussion of the target spin asymmetries in pion electroproduction and Sect. 4 to those for vector mesons and the GPD E . Finally, in Sect. 5, a summary is presented.

2 Target asymmetries

The electroproduction cross sections measured with a transversely or longitudinally polarized target consist of many terms, each can be projected out by a $\sin \varphi$ or $\cos \varphi$ moment where φ is a linear combination of ϕ , the azimuthal angle between the lepton and the hadron plane and ϕ_s , the orientation of the target spin vector [8]. In Tab. 1 the features of some of these moments are displayed. As the dominant interference terms reveal the target asymmetries provide detailed information on the $\gamma^* p \rightarrow MB$ amplitudes and therefore on the underlying dynamics that generates them.

A number of these moments have been measured recently. A particularly striking result is the $\sin \phi_s$ moment which has been measured by the HERMES collaboration for π^+ electroproduction [9]. The data on this moment, shown in Fig. 1, exhibit a mild t -dependence and do not show any indication for a turnover towards zero for $t' \rightarrow 0$. Inspection of Tab. 1 reveals that this behavior of $A_{UT}^{\sin \phi_s}$ at small $-t'$ requires a contribution from the interference term $\text{Im}[\mathcal{M}_{0-,++}^* \mathcal{M}_{0+,0+}]$. Both the contributing amplitudes are helicity non-flip ones and are therefore not forced to vanish in the forward direction by angular momentum conservation. Thus, we see that for pion electroproduction there are strong contributions from $\gamma_T^* \rightarrow \pi$ transitions. The underlying dynamical mechanism for such transitions will be discussed in Sect. 3.

For ρ^0 production the $\sin(\phi - \phi_s)$ moment has been measured by HERMES [10] and COMPASS [11]; the latter data being still preliminary. The HERMES data are shown in Fig. 2. In the handbag approach $A_{UT}^{\sin(\phi - \phi_s)}$ can also be expressed by an interference term of the convolutions of the GPDs H and E with hard scattering kernels

$$A_{UT}^{\sin \phi - \phi_s} \sim \text{Im}\langle E \rangle^* \langle H \rangle \quad (1)$$

instead of the helicity amplitudes. Given that H is known from the analysis of the ρ^0 and ϕ cross sections and SDMEs, A_{UT} provides information on E [12]. In order to calculate this

for the squared invariant momentum transfer $-t' = -t + t_0 \lesssim 0.6 \text{ GeV}^2$ where t_0 is the value of t for forward scattering. This analysis fixes the GPD H for quarks and gluons quite well. The other GPDs do practically not contribute to the cross sections and SDMEs at small skewness.

As mentioned spin effects in hard exclusive meson electroproduction will be briefly reviewed and their implications on the handbag approach and above all for the determination of the GPDs, discussed. In Sect. 2 the role of target spin asymmetries in meson electroproduction is examined. Sect. 3 is devoted to a discussion of the target spin asymmetries in pion electroproduction and Sect. 4 to those for vector mesons and the GPD E . Finally, in Sect. 5, a summary is presented.

observable	dominant interf. term	amplitudes	low t' behavior
$A_{UT}^{\sin(\phi-\phi_s)}$	LL	$\text{Im}[\mathcal{M}_{0-,0+}^* \mathcal{M}_{0+,0+}]$	$\propto \sqrt{-t'}$
$A_{UT}^{\sin(\phi_s)}$	LT	$\text{Im}[\mathcal{M}_{0-,++}^* \mathcal{M}_{0+,0+}]$	const.
$A_{UT}^{\sin(2\phi-\phi_s)}$	LT	$\text{Im}[\mathcal{M}_{0\mp,-+}^* \mathcal{M}_{0\pm,0+}]$	$\propto t'$
$A_{UT}^{\sin(\phi+\phi_s)}$	TT	$\text{Im}[\mathcal{M}_{0-,++}^* \mathcal{M}_{0+,++}]$	$\propto \sqrt{-t'}$
$A_{UT}^{\sin(2\phi+\phi_s)}$	TT	$\propto \sin \theta_\gamma$	$\propto t'$
$A_{UT}^{\sin(3\phi-\phi_s)}$	TT	$\text{Im}[\mathcal{M}_{0-, -+}^* \mathcal{M}_{0+, -+}]$	$\propto (-t')^{(3/2)}$
$A_{UL}^{\sin(\phi)}$	LT	$\text{Im}[\mathcal{M}_{0-,++}^* \mathcal{M}_{0-,0+}]$	$\propto \sqrt{-t'}$

Table 1: Features of the asymmetries for transversally and longitudinally polarized targets. The angle θ_γ describes the rotation in the lepton plane from the direction of the incoming lepton to the virtual photon one; it is very small.

target asymmetry E is needed. What is known about the GPD E will be recapitulated in Sect. 4.

3 Target spin asymmetries in π^+ production

In Ref. [13] electroproduction of positively charged pions has been investigated in the same handbag approach as applied to vector meson production [7]. To the asymptotically leading amplitudes for longitudinally polarized photons the GPDs $\tilde{H}$ and $\tilde{E}$ contribute in the isovector combination

$$\tilde{F}^{(3)} = \tilde{F}_v^u - \tilde{F}_v^d. \quad (2)$$

instead of H and E for vector mesons. In deviation to work performed in collinear approximation the full electromagnetic form factor of the pion as measured by the $F_\pi - 2$ collaboration [14] is naturally taken into account ¹ (see also the recent work by Bechler and Mueller [15]). The GPDs $\tilde{H}$ and $\tilde{E}$ are again constructed with the help of double distributions with the forward limit of $\tilde{H}$ being the polarized parton distributions while that of $\tilde{E}$ is parameterized analogously to the familiar parton distributions

$$\tilde{e}^u = -\tilde{e}^d = \tilde{N}_e x^{-0.48} (1-x)^5, \quad (3)$$

with $\tilde{N}_e$ fitted to experiment.

As is mentioned in Sect. 2 experiment requires a strong contribution from the helicity-non-flip amplitude $\mathcal{M}_{0-,++}$ which does not vanish in the forward direction. How can this amplitude be modeled in the frame work of the handbag approach? From the usual helicity non-flip GPDs $H, E, \dots$ one obtains a contribution to $\mathcal{M}_{0-,++}$ that vanishes $\propto t'$ if it is non-zero at all. However, there is a second set of GPDs, the helicity-flip or transversity

¹As compared to other work $\tilde{E}$ contains only the non-pole contribution.

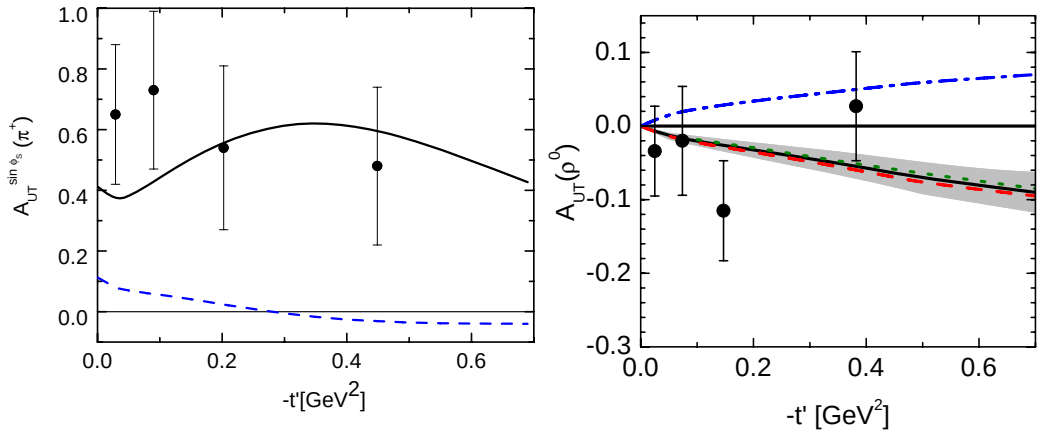


Figure 2: The $\sin \phi_s$ moment for a transversely polarized target at $Q^2 \simeq 2.45 \text{ GeV}^2$ and $W = 3.99 \text{ GeV}$ for π^+ production. The predictions from the handbag approach of Ref. [13] are shown as a solid line. The dashed line is obtained disregarding the twist-3 contribution. Data are taken from Ref. [9].

Figure 3: The asymmetry $A_{UT}^{\sin(\phi-\phi_s)}$ for ρ^0 production at $W = 5 \text{ GeV}$ and $Q^2 = 2 \text{ GeV}^2$. Data taken from Ref. [10]. The lines represent the results presented in Ref. [12]. For further notations see text and Ref. [12].

ones $H_T, E_T, \dots$ [16,17]. As inspection of Fig. 1 where the helicity configuration of the process is specified, reveals the proton-parton vertex is of non-flip nature in this case and, hence, is not forced to vanish in the forward direction by angular momentum conservation. One also sees from Fig. 1, that the helicity configuration of the subprocess is the same as for the full amplitude. Therefore, also the subprocess amplitude has not to vanish in the forward direction and so the full amplitude. The prize to pay is that quark and antiquark forming the pion have the same helicity. Therefore, the twist-3 pion wave function is needed instead of the familiar twist-2 one. The dynamical mechanism building up the amplitude $\mathcal{M}_{0-,++}$ is so of twist-3 order. This mechanism has been first proposed in Ref. [18] for photo- and electroproduction of mesons where $-t$ is considered as the large scale [19].

In Ref. [13] the twist-3 pion wave function is taken from Ref. [20] with the three-particle Fock component neglected. This wave function, still containing a pseudoscalar and a tensor component, is proportional to the parameter $\mu_\pi = m_\pi^2/(m_u + m_d) \simeq 2 \text{ GeV}$ at the scale of 2 GeV as a consequence of the divergency of the axial-vector current (m_u and m_d are current quark masses). It is further assumed that the dominant transversity GPD is H_T while the other three can be neglected. The forward limit of H_T^a is the transversity distribution $\delta^a(x)$ which has been determined in [21] in an analysis of data on the asymmetries in semi-inclusive electroproduction of charged pions measured with a transversely polarized target. Using these results for $\delta^a(x)$ the GPDs H_T^a have been modeled in a manner analogously to that of the other GPDs (see Eq. (5)) ².

It is shown in Ref. [13] that with the described model GPDs, the π^+ cross sections as measured by HERMES [22] are nicely fitted as well as the transverse target asymme-

²While the relative signs of δ^u and δ^d is fixed in the analysis performed in Ref. [21] the absolute sign is not. Here, in π^+ electroproduction a positive δ^u which goes along with a negative δ^d is required by the signs of the target asymmetries.

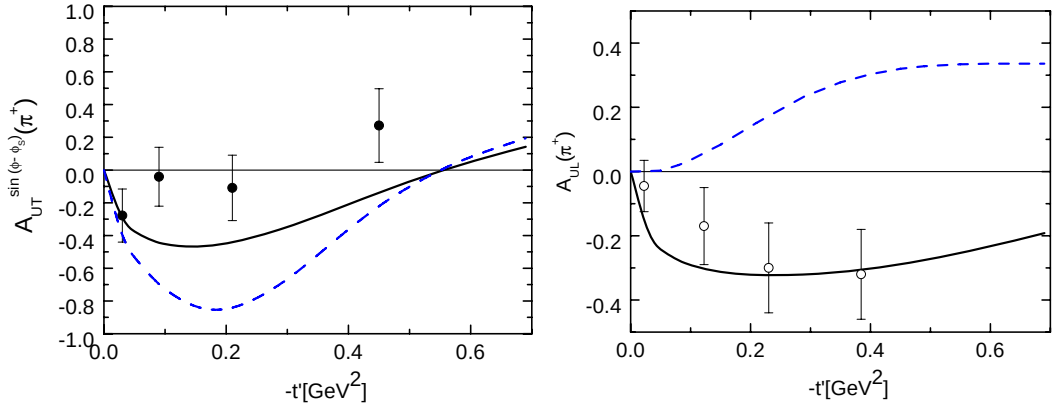


Figure 4: Left: Predictions for the $\sin(\phi - \phi_s)$ moment at $Q^2 = 2.45 \text{ GeV}^2$ and $W = 3.99 \text{ GeV}$ shown as solid lines [13]. The dashed line represents the longitudinal contribution to the $\sin(\phi - \phi_s)$ moment. Data are taken from [9].

Figure 5: Right: The asymmetry for a longitudinally polarized target at $Q^2 \simeq 2.4 \text{ GeV}^2$ and $W \simeq 4.1 \text{ GeV}$. The dashed line is obtained disregarding the twist-3 contribution. Data are taken from [23].

tries [9]. This can be seen for $A_{UT}^{\sin \phi_s}$ from Fig. 1. Also the $\sin(\phi - \phi_s)$ moment which is dominantly fed by an interference term of the two amplitudes for longitudinally polarized photons (see Tab. 1), is fairly well described as is obvious from Fig. 4. Very interesting is also the asymmetry for a longitudinally polarized target which is dominated by the interference term between $\mathcal{M}_{0-,++}$ which comprises the twist-3 effect, and the nucleon helicity-flip amplitude for $\gamma_L^* \rightarrow \pi$ transition, $\mathcal{M}_{0-,0+}$. Results for $A_{UL}^{\sin \phi}$ are displayed in Fig. 5 and compared to the data [23]. Also in this case good agreement between theory and experiment is to be noticed. In both the cases, $A_{UT}^{\sin \phi_s}$ and $A_{UL}^{\sin \phi}$, the prominent role of the twist-3 mechanism is clearly visible. Switching it off one obtains the dashed lines which are significantly at variance with experiment. In this case the transverse amplitudes are only fed by the pion-pole contribution. The other transverse target asymmetries quoted in Tab. 1 are predicted to be small in absolute value which is in agreement with experiment [9]. Thus, in summary, there is strong evidence for transversity in hard exclusive pion electroproduction. It can be regarded as a non-trivial result that the transversity distributions determined from data on inclusive pion production lead to a transversity GPD which is nicely in agreement with target asymmetries measured in exclusive pion electroproduction.

It is to be stressed that information on the amplitude $\mathcal{M}_{0-,++}$ can also be obtained from the asymmetries measured with a longitudinally polarized beam or with a longitudinally polarized beam and target. The first asymmetry, $A_{LU}^{\sin \phi}$, is dominated by the same interference term as $A_{UL}^{\sin \phi}$ but diluted by the factor $\sqrt{(1-\varepsilon)/(1+\varepsilon)}$. Also the second asymmetry, $A_{LL}^{\cos \phi}$, is dominated by the interference term $\mathcal{M}_{0-,++}^* \mathcal{M}_{0-,0+}$. However, in this case its real part occurs. For HERMES kinematics it is predicted to be rather large and positive at small $-t'$ and changes sign at $-t' \simeq 0.4 \text{ GeV}^2$ [13]. A measurement of these asymmetries would constitute a serious check of the twist-3 effect.

Although the main purpose of the work presented in Ref. [13] is focused on the analysis of the HERMES data one may also be interested in comparing this approach with the Jefferson Lab data on the cross sections [14]. With the GPDs $\tilde{H}$, $\tilde{E}$ and H_T in their present

form the agreement with these data is reasonable for the transverse cross section while the longitudinal one is somewhat too small. It is however to be stressed that the approach advocated for in Refs. [7, 13, 12] is designed for small skewness. At larger values of it the parameterizations of the GPDs are perhaps too simple and may require improvements. It is also important to realize that the GPDs are probed by the HERMES, COMPASS and HERA data only at x less than about 0.6. One may therefore change the GPDs at large x to some extent without changing the results for cross sections and asymmetries in the kinematical region of small skewness. For Jefferson Lab kinematics, on the other hand, such changes of the GPDs may matter.

4 The GPD E

In Ref. [24] the electromagnetic form factors of the proton and neutron have been utilized in order to determine the zero-skewness GPDs for valence quarks through the sum rules which for the case of the Pauli form factor, reads

$$F_2^{p(n)} = \int_0^1 dx \left[e_{u(d)} E_v^u(x, \xi = 0, t) + e_{d(u)} E_v^d(x, \xi = 0, t) \right]. \quad (4)$$

In order to determine the GPDs from the integral a parameterization of the GPD is required for which the ansatz

$$E_v^a(x, 0, t) = e_v^a(x) \exp \left[t(\alpha'_v \ln(1/x) + b_e^a) \right] \quad (5)$$

is made in a small $-t$ approximation [24]. The forward limit of E is parameterized analogously to that of the usual parton distributions:

$$e_v^a = N_a x^{\alpha_v(0)} (1-x)^{\beta_v^a}, \quad (6)$$

where $\alpha_v(0)$ ($\simeq 0.48$) is the intercept of a standard Regge trajectory and α'_v in Eq. (5) its slope. The normalization N_a is fixed from the moment

$$\kappa^a = \int dx E_v^a(x, \xi, t = 0), \quad (7)$$

where κ^a is the contribution of flavor- a quarks to the anomalous magnetic moments of the proton and neutron ($\kappa^u = 1.67$, $\kappa^d = -2.03$). A best fit to the data on the nucleon form factors provides the powers $\beta_v^u = 4$ and $\beta_v^d = 5.6$. However, other powers are not excluded in the 2004 analysis presented in [24]; the most extreme set of powers, still in agreement with the form factor data, is $\beta_v^u = 10$ and $\beta_v^d = 5$. The analysis performed in [24] should be repeated since new form factor data are available from Jefferson Lab, e.g. G_E^n and G_M^n are now measured up to $Q^2 = 3.5$ and 5.0 GeV², respectively [25, 26]. These new data seem to favor $\beta_v^u < \beta_v^d$. The zero-skewness GPDs E_v are used as input to a double distribution from which the valence quark GPDs for non-zero skewness are constructed [12].

In [12], following Diehl and Kugler [27], E for gluons and sea quarks has been estimated from positivity bounds and a sum rule for the second moments of E which follows from a combination of Ji's sum rule [28] and the momentum sum rule of deep inelastic

lepton-nucleon scattering. It turned out that the valence quark contribution to that sum rule is very small, in particular if $\beta_v^u < \beta_v^d$, with the consequence of an almost exact cancellation of the gluon and sea quark moments. The GPDs E^g and E^{sea} are parameterized analogously to E_v , see Eqs. (5), (6). The normalization of E^{sea} is fixed by assuming that an appropriate positivity bound (see Refs. [29, 30]) is saturated while that of E^g is determined from the sum rule. Several variants of E have been exploited in [12] in a calculation of $A_{UT}^{\sin(\phi-\phi_s)}$ within the handbag approach. The results for a few variants are compared to the HERMES data on ρ^0 production [10] in Fig. 2. Agreement between theory and experiment is to be noted. Similar agreement is obtained for the preliminary COMPASS data [11]. Combining both the experiments a negative value of $A_{UT}^{\sin(\phi-\phi_s)}$ for ρ^0 production is favored in agreement with the theoretical results obtained in [12], only the extreme variant $\beta_v^u = 10$ and $\beta_v^d = 5$ (dashed-dotted line in Fig. 2) seems to be ruled out. In [12] predictions for ω , ρ^+ , K^{*0} and ϕ productions are also given. Their comparison with forthcoming data from HERMES and COMPASS may lead to a fair determination of the GPD E .

With E at hand one may exploit Ji's sum rule for the parton angular momenta. At zero skewness the sum rule reads

$$\langle J^a \rangle = \frac{1}{2} [q_{20}^a + e_{20}^a], \quad \langle J^g \rangle = \frac{1}{2} [q_{20}^g + e_{20}^g]. \quad (8)$$

From a variant with $\beta_v^u = 4$, $\beta_v^d = 5.6$ and neglected E^g and E^{sea} (solid line in Fig. 2) for instance one obtains

$$\langle J^u \rangle = 0.250, \quad \langle J^d \rangle = 0.020, \quad \langle J^s \rangle = 0.015, \quad \langle J^g \rangle = 0.214, \quad (9)$$

at the scale of 4 GeV². The angular momenta sum up to $\simeq 1/2$, the spin of the proton. A very characteristic stable pattern is obtained in Ref. [12]: For all variants investigated, J^u and J^g are large while the other two angular momenta are very small. The angular momenta of the valence quarks are $\langle J_v^u \rangle = 0.222$ and $\langle J_v^d \rangle = -0.015$. These values are identical to the results quoted in Ref. [24] (for variant 1). They are in agreement with a recent lattice result [31].

5 Summary

Recent measurements of single spin asymmetries in hard meson electroproduction has been reviewed. The data clearly show that a leading-twist calculation of meson electroproduction within the handbag approach is insufficient. They demand higher-twist and/or power corrections which manifest themselves through substantial contributions from $\gamma_T^* \rightarrow V, P$ transitions.

A most striking effect is the target asymmetry $A_{UT}^{\sin \phi_s}$ in π^+ electroproduction. The interpretation of this effect requires a large contribution from the helicity non-flip amplitude $\mathcal{M}_{0-,++}$. Within the handbag approach such a contribution is generated by the helicity-flip or transversity GPDs in combination with a twist-3 pion wave function [13]. This explanation establishes an interesting connection to transversity parton distributions measured in inclusive processes. Further studies of transversity in exclusive reactions are certainly demanded. For instance, data on the asymmetries obtained with a longitudinally polarized beam and with likewise polarized beam and target would be very helpful

in settling this dynamical issue. Good data on π^0 electroproduction would also be highly welcome. They would not only allow for an additional test of the twist-3 mechanism but also give the opportunity to verify the model GPDs $\tilde{H}$ and $\tilde{E}$ as used in Ref. [13].

One may wonder whether the twist-3 mechanism does not apply to vector-meson electroproduction as well and offers an explanation of the experimentally observed $\gamma_T^* \rightarrow V_L$ transitions seen for instance in the SDME r_{00}^{05} . It however turned out that this effect is too small in comparison to the data. The reason is that instead of the parameter μ_π the mass of the vector meson sets the scale of the twist-3 effect. This amounts to a reduction by about a factor of three. Further suppression comes from the unfavorable flavor combination of H_T occurring for uncharged vector mesons, e.g. $e_u H_T^u - e_d H_T^d$ for ρ^0 production instead of $H_T^u - H_T^d$ for π^+ production. Perhaps the gluonic GPD H_T^g may lead to a larger effect.

From the small value of the ratio of the longitudinal and transverse electroproduction cross sections for ρ^0 and ϕ mesons it also clear that the transitions from transversely polarized virtual photons to likewise polarized vector mesons are large too. In the handbag approach advocated in [7] such transitions are also well described. The infrared divergence occuring in collinear approximation is regularized by the quark transverse momenta in the modified perturbative approach.

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SPIN CORRELATIONS OF THE ELECTRON AND POSITRON IN THE TWO-PHOTON PROCESS $\gamma\gamma \rightarrow e^+e^-$

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Abstract

The spin structure of the process of electron-positron pair production by two photons $\gamma\gamma \rightarrow e^+e^-$ is theoretically investigated. It is shown that, if the primary photons are unpolarized, the final electron and positron are unpolarized as well but their spins are strongly correlated. Explicit expressions for the components of the correlation tensor of the final (e^+e^-) system are derived, and the relative fractions of singlet and triplet states of the (e^+e^-) pair are found. It is demonstrated that in the process $\gamma\gamma \rightarrow e^+e^-$ one of the incoherence inequalities of the Bell type for the correlation tensor components is always violated and, thus, spin correlations of the electron and positron in this process have the strongly pronounced quantum character.

1. Let us consider the process of electron-positron pair production by two photons, $\gamma\gamma \rightarrow e^+e^-$. The spin state of the electron-positron system is described, in the general case, by the two-particle spin density matrix:

$$\hat{\rho}^{(e^-e^+)} = \frac{1}{4} \left[\hat{I}^{(e^-)} \otimes \hat{I}^{(e^+)} + \hat{I}^{(e^-)} \otimes (\hat{\sigma}^{(e^+)} \mathbf{P}^{(e^+)}) + (\hat{\sigma}^{(e^-)} \mathbf{P}^{(e^-)}) \otimes \hat{I}^{(e^+)} + \sum_{i=1}^3 \sum_{k=1}^3 T_{ik} \hat{\sigma}_i^{(e^-)} \otimes \hat{\sigma}_k^{(e^+)} \right], \quad (1)$$

where $\hat{I}$ is the two-row unit matrix, $\mathbf{P}^{(e^-)}$ and $\mathbf{P}^{(e^+)}$ are the polarization vectors of the electron and positron, respectively, T_{ik} – components of the correlation tensor ($T_{ik} = \langle \hat{\sigma}_i^{(e^-)} \otimes \hat{\sigma}_k^{(e^+)} \rangle$, $i, k = \{1, 2, 3\} = \{x, y, z\}$). In the absence of correlations, we have: $T_{ik} = P_i^{(e^-)} P_k^{(e^+)}$.

The process $\gamma\gamma \rightarrow e^+e^-$ is described by two well-known Feynman diagrams [1]. Within the first nonvanishing approximation over the electromagnetic constant (Born approximation), in case of unpolarized primary photons the final electron and positron prove to be unpolarized as well, but their spins are correlated.

Thus, in the above formula for the spin density matrix $\hat{\rho}^{(e^+e^-)}$ (1)

$$\mathbf{P}^{(e^-)} = \mathbf{P}^{(e^+)} = 0.$$

The components of the correlation tensor of the electron-positron pair, generated in the interaction of unpolarized γ quanta, may be calculated by applying the results of paper [2]. Finally we obtain the following expressions:

$$T_{zz} = 1 - \frac{2(1 - \beta^2)[\beta^2(1 + \sin^2 \theta) + 1]}{1 + 2\beta^2 \sin^2 \theta - \beta^4 - \beta^4 \sin^4 \theta}, \quad (2)$$

$$T_{yy} = \frac{(1 - \beta^2) [\beta^2 (1 + \sin^2 \theta) - 1] - \beta^2 \sin^4 \theta}{1 + 2 \beta^2 \sin^2 \theta - \beta^4 - \beta^4 \sin^4 \theta}, \quad (3)$$

$$T_{xx} = \frac{(1 - \beta^2) [\beta^2 (1 + \sin^2 \theta) - 1] + \beta^2 \sin^4 \theta}{1 + 2 \beta^2 \sin^2 \theta - \beta^4 - \beta^4 \sin^4 \theta}. \quad (4)$$

Here the axis z is aligned along the positron momentum in the c.m. frame, the axis x lies in the reaction plane and the axis y is directed along the normal to the reaction plane ; $\beta = \frac{v}{c}$, $\mathbf{v}$ is the positron velocity in the c.m. frame ; $1 - \beta^2 = \frac{m_e c^2}{E_+}$, where E_+ is the positron (or electron) energy in the c.m. frame ; θ is the angle between the positron momentum and the momentum of one of the photons in the c.m. frame.

Meantime, the differential cross section of the process $\gamma\gamma \rightarrow e^+e^-$ in the c.m. frame has the following form [1,2]:

$$\frac{d\sigma}{d\Omega} = r_0^2 \frac{1 - \beta^2}{4} \beta \left[\frac{1 + 2 \beta^2 \sin^2 \theta - \beta^4 - \beta^4 \sin^4 \theta}{(1 - \beta^2 \cos^2 \theta)^2} \right], \quad (5)$$

where $r_0 = \frac{e^2}{m_e c^2}$.

The “trace” of the correlation tensor of the final (e^+e^-) pair is determined by the formula:

$$T = T_{xx} + T_{yy} + T_{zz} = 1 - \frac{4(1 - \beta^2)}{1 + 2 \beta^2 \sin^2 \theta - \beta^4 - \beta^4 \sin^4 \theta}. \quad (6)$$

In doing so, the relative fraction of the triplet states is as follows [3]:

$$W_t = \frac{T + 3}{4} = 1 - \frac{1 - \beta^2}{1 + 2 \beta^2 \sin^2 \theta - \beta^4 - \beta^4 \sin^4 \theta}, \quad (7)$$

and the relative fraction of the singlet state (total spin $S = 0$) equals

$$W_s = \frac{1 - T}{4} = 1 - W_t = \frac{1 - \beta^2}{1 + 2 \beta^2 \sin^2 \theta - \beta^4 - \beta^4 \sin^4 \theta}. \quad (8)$$

At $\beta \ll 1$ we have: $W_t \approx 0$, $W_s \approx 1$, i.e. the (e^+e^-) pair is generated in the singlet state only. On the contrary, under the conditions $\beta \approx 1$, and $2 \sin^2 \theta - \sin^4 \theta > 1 - \beta^2$ we obtain: $W_t \approx 1$, $W_s \approx 0$ (almost pure triplet state of the pair).

2. It should be noted that the charge parity C of the system of two photons is positive and it is conserved in the process $\gamma\gamma \rightarrow e^+e^-$. This means that electron-positron pairs generated in this process have the positive charge parity ($C = 1$). According to the well-known formula for charge parity of a fermion-antifermion system, $C = (-1)^{L+S}$, where L is the orbital momentum and S is the total spin. Thus, in the process $\gamma\gamma \rightarrow e^+e^-$, singlet electron-positron pairs ($S = 0$) are produced in the states with even orbital and total angular momenta ($L = J$) and negative space parity $P = (-1)^{L+1}$. Due to the conservation of total angular momentum and space parity in the process $\gamma\gamma \rightarrow e^+e^-$, the contribution into generation of singlet (e^+e^-) pairs is provided only by states of two photons with even total angular momenta and negative space parity, which, according to Bose symmetry, are antisymmetric over the polarizations of two photons (they correspond to the total spin of

the two-photon system equaling 1). If both the photons are unpolarized, then the relative fraction of such photon pairs amounts to 1/4.

Meantime, triplet electron-positron pairs ($S = 1$) are produced, in the process $\gamma\gamma \rightarrow e^+e^-$, in the states with odd orbital momenta and positive space parity. In doing so, total angular momenta may take both even and odd values. The contribution into generation of triplet (e^+e^-) pairs is provided only by states of two photons with positive space parity, being symmetric over polarizations, which correspond to the total spins of the two-photon system equaling 0 and 2. If both the photons are unpolarized, then the relative fraction of such photon pairs amounts to 3/4.

3. Now let us consider the particular cases $\theta = 0$ and $\theta = \pi$. In accordance with the formula for the differential cross section $\frac{d\sigma}{d\Omega}$ (5), here we obtain :

$$\frac{d\sigma}{d\Omega} = \frac{1}{4} r_0^2 \beta (1 + \beta^2). \quad (9)$$

In the ultrarelativistic limit formula (9) turns to $\frac{d\sigma}{d\Omega} = \frac{r_0^2}{2}$.

According to the general expressions (2)–(4) for the correlation tensor components, at $\theta = 0$ and $\theta = \pi$ we have:

$$T_{zz} = 1 - \frac{2(1 + \beta^2)(1 - \beta^2)}{1 - \beta^4} = -1; \quad T_{xx} = T_{yy} = -\frac{1 - \beta^2}{1 + \beta^2}. \quad (10)$$

In doing so, the “trace” of the correlation tensor (6) takes the value

$$T = 1 - \frac{4}{1 + \beta^2} = -\frac{3 - \beta^2}{1 + \beta^2}, \quad (11)$$

and the relative fractions of the triplet states W_t (7) and the singlet state W_s (8) amount to

$$W_t = \frac{T + 3}{4} = \frac{\beta^2}{1 + \beta^2}, \quad W_s = \frac{1 - T}{4} = \frac{1}{1 + \beta^2}. \quad (12)$$

At nonrelativistic velocities $W_t \approx 0$, $W_s \approx 1$, in accordance with the general case ; meantime, in the ultrarelativistic limit ($\beta \rightarrow 1$) we have: $W_t = W_s = \frac{1}{2}$.

Let us remark that in the process $\gamma\gamma \rightarrow e^+e^-$ we observe the violation of the “incoherence” inequalities, established previously at the classical level, according to which for incoherent “classical” mixtures of factorizable two-particle spin states the sum of any two (or three) diagonal components of the correlation tensor cannot exceed unity [3]. Indeed, at $\theta = 0$ and $\theta = \pi$, in particular, we obtain:

$$|T_{zz} + T_{xx}| = |T_{zz} + T_{yy}| = \frac{2}{1 + \beta^2} > 1, \quad (13)$$

since $\beta < 1$.

Thus, the spin correlations of the final electron and positron in the considered process have the strongly pronounced quantum character.

Finally, it should be noted that analogous results hold as well for the processes of generation of a muon pair and a τ lepton pair by two photons : $\gamma\gamma \rightarrow \mu^+\mu^-$, $\gamma\gamma \rightarrow \tau^+\tau^-$, which become possible at considerably higher energies.

4. Summary

1. The theoretical investigation of spin structure of the process $\gamma\gamma \rightarrow e^+e^-$ is performed. It is shown that, if the primary photons are unpolarized, the final electron and positron are not polarized as well but their spins are strongly correlated.

2. Explicit expressions for the components of correlation tensor of the final (e^+e^-) system are derived, and the relative fractions of singlet and triplet states of the (e^+e^-) pair are found.

3. It is demonstrated that in the process $\gamma\gamma \rightarrow e^+e^-$ the “incoherence” inequalities for the correlation tensor components may be violated.

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ANOTHER NEW TRACE FORMULA FOR THE COMPUTATION OF THE FERMIONIC LINE

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Abstract

We reexpress the fermion line as a trace over spinor indices but using the spin formalism. Our spin formulation is appropriate to the computation of dipole moments, electric or magnetic for which we have developed such formalism. Another formalism based on helicity is compared to our. We just confirm here that both formalisms are equivalent. We test the power of our trace formula by analytically computing the quark dipole magnetic moment in few lines compared to the direct computation using spinor components we made in a previous work. The trace formulation has the advantage of making the amplitudes or even the squared amplitudes easily computable either analytically or symbolically.

1 The amplitude as a trace.

Cross- sections and lifetimes being the squared modulus of amplitudes are traces

$$Tr(\bar{f} \frac{\not{p}' + m}{2m} \frac{1 + \gamma_5 \not{s}'}{2} f \frac{\not{p} + m}{2m} \frac{1 + \gamma_5 \not{s}}{2}) \quad (1)$$

Observables such as form factors, magnetic (electric) dipole moments etc, are probability amplitudes and are usually not expressed as traces (momenta are not paired to form projectors). In this paper we propose to rewrite the amplitude as a trace

$$\begin{aligned} \bar{u}_{\alpha'}(p', s') f_{\alpha'\alpha}(s', s) u_{\alpha}(p, s) &= f_{\alpha'\alpha}(s', s) u_{\alpha}(p, s) \bar{u}_{\alpha'}(p', s') \\ &= Tr(f \rho), \quad \rho_{\alpha\alpha'} = u_{\alpha}(p, s) \bar{u}_{\alpha'}(p', s') \end{aligned} \quad (2)$$

The amplitude is then rewritten in term of the generalized spin density matrix ρ . To find ρ we directly relate it to its form in the rest frame $\rho|_{rf}$; calculate the latter then boost it again to the initial frame.

$$\begin{aligned} u(k, s) &= \exp(-\frac{\omega}{2} \gamma_0 \frac{\vec{\gamma} \cdot \vec{k}}{|\vec{k}|}) \begin{pmatrix} \chi_s \\ 0 \end{pmatrix}, \quad \bar{u}(k, s) = \begin{pmatrix} \chi_s^\dagger & 0 \end{pmatrix} \exp(-\frac{\omega}{2} \gamma_0 \frac{\vec{\gamma} \cdot \vec{k}}{|\vec{k}|}) \gamma_0 \\ \rho|_{rf} &= \chi_s(\vec{\zeta}) \tilde{\chi}_{s'}^\dagger(\vec{\zeta}') = (\frac{1+s\vec{\zeta} \cdot \vec{\sigma}}{2}) (\frac{1+s'\vec{\zeta}' \cdot \vec{\sigma}}{2}) \end{aligned} \quad (3)$$

After lengthy but straightforward calculations we get the expression for ρ :

$$\begin{aligned} \rho(k, s, k', s') &= (\frac{\not{k} + m}{2m}) (\frac{1 + \gamma^5 \not{s}}{2}) \Re (\frac{\not{k}' + m}{2m}) (\frac{1 + \gamma^5 \not{s}'}{2}) \\ \Re &= \sqrt{\frac{mm'}{(k_0 + m)(k'_0 + m')}} (1 + \gamma_0) \end{aligned} \quad (4)$$

This expression can further be arranged in a covariant form which in addition does not depend on specific representations of the Dirac matrices. To this end we introduce the additional time-like vector $p = (1, \vec{0})$ and write the final form of $\Re$ as.

$$\Re = \sqrt{\frac{mm'}{(pk+m)(pk'+m')}}(p+1)$$

Our trace formulation of the amplitude in the spin framework is well adapted to the computation of the dipole magnetic or electric moments of the quarks inside baryons compared to the calculation of the same quantities using the explicit spinor components [1][2]. There is however a similar trace formulation[3] but in the helicity framework. Such a formulation although competitive to our is not adequate to dipole moment calculations as the latter ones are interpreted at the end in terms of longitudinal and transverse spin structure functions. We can of course retrieve all results of the helicity formalism as particular cases of the spin formalism.

2 Application: The convection current part of the quark dipole magnetic moment.

By using the Gordon decomposition we divide the dipole magnetic moment expression into two terms, one is the convection current part and the other is the spin part. In the application of our trace formalism we only compute the convection part for illustration of the power of the formulation. The convection part has the form:

$$\int \left[Tr \vec{\nabla} \rho \right]_{\vec{q}=0} \times \vec{k} \frac{d^3 k}{(2\pi)^3} \quad (5)$$

In the above calculation we use the identities $\vec{\nabla}_k k_0 = \frac{\vec{k}}{k_0}$ and $\vec{\nabla}_k |\vec{k}| = \frac{\vec{k}}{|\vec{k}|}$ to eliminate all differentiation leading to terms proportional to $\vec{k}$ as we have to take the cross product with $\vec{k}$ in (5). In particular any differentiation of the spin variable is eliminated as s is function of $\vec{\xi}$ and of the boost parameter $\omega = -\tanh^{-1}(\frac{|\vec{k}|}{k_0})$, idem for the primed spin. With this remark only the variable $\vec{k}$ and $\vec{k}'$ are differentiated and we get

$$\begin{aligned} & -\frac{1}{2} \frac{1}{(k_0+m)} \int (\vec{k} \times Tr \left[(\frac{\vec{k}+m}{2m}) (\frac{1+\gamma_5 \not{s}}{2}) (\gamma_0 \vec{\gamma}) \right]) \frac{d^3 k}{(2\pi)^3} \\ & = \frac{1}{8m} \frac{1}{(k_0+m)} \int (\vec{k} \times Tr \gamma_5 \not{k} \not{s} \gamma_0 \vec{\gamma}) \frac{d^3 k}{(2\pi)^3} \\ & = -\frac{1}{2m} \frac{1}{(k_0+m)} \int |\vec{k}|^2 \vec{s}_\perp \frac{d^3 k}{(2\pi)^3} \\ & = -\frac{x}{(x+1)} \int \frac{|\vec{k}|^2}{2m^2} \vec{s}_\perp \frac{d^3 k}{(2\pi)^3} \end{aligned} \quad (6)$$

In the equation above, we approximate the factor k_0 by the average value of the relativistic quark energy inside the nucleon and keep the notation k_0 to designate the average value. Inserting the missing factors $\mu = \frac{Q}{2m}$ and $\sqrt{\frac{m}{k_0}}$ for each spinor u in (5) to complete the definition of the magnetic moment and define the parameter $x = \frac{m}{k_0}$ we get the convection current part in terms of the longitudinal spin $\vec{S}$ and the transverse spin $\vec{\delta}$:

$$\begin{aligned} & \frac{x\mu}{(x+1)} (\vec{S} - \frac{\vec{\delta}}{2x}) \\ & \mu = \frac{Q}{2m} \end{aligned} \quad (7)$$

Compare the gain of time and effort in obtaining this contribution with the lengthy method that used the components of the Dirac spinor in one of our previous work [1]

3 Conclusion.

Dipole moments, electric or magnetic are naturally expressed in terms of the spin(longitudinal and transverse) of the quarks inside the parent baryons. Our trace formulation of the amplitude in the spin framework (as opposed to helicity formulation) is appropriate to such computations. On the other hand our formulation has the advantage of making the amplitudes or even the squared amplitudes easily computable either analytically or symbolically.

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ANALYTIC PERTURBATION THEORY AND PROTON SPIN STRUCTURE FUNCTION g_1^p

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Abstract

The interplay between higher orders of the perturbative QCD (pQCD) expansion and higher twist contributions in the analysis of recent Jefferson Lab (JLab) data on the lowest moment of the spin-dependent proton $\Gamma_1^p(Q^2)$ at $0.05 < Q^2 < 3 \text{ GeV}^2$ is studied. We demonstrate that the values of the higher twist coefficients $\mu_{2k}^{p,n}$ extracted from the data by using the singularity-free analytic perturbation theory provide a better convergence of the higher-twist series than with the standard pQCD. From the high-precision proton data, we extract the value of the singlet axial charge $a_0(1 \text{ GeV}^2) = 0.33 \pm 0.05$. We observe a slow Q^2 -dependence of fitted values of the twist coefficient μ_4 and a_0 when going to lower energy scales, which can be explained by the RG evolution of $\mu_4(Q^2)$ and $a_0(Q^2)$. As the main result, a good quantitative description of all the JLab data sets down to $Q \simeq 350 \text{ MeV}$ is achieved.

The spin structure of the nucleon remains to be the essential problem of nonperturbative QCD and hadronic physics [1]. Higher-twist parameters (known also as the colour polarizabilities) are important ingredients of the nucleon spin structure. Their extraction from experimental studies is relatively complicated as they are most pronounced at low momentum transfer $Q^2 \lesssim 1 \text{ GeV}^2$. Although very accurate Jefferson Lab (JLab) data on the lowest moments of the spin-dependent proton and neutron structure functions $\Gamma_1^{p,n}(Q^2)$ in the range $0.05 < Q^2 < 3 \text{ GeV}^2$ are now available [2], higher twist contributions are shadowed by Landau singularities of QCD coupling. As it was shown in Ref. [3], this problem may be solved by the use of singularity-free couplings which allowed a quite accurate extraction of higher twist and a fairly well description of data down to rather low $Q \gtrsim \Lambda_{QCD}$.

In this report, we extend and generalize the analysis started in Ref. [3] by considering the singlet channel involving the $\Gamma_1^p(Q^2)$ provided by the JLab data. This allows, in particular, to determine the singlet axial charge a_0 , which in the quark-parton model is identified with the total spin carried by quarks in the proton. For this purpose, we perform the global analysis of the JLab data on $\Gamma_1^p(Q^2)$ using the advantages of the Analytic Perturbation Theory (APT) [4], and extract a_0 , as well as the coefficient μ_4^p of the $1/Q^2$ subleading twist-4 term, which contains information on quark-gluon correlations in nucleons.

In the following we use the total expression for the perturbative part of $\Gamma_1^{p,n}(Q^2)$ including the higher twist contributions

$$\Gamma_1^{p,n}(Q^2) = \frac{1}{12} \left[\left(\pm a_3 + \frac{1}{3} a_8 \right) E_{NS}(Q^2) + \frac{4}{3} a_0 E_S(Q^2) \right] + \sum_{i=2}^{\infty} \frac{\mu_{2i}^{p,n}(Q^2)}{Q^{2i-2}}, \quad (1)$$

where the triplet and octet axial charges are $a_3 \equiv g_A = 1.267 \pm 0.004$ and $a_8 = 0.585 \pm 0.025$, respectively, and E_S and E_{NS} are the singlet and nonsinglet Wilson coefficients, respectively, calculated as series in powers of α_s :

$$E_{NS}(Q^2) = 1 - \frac{\alpha_s}{\pi} - 3.558 \left(\frac{\alpha_s}{\pi} \right)^2 - 20.215 \left(\frac{\alpha_s}{\pi} \right)^3 - O(\alpha_s^4), \quad (2)$$

$$E_S(Q^2) = 1 - \frac{\alpha_s}{\pi} - 1.096 \left(\frac{\alpha_s}{\pi} \right)^2 - O(\alpha_s^3). \quad (3)$$

The infrared behavior of the strong coupling is crucial for the extraction of the nonperturbative information from the low energy data. The moments of the structure functions are analytic functions in the complex Q^2 -plane with a cut along the negative real axis. In contrast to the standard perturbation theory, the APT method supports required analytic properties of the nucleon spin sum rules $\Gamma_1^{p,n}(Q^2)$ (for details, see Ref. [5]).

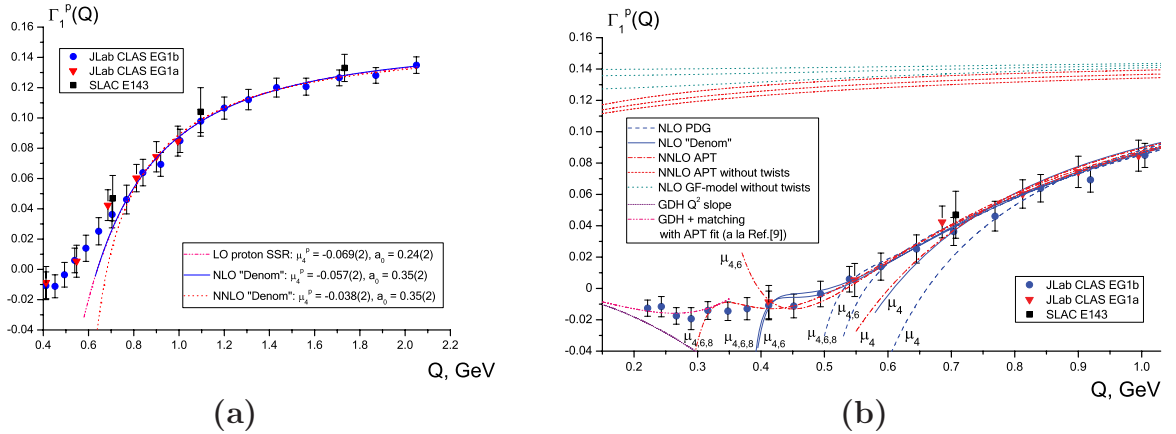


Figure 1: **(a)** Best fits of JLab and SLAC data on proton spin sum rule $\Gamma_1^p(Q^2)$ calculated using the PT in various loop orders with fixed $Q_{min} = 0.8$ GeV. **(b)** Best 1,2,3-parametric fits of the JLab and SLAC data on proton spin sum rule $\Gamma_1^p(Q^2)$ calculated with different models of running coupling.

In Fig. 1a, we show fits of proton spin sum rule data in different orders of perturbation theory taking only into account the μ_4 -term. One can see there that the higher-loop contributions are effectively “absorbed” into the value of μ_4 which decreases in magnitude with increasing loop order while all the fitting curves are very close to each other. This exhibits a kind of duality between higher orders of PT and HT terms, moving the pQCD frontier between the PT and HT contribution to lower Q values in both nonsinglet and singlet channels. At the same time, the value of a_0 is quite stable in higher loop orders.

In Fig. 1b, we show best fits of the combined data set for the function $\Gamma_1^p(Q^2)$ (the data uncertainties are statistical only) in the standard PT and in the APT approaches. One can see that the perturbative parts of $\Gamma_1^p(Q^2)$ calculated in the APT and in the so-called “glueball-freezing” model [6] are close to each other down to $Q \sim \Lambda$.

We test a separation of perturbative and nonperturbative physics and perform a systematic comparison of the extracted values of the higher twist terms in different versions of perturbation theory. In Table 1, we present the combined fit results of the proton $\Gamma_1^p(Q^2)$ data (elastic contribution excluded) in APT and conventional PT. One can see there is noticeable sensitivity of the extracted a_0 and μ_4 w.r.t. the minimal fitting scale Q_{min}^2 variations, which may be (at least, partially) compensated by their RG $\log Q^2$ -evolution. For completeness, we included in Table 1 APT fits for $a_0(Q_0^2)$ and $\mu_4(Q_0^2)$ taking into

account their RG evolution. As a result, we extract the value of the singlet axial charge $a_0(1 \text{ GeV}^2) = 0.33 \pm 0.05$. This value is very close to the corresponding COMPASS 0.35 ± 0.06 [7] and HERMES 0.35 ± 0.06 [8] results.

Table 1: Combined fit results of the proton $\Gamma_1^p(Q^2)$ data (elastic contribution excluded). APT fit results a_0 and $\mu_{4,6,8}^{APT}$ (at the scale $Q_0^2 = 1 \text{ GeV}^2$) are given without and with taking into account the RG Q^2 -evolution of $a_0(Q^2)$ and $\mu_4^{APT}(Q^2)$. The minimal borders of fitting domains in Q^2 are settled from the *ad hoc* restriction $\chi^2 \leq 1$ and monotonous behavior of the resulting fitted curves.

Method	$Q_{min}^2, \text{ GeV}^2$	a_0	μ_4/M^2	μ_6/M^4	μ_8/M^6
PT	0.59	0.33(3)	-0.050(4)	0	0
	0.35	0.43(5)	-0.087(9)	0.024(5)	0
	0.29	0.37(5)	-0.060(15)	-0.001(8)	0.006(5)
APT (no evolution)	0.47	0.35(4)	-0.054(4)	0	0
	0.17	0.39(3)	-0.069(4)	0.0081(8)	0
	0.10	0.43(3)	-0.078(4)	0.0132(9)	-0.0007(5)
APT (with evolution)	0.47	0.33(4)	-0.051(4)	0	0
	0.17	0.31(3)	-0.059(4)	0.0098(8)	0
	0.10	0.32(4)	-0.065(4)	0.0146(9)	-0.0006(5)

In order to see how the Q_{min}^2 scale and fit results for the μ -terms change with varying Λ_{QCD} , we have performed three different NLO fits with $\Lambda_{QCD} = 300, 400, 500 \text{ MeV}$. It turns out that the term μ_4 is quite sensitive to the Landau singularity position, and its value noticeably increases with increasing Λ_{QCD} . The APT is free of such a problem thus providing a reliable tool of investigating the behavior of higher twist terms extracted directly from the low-energy data.

In the APT approach the convergence of both the higher orders and higher twist series is much better. In both the nonsinglet and singlet case, while the twist-4 term happened to be larger in magnitude in the APT than in the conventional PT, the subsequent terms are essentially smaller and quickly decreasing (as the APT absorbs some part of nonperturbative dynamics described by higher twists). This is the main reason of the shift of the pQCD frontier to lower Q values. A satisfactory description of the proton spin sum rule data down to $Q \sim \Lambda_{QCD} \simeq 350 \text{ MeV}$ was achieved (see Fig. 1b). In a sense, this could be natural if the main reason of such a success was the disappearance of unphysical singularities. Note that the data at very low $Q \sim \Lambda_{QCD}$ are usually dropped from the analysis of a_0 and higher-twist term in the standard PT analysis because of Landau singularities. The compatibility of our results for a_0 and previous results [7, 8] demonstrates the universality of the nucleon spin structure at large and low Q^2 scales.

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TRANSVERSE MOMENTUM DEPENDENT PARTON DISTRIBUTIONS AND AZIMUTHAL ASYMMETRIES IN LIGHT-CONE QUARK MODELS

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Abstract

We review the information on the spin and orbital angular momentum structure of the nucleon encoded in the T-even transverse momentum dependent parton distributions within light-cone quark models. Model results for azimuthal spin asymmetries in semi-inclusive lepton-nucleon deep-inelastic scattering are discussed, showing a good agreement with available experimental data and providing predictions to be further tested by future CLAS, COMPASS and HERMES data.

1 TMDs and Light-Cone CQMs

A convenient framework for the analysis of hadronic states is quantization on the light-cone. The proton state, for example, can be represented as a superposition of light-cone wave functions (LCWFs), one for each of the Fock components (qqq), ($qqq\bar{q}$), ... of the nucleon state. This light-cone representation has a number of simplifying properties [1]. In particular it allows one to describe the hadronic matrix elements which parametrize the soft-contribution in inclusive and exclusive reactions in terms of overlap of LCWFs with different parton configurations. In principle, there is an infinite number of LCWFs in such an expansion. However, there are many situations where one can confine the analysis to the contribution of the Fock components with a few partons. For example, light-cone models limited to the minimal Fock-space configuration of valence quarks are able to reproduce the main features of the hadron electromagnetic form factors [2] as well to account for the behaviour of the hadron structure functions in deeply inelastic processes at large values of the Bjorken variable x [3, 4].

Here the LCWFs of constituent quark models (CQMs) will be used to describe transverse momentum dependent parton distributions (TMDs) which are a natural extension of standard parton distributions from one to three-dimensions in momentum space. In particular, to disentangle the spin-spin and spin-orbit quark correlations encoded in the different TMDs, we expand the three-quark LCWF in a basis of eigenstates of orbital angular momentum. In the light-cone gauge $A^+ = 0$, such an expansion involves six independent amplitudes corresponding to the different combinations of quark helicity and orbital angular momentum. Explicit expressions for the light-cone amplitudes have been obtained in Ref. [5] representing the light-cone spinors of the quarks through the unitary

Melosh rotations which boost the rest-frame spin into the light-cone. Furthermore, assuming SU(6) symmetry, the light-cone amplitudes have a particularly simple structure, with the spin and isospin dependence factorized from a momentum-dependent function which is spherically symmetric. Under this assumption the orbital angular momentum content of the wave function is fully generated by the Melosh rotations and therefore matches the analytical structure expected from model-independent arguments [6]. The model dependence enters the choice of the momentum-dependent part of the LCWF, as obtained, for example, from the eigenvalue equation of the Hamiltonian with a specific potential model for the bound state of the three quarks. Using a more phenomenological description, we choose this part by assuming a specific functional form with parameters fitted to hadronic structure constants. This is the strategy adopted also in Ref. [7] through a fit of the LCWF to the anomalous magnetic moments of the nucleon. The same wave function was also used to predict many other hadronic properties [8], providing a good description of available experimental data, and being able to capture the main features of hadronic structure functions, like parton distributions [4], generalized parton distributions (GPDs) [3] and TMDs [5].

The eight leading twist TMDs, f_1 , $f_{1T}^\perp$, g_1 , g_{1T} , $g_{1L}^\perp$, h_1 , $h_{1T}^\perp$, $h_{1L}^\perp$, and $h_1^\perp$, are defined in terms of the same quark correlation functions entering the definition of ordinary parton distributions, but without integration over the transverse momentum. Among them, the Boer-Mulders $h_1^\perp$ [9] and the Sivers $f_{1T}^\perp$ [10] functions are T-odd, i.e. they change sign under “naive time reversal”, which is defined as usual time reversal, but without interchange of initial and final states. Since non-vanishing T-odd TMDs require gauge boson degrees of freedom which are not taken into account in our light-cone quark model, our model results will be discussed only for the T-even TMDs.

Projecting the correlator for quarks of definite longitudinal (s_L) or transverse ($\mathbf{s}_T$) polarizations, one obtains in nucleon states described by the polarization vector $\mathbf{S} = (s_L, \mathbf{S}_T)$ the following spin densities in the momentum space

$$\tilde{\rho}(x, \mathbf{k}_T^2, s_L, \mathbf{S}) = \frac{1}{2} \left[f_1 + S_T^i \epsilon^{ij} k_T^j \frac{1}{m} f_{1T}^\perp + s_L s_L g_{1L} + s_L S_T^i k_T^i \frac{1}{m} g_{1T} \right], \quad (1)$$

$$\begin{aligned} \tilde{\rho}(x, \mathbf{k}_T^2, \mathbf{s}_T, \mathbf{S}_T) = & \frac{1}{2} \left[f_1 + S_T^i \epsilon^{ij} k_T^j \frac{1}{m} f_{1T}^\perp + s_T^i \epsilon^{ij} k_T^j \frac{1}{m} h_1^\perp + s_T^i S_T^i h_1 \right. \\ & \left. + s_T^i (2k_T^i k_T^j - \mathbf{k}_T^2 \delta^{ij}) S_T^j \frac{1}{2m^2} h_{1T}^\perp + s_L s_T^i k_T^i \frac{1}{m} h_{1L}^\perp \right], \quad (2) \end{aligned}$$

where the distribution functions depend on x and $\mathbf{k}_T^2$. The unpolarized TMD f_1 , the helicity TMD g_{1L} , and the transversity TMD h_1 in Eqs. (1) and (2) correspond to monopole distributions in the momentum space for unpolarized, longitudinally and transversely polarized nucleon, respectively. They can be obtained from the overlap of LCWFs which are diagonal in the orbital angular momentum, but probe different transverse momentum and helicity correlations of the quarks inside the nucleon. All the other TMDs require a transfer of orbital angular momentum between the initial and final state. In particular, g_{1T} and $h_{1L}^\perp$ correspond to quark densities with specular configurations for the quark and nucleon spin: g_{1T} describes longitudinally polarized quarks in a transversely polarized nucleon, while $h_{1L}^\perp$ gives the distribution of transversely polarized quarks in longitudinally polarized nucleon. Therefore, g_{1T} requires helicity flip of the nucleon which is not

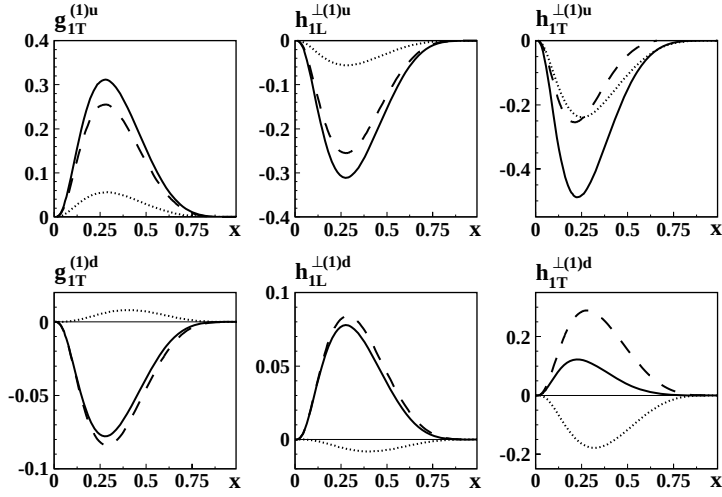


Figure 1: Transverse moments of TMDs as function of x for up (upper panels) and down (lower panels) quark. The solid curves show the total results, sum of the partial wave contributions. In the case of $g_{1T}^{(1)}$ and $h_{1L}^{\perp(1)}$ the dashed and dotted curves give the results from the S-P and P-D interference terms, respectively. In the case of $h_{1T}^{\perp(1)}$, the dashed curve is the result from P-wave interference, and the dotted curve is due to the interference of S and D waves.

compensated by a change of the quark helicity, and viceversa $h_{1L}^{\perp}$ involves helicity flip of the quarks but is diagonal in the nucleon helicity. As a result, in both cases, the LCWFs of the initial and final states differ by one unit of orbital angular momentum and the associated spin distributions have a dipole structure. Finally, for transverse polarizations in perpendicular directions of both the quarks and the nucleon, one has a quadrupole distribution with strength given by $h_{1T}^{\perp}$. In this case, the nucleon helicity flips in the direction opposite to the quark helicity, with a mismatch of two units for the orbital angular momentum of the initial and final LCWFs.

In Fig. 3 is shown the interplay between the different partial-wave contributions to the transverse moments $g_{1T}^{(1)}$, $h_{1L}^{\perp(1)}$ and $h_{1T}^{\perp(1)}$, defined as $g_{1T}^{(1)}(x) = \int d^2\mathbf{k}_T (\mathbf{k}_T^2/2m^2) g_{1T}(x, \mathbf{k}_T^2)$, etc. While the first two functions $g_{1T}^{(1)}$ and $h_{1L}^{\perp(1)}$ are dominated by the contribution due to P-wave interference, in the case of $h_{1T}^{\perp(1)}$ the contribution from the D wave is amplified through the interference with the S wave. The total results for up and down quarks obey the SU(6) isospin relation, i.e. the functions for up quarks are four times larger than for down quark and with opposite sign. This does not apply to the partial-wave contributions, as it is evident in particular for the terms containing D-wave contributions.

Among the distributions in Eqs. (1) and (2), the dipole correlations related to g_{1T} and $h_{1L}^{\perp}$ have characteristic features of intrinsic transverse momentum, since they are the only ones which have no analog in the spin densities related to the GPDs in the impact parameter space [11, 12]. The results in the light-cone quark model of Ref. [5] for the densities with longitudinally polarized quarks in a transversely polarized proton are shown in Fig. 2. The sideways shift in the positive (negative) x direction for up (down) quark due to the dipole term $\propto s_L S_T^i k_T^i \frac{1}{m} g_{1T}$ is sizable, and corresponds to an average deformation $\langle \mathbf{k}_x^u \rangle = 55.8$ MeV, and $\langle \mathbf{k}_x^d \rangle = -27.9$ MeV. The dipole distortion $\propto S_L s_T^i k_T^i \frac{1}{m} h_{1L}^{\perp}$ in the case of transversely polarized quarks in a longitudinally polarized proton is equal but with opposite sign, since in our model $h_{1L}^{\perp} = -g_{1T}$. (Also other quark model relations among

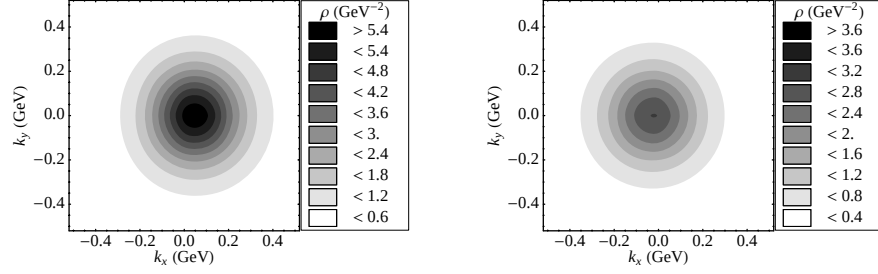


Figure 2: Quark densities in the $\mathbf{k}_T$ plane for longitudinally polarized quarks in a transversely polarized proton for up (left panel) and down (right panel) quark.

TMDs [13] are satisfied in our model, see [5].) These model results are supported from a recent lattice calculation [14,15] which gives, for the density related to g_{1T} , $\langle \mathbf{k}_x^u \rangle = 67(5)$ MeV, and $\langle \mathbf{k}_x^d \rangle = -30(5)$ MeV. For the density related to $h_{1L}^\perp$, they also find shifts of similar magnitude but opposite sign: $\langle \mathbf{k}_x^u \rangle = -60(5)$ MeV, and $\langle \mathbf{k}_x^d \rangle = 15(5)$ MeV.

2 Results for azimuthal SSAs

In Ref. [16] the present results for the T-even TMDs were applied to estimate azimuthal asymmetries in SIDIS, discussing the range of applicability of the model, especially with regard to the scale dependence of the observables and the transverse-momentum dependence of the distributions. Here we review the results for the Collins asymmetry $A_{UT}^{\sin(\phi+\phi_S)}$ and for $A_{UT}^{\sin(3\phi-\phi_S)}$, due to the Collins fragmentation function and to the chirally-odd TMDs h_1 , and $h_{1T}^\perp$, respectively. In both cases, we use the results extracted in [17] for the Collins function. In the denominator of the asymmetries we take f_1 from [18] and the unpolarized fragmentation function from [19], both valid at the scale $Q^2 = 2.5 \text{ GeV}^2$. In Fig. 2 the results for the Collins asymmetry in DIS production of charged pions off proton and deuterium targets are shown as function of x . The model results for h_1 evolved from the low hadronic scale of the model to $Q^2 = 2.5 \text{ GeV}^2$ ideally describe the HER-

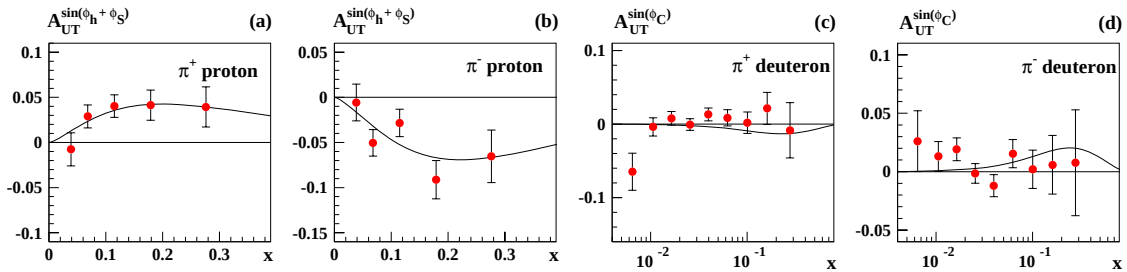


Figure 3: The single-spin asymmetry $A_{UT}^{\sin(\phi_h+\phi_S)} \equiv -A_{UT}^{\sin\phi_C}$ in DIS production of charged pions off proton and deuterium targets, as function of x . The theoretical curves are obtained on the basis of the light-cone CQM predictions for $h_1(x, Q^2)$ from Ref. [4,5]. The (preliminary) proton target data are from HERMES [20], the deuterium target data are from COMPASS [21].

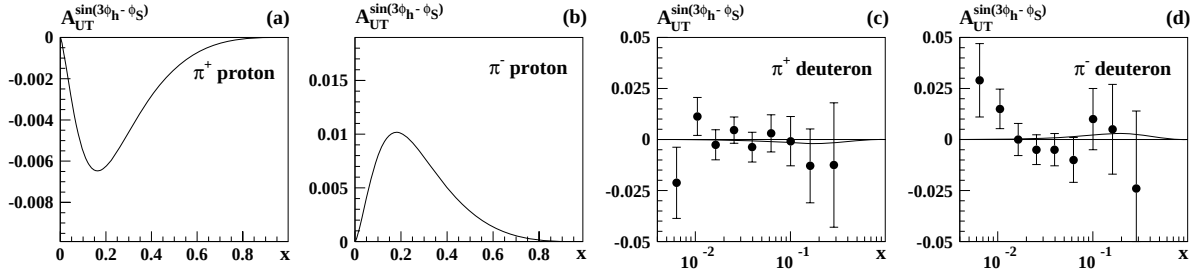


Figure 4: The single-spin asymmetry $A_{UT}^{\sin(3\phi_h - \phi_S)}$ in DIS production of charged pions off proton and deuterium targets, as function of x . The theoretical curves are obtained by evolving the light-cone CQM predictions for $h_{1T}^{\perp(1)}$ of Ref. [5] to $Q^2 = 2.5 \text{ GeV}^2$, using the h_1 evolution pattern. The preliminary COMPASS data are from Ref. [24].

MES data [20] for a proton target (panels (a) and (b) of Fig. 2). This is in line with the favourable comparison between our model predictions [4] and the phenomenological extraction of the transversity and the tensor charges in Ref. [22]. Our results are compatible also with the COMPASS data [21] for a deuteron target (panels (c) and (d) of Fig. 2) which extend down to much lower values of x .

In the case of the asymmetry $A_{UT}^{\sin(3\phi - \phi_S)}$ we face the question how to evolve $h_{1T}^{\perp(1)}$ from the low scale of the model to the relevant experimental scale. Since exact evolution equations are not available in this case, we “simulate” the evolution of $h_{1T}^{\perp(1)}$ by evolving it according to the transversity-evolution pattern. Although this is not the correct evolution pattern, it may give us a rough insight on the possible size of effects due to evolution (for a more detailed discussion we refer to [16]). The evolution effects give smaller asymmetries in absolute value and shift the peak at lower x values in comparison with the results obtained without evolution. The results shown in Fig. 4 are also much smaller than the bounds allowed by positivity, $|h_{1T}^{\perp(1)}| \leq \frac{1}{2} (f_1(x) - g_1(x))$, and constructed using parametrizations of the unpolarized and helicity distributions at $Q^2 = 2.5 \text{ GeV}^2$. Precise measurements in range $0.1 \lesssim x \lesssim 0.6$ are planned with the CLAS 12 GeV upgrade [23] and will be able to discriminate between these two scenarios. There exist also preliminary deuteron target data [24] which are compatible, within error bars, with the model predictions both at the hadronic and the evolved scale.

Acknowledgments

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COLOUR MODIFICATION OF FACTORISATION IN SINGLE-SPIN ASYMMETRIES

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Abstract

We discuss the way in which factorisation is partially maintained but nevertheless modified by process-dependent colour factors in hadronic single-spin asymmetries. We also examine QCD evolution of the twist-three gluonic-pole strength defining an effective T-odd Sivers function in the large- x limit, where evolution of the T-even transverse-spin DIS structure function g_2 is known to be multiplicative.

1 Preamble

1.1 Motivation

Single-spin asymmetries (SSA's) have long been something of an enigma in high-energy hadronic physics. Prior to the first experimental studies, hadronic SSA's were predicted to be very small for a variety of reasons. Experimentally, however, they turn out to be large (up to the order of 50% and more) in many hadronic processes. It was also long held that such asymmetries should eventually vanish with growing energy and/or p_T . Again, however, the SSA's so far observed show no signs of high-energy suppression.

1.2 SSA Basics

Typically, SSA's reflect spin-momenta correlations of the form $\mathbf{s} \cdot (\mathbf{p} \wedge \mathbf{k})$, where $\mathbf{s}$ is some particle polarisation vector, while $\mathbf{p}$ and $\mathbf{k}$ are initial/final particle/jet momenta. A simple example might be: $\mathbf{p}$ the beam direction, $\mathbf{s}$ the target polarisation (transverse therefore with respect to $\mathbf{p}$) and $\mathbf{k}$ the final-state particle direction (necessarily then out of the $\mathbf{p}$ - $\mathbf{s}$ plane). Polarisation involved in SSA's must usually thus be transverse (although there are certain special exceptions).

It is more convenient to use an helicity basis via the transformation

$$|\uparrow/\downarrow\rangle = \frac{1}{\sqrt{2}} [|+\rangle \pm i |-\rangle]. \quad (1)$$

A transverse-spin asymmetry then takes on the (schematic) form

$$\mathcal{A}_N \sim \frac{\langle \uparrow | \uparrow \rangle - \langle \downarrow | \downarrow \rangle}{\langle \uparrow | \uparrow \rangle + \langle \downarrow | \downarrow \rangle} \sim \frac{2\Im \langle + | - \rangle}{\langle + | + \rangle + \langle - | - \rangle}. \quad (2)$$

The appearance of both $|+\rangle$ and $|-\rangle$ in the numerator signals the presence of a helicity-flip amplitude. The precise form of the numerator implies interference between two different helicity amplitudes: one helicity-flip and one non-flip, with a relative phase difference (the imaginary phase implying naïve T-odd processes).

Early on Kane, Pumplin and Repko [1] realised that in the massless (or high-energy) limit and the Born approximation a gauge theory such as QCD cannot furnish either requirement: for a massless fermion, helicity is conserved and tree-diagram amplitudes are always real. This led to the now infamous statement [1]: “... *observation of significant polarizations in the above reactions would contradict either QCD or its applicability.*”

It therefore caused much surprise and interest when large asymmetries were found; QCD nevertheless survived! Efremov and Teryaev [2] soon discovered one way out within the context of perturbative QCD. Consideration of the three-parton correlators involved in, *e.g.* g_2 , leads to the following crucial observations: the relevant mass scale is not that of the current quark, but of the hadron and the pseudo-two-loop nature of the diagrams can generate an imaginary part in certain regions of partonic phase space [3].

It took some time, however, before real progress was made and the richness of the newly available structures was fully exploited—see [4]. Indeed, it turns out that there are a variety of mechanisms that can generate SSA’s:

- Transversity: this correlates hadron helicity flip to quark flip. Chirality conservation, however, requires another T-odd (distribution or fragmentation) function.
- Internal quark motion: the transverse polarisation of a quark may be correlated with its own transverse momentum. This corresponds to the Sivers function [5] and requires orbital angular momentum together with soft-gluon exchange.
- Twist-3 transverse-spin dependent three-parton correlators (*cf.* g_2): here the pseudo two-loop nature provides effective spin flip (via the extra parton) and also the required imaginary part (via pole terms).

The second and third mechanisms turn out to be related.

2 Single-Spin Asymmetries

2.1 Single-Hadron Production

As a consequence of the multiplicity of underlying mechanisms, there are various types of distribution and fragmentation functions that can be active in generating SSA’s (even competing in the same process):

- higher-twist distribution and fragmentation functions,
- k_T -dependent distribution and fragmentation functions,
- interference fragmentation functions,
- higher-spin functions, *e.g.* vector-meson fragmentation functions.

Consider then hadron production with one initial-state, transversely polarised hadron:

$$A^\uparrow(P_A) + B(P_B) \rightarrow h(P_h) + X, \quad (3)$$

where hadron A is transversely polarised while B is not. The unpolarised (or spinless) hadron h is produced at large transverse momentum $\mathbf{P}_{hT}$ and PQCD is thus applicable. Typically, A and B are protons while h may be a pion or kaon *etc.*

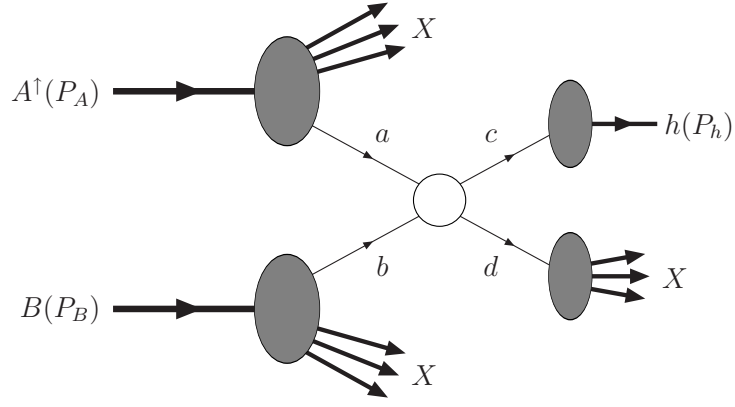


Figure 1: Factorisation in single-hadron production with a transversely polarised hadron.

The following SSA may then be measured:

$$A_T^h = \frac{d\sigma(\mathbf{S}_T) - d\sigma(-\mathbf{S}_T)}{d\sigma(\mathbf{S}_T) + d\sigma(-\mathbf{S}_T)}. \quad (4)$$

Assuming standard factorisation to hold, the differential cross-section for such a process may be written formally as (*cf.* Fig. 1)

$$d\sigma = \sum_{abc} \sum_{\alpha\alpha'\gamma\gamma'} \rho_{\alpha'\alpha}^a f_a(x_a) \otimes f_b(x_b) \otimes d\hat{\sigma}_{\alpha\alpha'\gamma\gamma'} \otimes \mathcal{D}_{h/c}^{\gamma'\gamma}(z), \quad (5)$$

where f_a (f_b) is the density of parton type a (b) inside hadron A (B), $\rho_{\alpha\alpha'}$ is the spin density matrix for parton a , $\mathcal{D}_{h/c}^{\gamma'\gamma}$ is the fragmentation matrix for parton c into the final hadron h and $d\hat{\sigma}_{\alpha\alpha'\gamma\gamma'}$ is the partonic cross-section:

$$\left(\frac{d\hat{\sigma}}{d\hat{t}}\right)_{\alpha\alpha'\gamma\gamma'} = \frac{1}{16\pi\hat{s}^2} \frac{1}{2} \sum_{\beta\delta} \mathcal{M}_{\alpha\beta\gamma\delta} \mathcal{M}_{\alpha'\beta\gamma'\delta}^*, \quad (6)$$

where $\mathcal{M}_{\alpha\beta\gamma\delta}$ is the amplitude for the hard partonic process, see Fig. 2.

The off-diagonal elements of $\mathcal{D}_{h/c}^{\gamma'\gamma}$ vanish for an unpolarised produced hadron; *i.e.*, $\mathcal{D}_{h/c}^{\gamma'\gamma} \propto \delta_{\gamma\gamma'}$. Helicity conservation then implies $\alpha = \alpha'$, so that there can be no dependence on the spin of hadron A and all SSA's must vanish. To avoid such a conclusion, either intrinsic quark transverse motion, or higher-twist effects must be invoked.

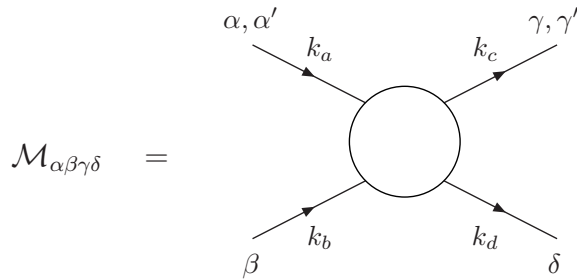


Figure 2: The hard partonic amplitude, $\alpha\beta\gamma\delta$ are Dirac indices.

2.2 Intrinsic Transverse Motion

Quark intrinsic transverse motion can generate SSA's in three essentially different ways (all necessarily T -odd effects):

1. $\mathbf{k}_T$ in hadron A requires $f_a(x_a)$ to be replaced by $\mathcal{P}_a(x_a, \mathbf{k}_T)$, which may then depend on the spin of A (distribution level);
2. $\boldsymbol{\kappa}_T$ in hadron h allows $\mathcal{D}_{h/c}^{\gamma'}$ to be non-diagonal (fragmentation level);
3. $\mathbf{k}'_T$ in hadron B requires $f_b(x_b)$ to be replaced by $\mathcal{P}_b(x_b, \mathbf{k}'_T)$ —the spin of b in the unpolarised B may then couple to the spin of a (distribution level).

The three corresponding mechanisms are: 1. the Sivers effect [5]; 2. the Collins effect [6]; 3. an effect studied by Boer [7] in Drell–Yan processes. Note that all such intrinsic- $\mathbf{k}_T$, $-\boldsymbol{\kappa}_T$ or $-\mathbf{k}'_T$ effects are T -odd; *i.e.*, they require ISI or FSI. Note too that when transverse parton motion is included, the QCD factorisation theorem is not completely proven, but see [8].

Assuming factorisation to be valid, the cross-section is

$$E_h \frac{d^3\sigma}{d^3\mathbf{P}_h} = \sum_{abc} \sum_{\alpha\alpha'\beta\beta'\gamma\gamma'} \int dx_a dx_b d^2\mathbf{k}_T d^2\mathbf{k}'_T \frac{d^2\boldsymbol{\kappa}_T}{\pi z} \times \mathcal{P}_a(x_a, \mathbf{k}_T) \rho_{\alpha'\alpha}^a \mathcal{P}_b(x_b, \mathbf{k}'_T) \rho_{\beta'\beta}^b \left(\frac{d\hat{\sigma}}{d\hat{t}} \right)_{\alpha\alpha'\beta\beta'\gamma\gamma'} \mathcal{D}_{h/c}^{\gamma'\gamma}(z, \boldsymbol{\kappa}_T), \quad (7)$$

where again

$$\left(\frac{d\hat{\sigma}}{d\hat{t}} \right)_{\alpha\alpha'\beta\beta'\gamma\gamma'} = \frac{1}{16\pi\hat{s}^2} \sum_{\beta\delta} \mathcal{M}_{\alpha\beta\gamma\delta} \mathcal{M}_{\alpha'\beta'\gamma'\delta}^*. \quad (8)$$

The Sivers effect relies on T -odd k_T -dependent distribution functions and predicts an SSA of the form

$$E_h \frac{d^3\sigma(\mathbf{S}_T)}{d^3\mathbf{P}_h} - E_h \frac{d^3\sigma(-\mathbf{S}_T)}{d^3\mathbf{P}_h} = |\mathbf{S}_T| \sum_{abc} \int dx_a dx_b \frac{d^2\mathbf{k}_T}{\pi z} \Delta_0^T f_a(x_a, \mathbf{k}_T^2) f_b(x_b) \frac{d\hat{\sigma}(x_a, x_b, \mathbf{k}_T)}{d\hat{t}} D_{h/c}(z), \quad (9)$$

where $\Delta_0^T f$ (related to $f_{1T}^\perp$) is a T -odd distribution.

2.3 Higher Twist

Efremov and Teryaev [2] showed that in QCD non-vanishing SSA's can also be obtained by invoking higher twist and the so-called gluonic poles in diagrams involving qqg correlators. Such asymmetries were later evaluated in the context of QCD factorisation by Qiu and Sterman, who studied both direct-photon production [4] and hadron production [9]. This program has been extended by Kanazawa and Koike [10] to the chirally-odd contributions.

The various possibilities are:

$$\begin{aligned}
d\sigma = \sum_{abc} \Big\{ & G_F^a(x_a, y_a) \otimes f_b(x_b) \otimes d\hat{\sigma} \otimes D_{h/c}(z) \\
& + \Delta_T f_a(x_a) \otimes E_F^b(x_b, y_b) \otimes d\hat{\sigma}' \otimes D_{h/c}(z) \\
& + \Delta_T f_a(x_a) \otimes f_b(x_b) \otimes d\hat{\sigma}'' \otimes D_{h/c}^{(3)}(z) \Big\}. \quad (10)
\end{aligned}$$

The first term represents the chirally-even three-parton correlator pole mechanism, as proposed in [2] and studied in [4, 9]; the second contains transversity and corresponds to the chirally-odd contribution analysed in [10]; and the third also contains transversity but requires a twist-3 fragmentation function $D_{h/c}^{(3)}$.

2.4 Phenomenology

Anselmino *et al.* [11] have compared data with various models inspired by the previous possible (k_T -dependent) mechanisms and find good descriptions although they were not able to differentiate between contributions. The results of Qiu and Sterman [4] (based on three-parton correlators) also compare well but are rather complex. However, the twist-3 correlators (as in g_2) obey constraining relations with k_T -dependent densities and also exhibit a novel factorisation property and thus simplification, to which we now turn.

2.5 Pole Factorisation

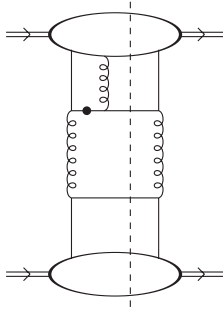


Figure 3: Example of a propagator pole in a three-parton diagram.

Efremov and Teryaev [2] noticed that certain twist-3 diagrams involving three-parton correlators can supply the necessary imaginary part via a pole term while spin-flip is implicit, due to the gluon. The standard propagator prescription (denoted $\bullet$ in Fig. 3, with momentum k),

$$\frac{1}{k^2 \pm i\epsilon} = \mathbf{P} \frac{1}{k^2} \mp i\pi\delta(k^2). \quad (11)$$

leads to an imaginary contribution for $k^2 \rightarrow 0$. A gluon with momentum $x_g p$ inserted into an (initial or final) external line p' sets $k = p' - x_g p$ and thus as $x_g \rightarrow 0$ we have that $k^2 \rightarrow 0$. The gluon vertex may then be factored out together with the quark propagator and pole, see Fig. 4. Such factorisation can be performed systematically for all poles (gluon and fermion): on all external legs with all insertions [12]. The structures

$$\text{Pole Part} = -i\pi \frac{p' \cdot \xi}{p' \cdot p} \times \text{Diagram}$$

Figure 4: An example of pole factorisation: p is the incoming proton momentum, p' the outgoing hadron and ξ is the gluon polarisation vector (lying in the transverse plane).

are still complex: for a given correlator there are many insertions, leading to different signs and momentum dependence.

The colour structures of the various diagrams (with different possible types of soft insertions) are also different (we shall examine this question shortly). In all the cases we examined it turns out that just one diagram dominates in the large- N_c limit, see Fig. 5. All other insertions into external (on-shell) legs are relatively suppressed by $1/N_c^2$. This has all been examined in detail by Ramilli (Insubria U. Masters thesis [13]): the leading diagrams provide a good approximation. The analysis has yet to be repeated for all the other possible twist-3 contributions (*e.g.* also in fragmentation).

A question immediately arises: could there be any direct relationship between the twist-3 and k_T -dependent mechanisms? It might be hoped that, via the equations of motion *etc.*, unique predictions for single-spin azimuthal asymmetries could be obtained by linking the (*e.g.* Siverson- or Collins-like) k_T -dependent mechanisms to the (Efremov–Teryaev) higher-twist three-parton mechanisms. An early attempt was made by Ma and Wang [14] for the Drell–Yan process, but the predictions were found not to be unique. On the other hand, Ji *et al.* [15] have since also examined the relationship between the k_T -dependent and higher-twist mechanisms by matching the two in an intermediate k_T region of common validity.

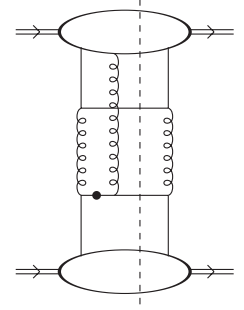


Figure 5: Example of a dominant propagator pole diagram.

3 More on Multiparton Correlators

3.1 Colour Modification

In [16] we provided an *a posteriori* proof of the relation between twist-3 and k_T -dependence. The starting point is a factorised formula for the Siverson function:

$$d\Delta\sigma \sim \int d^2k_T dx f_S(x, k_T) \epsilon^{\rho s P k_T} \text{Tr}[\gamma_\rho H(xP, k_T)] . \quad (12)$$

Expanding the subprocess coefficient function H in powers of k_T and keeping the first non-vanishing term leads to

$$\sim \int d^2k_T dx f_S(x, k_T) k_T^\alpha \epsilon^{\rho s P k_T} \text{Tr}\left[\gamma_\rho \frac{\partial H(xP, k_T)}{\partial k_T^\alpha}\right]_{k_T=0} . \quad (13)$$

By exploiting various identities and the fact that there are other momenta involved, this can be rearranged into the following form:

$$d\Delta\sigma \sim M \int dx f_S^{(1)}(x) \epsilon^{\alpha s P n} \text{Tr}\left[\not{p} \frac{\partial H(xP, k_T)}{\partial k_T^\alpha}\right]_{k_T=0} , \quad (14)$$

where

$$f_S^{(1)}(x) = \int d^2k_T f_S(x, k_T) \frac{k_T^2}{2M^2} . \quad (15)$$

The final expression coincides with the master formula of Koike and Tanaka [17] for twist-3 gluonic poles in high- p_T processes. The Sivvers function can thus be identified with the gluonic-pole strength $T(x, x)$ multiplied by a process-dependent colour factor.

The sign of the Sivvers function depends on which of ISI or FSI is relevant:

$$f_S^{(1)}(x) = \sum_i C_i \frac{1}{2M} T(x, x), \quad (16)$$

where C_i is a colour factor defined relative to an Abelian subprocess. Emission of an extra hard gluon is needed to generate high p_T and, according to the process under consideration, FSI may occur *before* or *after* this emission, leading to different colour factors. In this sense, factorisation is broken in SIDIS, albeit in a simple and accountable manner. Figure 6 depicts the application of this relation to high- p_T SIDIS.

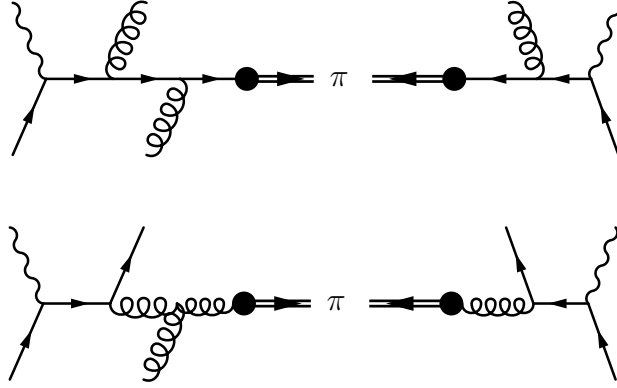


Figure 6: Twist-3 SIDIS π production via quark and gluon fragmentation.

3.2 Asymptotic Behaviour

The relation between gluonic poles (*e.g.* the Sivvers function, and T-even transverse-spin effects, *e.g.* g_2 [18–22]) still remains unclear. Although there are model-based estimates and approximate sum rules, the compatibility of general twist-3 evolution with dedicated studies of gluonic-pole evolution (at LO [23, 24] and at NLO [25]) is still unproven.

In the large- x limit the evolution equations for g_2 diagonalise in the double-moment arguments [26]. For the Sivvers function and gluonic poles, this is the important kinematical region [4]. The gluonic-pole strength $T(x)$, corresponds to a specific matrix element [4]. It is also the residue of a general qqg vector correlator $b_V(x_1, x_2)$ [27]:

$$b_V(x_1, x_2) = \frac{T(\frac{x_1+x_2}{2})}{x_1 - x_2} + \text{regular part}, \quad (17)$$

which is defined as

$$b_V(x_1, x_2) = \frac{i}{M} \int \frac{d\lambda_1 d\lambda_2}{2\pi} e^{i\lambda_1(x_1-x_2)+i\lambda_2 x_2} \epsilon^{\mu s p_1 n} \langle p_1, s | \bar{\psi}(0) \not{n} D_\mu(\lambda_1) \psi(\lambda_2) | p_1, s \rangle. \quad (18)$$

There is though one other correlator, projected onto an axial Dirac matrix:

$$b_A(x_1, x_2) = \frac{1}{M} \int \frac{d\lambda_1 d\lambda_2}{\pi} e^{i\lambda_1(x_1-x_2)+i\lambda_2 x_2} \langle p_1, s | \bar{\psi}(0) \not{n} \gamma^5 s \cdot D(\lambda_1) \psi(\lambda_2) | p_1, s \rangle. \quad (19)$$

This last is required to complete the description of transverse-spin effects, in both SSA's and g_2 . The two correlators have opposite symmetry properties for $x_1 \leftrightarrow x_2$:

$$b_A(x_1, x_2) = b_A(x_2, x_1), \quad b_V(x_1, x_2) = -b_V(x_2, x_1), \quad (20)$$

determined by T invariance. In both DIS and SSA's just one combination appears [19]:

$$b_-(x_1, x_2) = b_A(x_2, x_1) - b_V(x_1, x_2). \quad (21)$$

The QCD evolution equations [20–22] are best expressed in terms of another quantity, which is determined by matrix elements of the gluon field strength:

$$Y(x_1, x_2) = (x_1 - x_2) b_-(x_1, x_2). \quad (22)$$

It should be safe to assume that $b_-(x_1, x_2)$ has no double pole and thus

$$T(x) = Y(x, x). \quad (23)$$

Evolution is easier studied for Mellin-moment, implying double moments of $Y(x, y)$:

$$Y^{mn} = \int dx dy x^m y^n Y(x, y), \quad (24)$$

where the allowed regions are $|x|, |y|$ and $|x - y| < 1$ (recall that negative values indicate antiquark distributions). We wish to examine the behaviour for x and y both close to unity and therefore close to each other. The gluonic pole thus provides the dominant contribution:

$$\lim_{x, y \rightarrow 1} Y(x, y) = T\left(\frac{x+y}{2}\right) + O(x - y). \quad (25)$$

In this approximation (now large m, n) the leading-order evolution equations simplify:

$$\frac{d}{ds} Y^{nn} = 4 \left(C_F + \frac{C_A}{2} \right) \ln n Y^{nn}, \quad (26)$$

where the evolution variable is $s = \beta_0^{-1} \ln \ln Q^2$. In terms of $T(x)$ this is

$$\dot{T}(x) = 4 \left(C_F + \frac{C_A}{2} \right) \int_x^1 dz \frac{(1-z)}{(1-x)} \frac{1}{(z-x)_+} T(z), \quad (27)$$

which is similar to the unpolarised case, but differs by a colour factor $(C_F + C_A/2)$ and a softening factor $(1-z)/(1-x)$.

The extra piece in the colour factor $(C_A/2)$ *vis-a-vis* the unpolarised case (just C_F) reflects the presence of a third active parton—the gluon. That is, the pole structure of three-parton kernels is identical, but the effective colour charge of the extra gluon is $C_A/2$. The softening factor is inessential to the asymptotic solution, it merely implies standard evolution for the function $f(x) = (1-x)T(x)$. For an initial $f(x, Q_0^2) = (1-x)^a$, the asymptotic solution [28] is the same but modified by $a \rightarrow a(s)$, with

$$a(s) = a + 4 \left(C_F + \frac{C_A}{2} \right) s. \quad (28)$$

For $T(x)$, a shifts to $a - 1$ and the evolution modification is identical; the spin-averaged asymptotic solutions are thus also valid for $T(x)$. This large- x limit of the evolution agrees, barring the colour factor itself, with studies of gluonic-pole evolution [23–25]. We should note here that Braun *et al.* [29] have found errors in [23] and [25]; our results for the logarithmic term are, however, unaffected.

4 Summary and Conclusions

Viewing the Sivers function as an effective twist-3 gluonic-pole contribution [16], it is seen to be process dependent: besides a sign (ISI *vs.* FSI), there is a process-dependent colour factor. This factor is determined by the colour charge of the initial and final partons. It generates the sign difference between SIDIS and Drell–Yan at low p_T , but in hadronic reactions at high p_T it is more complicated. Such a picture is complementary to the matching in the region of common validity. The matching between various p_T regions now takes the form of a p_T -dependent colour factor. It also lends some justification to the possibility of global Sivers function fits [30].

We have shown that generic twist-3 evolution is applicable to the Sivers function. Its effective nature allows us to relate the evolution of T-odd (Sivers function) and T-even (gluonic pole) quantities. A vital ingredient here is the large- x approximation, where gluonic-poles dominate and the evolution simplifies. The Sivers-function evolution is then multiplicative and described by a colour-factor modified twist-2 spin-averaged kernel [31].

Acknowledgments

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IMPACT PARAMETER REPRESENTATION AND x -DEPENDENCE OF THE TRANSVERSITY SPIN STRUCTURE OF THE NUCLEONS

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Abstract

In the frame work of the our model of t -dependence of generalized parton distributions we obtain its impact parameter representation. On this basis we calculate and compare with the results of other models the transversity spin structure of the proton and neutron. We calculate as function of x the unpolarized density, the density of unpolarized quarks in the transversely polarized nucleon.

1 Introduction

The electromagnetic hadrons form factors are related to the first moments of the Generalized Parton distributions (GPDs) [1, 3, 2]

$$F_1^q(t) = \int_0^1 dx \mathcal{H}^q(x, t), \quad F_2^q(t) = \int_0^1 dx \mathcal{E}^q(x, t). \quad (1)$$

Taking the matrix elements of energy-momentum tensor $T_{\mu\nu}$ instead of the electromagnetic current J^μ one can obtain the gravitational form factors of quarks which are related to the second moments. For $\xi = 0$ one has

$$\int_0^1 dx x \mathcal{H}_q(x, t) = A_q(t); \quad \int_0^1 dx x \mathcal{E}_q(x, t) = B_q(t). \quad (2)$$

Non-forward parton densities also provide information about the distribution of the parton in impact parameter space [4] which is connected with t -dependence of *GPDs*. Now we cannot obtain this dependence from the first principles, but it must be obtained from the phenomenological description with *GPDs* of the nucleon electromagnetic form factors.

Note that in [5, 6] it was shown that at large $x \rightarrow 1$ and momentum transfer the behavior of GPDs requires a larger power of $(1 - x)$ in the t -dependent exponent with $n \geq 2$. It was noted that $n = 2$ naturally leads to the Drell-Yan-West duality between parton distributions at large x and the form factors.

In [7], a simple ansatz was proposed which will be good for describing the form factors of the proton and neutron by taking into account a number of new data that have appeared in the last years. We choose the t -dependence of GPDs in the form

$$\mathcal{H}^q(x, t) = q(x) \exp[a_+ \frac{(1-x)^2}{x^m} t]; \quad \mathcal{E}^q(x) = \frac{k_q}{N_q} (1-x)^{\kappa_1} q(x) \exp[a_- \frac{(1-x)^2}{x^m} t], \quad (3)$$

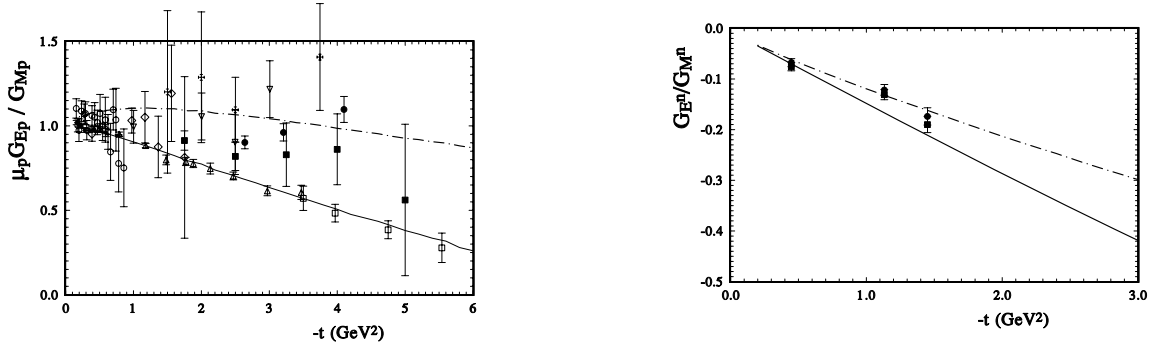


Figure 1: a)[left] $\mu_p G_E^p / G_M^p$ (hard and dot-dashed lines correspond to variant (I) and (II)); b)[right] G_E^n / G_M^n (hard and dot-dashed lines corresponds to variant (II) and (I));

with the standard normalization of the form factors $\kappa_1 = 1.53$ and $\kappa_2 = 0.31$. The size of the parameter $m = 0.4$ was determined by the low t experimental data; the free parameters $a_{\pm}$ (a_+ - for $\mathcal{H}$ and a_- - for $\mathcal{E}$) were chosen to reproduce the experimental data in a wide t region. The $q(x)$ was taken from the MRST2002 global fit [10] with the scale $\mu^2 = 1 \text{ GeV}^2$. In all our calculations we restrict ourselves, as in other works, only to the contributions of u and d quarks and the terms in $\mathcal{H}^q$ and $\mathcal{E}^q$.

2 Proton and neutron electromagnetic form factors

The proton Dirac form factor calculated in [7, 9] reproduces sufficiently well the behavior of experimental data not only at high t but also at low t . Our description of the ratio of the Pauli to the Dirac proton form factors and the ratio of G_E^p / G_M^p shows (see Fig.1a) that in our model we can obtain the results of both the methods (Rosenbluth and Polarization) by changing the slope of $\mathcal{E}$ [9]. Based on the model developed for proton the neutron form factors are calculated too. To do this the isotopic invariance can be used to change the proton GPDs to neutron GPDs. The calculations of G_E^n and G_M^n (see Fig.1b) show that the variant which describes the polarization data is in better agreement with the experimental data that is coincides with our calculations.

3 Transversety asymmetry of the gravitational form factors

The representation (2) combined with our model (we use here the first variant of parameters describing the experimental data obtained by the polarization method) allows to calculate the gravitational form factors of valence quarks and their contribution (being just their sum) to gravitational form factors of nucleon [8, 9]. Note that nonperturbative analysis within the framework of the lattice OCD indicates that the net quark contribution to the anomalous gravimagnetic moment $B_{u+d}(0)$ is close to zero [11, 12]. Let us examine the distribution of matter (that is, gravitational charge) in the polarized nucleon. For that purpose we generalize (4) in a straightforward way and introduce the

gravitomagnetic transverse density

$$\rho_T^{Gr}(\vec{b}) = \frac{1}{2\pi} \int_{-\infty}^{\infty} dq \, q \, J_0(qb) A(q^2) + \sin(\phi) \frac{1}{2\pi} \int_0^{\infty} dq \frac{q^2}{2M_N} J_1(qb) B(q^2). \quad (4)$$

Our calculations of the corresponding three dimension plots over the impact parameter b_x and b_y are shown in fig.2 for the total and the separate parts of the u and d quarks and for different values of x . One can see that the asymmetry at small x is quite small and its maximal values are concentrated almost at the same distances from the center of nucleon as the transverse charge density of the proton.

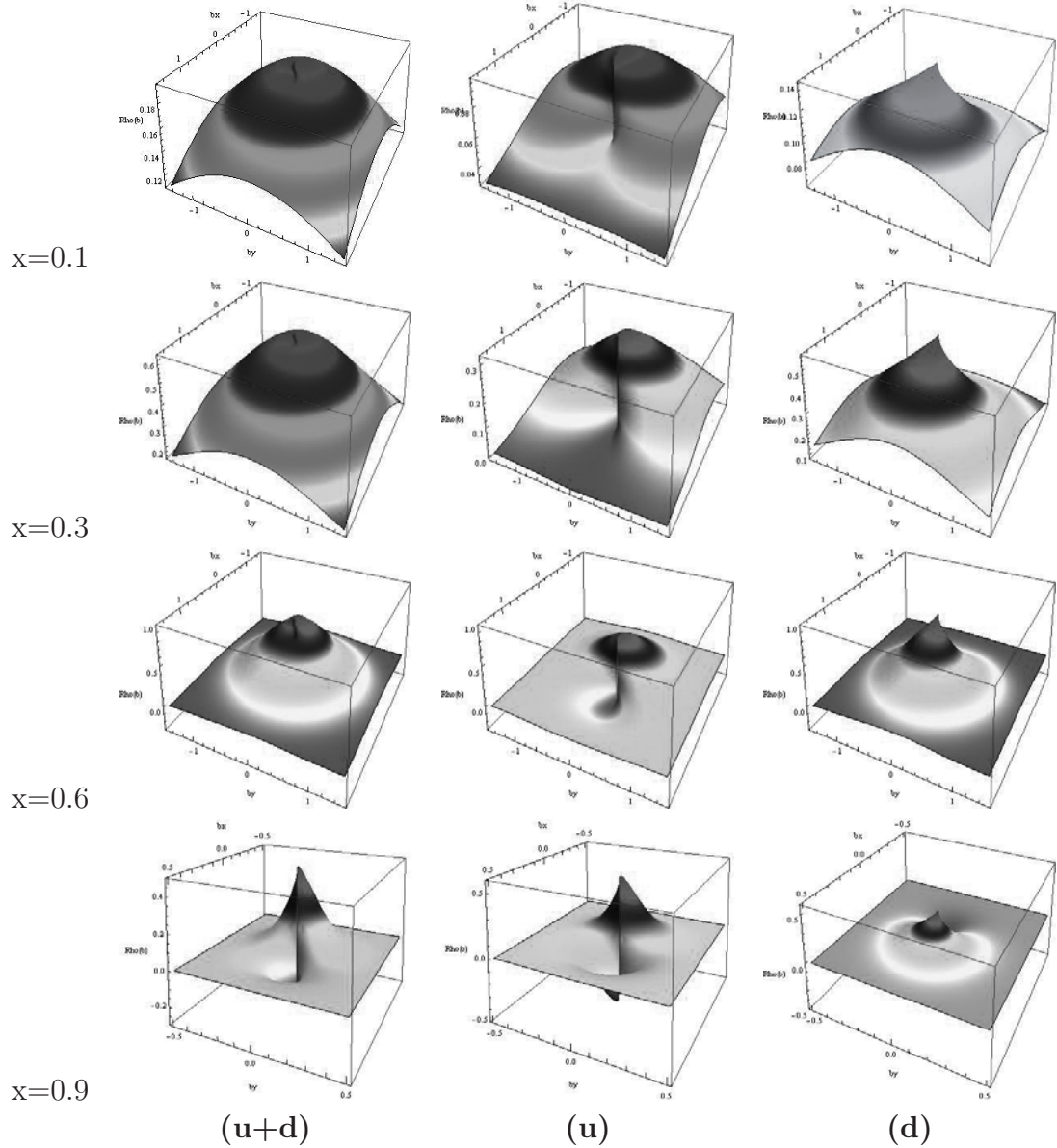


Figure 2: (a) The graviton density of transverse polarized the nucleon.

4 Conclusion

We introduced a simple new form of the t -dependence of GPDs. It satisfies the conditions of the non-factorization, condition on the power of $(1-x)^n$ in the exponential form of the t -dependence. With this simple form we obtained a good description of the proton electromagnetic Sachs form factors. Using the isotopic invariance we obtained good descriptions of the neutron Sachs form factors without changing any parameters.

On the basis of our results we calculated the contribution of the u and d quarks to the gravitational form factor of the nucleons. The obtained transverse asymmetry determined by the contribution of the u -quark distribution. The relative asymmetry of the sum of the u and d quarks contributions is very small at small x and it is large at large x .

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POLARIZED PARTON DISTRIBUTIONS FROM NLO QCD ANALYSIS OF WORLD DIS AND SIDIS DATA

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Abstract

The combined analysis of polarized DIS and SIDIS data is performed in NLO QCD. The new parameterization on polarized PDFs is constructed. The especial attention is paid to the impact of novel SIDIS data on the polarized distributions of light sea and strange quarks.

Since the observation of the famous spin crisis in 1987 one of the most intriguing and still unsolved problems of the modern high energy physics is the nucleon spin puzzle. The key component of this problem, which attracted the great both theoretical and experimental efforts during many years is the finding of the polarized parton distributions functions (PDFs) in nucleon.

The basic process which enables us to solve these problems is the process of semi-inclusive DIS (SIDIS). However, until recently the quality of the polarized SIDIS data was rather poor, so that its inclusion in the analysis did not helped us to solve the main task of SIDIS measurements: to extract the polarized sea and valence PDFs of all active flavors. Only in 2004 the first polarized SIDIS data with the identification of produced hadrons (pions and kaons) were published [1]. These data were included in the global QCD analysis in Ref. [2]. Recently, the new data on the SIDIS asymmetries $A_d^{\pi^\pm}, A_d^{K^\pm}$ were published [3] by the COMPASS collaboration. It is of importance that this data cover the most important and badly investigated low x region. In this paper we include this data in the new global QCD analysis of all existing polarized DIS and SIDIS data. The elaborated parameterization on the polarized PDFs in some essential points differs from the parameterization of Ref. [4] (see below).

We tried to be as close as possible to our previous NLO QCD analysis of pure inclusive DIS data [5]. Namely, data we parameterize the singlet $\Delta\Sigma$ and two nonsinglet $\Delta q_3, \Delta q_8$ combinations at the initial scale $Q_0^2 = 1 \text{ GeV}^2$ in the common form (which is used also for ΔG and $\Delta\bar{u}, \Delta\bar{d}$ distributions)

$$\Delta q = \eta \frac{x^\alpha (1-x)^\beta (1 + \gamma x + \delta \sqrt{x})}{\int_0^1 x^\alpha (1-x)^\beta (1 + \gamma x + \delta \sqrt{x}) dx}, \quad (1)$$

Then, the quantities $\Delta u + \Delta\bar{u}, \Delta d + \Delta\bar{d}, \Delta s = \Delta\bar{s}$ are determined as: $\Delta u + \Delta\bar{u} = \frac{1}{6}(3\Delta q_3 + \Delta q_8 + 2\Delta\Sigma)$, $\Delta d + \Delta\bar{d} = \frac{1}{6}(3\Delta q_3 - \Delta q_8 + 2\Delta\Sigma)$, $\Delta s = \Delta\bar{s} = \frac{1}{6}(\Delta\Sigma - \Delta q_8)$. Further, to properly describe the SIDIS data we, besides $\Delta\Sigma, \Delta q_3$ and Δq_8 , parametrize the sea PDFs of u and d flavors. For the DIS sector we introduced the additional factors $\gamma_{\Delta q_3} x, \gamma_{\Delta q_8} x$ for Δq_3 and Δq_8 to provide the better flexibility of the parametrizations required by the inclusion of SIDIS data. Besides, we introduce the additional factors

$\delta_{\Delta q_8} \sqrt{x}$ for Δq_8 and $\gamma_{\Delta G} x$ for ΔG . to provide the possibility of sign-changing scenarios for Δs and ΔG , respectively.

We analyze the inclusive A_1 and semi-inclusive A_1^h asymmetries. We work in $\overline{MS}$ factorization scheme. We use here the latest NLO parameterization on fragmentation functions from Ref. [6]. Calculating F_2 and F_2^h we use parameterization for R from [7] and the recent NLO parameterization on unpolarized PDFs from Ref. [8].

For our analysis we collected all accessible in literature polarized DIS and SIDIS data. We include the inclusive proton, deuteron and neutron data by SMC, E143, E155, E154, COMPASS, HERMES, CLAS. The semi-inclusive data are collected by SMC, HERMES and COMPASS. We include also the latest COMPASS data from Ref [3]. In total we have 232 points for the inclusive polarized DIS and 202 points for semi-inclusive polarized DIS. For 16 fit parameters $\chi_0^2|_{inclusive} = 188.4$ and $\chi_0^2|_{semi-inclusive} = 194.8$ for DIS and SIDIS data, while $\chi_0^2|_{total} = 383.9$ for the full set of data (434 points). Thus, one can conclude that the fit quality is quite good: $\chi_0^2/D.O.F. \simeq 0.84$.

The optimal values of our fit parameters are presented in Table 1. Certainly, the

Table 1: Optimal values of the global fit parameters at the initial scale $Q_0^2 = 1 \text{ GeV}^2$.

	$\Delta\Sigma$	Δq_3	Δq_8		ΔG	$\Delta\bar{u}$	Δd
α	1.0227	-0.6342	-0.7916	α	0.9040	-0.3506	0.2802
β	3.3891	3.1418	$= \beta_{\Delta q_3}$	β	$= \beta_{\Delta\bar{u}}$	15.0 (fixed)	$= \beta_{\Delta\bar{u}}$
γ	0.0 (fixed)	23.9180	36.8400	γ	-5.6703	0.0 (fixed)	0.0 (fixed)
δ	0.0 (fixed)	0.0 (fixed)	-13.7480	δ	0.0000 (fixed)	0.0 (fixed)	0.0 (fixed)
η	0.3846	1.2660	0.6170	η	-0.1828	0.0672	-0.0792

construction of the best fit should be accompanied by the reliable method of uncertainties estimation. We choose the modified Hessian method [9], [10] which well works (as well as the Lagrange multipliers method – see [4] and references therein) even in the case of deviation of χ^2 profile from the quadratic parabola, and was successfully applied in a lot of physical tasks. Besides, very important question arises about choice of $\Delta\chi^2$ determining the uncertainty size. The standard choice is $\Delta\chi^2 = 1$, just as we did before in Ref. [5]. However, the such choice of $\Delta\chi^2$ can lead to underestimation of uncertainties. The alternative choice of $\Delta\chi^2$ is based on the equation for the cumulative χ^2 distribution. In our case (16 parameters) the $\Delta\chi^2$ value is equal to 18.065. We calculate the uncertainties for both $\Delta\chi^2 = 1$ and $\Delta\chi^2 = 18.065$ options. The first moments of PDFs together with their uncertainties are presented in Table 2.

Let us now discuss the obtained parameterization. First, one can see that the results on the first moments $\Delta_1\Sigma \equiv \eta_{\Delta\Sigma}$ and $\Delta_1G \equiv \eta_{\Delta G}$ are very close to the respective results (scenario with $\Delta G < 0$) obtained in Ref. [5] in the case of pure inclusive DIS. What concerns nonsinglet combinations $\Delta q_3(x)$ and $\Delta q_8(x)$, we see (compare Table 1 with Table 3 from Ref. [5]) that, even despite the small difference in initial scales ($Q_0^2 = 1 \text{ GeV}^2$ here and $Q_0^2 = 3 \text{ GeV}^2$ in Ref. [5]), the values of $\Delta q_3(x)$ are very similar for both parameterizations. At the same time, the values of $\Delta q_8(x)$ differ much more essentially, which occurs due to the impact (first of all on Δs – see below) of the additional SIDIS data.

Let us now discuss the impact of SIDIS data on the polarized strangeness in nucleon. It is of importance because, after the appearance of the first results on Δs extraction

Table 2: Estimations of the uncertainties on the first moments of polarized PDFs for two options of $\Delta\chi^2$ choice.

	$\Delta\chi^2 = 1$	$\Delta\chi^2 = 18.065$
$\Delta\Sigma$	$0.3846^{+0.0050}_{-0.0122}$	$0.3846^{+0.0342}_{-0.0389}$
$\Delta u + \Delta\bar{u}$	$0.8640^{+0.0028}_{-0.0049}$	$0.8640^{+0.0114}_{-0.0084}$
$\Delta d + \Delta\bar{d}$	$-0.4020^{+0.0028}_{-0.0048}$	$-0.4020^{+0.0115}_{-0.0130}$
$\Delta s = \Delta\bar{s}$	$-0.0387^{+0.0014}_{-0.0024}$	$-0.038738^{+0.0061}_{-0.0065}$
ΔG	$-0.1828^{+0.0720}_{-0.1090}$	$-0.1828^{+0.1693}_{-0.2831}$
$\Delta\bar{u}$	$0.0672^{+0.0263}_{-0.0270}$	$0.0672^{+0.06483}_{-0.0737}$
$\Delta\bar{d}$	$-0.0792^{+0.0191}_{-0.0238}$	$-0.0792^{+0.0795}_{-0.0830}$

from SIDIS data performed by HERMES [1], we met the puzzle with the positive Δs in the middle x HERMES region $0.023 < x < 0.6$, while the total moment $\Delta_1 s$ definitely should be negative in accordance with the sum rule for Δq_8 . Thus, to satisfy this requirement, $\Delta s(x)$ should possess the compensating negative behavior in the inaccessible for HERMES low x region $0 < x < 0.023$, i.e. the sign-changing scenario for Δs should be realized. Looking at Fig. 1 we see that this is indeed the case and we produce the best fit namely for the sign-changing Δs scenario, as well as in Ref. [4]. However, while within parametrization [4] Δs changes the sign one time, within our parametrization Δs changes the sign twice. It seems that this distinction occurs due to the inclusion of the latest COMPASS semi-inclusive data [3], which allow to better fix¹ Δs shape. This is illustrative to compare the NLO results on Δs obtained here with the respective results of direct Δs extraction in LO by COMPASS – see Fig. 1. We see very similar Δs behavior in both cases.

As it was discussed above, the exclusive task for SIDIS is the finding of sea $\Delta\bar{q}$ (and, consequently, valence $\Delta q_V = \Delta q - \Delta\bar{q}$) PDFs in nucleon. Comparing the central values of $\Delta\bar{u}$ and $\Delta\bar{d}$ with [4] we find that they are quite similar. At the same time, there are also some distinctions. The point is that recently, analyzing the SIDIS data on $h^\pm$ production COMPASS obtained rather surprising result [11] that the sum $[\Delta_1\bar{u} + \Delta_1\bar{d}](Q^2 = 10 \text{ GeV}^2)$ is just zero within the errors (see Table 2 in Ref. [11])

$$[\Delta_1\bar{u} + \Delta_1\bar{d}]_{\text{COMPASS}} = 0.0 \pm 0.04 \pm 0.03. \quad (2)$$

This result was confirmed in the subsequent COMPASS paper [3], where sum $\Delta\bar{u}(x, Q^2 = 3 \text{ GeV}^2) + \Delta\bar{d}(x, Q^2 = 3 \text{ GeV}^2)$ of the local PDFs was extracted from the measured asymmetries $A_{1d}, A_{1d}^{\pi^\pm}, A_{1d}^{K^\pm}$ in the region $0.004 < x < 0.3$ (see Fig. 4 in Ref. [3]) and occurs to be about zero in the whole this region (the central values occur in both positive and negative vicinities of zero).

Thus, at least in the leading order (COMPASS analysis) the sum $\Delta\bar{u}(x) + \Delta\bar{d}(x)$ is about zero in the region $3 \text{ GeV}^2 < Q^2 < 10 \text{ GeV}^2$, which sheds new light on our

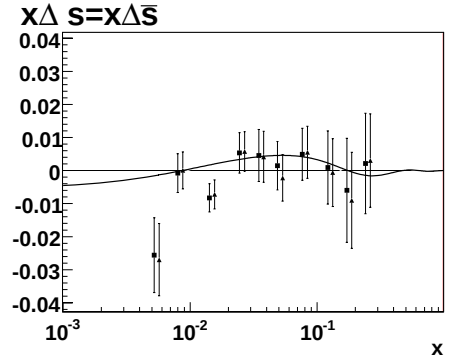


Figure 1: Obtained NLO parameterization on Δs (solid line) in comparison with the COMPASS results [3] obtained in LO QCD (points with error bars), $Q^2 = 3 \text{ GeV}^2$.

¹Notice that after inclusion of the latest COMPASS data to our fit $\chi^2/D.O.F.$ value becomes small only if we allow Δs to change the sign twice (due to the additional parameter $\delta_{\Delta q_8} \sqrt{x}$).

understanding of polarized light quark sea. Namely, the sea is extremely asymmetric ($\Delta\bar{u} \simeq -\Delta\bar{d}$), on the contrary to the assumption of symmetric sea scenario $\Delta\bar{u}(x, Q_0^2) = \Delta\bar{d}(x, Q_0^2)$, applied in the practically all existing parameterizations based on the pure inclusive DIS data analysis. Our analysis shows that the sum $\Delta\bar{u} + \Delta\bar{d}$ is very small quantity in NLO QCD too: $[\Delta_1\bar{u} + \Delta_1\bar{d}](Q^2 = 1 \text{ GeV}^2) = -0.01_{-0.02}^{+0.01}$.

In conclusion, the new combined analysis of polarized DIS and SIDIS data in NLO QCD is presented. The impact of modern SIDIS data on polarized PDFs is studied, which is of especial importance for the light sea quark PDFs and strangeness in nucleon. The obtained results are in agreement with the latest direct leading order COMPASS analysis of SIDIS asymmetries [3] as well as with the recent global fit analysis in NLO QCD of Ref. [4], where the SIDIS data were also applied. Nevertheless, there also some distinctions concerning, first of all, the polarized quark sea peculiarities. At the same time, the quality of SIDIS data is still not sufficient to make the eventual conclusions about the quantities influenced mainly by SIDIS.

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POLARIZED PDFs: REMARKS ON METHODS OF QCD ANALYSIS OF THE DATA

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Abstract

The results on polarized parton densities (PDFs) obtained using different methods of QCD analysis of the present polarized DIS data are discussed. It is pointed out that the precise data in the *preasymptotic* region require a more careful matching of the QCD predictions to the data in this region in order to determine the polarized PDFs correctly.

In this talk we discuss how the results on polarized PDFs depend on the method used in the QCD analysis, accounting or not for the kinematic and dynamic $1/Q^2$ corrections to the spin structure function g_1 .

One of the peculiarities of polarized DIS is that more than a half of the present data are at moderate Q^2 and W^2 ($Q^2 \sim 1 - 4 \text{ GeV}^2$, $4 \text{ GeV}^2 < W^2 < 10 \text{ GeV}^2$), or in the so-called *preasymptotic* region. So, in contrast to the unpolarized case, this region cannot be excluded from the analysis, and the role of the $1/Q^2$ terms (*kinematic* - γ^2 factor, target mass corrections, and *dynamic* - higher twist corrections to the spin structure function g_1) in the determination of the polarized PDFs has to be investigated.

The best manner to determine the polarized PDFs is to perform a QCD fit to the data on g_1/F_1 , which can be obtained if both the $A_{||}$ and $A_{\perp}$ asymmetries are measured. The data on the photon-nucleon asymmetry A_1 are not suitable for the determination of PDFs because the structure function g_2 is not well known in QCD and the approximation

$$(A_1)^{theor} = g_1/F_1 - \gamma^2 g_2/F_1 \approx (g_1/F_1)^{theor} \quad (1)$$

used by some of the groups in the *preasymptotic* region, where γ^2 cannot be neglected, is *not* reasonable.

In QCD, one can split g_1 and F_1 into leading (LT) and dynamical higher twist (HT) pieces

$$g_1 = (g_1)_{LT,TMC} + (g_1)_{HT}, \quad F_1 = (F_1)_{LT,TMC} + (F_1)_{HT}. \quad (2)$$

In the LT pieces in Eq. (2) the calculable target mass corrections (TMC) are included

$$\begin{aligned} g_1(x, Q^2)_{LT,TMC} &= g_1(x, Q^2)_{LT} + g_1(x, Q^2)_{TMC}, \\ F_1(x, Q^2)_{LT,TMC} &= F_1(x, Q^2)_{LT} + F_1(x, Q^2)_{TMC}. \end{aligned} \quad (3)$$

They are inverse powers of Q^2 kinematic corrections, which, however, effectively belong to the LT part of g_1 . Then, approximately

$$\frac{g_1}{F_1} \approx \frac{(g_1)_{\text{LT}}}{(F_1)_{\text{LT}}} \left[1 + \frac{(g_1)_{\text{TMC+HT}}}{(g_1)_{\text{LT}}} - \frac{(F_1)_{\text{TMC+HT}}}{(F_1)_{\text{LT}}} \right]. \quad (4)$$

Note that the LT pieces $(g_1)_{\text{LT}}$ and $(F_1)_{\text{LT}}$ are expressed in terms of the polarized and unpolarized PDFs, respectively. In what follows only the first terms in the TMC and HT expansions will be considered

$$\begin{aligned} g_1(x, Q^2)_{\text{TMC}} &= M^2/Q^2 g_1^{(1)}(x, Q^2)_{\text{TMC}} + \mathcal{O}(M^4/Q^4), \\ F_1(x, Q^2)_{\text{TMC}} &= M^2/Q^2 F_1^{(1)}(x, Q^2)_{\text{TMC}} + \mathcal{O}(M^4/Q^4); \end{aligned} \quad (5)$$

$$g_1(x, Q^2)_{\text{HT}} = h^{g_1}(x)/Q^2 + \mathcal{O}(\Lambda^4/Q^4), \quad 2xF_1(x, Q^2)_{\text{HT}} = h^{2xF_1}(x)/Q^2 + \mathcal{O}(\Lambda^4/Q^4). \quad (6)$$

There are essentially two methods to fit the data - taking or NOT taking into account the HT corrections to g_1 . According to the first [1], the data on g_1/F_1 have been fitted including the contribution of the first term $h(x)/Q^2$ in $(g_1)_{\text{HT}}$ and using the experimental data for the unpolarized structure function F_1

$$\left[\frac{g_1(x, Q^2)}{F_1(x, Q^2)} \right]_{\text{exp}} \Leftrightarrow \frac{g_1(x, Q^2)_{\text{LT,TMC}} + h^{g_1}(x)/Q^2}{F_1(x, Q^2)_{\text{exp}}} \quad (\text{Method I}). \quad (7)$$

According to the second approach [2] only the LT terms for g_1 and F_1 in (2) have been used in the fit to the g_1/F_1 data

$$\left[\frac{g_1(x, Q^2)}{F_1(x, Q^2)} \right]_{\text{exp}} \Leftrightarrow \frac{g_1(x, Q^2)_{\text{LT}}}{F_1(x, Q^2)_{\text{LT}}} \quad (\text{Method II}). \quad (8)$$

It is obvious that the two methods are equivalent in the pure DIS region where TMC and HT can be ignored. To be equivalent in the *preasymptotic* region requires a cancellation between the ratios $(g_1)_{\text{TMC+HT}}/(g_1)_{\text{LT}}$ and $(F_1)_{\text{TMC+HT}}/(F_1)_{\text{LT}}$ in (4). Then $(g_1)_{\text{LT}}$ obtained from the best fit to the data will coincide within the errors independently of the method which has been used. In Fig. 1 these ratios based on our results on target mass [3]

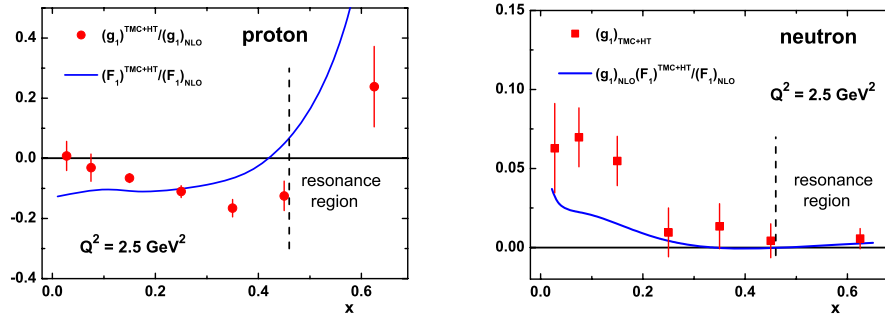


Figure 1: Comparison of the ratios of (TMC+HT)/LT for g_1 and F_1 structure functions for proton and neutron targets (see the text).

and higher twist [4] corrections to g_1 , and the results [5] on the unpolarized structure function F_1 are presented. Note that for the neutron target, $(g_1)_{\text{TMC+HT}}$ is compared with $(g_1)_{\text{LT}} \frac{(F_1)_{\text{TMC+HT}}}{(F_1)_{\text{LT}}}$ because of a node-type behaviour of $(g_1)_{\text{LT}}$. Also, LT means the NLO QCD approximation for both, g_1 and F_1 . As seen from Fig. 1, (TMC+HT) corrections to g_1 and F_1 in the ratio g_1/F_1 do not cancel and ignoring them using the second method is incorrect and will impact on the determination of the polarized PDFs.

Modifications of the second method of analysis in which F_1 is treated in different way are also presented in the literature [6].

How the results on NLO($\overline{\text{MS}}$) polarized PDFs depend on the method used for their determination is discussed in detail in Ref. [7]. To illustrate this dependence here we will compare the NLO LSS'06 set of polarized parton densities [4] determined by Method I with those obtained recently by DSSV [8] using Method II. As expected, the curves corresponding to the ratios $g_1^{\text{tot}}(\text{LSS})/(F_1)_{\text{exp}}$ and $g_1(\text{DSSV})_{\text{NLO}}/F_1(\text{MRST})_{\text{NLO}}$ practically coincide although different expressions were used for g_1 and F_1 in the fit. In Fig. 2 the LSS and DSSV LT(NLO) pieces of g_1 are compared for a proton target. Surprisingly they coincide for $x > 0.1$ although the HT corrections, taken into account in LSS'06 and ignored in the DSSV analysis, do NOT cancel in the ratio g_1/F_1 in this region, as has already been discussed above. The understanding of this puzzle is connected with the fact that in the DSSV fit to all available g_1/F_1 data a factor $(1 + \gamma^2)$ was introduced on the RHS of Eq. (8)

$$\left[\frac{g_1(x, Q^2)}{F_1(x, Q^2)} \right]_{\text{exp}} \Leftrightarrow \frac{g_1(x, Q^2)_{\text{LT}}}{(1 + \gamma^2) F_1(x, Q^2)_{\text{LT}}} . \quad (9)$$

There is no rational explanation for such a correction. The authors point out [9] that it is impossible to achieve a good description of the g_1/F_1 data, especially of the CLAS ones [10], without this correction.

It turns out empirically that the $1/(1+\gamma^2)$ factor accidentally more or less accounts for the TM and HT corrections to g_1 and F_1 in the ratio g_1/F_1 . The relation

$$1 + \frac{(g_1)_{\text{TMC+HT}}}{(g_1)_{\text{LT}}} - \frac{(F_1)_{\text{TMC+HT}}}{(F_1)_{\text{LT}}} \approx \frac{1}{(1 + \gamma^2)} \quad (10)$$

is satisfied with an accuracy between 4% and 18% for the CLAS proton data ($x > 0.1$ and Q^2 between 1 and 4 GeV^2). That is the reason why the LT(NLO) pieces of g_1^p obtained by LSS and DSSV are in a good agreement for $x > 0.1$ (see fig. 2). It

is important to mention that introducing the $(1 + \gamma^2)$ factor does not help at $x < 0.2$ because γ^2 is small in this region and cannot mimic the difference between TM and HT corrections to g_1 and F_1 (LHS of Eq. (10)) for proton as well as for neutron target (see Fig. 1). That is why the LT(NLO) pieces of g_1^p obtained by LSS and DSSV groups differ in this region - the smaller x is, the greater is the difference (see Fig. 2). To summarize: It is impossible to describe the precise data in the preasymptotic region like the CLAS data using the second method. Its empirical modification by introducing the $(1 + \gamma^2)$ factor

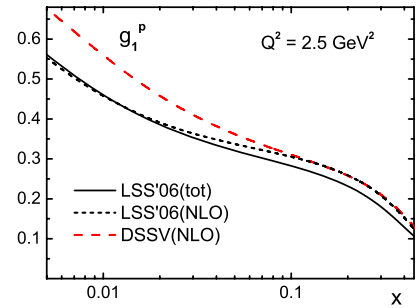


Figure 2: Comparison between LT pieces $g_1(\text{LSS})_{\text{NLO}}$ and $g_1(\text{DSSV})_{\text{NLO}}$.

accounts approximately for the TM and HT effects, but only in the x region: $x > 0.1$, 0.2 for proton and neutron targets, respectively.

The NLO($\overline{\text{MS}}$) PDFs determined from LSS'06 and DSSV analyses are compared in Fig. 3. The AAC'08 PDFs [11] using a modification of the second method of analysis are also presented. Note that all these analyses include the precise CLAS data in the preasptotic region. The results are presented for the sums $(\Delta u + \Delta \bar{u})$ and $(\Delta d + \Delta \bar{d})$ because they can only be separated using SIDIS data (the DSSV analysis). Although the first moments obtained for the PDFs are almost identical, the polarized quark densities themselves are different, especially $\Delta \bar{s}(x)$ (in all the analyses $\Delta s(x) = \Delta \bar{s}(x)$ is assumed).

Let us discuss the impact of HT effects on $(\Delta u + \Delta \bar{u})$ and $(\Delta d + \Delta \bar{d})$ parton densities which should be well determined from the inclusive DIS data. $(\Delta u + \Delta \bar{u})$ extracted by LSS and DSSV are well consistent. As was discussed above the HT effects for the proton target are effectively accounted for in the DSSV analysis for $x > 0.1$ by the introduction of the $(1 + \gamma^2)$ factor. However, this factor cannot account for the HT effects for a neutron target at $x < 0.2$ (see Fig. 1) and the impact of higher twist on $(\Delta d + \Delta \bar{d})$ determined by DSSV is demonstrated in Fig. 3. The positive HT effects are absorbed into $(\Delta d + \Delta \bar{d})_{\text{DSSV}}$ and it is thus less negative in this region. The influence of HT effects at small x (not accounted for by the DSSV group) on both parton densities is not sizable because of two reasons: first, their values are small and second, the data in this region are not precise enough to indicate the impact of higher twist on their values. The impact of HT effects on both $(\Delta u + \Delta \bar{u})_{\text{AAC}}$ and $(\Delta d + \Delta \bar{d})_{\text{AAC}}$ is larger because the AAC Collaboration has not taken them into account at all. The difference between the strange sea densities $\Delta \bar{s}(x, Q^2)_{\text{LSS}}$ and $\Delta \bar{s}(x, Q^2)_{\text{AAC}}$ for $x > 0.1$ is due to the different positivity conditions which have been used by the two groups (see the unpolarized $xs(x)$ in Fig. 3).

In contrast to a negative $\Delta \bar{s}(x, Q^2)$ obtained in all analyses of inclusive DIS data, the DSSV global analysis yields a changing in sign $\Delta \bar{s}(x, Q^2)$: positive for $x > 0.03$ and negative for small x. Its first moment is negative (practically fixed by the SU(3)

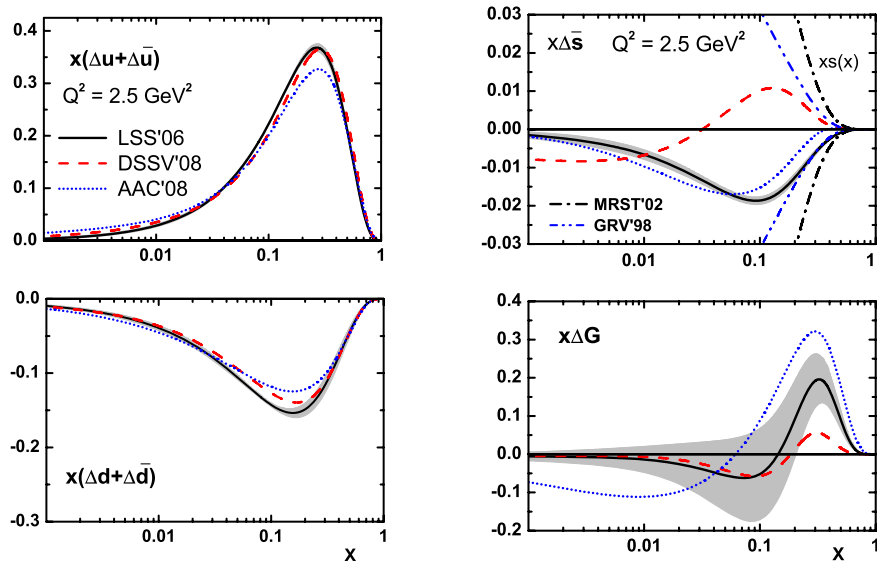


Figure 3: Comparison between LSS'06, DSSV and AAC'08 NLO PDFs in ($\overline{\text{MS}}$) scheme.

symmetric value of a_8) and almost identical with that obtained in the inclusive DIS analyses. It was shown [12] that the determination of $\Delta\bar{s}(x)$ from SIDIS strongly depends on the fragmentation functions (FFs) and the new FFs [13] are crucially responsible for the unexpected behavior of $\Delta\bar{s}(x)$. So, obtaining a final and unequivocal result for $\Delta\bar{s}(x)$ remains a challenge for further research on the internal spin structure of the nucleon.

In conclusion, the fact that more than a half of the present polarized DIS data are in the preasymptotic region makes the QCD analysis of the data more complex and difficult. In contrast to the unpolarized case, the $1/Q^2$ terms (*kinematic* - target mass corrections, and *dynamic* - higher twist corrections to the spin structure function g_1) cannot be ignored, and their role in determining the polarized PDFs is important. Sets of polarized PDFs extracted from the data using different methods of QCD analysis, accounting or not accounting for the kinematic and dynamic $1/Q^2$ corrections, are considered. The impact of higher twist effects on the determination of the parton densities is demonstrated. It is pointed out that the very accurate DIS data in the preasymptotic region require a more careful matching of QCD to the data in order to extract the polarized PDFs correctly.

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POTENTIAL FOR A NEW MUON $g-2$ EXPERIMENT

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Abstract

A new experiment to measure the muon $g - 2$ factor is proposed. We suppose the sensitivity of this experiment to be about 0.03 ppm. The developed experiment can be performed on an ordinary storage ring with a noncontinuous field created by usual magnets. When the total length of straight sections of the ring is appropriate, the spin rotation frequency becomes almost independent of the particle momentum. In this case, a high-precision measurement of an average magnetic field can be carried out with polarized proton beams. A muon beam energy can be arbitrary. Possibilities to avoid betatron resonances are analyzed.

Measurement of the anomalous magnetic moment of the muon is very important because it can in principle bring a discovery of new physics. Experimental data dominated by the BNL E821 experiment, $a_{\mu\pm}^{exp} = 116592080(63) \times 10^{-11}$ (0.54 ppm), are not consistent with the theoretical result, $a_{\mu\pm}^{the} = 116591790(65) \times 10^{-11}$, where $a = (g - 2)/2$. The discrepancy is 3.2σ : $a_{\mu\pm}^{exp} - a_{\mu\pm}^{the} = +290(90) \times 10^{-11}$ [1]. In this situation, the existence of the inconsistency should be confirmed by new experiments. The past BNL E821 experiment [2] was based on the use of electrostatic focusing at the “magic” beam momentum $p_m = mc/\sqrt{a}$ ($\gamma_m = \sqrt{1 + 1/a} \approx 29.3$). An upgraded (but not started up) experiment, E969 [3], with goals of $\sigma_{syst} = 0.14$ ppm and $\sigma_{stat} = 0.20$ ppm is based on the same principle.

Since the muon $g - 2$ experiment is very important, a search for new methods of its performing is necessary. One of new methods has been proposed by Farley [4]. Its main distinctions from the usual $g - 2$ experiments are *i*) noncontinuous magnetic field which is uniform into circular sectors, *ii*) edge focusing, and *iii*) measurement of an average magnetic field with polarized proton beams instead of protons at rest. A chosen energy of muons can be different from the “magic” energy. Its increasing prolongs the lab lifetime of muons. As a result, a measurement of muon $g - 2$ at the level of 0.03 ppm appears feasible [4].

In the present work, we develop the ideas by Farley. We adopt his propositions to measure the average magnetic field with polarized proton beams and to use a ring with a noncontinuous field for keeping the spin rotation frequency independent of the particle momentum. We also investigate the most interesting case when the beam energy can be arbitrary. However, we propose to perform the high-precision muon $g - 2$ experiment on an ordinary storage ring with a noncontinuous field created by usual magnets. We show that the independence of the spin rotation frequency on the particle momentum can be reached not only in a continuous uniform magnetic field [2, 3] and a noncontinuous and locally uniform one [4] but also in a usual storage ring with a noncontinuous and nonuniform magnetic field. In the last case, the total length of straight sections of the

ring should be appropriate. We also analyze possibilities to avoid the betatron resonance $\nu_x = 1$ (ν_x is the horizontal tune).

The system of units $\hbar = c = 1$ is used.

Let us consider spin dynamics in a usual storage ring with a noncontinuous magnetic field and magnetic focusing. The general equation for the angular velocity of spin precession in the cylindrical coordinates is given by (see Ref. [5])

$$\boldsymbol{\omega}^{(a)} = -\frac{e}{m} \left\{ a\mathbf{B} - \frac{a\gamma}{\gamma+1} \boldsymbol{\beta}(\boldsymbol{\beta} \cdot \mathbf{B}) + \left(\frac{1}{\gamma^2-1} - a \right) (\boldsymbol{\beta} \times \mathbf{E}) + \frac{1}{\gamma} \left[\mathbf{B}_{\parallel} - \frac{1}{\beta^2} (\boldsymbol{\beta} \times \mathbf{E})_{\parallel} \right] \right\}. \quad (1)$$

Eq. (1) is useful for analytical calculations of spin dynamics with allowance for field misalignments and beam oscillations. This equation does not contain terms depending on an electric dipole moment and other small terms. The sign $\parallel$ denotes a horizontal projection for any vector. The vertical magnetic field, B_z , and the radial electric one, E_ρ , are the main fields in the muon $g-2$ experiment. The spin precession caused by these fields is defined by

$$\omega_z^{(a)} = -\frac{e}{m} \left[aB_z - \left(\frac{1}{\gamma^2-1} - a \right) \beta_\phi E_\rho \right]. \quad (2)$$

Let $\Omega^{(a)}$ denotes the average value of $\omega^{(a)}$. The spin coherence is kept when

$$\frac{d\Omega_z^{(a)}}{dp} = 0. \quad (3)$$

For a storage ring with noncontinuous fields, the quantities B_z and E_ρ should be averaged.

Condition (3) can be satisfied for ordinary storage rings with usual magnets (Fig. 1). If the field created by the magnets is given by $B_z(\rho) = \text{const} \cdot \rho^{-n}$, the field index and betatron tunes into bending sections are equal to

$$n = -\frac{R_0}{B_0} \left(\frac{\partial B_z}{\partial \rho} \right)_{\rho=R_0}, \quad \nu_x = \sqrt{1-n}, \quad \nu_z = \sqrt{n},$$

where $B_0 \equiv B_z(R_0)$, $x = \rho - R_0$. Average angular velocity of spin precession is given by

$$\Omega^{(a)} = \frac{\omega_z^{(a)} \pi \rho}{\pi \rho + L} = -\frac{\pi e a \rho B_z(\rho)}{m(\pi \rho + L)}, \quad (4)$$

where L is a half of the total length of the straight sections (Fig. 1). The muon anomaly is given by

$$a_\mu = \frac{g_p - 2}{2} \frac{m_\mu}{m_p} \frac{\Omega_\mu^{(a)}}{\Omega_p^{(a)}}, \quad (5)$$

where the fundamental constants g_p and m_μ/m_p are measured with a high precision. The magnetic field is the same for muons and protons when they move on the same trajectory. In this case, their momenta coincide.

When the momentum increases ($p > p_0$), the magnetic field becomes weaker, but the time of flight in the magnetic field becomes longer. Condition (3) is satisfied when

$$L = L_0 = \frac{n}{1-n} \pi R_0. \quad (6)$$

where R_0 corresponds to p_0 and B_0 . In this case

$$R_0 = \frac{p_0}{|e|B_0}, \quad \Omega^{(a)} = \Omega_0^{(a)} = -(1-n) \frac{eaB_0}{m} \quad (7)$$

and the following relation takes place:

$$\frac{\Delta C}{C_0} = \frac{\Delta p}{p_0} = (1-n) \frac{x}{R_0},$$

where C is the orbit circumference. As a result, the momentum compaction factor is

$$\alpha = \frac{\Delta C/C_0}{\Delta p/p_0} = 1. \quad (8)$$

The real value of the length of the straight sections, L , can slightly differ from L_0 . The difference between the real and nominal values of the average angular velocity of spin rotation is given by

$$\frac{\Omega^{(a)} - \Omega_0^{(a)}}{\Omega_0^{(a)}} = n \cdot \frac{L - L_0}{L_0} \cdot \frac{p - p_0}{p_0} - \frac{n}{2(1-n)} \cdot \frac{(p - p_0)^2}{p_0^2}. \quad (9)$$

It is important that Eq. (9) does not depend explicitly on B . The first term in the r.h.s. of this equation disappears if we *define* $L_0 = L$. In this case, p_0 is the vertex of a parabola in the momentum space. To find p_0 and adjust the ring lattice, one can make measurements with proton beams. Three measurements with different values of p are sufficient. The average proton momentum can be kept with radio frequency (RF) cavities put into straight sections of the ring. The longitudinal electric field in the cavities does not influence the spin dynamics.

Condition (3) leading to Eq. (8) should not be exactly satisfied. It can be shown that the relation $\alpha = 1$ leads to the betatron resonance $\nu_x = 1$ which results in zeroth frequency of horizontal coherent betatron oscillation (CBO) of the beam as a whole and a loss of the beam [6]. Therefore, the total length of straight sections should slightly differ from L_0 so that the CBO tune would be small but nonzero:

$$\nu_{CBO} \equiv |1 - \nu_x| = \left| 1 - \sqrt{1 + \frac{L - L_0}{L_0} n} \right| \ll 1. \quad (10)$$

We expect that the CBO tune about 0.01 is sufficient to keep the beam. In this case, the appropriate choice of the total length of straight sections $(L - L_0)n/L_0 \sim 0.01$ reduces the dependence of the spin rotation frequency on the beam momentum by two orders of magnitude. As a result, the use of proton beams for measuring the average magnetic field becomes quite possible.

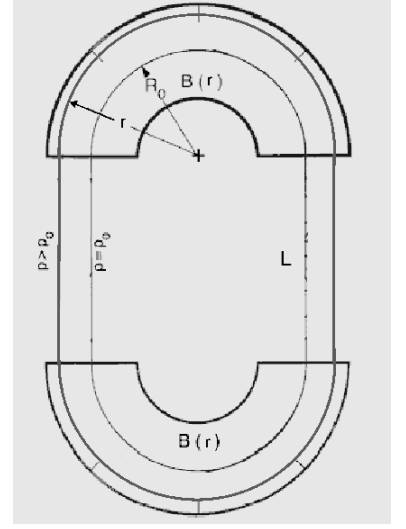


Figure 1: The storage ring.

Experimental details depend on the beam momentum. If it is higher than in the completed experiment, the muon lifetime in the laboratory frame increases and the RF cavities may be helpful not only for protons but also for muons to keep the spin coherence. Otherwise, the use of low muon momentum (~ 0.3 GeV/ c) and much higher statistics [7] may even be more preferable. In this case, the RF cavities are unnecessary for muons.

The problem of systematical errors is very important. A noncontinuous vertical magnetic field leads to a longitudinal magnetic field. It was asserted in Ref. [8] that this circumstance causes “the need to know $\int \mathbf{B} \cdot d\mathbf{l}$ for the muons to a precision of 10 ppb”. However, we should take into account that the longitudinal magnetic field cannot be neglected only on small segments of the beam trajectory near edges of magnets. As a result, the above estimate should be decreased by about 3 orders of magnitude. In addition, the field of magnets is well-known and $\oint \mathbf{B} \cdot d\mathbf{l} = 0$. All needed corrections can be properly calculated with spin tracking. The corrections and systematical errors in the proposed experiment and the Farley’s one are very similar. Therefore, we suppose the sensitivity of the proposed experiment and the Farley’s one to be approximately the same (about 0.03 ppm). Since the developed experiment can be carried out with usual magnets and does not need strong efforts for adjusting the magnetic field, it must be comparatively cheap.

As the theoretical predictions and the experimental data do not agree, performing new experiments based on different ring lattices is necessary. Such experiments will be very important even if they will not provide better precision as compared with the usual $g-2$ experiments [2, 3].

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EVOLUTION EQUATIONS FOR TRUNCATED MELLIN MOMENTS OF THE PARTON DENSITIES

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Abstract

Evolution equations for truncated Mellin moments of the parton distributions in the nucleon are presented. In this novel approach the equations are diagonal and exact for each n th moment and for every truncation point $x_0 \in (0; 1)$. They have the same form as those for the partons themselves. The modified splitting function for n th truncated moment $P'(n, x)$ is equal to $x^n P(x)$, where $P(x)$ is DGLAP splitting function for the partons. The evolution equations for truncated moments can be used in different approximations: LO, NLO etc. and for polarised as well as unpolarised densities. Presented approach can be an additional useful tool in the PQCD analysis of the nucleon structure functions.

1 Introduction

Unpolarised and polarised parton distribution functions as well as their Mellin moments are nowadays the subject of intensive theoretical and experimental studies. Usually, the central role in the perturbative QCD analysis is played by the parton distribution functions (PDFs), which obey the well-known DGLAP evolution equations [1]. However recently, moments have become a great of importance and now they are also a powerful tool in studying the structure functions of the nucleon. An approach based on the moments has many advantages. Moments directly refer to sum rules - fundamental relations in QCD. They provide knowledge about contributions to momentum or spin of the nucleon coming from quarks and gluons. This is essential in resolving the ‘spin puzzle’. We propose a novel - truncated Mellin moments approach (TMMA) with evolution equations, which is a generalization of the full (untruncated) moments case and additionally offers new possibilities. From the one hand, TMMA enables one directly to study the evolution of physical values, and from the other hand, allows to avoid uncertainties from non-available experimentally x -regions. The idea of truncated moments of parton densities was introduced in [2]. The authors obtained non-diagonal evolution equations, where each n th truncated moment couples to all higher moments. Then, truncated moments were applied in $\ln^2 x$ approximation, where we obtained the diagonal solutions [3], and also in NLO analysis of SIDIS data with use of polynomial expansion methods [4]. In [5], [6] we derived within DGLAP approach diagonal and exact evolution equations for truncated moments in the case of single and double truncation as well. Here we present this promising treatment, its advantages, possible applications and perspectives.

2 Evolution equations for truncated moments

Truncated Mellin moments of the parton densities obey the DGLAP-like evolution equations [5] (for clarity we show only the nonsinglet part):

$$\frac{d\bar{q}_n(x_0, Q^2)}{d \ln Q^2} = \frac{\alpha_s(Q^2)}{2\pi} (P' \otimes \bar{q}_n)(x_0, Q^2), \quad (1)$$

where $q(x, Q^2)$ is the parton distribution function and $\bar{q}_n(x_0, Q^2)$ denotes its n th moment truncated at x_0 :

$$\bar{q}_n(x_0, Q^2) = \int_{x_0}^1 dx x^{n-1} q(x, Q^2). \quad (2)$$

A role of the splitting function plays

$$P'(n, z) = z^n P(z). \quad (3)$$

Also the double truncated moments

$$\bar{q}_n(x_{min}, x_{max}, Q^2) = \int_{x_{min}}^{x_{max}} dx x^{n-1} q(x, Q^2) \quad (4)$$

fulfill the DGLAP-type evolution [6]. The main advantage of this approach is that the equations are exact and diagonal (no mixing between moments of different orders) and hence they can be solved with use of standard methods of solving the DGLAP equations. Furthermore, dealing with truncated moments, one can use a wide range of experimental data in terms of smaller number of free parameters. In this way, no assumptions on the shape of parton distributions are needed. TMMA enables also one to avoid uncertainties from the unmeasurable very small $x \rightarrow 0$ and high $x \rightarrow 1$ regions and can be used for different approximations: LO, NLO etc. and in the polarised as well as unpolarised case.

3 Possible applications

TMMA, which refers to the physical values - moments (not to the parton densities), allows one to study directly their evolution and the scaling violation. The solutions for truncated moments can be also used in the determination of the parton distribution functions via differentiation

$$q(x, Q^2) = -x^{1-n} \frac{\partial \bar{q}_n(x, Q^2)}{\partial x}, \quad (5)$$

In order to reconstruct initial parton densities at scale Q_0^2 from their truncated moments, given e.g. by experimental data at scale Q^2 , we evolve moments between these two scales (from Q^2 to Q_0^2) and then perform the final fit of free parameters - for details see [6].

We tested this method on the nonsinglet function parametrised in a general form

$$q(x, Q_0^2) = N(\alpha, \beta, \gamma) x^\alpha (1-x)^\beta (1+\gamma x), \quad (6)$$

and also on the original fits for HERMES and COMPASS data.

From Figs. 1-2 one can see satisfactory agreement, even for the limited x -range of the data. Due to its large- x sensitivity, the second moment can be used in the precise final reconstruction of the parton density.

4 Perspectives

Evolution equations for the truncated Mellin moments of the parton densities can be a promising tool in the QCD analysis of the nucleon structure functions. Here we list a few of the valuable future applications of TMMA:

- Studying the fundamental properties of nucleon structure, concerning moments of F_1 , F_2 and g_1 . These are: momentum fraction carried by quarks, quark helicities

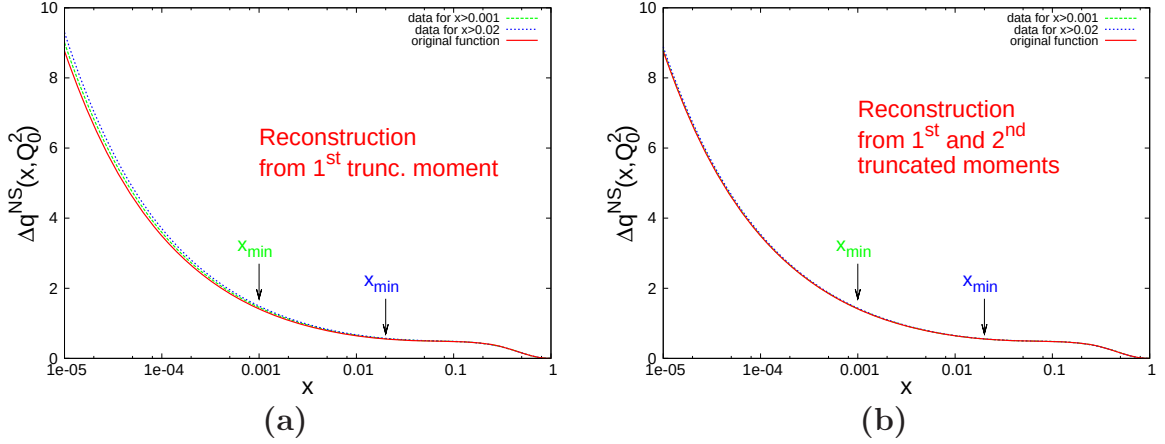


Figure 1: **(a)** Reconstruction of the initial parton density at $Q_0^2 = 1 \text{ GeV}^2$ from the first truncated moment given at $Q^2 = 10 \text{ GeV}^2$ in the broad range $0.001 \leq x \leq 0.7$ (dashed) and in the limited one $0.02 \leq x \leq 0.7$ (dotted). Comparison with the original function (6) (solid), where $\alpha = -0.4$ reflects the theoretical knowledge on the small- x behaviour - see e.g. [7]. **(b)** Similarly as in **(a)** but after additional precise reconstruction of the parameter γ (6) from the second truncated moment.

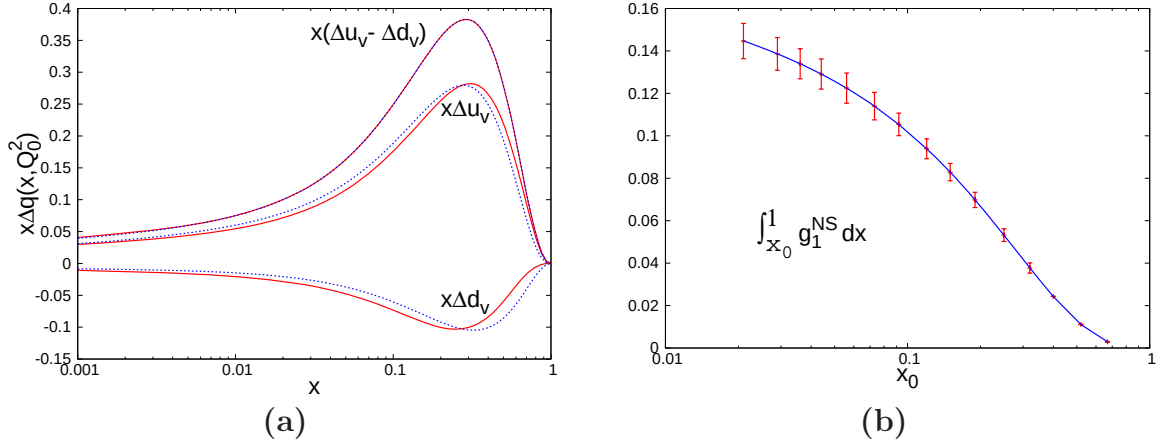


Figure 2: **(a)** Reconstruction of the initial parton densities from HERMES data for the first truncated moment of g_1^{NS} at $Q^2 = 5 \text{ GeV}^2$ [9] (dotted) in comparison with original BB fit [8]. Plots for $x(\Delta u_v - \Delta d_v)$ overlap each other. **(b)** The first truncated moment of the nonsinglet polarised structure function g_1^{NS} vs the truncation point x_0 , calculated from the reconstructed fit (solid). $Q^2 = 5 \text{ GeV}^2$. Comparison with HERMES data [9] based on BB fit [8] (points with error bars).

contributions to the spin of nucleon and, what is particularly important, estimation of the polarised gluon contribution ΔG from COMPASS and RHIC data.

- Determination of Higher Twist (HT) effects from moments of g_2 , which will be measured at JLab. This is possible via generalization of Wandzura-Wilczek relation [10] for truncated moments and test of Burkhardt-Cottingham [11] and Efremov-Leader-Teryaev [12] sum rules. HT corrections can provide information on the quark-hadron duality.
- Predictions for generalized parton distributions (GPDs). Moments of GPDs can be related to the total angular momentum (spin and orbital) carried by various quark flavors. Measurements of DVCS, sensitive to GPDs, will be done at JLab. This would be an important step towards a full accounting of the nucleon spin.

Concluding, in light of the recent progress in experimental program, the comprehensive theoretical analysis of the structure functions and their moments is of a great importance.

I thank Anatoly Efremov for inviting me to this meeting and warm hospitality.

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NEW DEVELOPMENTS IN THE QUANTUM STATISTICAL APPROACH OF THE PARTON DISTRIBUTIONS

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Abstract

We briefly recall the main physical features of the parton distributions in the quantum statistical picture of the nucleon. Some predictions from a next-to-leading order QCD analysis are compared to recent experimental results.

A new set of parton distribution functions (PDF) was constructed in the framework of a statistical approach of the nucleon [1], which has the following characteristic features:

- For quarks (antiquarks), the building blocks are the helicity dependent distributions $q_{\pm}(x)$ ($\bar{q}_{\pm}(x)$) and we define $q(x) = q_+(x) + q_-(x)$ and $\Delta q(x) = q_+(x) - q_-(x)$ (similarly for antiquarks).
- At the initial energy scale taken at $Q_0^2 = 4\text{GeV}^2$, these distributions are given by the sum of two terms, a quasi Fermi-Dirac function and a helicity independent diffractive contribution, which leads to a universal behavior at very low x for all flavors.
- The flavor asymmetry for the light sea, *i.e.* $\bar{d}(x) > \bar{u}(x)$, observed in the data is built in. This is clearly understood in terms of the Pauli exclusion principle, based on the fact that the proton contains two u quarks and only one d quark.
- The chiral properties of QCD lead to strong relations between $q(x)$ and $\bar{q}(x)$. For example, it is found that the well established result $\Delta u(x) > 0$ implies $\Delta \bar{u}(x) > 0$ and similarly $\Delta d(x) < 0$ leads to $\Delta \bar{d}(x) < 0$.
- Concerning the gluon, the unpolarized distribution $G(x, Q_0^2)$ is given in terms of a quasi Bose-Einstein function, with only *one free parameter*, and for simplicity, one assumes zero gluon polarization, *i.e.* $\Delta G(x, Q_0^2) = 0$, at the initial energy scale Q_0^2 .
- All unpolarized and polarized light quark distributions depend upon *eight* free parameters, which were determined in 2002 (see Ref. [1]), from an NLO fit of a selected set of accurate DIS data.

More recently, new tests against experimental (unpolarized and polarized) data turned out to be very satisfactory, in particular in hadronic reactions, as reported in Refs. [2–4]. The statistical approach has been extended [5] to the interesting situation where the PDF have, in addition to the usual Bjorken x dependence, an explicit k_T transverse momentum

dependence (TMD) and this might be used in future calculations with no k_T integration. This will not be treated here for lack of time, although it is a very important topic, with a growing interest because it is now clear that several new phenomena are sensitive to TMD effects.

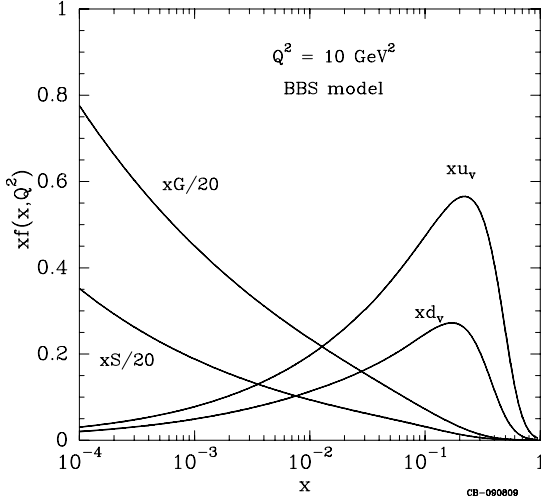


Figure 1: BBS predictions for various statistical parton distributions versus x , at $Q^2 = 10 \text{ GeV}^2$.

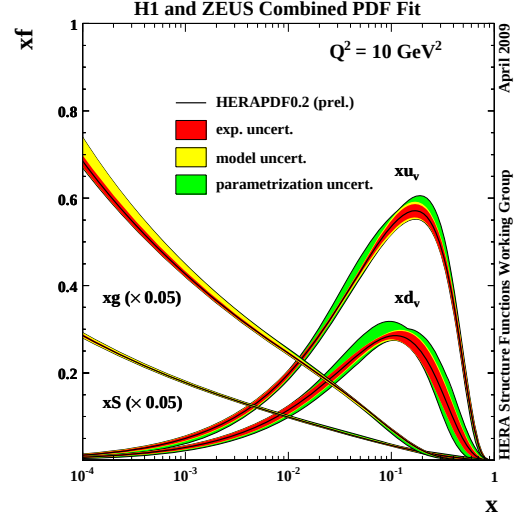


Figure 2: Parton distributions at $Q^2=10 \text{ GeV}^2$ as determined by the H1PDF fit, with different uncertainties (Taken from Ref. [6]).

We display on Fig. 1 the resulting unpolarized statistical PDF versus x at $Q^2=10 \text{ GeV}^2$, where xu_v is the u -quark valence, xd_v the d -quark valence, xG the gluon and xS stands for twice the total antiquark contributions, *i.e.* $xS(x) = 2x(\bar{u}(x) + \bar{d}(x) + \bar{s}(x)) + \bar{c}(x)$. Note that xG and xS are downscaled by a factor 0.05. They can be compared with the parton distributions as determined by the H1PDF 2009 QCD NLO fit, shown in Fig. 2, and the agreement is rather good. The results are based on recent ep collider data combined with previously published data and the accuracy is typically in the range of 1.3 - 2 %.

In the statistical approach the unpolarized gluon distribution has a very simple expression given by $xG(x, Q_0^2) = \frac{A_G x^{b_G}}{\exp(x/\bar{x}) - 1}$, where $\bar{x} = 0.099$ is the universal temperature, $A_G = 20.53$ is determined by the momentum sum rule and $b_G = 0.90$ is the only free parameter. It is consistent with the available data, mostly coming indirectly from the QCD Q^2 evolution of $F_2(x, Q^2)$, in particular in the low x region. However it is known that ep DIS cross section is characterized by two independent structure functions, $F_2(x, Q^2)$ and the longitudinal structure function $F_L(x, Q^2)$. For low Q^2 , the contribution of the later to the cross section at HERA is only sizeable at x smaller than approximately 10^{-3} and in this domain the gluon density dominates over the sea quark density. More precisely, it was shown that using some approximations [7] one has, $xG(x, Q^2) = \frac{3}{10} 5.9 [\frac{3\pi}{2\alpha_s} F_L(0.4x, Q^2) - F_2(0.8x, Q^2)] \simeq \frac{8.3}{\alpha_s} F_L(0.4x, Q^2)$. Before HERA was shut down, a dedicated run period, with reduced proton beam energy, was approved, allowing H1 to collect new results on F_L . We show on Fig. 3 the expectations of the

statistical approach compared to the new data, whose precision is reasonable. The trend and the magnitude of the prediction are in fair agreement with the data, so this is another test of the good predictive power of our theoretical framework.

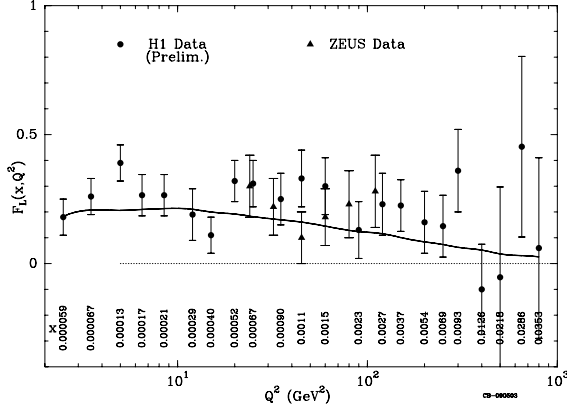


Figure 3: The longitudinal proton structure function $F_L(x, Q^2)$ averaged in x at given values of Q^2 . Data from [8,9] compared to the BBS theoretical prediction.

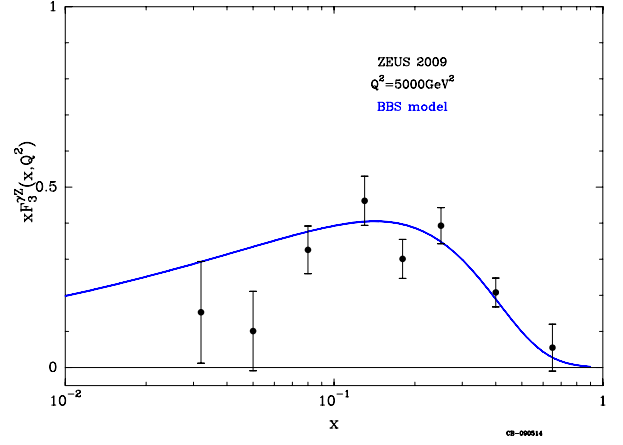


Figure 4: The interference term $xF_3^{\gamma Z}$ extracted in $e^\pm p$ collisions at HERA. Data from [10] compared to the BBS prediction.

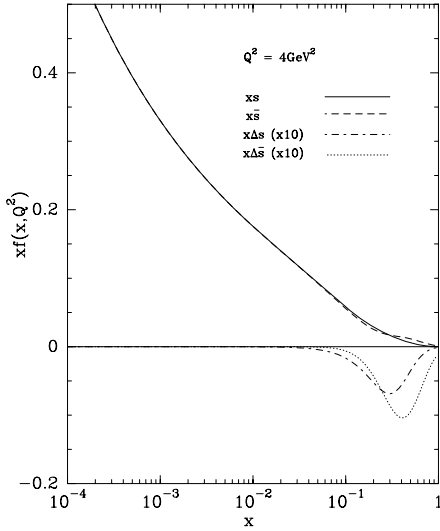


Figure 5: Strange quark and antiquark distributions determined at NLO.

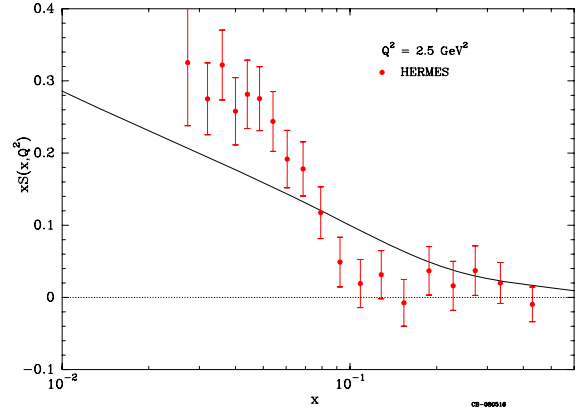


Figure 6: The strange parton distribution $xS(x) = xs(x) + x\bar{s}(x)$ at $Q^2=2.5 \text{ GeV}^2$. The theoretical prediction from [11] is compared to data from [12].

One can also test the behavior of the interference term between the photon and the Z exchanges, which can be isolated in neutral current $e^\pm p$ collisions at high Q^2 . We have to a good approximation, if sea quarks are ignored, $xF_3^{\gamma Z} = \frac{x}{3}(2u_v + d_v)$ and the comparison between data and prediction is displayed in Fig. 4.

Concerning the strange quark and antiquark distributions, a careful analysis of the NuTeV CCFR data led us to the conclusion that $s(x, Q^2) \neq \bar{s}(x, Q^2)$ and the corresponding polarized distributions are unequal, small and negative [11], as shown in Fig. 5. The rapid rise one observes in the small x region is compatible with the data from Hermes, as shown in Fig. 6.

Finally let us recall that the subject of quark and antiquark transversity distribution in the proton is also a very interesting topic. By studying a possible connection between helicity and transversity, we have proposed a simple toy model (see Ref. [13]).

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POLARIZED HADRON STRUCTURE IN THE VALON MODEL AND THE NUCLEON AXIAL COUPLING CONSTANTS: a_3 AND a_8

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Abstract

We have utilized the concept of valon model to calculate the spin structure function of the nucleon. The structure of the valon itself developed through the perturbative dressing of a valence quark in QCD, which is independent of the hosting hadron. In this paper we calculate the nucleon axial coupling constants, a_3 and a_8 in the valon framework. We compare the results with experimental data and find good agreement between them.

1 Polarized hadron structure in the Valon model

One of the fundamental properties of the nucleon is its spin. The spin structure of nucleon has been the subject of heated debates over the past twenty years. The key question is that how the spin of the nucleon is distributed among its constituent partons. Any determination of parton distribution function in a nucleon in the framework of QCD always involves some model-dependent procedures.

We utilized the valon model [1] to study polarized nucleon structure. In the valon model, it is assumed that a nucleon is composed of three dressed valence quarks: the valons. Each valon has its own internal structure which can be probed at high enough Q^2 . At low Q^2 , a valon behaves as a valence quark. The valons play a role in scattering problems as the constituents do in bound state problems. It is assumed that the valons stand at a level in between hadrons and partons and that the valon distributions are independent of the probe or Q^2 . In this representation a valon is viewed as a cluster of its own partons. The evolution of the parton distributions in a hadron is effected through the evolution of the valon structure, as the higher resolution of a probe reveals the parton content of the valon. The valon model is essentially a two component model. In this framework, the helicity distributions of various partons in a hadron are given by:

$$\delta q_i^h(x, Q^2) = \sum \int_x^1 \frac{dy}{y} \delta G_{valon}^h(y) \delta q_i^{valon}\left(\frac{x}{y}, Q^2\right) \quad (1)$$

where $\delta G_{valon}^h(y)$ is the helicity distribution of the valon in the hosting hadron. It is related to unpolarized valon distribution by:

$$\delta G_j(y) = \delta F_j(y) G_j(y). \quad (2)$$

$G_j(y)$ are the unpolarized valon distributions, where j refers to U and D type valons [2]. The functions $\delta F_j(y)$ are given in [3]. $\delta q_i^{valon}(z = x/y, Q^2)$ is the polarized parton distribution in the valon. Polarized parton distributions inside the valon are evaluated according to the DGLAP evolution equation subject to physically sensible initial conditions [3]. In the valon model the hadron structure is obtained by the convolution of valon structure and its distribution inside the hadron. Having specified the various components that contribute to the spin of a valon, we now turn to the polarized hadron structure, which is obtained by a convolution integral as follows:

$$g_1^h(x, Q^2) = \sum_{valon} \int_x^1 \frac{dy}{y} \delta G_{valon}^h(y) g_1^{valon}\left(\frac{x}{y}, Q^2\right). \quad (3)$$

The valon structure is generated by perturbative dressing in QCD. In such processes with massless quarks, helicity is conserved and therefore, the hard gluons can not induce sea quark polarization perturbatively. According to this description, it turns out that sea polarization is consistent with zero [3, 4]. This finding is supported by HERMES experiment and by the newly released data from COMPASS experiment [5–7]. There is an excellent agreement between the model predictions with the experimental data for spin structure functions. A sample is given in Figure 1.

2 Axial coupling constants: a_3 and a_8

The axial coupling constant can be write as a combination of PPDFs as follows:

$$\begin{aligned} g_A^3 &\equiv \langle p, s | \bar{u} \gamma_\mu \gamma_5 u - \bar{d} \gamma_\mu \gamma_5 d | p, s \rangle s^\mu = \Delta u(Q^2) - \Delta d(Q^2), \\ g_A^8 &\equiv \langle p, s | \bar{u} \gamma_\mu \gamma_5 u + \bar{d} \gamma_\mu \gamma_5 d - 2\bar{s} \gamma_\mu \gamma_5 s | p, s \rangle s^\mu = \Delta u(Q^2) + \Delta d(Q^2) - 2\Delta s(Q^2). \end{aligned} \quad (4)$$

In the model described in the section 1, we showed that [3, 4], the calculations in the framework of NLO perturbation theory, shows that the sea quark contribution to the spin of proton is essentially consistent with zero, then we have:

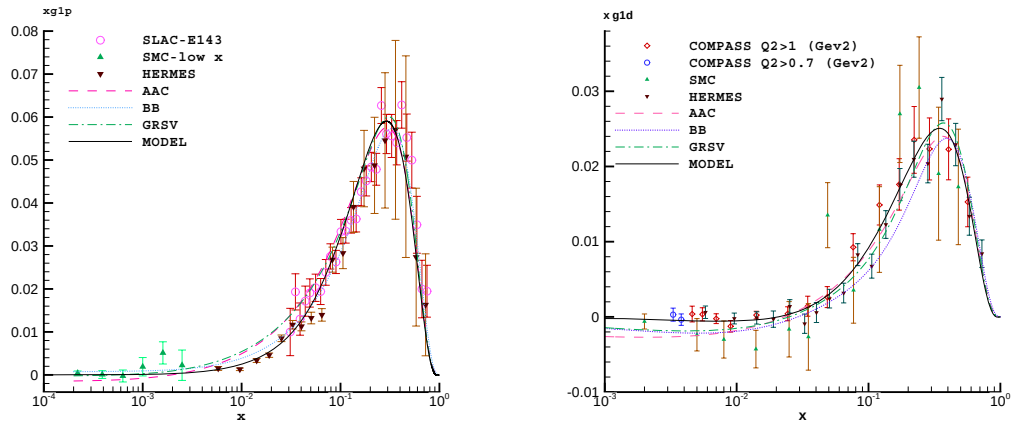


Figure 1: *Left:* Polarized proton structure function, xg_1^p at $Q^2 = 5 \text{ GeV}^2$. *Right:* Polarized deuteron structure function, xg_1^d at $Q^2 = 3 \text{ GeV}^2$. The results from model compared with the global fits [9–11] and the experimental data [5, 7, 12–14]

$$\begin{aligned} g_A^3 &= \Delta u_v(Q^2) - \Delta d_v(Q^2), \\ g_A^8 &= \Delta u_v(Q^2) + \Delta d_v(Q^2). \end{aligned} \quad (5)$$

The experimental values are $g_A^3 = 1.2573 \pm 0.0028$ and $g_A^8 = 0.579 \pm 0.025$ [8]. From the stated values for Δu_v and Δd_v we obtain $g_A^3 = 1.240 - 1.253$ which accommodates the experimental value with an accuracy of 2%. The octet coupling constant, g_A^8 , is equivalent to valence quark polarization in our model and its value is vary between 0.39 - 0.41. this result is in very good agreement with the newly COMPASS results [7]. The value for a_8 is substantially different from 0.579 ± 0.025 , a result that seems confirmed by the emerging experiments. The COMPASS collaboration provided direct information about the valence quark polarization [7]:

$$\Gamma_v(x_{\min}) = \int_{x_{\min}}^{0.7} [\delta u_v(x) + \delta d_v(x)] dx, \quad (6)$$

$$\delta u_v(x) + \delta d_v(x) = \frac{36}{5} \frac{g_1^d(x)}{(1 - 1.5\omega_D)} - \left[2(\delta \bar{u}(x) + \delta \bar{d}(x)) + \frac{2}{5}(\delta s(x) + \delta \bar{s}(x)) \right]. \quad (7)$$

With precise measurement of $g_1^d(x)$, COMPASS obtained the valence quark polarization for the full measured range of x

$$\Gamma_v(0.006 < x < 0.7) = 0.40 \pm 0.07 \text{ (stat.)} \pm 0.06 \text{ (syst.)} \quad (8)$$

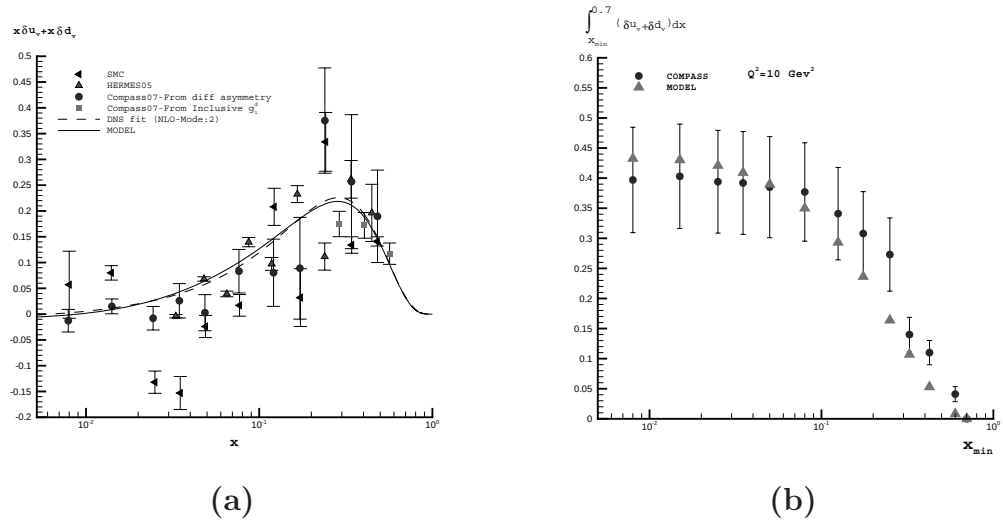


Figure 2: *Left:* $x\delta u_v(x) + x\delta d_v(x)$ at $Q^2 = 10 \left(\frac{\text{GeV}}{c}\right)^2$. *Right:* The integral of $\Delta u_v(x) + \Delta d_v(x)$ over the range $0.006 < x < 0.7$ measured by COMPASS coll. as the function of the low x limit of integration, $x_{\min}$ evaluated at $Q^2 = 10 \left(\frac{\text{GeV}}{c}\right)^2$. The results from the model calculation are compared with the experimental data and DNS global fit [15].

COMPASS collaboration although does not provide direct information about the sea polarization, but clearly indicates that their data is consistent with $\Delta \bar{u}_{sea} + \Delta \bar{d}_{sea} \sim 0$ and suggests that $\Delta \bar{u}_{sea}$ and $\Delta \bar{d}_{sea}$, if different from zero, must be of opposite sign

[7]. Newly released data from HERMES collaboration also shows that the strange sea quark polarization is very small and the helicity distribution is zero within experimental uncertainties [6]. The value that obtained by the HERMES collaboration for g_A^8 is $0.285 \pm 0.046(stat) \pm 0.057(sys)$ [6]. Accepting HERMES results and the data from COMPASS collaboration points to the fact that the role of sea polarization is marginal and consistent with zero. This fact is naturally explained in our model, which relies on the Next-to-Leading order calculations in perturbative QCD. This situation leaves us with the valence and gluon polarization and the orbital angular momentum in the nucleon. The gluon and orbital angular momentum contributions to the spin of proton is discussed in [3]. As a Conclusion, we calculate the axial coupling constant by using the results obtained for PPDFs and nucleon structure functions in the valon framework. In our model sea quark polarization is consistent with zero because we work in the framework of QCD and in these processes the sea quark polarization can not produces perturbatively. The results obtained for g_A^3 and g_A^8 have very good agreement with new released data from HERMES and COMPASS.

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AXIAL ANOMALIES, NUCLEON SPIN STRUCTURE AND HEAVY IONS COLLISIONS.

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Abstract

The axial (related to axial anomaly) and vector currents of heavy quarks are considered. The special attention is paid to the strangeness polarization mediated by gluon anomaly and treatment of the strange quarks in a heavy ones in a multiscale nucleon. It is shown that the straightforward modification of Heisenberg-Euler effective lagrangian allows to calculate the vector current of strange quarks and describes an analog of Chiral Magnetic Effect for strange and heavy quarks.

1. Introduction. The spin structure of nucleon is a major problem since EMC Spin crisis (puzzle) emerged in 80's.

The first observation [1] was related to the role of gluon anomaly which was interpreted as a (circular) gluon polarization. The extensive experimental investigations at HERMES, COMPASS and RHIC, however, did not find the significant polarization. Sometimes this is described as an "absence" of anomaly which is quite strange because of the fundamental character of this phenomenon. Because of this, I am going to discuss the manifestation of anomaly through strange quarks polarization mediated by very small polarizations of off-shell gluons.

2. Axial current of strange quarks, gluonic anomaly and strangeness polarization. The divergence of the singlet axial current contains a normal and an anomalous piece,

$$\partial^\mu j_{5\mu}^{(0)} = 2i \sum_q m_q \bar{q} \gamma_5 q - \left(\frac{N_f \alpha_s}{4\pi} \right) G_{\mu\nu}^a \tilde{G}^{\mu\nu, a}, \quad (1)$$

where N_f is the number of flavours. The two terms at the r.h.s. of the last equation are known to cancel in the limit of infinite quark mass. This is the so-called cancellation of physical and regulator fermions, related to the fact, that the anomaly may be regarded as a usual mass term in the infinite mass limit, up to a sign, resulting from the subtraction in the definition of the regularized operators.

Consequently, one should expect, that the contribution of infinitely heavy quarks to the first moment of g_1 is zero. This is exactly what happens [2] in a perturbative calculation of the triangle anomaly graph. One may wonder, what is the size of this correction for large, but finite masses and how does it compare with the purely perturbative result. To answer this question, one should calculate the r.h.s. of (1) for heavy fermions. The leading coefficient is of the order m^{-2} ,

$$\partial^\mu j_{5\mu}^c = \frac{\alpha_s}{48\pi m_c^2} \partial^\mu R_\mu \quad (2)$$

where

$$R_\mu = \partial_\mu \left(G_{\rho\nu}^a \tilde{G}^{\rho\nu,a} \right) - 4 (D_\alpha G^{\nu\alpha})^a \tilde{G}_{\mu\nu}^a . \quad (3)$$

The contribution [3] of heavy (say, charm) quarks to the nucleon forward matrix element is

$$\langle N(p, \lambda) | j_{5\mu}^{(c)}(0) | N(p, \lambda) \rangle = \frac{\alpha_s}{48\pi m_c^2} \langle N(p, \lambda) | R_\mu(0) | N(p, \lambda) \rangle \quad (4)$$

Note that the first term in R_μ does not contribute to the forward matrix element because of its gradient form, while the contribution of the second one is rewritten, by making use of the equation of motion, as matrix element of the operator

$$\begin{aligned} \langle N(p, \lambda) | j_{5\mu}^{(c)}(0) | N(p, \lambda) \rangle &= \frac{\alpha_s}{12\pi m_c^2} \langle N(p, \lambda) | g \sum_{f=u,d,s} \bar{\psi}_f \gamma_\nu \tilde{G}_\mu{}^\nu \psi_f | N(p, \lambda) \rangle \\ &\equiv \frac{\alpha_s}{12\pi m_c^2} 2m_N^3 s_\mu f_S^{(2)} , \end{aligned} \quad (5)$$

The parameter $f_S^{(2)}$ appears in calculations of the power corrections to the first moment of the singlet part of g_1 part of which is given by exactly the quark-gluon-quark matrix element we got. The non-perturbative calculations are resulting in the estimate $\bar{G}_A^c(0) = -\frac{\alpha_s}{12\pi} f_S^{(2)} \left(\frac{m_N}{m_c}\right)^2 \approx -5 \cdot 10^{-4}$. The seemingly naive application of this approach to the case of strange quarks was presented already ten years from now [3] giving for their contribution to the first moment of g_1 roughly $-5 \cdot 10^{-2}$, which is compatible with the experimental data which is preserved also now despite the problem of matching DIS and SIDIS analysis.

At that time the reason for such a success which was mentioned in [3] (and emerged due to discussion with Sergey Mikhailov who is participating in this meeting) was the possible applicability of a heavy quark expansions for strange quarks [4, 5] in the case of the vacuum condensates of heavy quarks. That analysis was also related to the anomaly equation for heavy quarks, however, for the trace anomaly, rather than the axial one.

The current understanding may also include what I would call "multiscale" picture of nucleon with (squared) strange quark mass (and Λ) being much smaller than that of nucleon but much larger than (genuine) higher twist parameter. Whether the ("semiclassical" as Maxim Polyakov calls it) smallness of higher twists holds for consecutive terms in the series of higher twists may be checked by use of very accurate JLAB data for g_1 . The simplest case of course is the non-singlet combination related to Bjorken Sum Rule. As higher twists are more pronounced at low Q^2 , one should take care on the Landau singularities which may be achieved by use of Analytic Perturbation Theory (the main experts in which are here) or Simonov's soft freezing. The result [6] looks like a first terms of converging series of higher twists compatible with semiclassical picture.

It is instructive to compare the physical interpretation of gluonic anomaly for massless and massive quarks. While in the former, most popular, case it corresponds to the circular polarization of on-shell gluons (recall, that it is rather small, according to various experimental data), in the case of massive quarks one deals with very small (because of small higher twist strength) correlation of nucleon polarization and a sort of polarization of *off-shell* gluons. As soon as strange quark mass is not very large, it partially compensated the smallness of higher twist and this gluon polarization is transmitted to the non-negligible strange quark polarization. The role of gluonic anomaly is therefore to produce the "anomaly-mediated" strangeness polarization.

3. Vector current of strange quarks and Chiral Magnetic Effect in Heavy Ions Collisions. The calculation of vector rather than axial current of heavy (and strange) quarks appears to be even easier. The answer is actually contained in the classical Heisenberg-Euler effective lagrangian for light-by light scattering. Calculating its variation with respect to the electromagnetic potential one immediately get the expression for the current. The transition from QED to QCD is performed by the substitution of three of quark-photon vertices by the quark-gluon ones. The C-parity ensures that the result contains only the symmetric $SU(3)$ structure d^{abc} and is proportional to Abelian one which was explored earlier [8]. The notion of this current allows to determine the strangeness contribution to the anomalous magnetic moment of the nucleon and to the mean square radius of the pion.

Another interesting manifestation of strange quark vector current emerges if one substitutes [7] the *two* rather than three quark-photon vertices by quark-gluon ones. The results describes the vector current induced by cooperative action of (two) gluonic and electromagnetic fields. Its physical realization corresponds, in particular, to heavy ions collisions where extremely strong magnetic fields are generated. The most interesting contributions comes from the term $(F\tilde{F})^2$ in Heisenberg-Euler lagrangian, leading to the current

$$j_\mu^s = \frac{7\pi\alpha\alpha_s}{45m_s^4} \tilde{F}_{\mu\nu} \partial^\nu (G\tilde{G}) \quad (6)$$

One may easily recognize here the analog of the famous Chiral Magnetic Effect (CME) (see [9] and Ref. therein). The correspondence is manifested by the substitution $\frac{1}{m_s^4} \partial^\nu (G\tilde{G}) \rightarrow \partial^\nu \int d^4z (G\tilde{G}) \rightarrow \partial^\nu \theta$

The latter requires the appearance of two scales, now in multiscale medium, (the ones corresponding to integration and taking the derivative) and leads to the derivative of topological field θ . The later property makes the interpretation of lattice simulations [10, 11] ambiguous. Note also that the calculated effect for heavy quarks is not directly related neither to topology nor to chirality. In particular, it is present also when all the three fields are electromagnetic ones, which may be of physical interest, as soon as the electromagnetic and chromodynamical fields are of the same order in heavy ions collisions.

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ORBITAL MOMENTUM EFFECTS DUE TO A LIQUID NATURE OF TRANSIENT STATE

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Abstract

It is argued that directed flow v_1 , the observable introduced for description of nucleus collisions, can be used for the detection of the nature of state of the matter in the transient state of hadron and nuclei collisions. We consider a possible origin of the directed flow in hadronic reactions as a result of rotation of the transient matter and trace analogy with nucleus collisions. Our proposal is that the presence of directed flow can serve as a signal that transient matter is in a liquid state.

Important tools in the studies of the nature of the new form of matter are the anisotropic flows which are the quantitative characteristics of the collective motion of the produced hadrons in the nuclear interactions. With their measurements one can obtain a valuable information on the early stages of reactions and observe signals of QGP formation. The experimental probes of collective dynamics in AA interactions, the momentum anisotropies v_n are defined by means of the Fourier expansion of the transverse momentum spectrum over the momentum azimuthal angle ϕ . The angle ϕ is the angle of the detected particle transverse momentum with respect to the reaction plane spanned by the collision axis z and the impact parameter vector $\mathbf{b}$ directed along the x axis. Thus, the anisotropic flows are the azimuthal correlations with the reaction plane. In particular, the directed flow is defined as

$$v_1(p_\perp) \equiv \langle \cos \phi \rangle_{p_\perp} = \langle p_x / p_\perp \rangle = \langle \hat{\mathbf{b}} \cdot \mathbf{p}_\perp / p_\perp \rangle \quad (1)$$

From Eq. (1) it is evident that this observable can be used for studies of multiparticle production dynamics in hadronic collisions provided that impact parameter $\mathbf{b}$ is fixed.

We assume that the origin of the transient state and its dynamics along with hadron structure can be related to the mechanism of spontaneous chiral symmetry breaking (χ SB) in QCD, which leads to the generation of quark masses and appearance of quark condensates. This mechanism describes transition of the current into constituent quarks. The gluon field is considered to be responsible for providing quarks with masses and its internal structure through the instanton mechanism of the spontaneous chiral symmetry breaking. Massive constituent quarks appear as quasiparticles, i.e. current quarks and the surrounding clouds of quark–antiquark pairs which consist of a mixture of quarks of the different flavors. Quark radii are determined by the radii of the surrounding clouds. Quantum numbers of the constituent quarks are the same as the quantum numbers of current quarks due to conservation of the corresponding currents in QCD.

Collective excitations of the condensate are the Goldstone bosons and the constituent quarks interact via exchange of the Goldstone bosons; this interaction is mainly due to

pion field. Pions themselves are the bound states of massive quarks. The interaction responsible for quark-pion interaction can be written in the form [1]:

$$\mathcal{L}_I = \bar{Q}[i\not{\partial} - M \exp(i\gamma_5 \pi^A \lambda^A / F_\pi)]Q, \quad \pi^A = \pi, K, \eta. \quad (2)$$

The interaction is strong, the corresponding coupling constant is about 4. The general form of the total effective Lagrangian ($\mathcal{L}_{QCD} \rightarrow \mathcal{L}_{eff}$) relevant for description of the non-perturbative phase of QCD includes the three terms [2]

$$\mathcal{L}_{eff} = \mathcal{L}_\chi + \mathcal{L}_I + \mathcal{L}_C.$$

Here $\mathcal{L}_\chi$ is responsible for the spontaneous chiral symmetry breaking and turns on first. The picture of a hadron consisting of constituent quarks embedded into quark condensate implies that overlapping and interaction of peripheral clouds occur at the first stage of hadron interaction. The interaction of the condensate clouds assumed to be of the shock-wave type, this condensate clouds interaction generates the quark-pion transient state. This mechanism is inspired by the shock-wave production process proposed by Heisenberg long time ago. At this stage, part of the effective lagrangian $\mathcal{L}_C$ is turned off (it is turned on again in the final stage of the reaction). Nonlinear field couplings transform then the kinetic energy to internal energy. As a result the massive virtual quarks appear in the overlapping region and transient state of matter is generated. This state consists of $\bar{Q}Q$ pairs and pions strongly interacting with quarks. This picture of quark-pion interaction can be considered as an origin for percolation mechanism of deconfinement resulting in the liquid nature of transient matter [3].

Part of hadron energy carried by the outer condensate clouds being released in the overlap region goes to generation of massive quarks interacting by pion exchange and their number was estimated as follows:

$$\tilde{N}(s, b) \propto \frac{(1 - \langle k_Q \rangle) \sqrt{s}}{m_Q} D_c^{h_1} \otimes D_c^{h_2} \equiv N_0(s) D_C(b), \quad (3)$$

where m_Q – constituent quark mass, $\langle k_Q \rangle$ – average fraction of hadron energy carried by the constituent valence quarks. Function D_c^h describes condensate distribution inside the hadron h and b is an impact parameter of the colliding hadrons. Thus, $\tilde{N}(s, b)$ quarks appear in addition to $N = n_{h_1} + n_{h_2}$ valence quarks.

The generation time of the transient state Δt_{tsq} in this picture obeys to the inequality

$$\Delta t_{tsq} \ll \Delta t_{int},$$

where Δt_{int} is the total interaction time. The newly generated massive virtual quarks play a role of scatterers for the valence quarks in elastic scattering; those quarks are transient ones in this process: they are transformed back into the condensates of the final hadrons.

Under construction of the model for elastic scattering it was assumed that the valence quarks located in the central part of a hadron are scattered in a quasi-independent way off the transient state with interaction radius of valence quark determined by its inverse mass:

$$R_Q = \kappa / m_Q. \quad (4)$$

The elastic scattering S -matrix in the impact parameter representation is written in the model in the form of linear fractional transform:

$$S(s, b) = \frac{1 + iU(s, b)}{1 - iU(s, b)}, \quad (5)$$

where $U(s, b)$ is the generalized reaction matrix, which is considered to be an input dynamical quantity similar to an input Born amplitude and related to the elastic scattering amplitude through an algebraic equation which enables one to restore unitarity. The function $U(s, b)$ is chosen in the model as a product of the averaged quark amplitudes

$$U(s, b) = \prod_{Q=1}^N \langle f_Q(s, b) \rangle \quad (6)$$

in accordance with assumed quasi-independent nature of the valence quark scattering. The essential point here is the rise with energy of the number of the scatterers like $\sqrt{s}$. The b -dependence of the function $\langle f_Q \rangle$ has a simple form $\langle f_Q(b) \rangle \propto \exp(-m_Q b/\xi)$.

These notions can be extended to particle production with account of the geometry of the overlap region and properties of the liquid transient state. Valence constituent quarks would excite a part of the cloud of the virtual massive quarks and those quark droplets will subsequently hadronize and form the multiparticle final state. This mechanism can be relevant for the region of moderate transverse momenta while the region of high transverse momenta should be described by the excitation of the constituent quarks themselves and application of the perturbative QCD to the parton structure of the constituent quark. The model allow to describe elastic scattering and the main features of multiparticle production. In particular, it leads to asymptotical dependencies

$$\sigma_{tot,el} \sim \ln^2 s, \quad \sigma_{inel} \sim \ln s, \quad \bar{n} \sim s^\delta. \quad (7)$$

The geometrical picture of hadron collision at non-zero impact parameters described above implies that the generated massive virtual quarks in overlap region will obtain large initial orbital angular momentum at high energies. The total orbital angular momentum can be estimated as follows

$$L(s, b) \simeq \alpha b \frac{\sqrt{s}}{2} D_C(b). \quad (8)$$

The parameter α is related to the fraction of the initial energy carried by the condensate clouds which goes to rotation of the quark system and the overlap region, which is described by the function $D_C(b)$, has an ellipsoidal form. It should be noted that $L \rightarrow 0$ at $b \rightarrow \infty$ and $L = 0$ at $b = 0$. At this point we would like to stress again on the liquid nature of transient state. Namely due to strong interaction between quarks in the transient state, it can be described as a quark-pion liquid. Therefore, the orbital angular momentum L should be realized as a coherent rotation of the quark-pion liquid as a whole in the xz -plane (due to mentioned strong correlations between particles presented in the liquid). It should be noted that for the given value of the orbital angular momentum L kinetic energy has a minimal value if all parts of liquid rotates with the same angular velocity. We assume therefore that the different parts of the quark-pion liquid in the overlap region indeed have the same angular velocity ω . In this model spin of the polarized hadrons has its origin in the rotation of matter hadrons consist of. In contrast, we assume rotation of

the matter during intermediate, transient state of hadronic interaction. Collective rotation of the strongly interacting system of the massive constituent quarks and pions is the main point of the proposed mechanism of the directed flow generation in hadronic and nuclei collisions. We concentrate on the effects of this rotation and consider directed flow for the constituent quarks supposing that directed flow for hadrons is close to the directed flow for the constituent quarks at least qualitatively. The assumed particle production mechanism at moderate transverse momenta is an excitation of a part of the rotating transient state of massive constituent quarks (interacting by pion exchanges) by the one of the valence constituent quarks with subsequent hadronization of the quark-pion liquid droplets. Due to the fact that the transient matter is strongly interacting, the excited parts should be located closely to the periphery of the rotating transient state otherwise absorption would not allow to quarks and pions to leave the region (quenching). The mechanism is sensitive to the particular rotation direction and the directed flow should have opposite signs for the particles in the fragmentation regions of the projectile and target respectively. It is evident that the effect of rotation (shift in p_x value) is most significant in the peripheral part of the rotating quark-pion liquid and is to be weaker in the less peripheral regions (rotation with the same angular velocity ω), i.e. the directed flow v_1 (averaged over all transverse momenta) should be proportional to the inverse depth Δl where the excitation of the rotating quark-pion liquid takes place. The geometrical picture of hadron collision has an apparent analogy with collisions of nuclei and it should be noted that the appearance of large orbital angular momentum should be expected in the overlap region in the non-central nuclei collisions. And then due to strongly interacting nature of the transient matter we assume that this orbital angular momentum realized as a coherent rotation of liquid. Thus, it seems that underlying dynamics could be similar to the dynamics of the directed flow in hadron collisions.

We can go further and extend the production mechanism from hadron to nucleus case also. This extension cannot be straightforward. First, there will be no unitarity corrections for the anisotropic flows and instead of valence constituent quarks, as a projectile we should consider nucleons, which would excite rotating quark liquid. Of course, those differences will result in significantly higher values of directed flow. But, the general trends in its dependence on the collision energy, rapidity of the detected particle and transverse momentum, should be the same. In particular, the directed flow in nuclei collisions as well as in hadron reactions will depend on the rapidity difference $y - y_{beam}$ and not on the incident energy. The mechanism therefore can provide a qualitative explanation of the incident-energy scaling of v_1 observed at RHIC [4].

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IDENTIFICATION OF EXTRA NEUTRAL GAUGE BOSONS AT THE ILC WITH POLARIZED BEAMS

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Abstract

The potential of the e^+e^- International Linear Collider to search for and distinguish between new neutral gauge bosons predicted within various classes of models with extended gauge sector in fermion pair production processes was examined. The presented analysis is based on the polarized differential distribution of fermions providing high sensitivity to Z' bosons of the studied processes. The discovery and identification reaches of the new neutral gauge bosons within the classes of E_6 and LR models and also ALR and SSM models were determined. It was shown that an important ingredient in this analysis is the electron (and to a lesser extent positron) beam polarization. The measurement of b - and c -quark pair production give important complementary information to those obtainable from pure leptonic processes.

Heavy resonances with mass around 1 TeV or higher are predicted by numerous New Physics (NP) scenarios, candidate solutions of conceptual problems of the standard model (SM). In particular, this is the case of models of gravity with extra spatial dimensions, grand-unified theories, electroweak models with extended spontaneously broken gauge symmetry, and supersymmetric (SUSY) theories with R -parity breaking ($\mathcal{R}_p$). These new heavy objects, or “resonances”, with mass $M \gg M_{W,Z}$, may be either produced or exchanged in reactions among SM particles at high energy colliders such as the LHC and ILC. A particularly interesting process to be studied in this regard at the LHC is the Drell-Yan (DY) dilepton production ($l = e, \mu$)

$$p + p \rightarrow l^+l^- + X, \quad (1)$$

where exchanges of the new particles can occur and manifest themselves as peaks in the (l^+l^-) invariant mass M . Once the heavy resonance is discovered at some $M = M_R$, further analysis is needed to identify the theoretical framework for NP to which it belongs. Correspondingly, for any NP model, one defines as *identification* reach the upper limit for the resonance mass range where it can be identified as the source of the resonance, against the other, potentially competitor scenarios, that can give a peak with same mass and same number of events under the peak.

On the basis of the lepton differential polar angle distribution, tests of the spin-2 of the Randall-Sundrum graviton excitation exchange in the process (1) at LHC, against the spin-1 hypothesis, have been performed [1] by using as basic observable an angular-integrated center-edge asymmetry, A_{CE} . The potential advantages of the asymmetry A_{CE}

to discriminate the spin-2 graviton resonance against the spin-1 hypothesis was discussed in Refs. [1, 2].

It would be interesting to compare the LHC results for the Z' identification reaches [1] with the foreseeable identification potential of the e^+e^- International Linear Collider (ILC). In this case, with the mass expected to lie well above the available c.m. energy, the Z' manifestations can be represented only by deviations of observables from the SM predictions, to be searched for in precision measurements of cross sections. For this discussion, we must utilize, as basic observables, the differential cross sections for the processes

$$e^+ + e^- \rightarrow f + \bar{f}, \quad f = e, \mu, \tau, c, b. \quad (2)$$

In a wide variety of electroweak theories, in particular those based on extended, spontaneously broken, gauge symmetries, the existence of one (or more) new neutral gauge bosons Z' is envisaged [3]. The three possible $U(1)$ Z' scenarios originating from the exceptional group E_6 spontaneous breaking are Z'_χ , Z'_ψ and Z'_η . Also, we consider the left-right model with Z'_{LR} originating from the breaking down of an $SO(10)$ grand-unification symmetry and the Z'_{ALR} predicted by the so-called “alternative” left-right scenario. The so-called sequential Z'_{SSM} , where the couplings to fermions are the same as those of the SM Z . Current Z' mass limits, from the Fermilab Tevatron collider, are in the range 500 – 900 GeV, depending on the model.

It is widely believed that new heavy vector bosons are “light” enough to be directly produced at the LHC by means of the DY mechanism and discovered through e.g. the leptonic decay mode. However, since the current limits on $M_{Z'}$ for the Z' models under study are well above the planned ILC energy of 0.5 TeV, it is expected to observe the virtual effects of new gauge bosons with the ILC at least at this energy option. A scenario we adopt in this section and the question we will address is that if the mass of a potential Z' is known from the LHC, whether the Z' model can be resolved at the ILC.

In the following, we will try to quantitatively discuss the above issues in the framework of the processes (2). In particular, our aim will be to assess the potential of electron and positron longitudinal polarization at the ILC, in enhancing the discovery reaches on Z' masses and distinction of Z' models. We will take as basic observables the polarized angular differential cross sections of the processes (2).

Expression of the polarized differential cross section for the process $e^+e^- \rightarrow f\bar{f}$ with $f \neq t$ can be found in [4]. We divide the angular range into bins. For Bhabha scattering, the angular range $|\cos\theta| < 0.90$ is divided into ten equal-size bins. Similarly, for annihilation into muon, tau and quark pairs we consider the analogous binning of the angular range $|\cos\theta| < 0.98$. For the Bhabha process, we combine the cross sections with the following initial electron and positron longitudinal polarizations: $(P^-, P^+) = (|P^-|, -|P^+|)$; $(-|P^-|, |P^+|)$; $(|P^-|, |P^+|)$; $(-|P^-|, -|P^+|)$. For the “annihilation” processes in Eq. (2), with $f \neq e, t$, we restrict ourselves to combining the $(P^-, P^+) = (|P^-|, -|P^+|)$ and $(-|P^-|, |P^+|)$ polarization configurations. Numerically, we take the “standard” envisaged values $|P^-| = 0.8$ and $|P^+| = 0.5$.

Regarding the ILC energy and time-integrated luminosity, for simplicity we assume the latter to be equally distributed among the different polarization configurations defined above. The explicit numerical results will refer to C.M. energy $\sqrt{s} = 0.5$ TeV with time-integrated luminosity $\mathcal{L}_{\text{int}} = 500 \text{ fb}^{-1}$. The assumed reconstruction efficiencies, that determine the expected statistical uncertainties, are 100% for e^+e^- final pairs; 95%

for final l^+l^- events ($l = \mu, \tau$); 35% and 60% for $c\bar{c}$ and $b\bar{b}$, respectively. The major systematic uncertainties are found to originate from uncertainties on beam polarizations and on the time-integrated luminosity: we assume $\delta P^-/P^- = \delta P^+/P^+ = 0.2\%$ and $\delta \mathcal{L}_{\text{int}}/\mathcal{L}_{\text{int}} = 0.5\%$, respectively.

As theoretical inputs, for the SM amplitudes we use the effective Born approximation [5]. Concerning the $\mathcal{O}(\alpha)$ QED corrections, the (numerically dominant) effects from initial-state radiation for Bhabha scattering and the annihilation processes in (2) are accounted for by a structure function approach including both hard and soft photon emission [6]. Effects of radiative flux return to the s -channel Z exchange are minimized by the cut $\Delta \equiv E_\gamma/E_{\text{beam}} < 1 - M_Z^2/s$ on the radiated photon energy, with $\Delta = 0.9$.

The expected sensitivity bounds on $M_{Z'}$ are assessed by means of “conventional” χ^2 analysis and by assuming a situation where no deviation from the SM predictions is observed within the experimental uncertainty. Accordingly, the corresponding limits on $M_{Z'}$ are determined by the condition $\chi^2(\mathcal{O}) \leq \chi^2_{\text{CL}}$, and we take $\chi^2_{\text{CL}} = 3.84$ for a 95% C.L. Also, in deriving limits on $M_{Z'}$ we combine all the final states of processes (2). In Table 1 we present the numerical results for the Z' sensitivity bounds obtained from the processes (2) at the ILC with $\sqrt{s} = 0.5$ TeV and $\mathcal{L}_{\text{int}} = 500 \text{ fb}^{-1}$.

Table 1: Discovery reaches (in TeV) on Z' bosons (at 95% C.L.) at the ILC with polarized (pol) and unpolarized (unp) beams at $\sqrt{s} = 0.5$ TeV and $\mathcal{L}_{\text{int}} = 500 \text{ fb}^{-1}$.

Model	Z'_ψ	Z'_η	Z'_χ	Z'_{LRS}	Z'_{SSM}	Z'_{ALR}
ILC unp	3.7	3.6	6.2	5.4	8.0	8.4
ILC pol	4.5	4.8	7.7	7.5	9.4	10.1

In distinction from the consideration above, where the “discovery reach” on $M_{Z'}$ was based on the assumption that no corrections are observed and the data are consistent with the SM predictions, we here make the hypothesis that a Z' signal is effectively observed (so that the SM is excluded at a certain C.L.) and the data is consistent with one of the Z' models. We want to assess the level at which this Z' model, that we call “true” model, can be expected to be distinguishable from the others, that may compete with it as sources of the deviations from the SM and that we call “tested” models, for any values of their mass $M_{Z'}$. Quantitatively, this amounts to determining the foreseeable “identification reach” on the “true” model.

In conclusion, we have explored in some detail how the Z' discovery reach at the ILC depends on the c.m. energy, on the available polarization, as well as on the model actually realized in Nature. The lower part of this range, up to $M_{Z'} \simeq 5$ TeV, will also be covered by the LHC, but the identification reach at the LHC is only up to $M_{Z'} < 2.2$ TeV.

In this LHC discovery range, the cleaner ILC environment, together with the availability of beam polarization, allow for an identification of the particular Z' version realized. Actually, this ILC identification range extends considerably beyond the LHC identification range. Specifically, the ILC with polarized beams at $\sqrt{s} = 0.5$ TeV allows to identify all considered Z' bosons if $M_{Z'} \lesssim 4.5$ TeV, substantially improving the LHC reach.

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QUARK INTRINSIC MOTION AND THE LINK BETWEEN TMDs AND PDFs IN COVARIANT APPROACH

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Abstract

The relations between TMDs and PDFs are obtained from the symmetry requirement: relativistic covariance combined with rotationally symmetric parton motion in the nucleon rest frame. This requirement is applied in the covariant parton model. Using the usual PDFs as an input, we are obtaining predictions for some polarized and unpolarized TMDs.

The transverse momentum dependent parton distribution functions (TMDs) [1,2] open the new way to more complete understanding of the quark-gluon structure of the nucleon. We studied this topic in our recent papers [3–5]. We have shown, that requirements of symmetry (relativistic covariance combined with rotationally symmetric parton motion in the nucleon rest frame) applied in the covariant parton model imply the relation between integrated unpolarized distribution function and its unintegrated counterpart. Obtained results are shortly discussed in the first part. Second part is devoted to the discussion of analogous relation valid for polarized distribution functions.

Unpolarized distribution function. In the covariant parton model we showed [6], that the parton distribution function $f_1^q(x)$ generated by the 3D distribution G_q of quarks reads:

$$f_1^q(x) = Mx \int G_q(p_0) \delta\left(\frac{p_0 + p_1}{M} - x\right) \frac{dp_1 d^2\mathbf{p}_T}{p_0} \quad (1)$$

and that this integral can be inverted

$$G_q\left(\frac{M}{2}x\right) = -\frac{1}{\pi M^3} \left(\frac{f_1^q(x)}{x}\right)'. \quad (2)$$

Further, due to rotational symmetry of the distribution G_q in the nucleon rest frame, the following relations for unintegrated distribution were obtained [5]:

$$f_1^q(x, \mathbf{p}_T) = MG_q\left(\frac{M}{2}\xi\right). \quad (3)$$

After inserting from Eq. (2) we get relation between unintegrated distribution and its integrated counterpart:

$$f_1^q(x, \mathbf{p}_T) = -\frac{1}{\pi M^2} \left(\frac{f_1^q(\xi)}{\xi}\right)'; \quad \xi = x \left(1 + \left(\frac{p_T}{Mx}\right)^2\right). \quad (4)$$

Now, using some input distributions $f_1^q(x)$ one can calculate transverse momentum distribution functions $f_1^q(x, \mathbf{p}_T)$. As the input we used the standard PDF parameterization [8] (LO at the scale 4GeV^2). The pictures of this distributions for u and d -quarks can be found in paper [5]. A part of this figures, but in different scale, is shown again in Fig 1. One can observe the following:

i) For fixed x the corresponding p_T -distributions are very close to the Gaussian distributions

$$f_1^q(x, p_T) \propto \exp\left(-\frac{p_T^2}{\langle p_T^2 \rangle}\right). \quad (5)$$

ii) The width $\langle p_T^2 \rangle = \langle p_T^2(x) \rangle$ depends on x . This result corresponds to the fact, that in our approach, due to rotational symmetry, the parameters x and p_T are not independent.

iii) Figures suggest the typical values of transversal momenta, $\langle p_T^2 \rangle \approx 0.01\text{GeV}^2$ or $\langle p_T \rangle \approx 0.1\text{GeV}$. These values correspond to the estimates based the analysis of the experimental data on structure function $F_2(x, Q^2)$ [5]. They are substantially lower, than the values $\langle p_T^2 \rangle \approx 0.25\text{GeV}^2$ or $\langle p_T \rangle \approx 0.44\text{GeV}$ following e.g. from the analysis of data on the Cahn effect [9] or HERMES data [10]. At the same time the fact, that the shape of obtained p_T -distributions (for fixed x) is close to the Gaussian, is remarkable. In fact, the Gaussian shape is supported by phenomenology.

Polarized distribution functions.

Relation between the distribution $g_1^q(x)$ and its unintegrated counterpart can be obtained in a similar way, however in general the calculation with polarized structure functions is slightly more complicated. First let us remind procedure for obtaining structure functions g_1, g_2 from starting distribution functions $G^\pm$ defined in [6], Sec. 2, see also the footnote there. In fact the auxiliary functions G_P, G_S are obtained in appendix of the paper [7]. If we assume that $Q^2 \gg 4M^2x^2$, then the approximations $|\mathbf{q}| \approx \nu$, $\frac{pq}{Pq} \approx \frac{p_0+p_1}{M}$ are valid and the equations (A1),(A2) can be with the use of (A3),(A4) rewritten as

$$G_X = \int \Delta G(p_0) w_X \delta\left(\frac{p_0+p_1}{M} - x\right) \frac{dp_1 d^2\mathbf{p}_T}{p_0}, \quad X = P, S \quad (6)$$

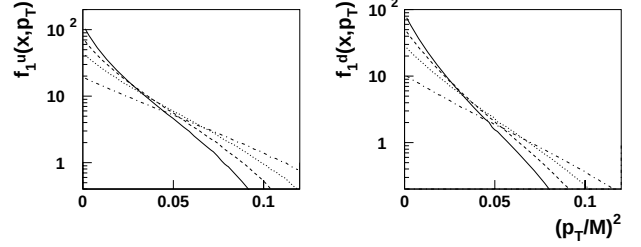


Figure 1: Transverse momentum dependent unpolarized distribution functions for u and d -quarks. Dependence on $(p_T/M)^2$ for $x = 0.15, 0.18, 0.22, 0.30$ is indicated by solid, dash, dotted and dash-dot curves.

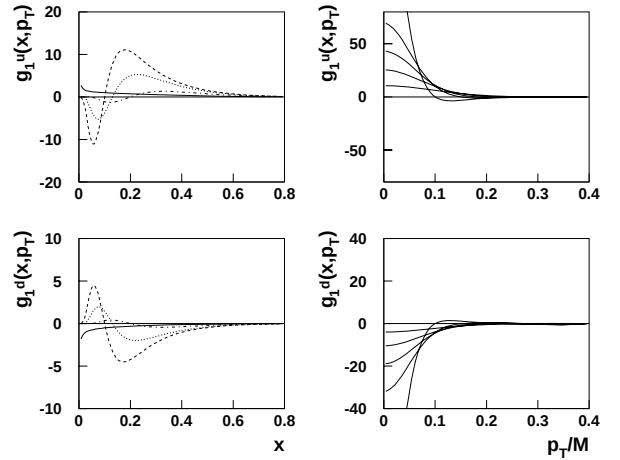


Figure 2: Transverse momentum dependent polarized distribution functions for u (upper figures) and d -quarks (lower figures). **Left part:** dependence on x for $p_T/M = 0.10, 0.13, 0.20$ is indicated by dash, dotted and dash-dot curves; solid curve corresponds to the integrated distribution $g_1^q(x)$. **Right part:** dependence on p_T/M for $x = 0.10, 0.15, 0.18, 0.22, 0.30$ from top to bottom for u -quarks, and symmetrically for d -quarks.

where

$$w_P = -\frac{m}{2M^2\nu} \frac{-p_1 \cos \omega + p_T \cos \varphi \sin \omega}{p_0 + m} \times \left(1 + \frac{1}{m} \left(p_0 - \frac{-p_1 - (-p_1 \cos \omega + p_T \cos \varphi \sin \omega) \cos \omega}{\sin^2 \omega} \right) \right), \quad (7)$$

$$w_S = \frac{m}{2M\nu} \left(1 + \frac{-p_1 \cos \omega + p_T \cos \varphi \sin \omega}{p_0 + m} \frac{1}{m} \times \left(-p_1 \cos \omega + p_T \cos \varphi \sin \omega - \frac{-p_1 - (-p_1 \cos \omega + p_T \cos \varphi \sin \omega) \cos \omega}{\sin^2 \omega} \cos \omega \right) \right). \quad (8)$$

Let us remark, that using the notation defined in [3] we can identify

$$-\cos \omega = S_L, \quad \sin \omega = S_T, \quad p_T \sin \omega \cos \varphi = \mathbf{p}_T \mathbf{S}_T, \quad (9)$$

which appear in definition of the TMDs [2]:

$$\frac{1}{2} \text{tr} [\gamma^+ \gamma_5 \phi^q(x, \mathbf{p}_T)] = S_L g_1^q(x, p_T) + \frac{\mathbf{p}_T \mathbf{S}_T}{M} g_{1T}^{\perp q}(x, p_T). \quad (10)$$

The expressions (7),(8) can be reordered in terms of powers of $\cos \varphi$:

$$w_P = -\frac{\cos \omega}{2M^2\nu} \left(\frac{p_T^2}{m + p_0} \cos^2 \varphi + (p_T \tan \omega + \frac{p_T}{m + p_0} p_1 (\tan \omega - \cot \omega)) \cos \varphi - p_1 \left(\frac{p_1}{m + p_0} + 1 \right) \right), \quad (11)$$

$$w_S = \frac{1}{2M\nu} \left(\frac{p_T^2}{m + p_0} \cos^2 \varphi - \frac{p_1 p_T \cot \omega}{m + p_0} \cos \varphi + m \right). \quad (12)$$

Now, in analogy with Eq.(46) in [7] we define (note that $Pq/qS = -M/\cos \omega$):

$$w_1 = M\nu \cdot w_S + \frac{M^2\nu}{\cos \omega} \cdot w_P, \quad w_2 = -\frac{M^2\nu}{\cos \omega} \cdot w_P, \quad (13)$$

which implies

$$g_k^q = \int \Delta G(p_0) w_k \delta \left(\frac{p_0 + p_1}{M} - x \right) \frac{dp_1 d^2 \mathbf{p}_T}{p_0}, \quad k = 1, 2. \quad (14)$$

After inserting from Eqs. (11),(12) definition (13) implies

$$w_1 = \frac{1}{2} \left(m + p_1 \left(1 + \frac{p_1}{m + p_0} \right) - p_T \tan \omega \left(1 + \frac{p_1}{m + p_0} \right) \cos \varphi \right), \quad (15)$$

$$w_2 = \frac{1}{2} \left(\frac{p_T^2}{m + p_0} \cos^2 \varphi + (p_T \tan \omega + \frac{p_T}{m + p_0} p_1 (\tan \omega - \cot \omega)) \cos \varphi - p_1 \left(\frac{p_1}{m + p_0} + 1 \right) \right), \quad (16)$$

$$w_1 + w_2 = \frac{1}{2} \left(\frac{p_T^2}{m + p_0} \cos^2 \varphi - \frac{p_1 p_T \cot \omega}{m + p_0} \cos \varphi + m \right). \quad (17)$$

Apparently, the terms proportional to $\cos \varphi$ disappear in the integrals (14) and the remaining terms give structure functions g_1, g_2 defined by Eqs. (15),(16) in [6].

Now, the integration over p_1 and further procedure can be done in a similar way as for unpolarized distribution. First, to simplify calculation, we assume $m \rightarrow 0$. For w_1 we get

$$g_1^q(x) = \frac{1}{2} \int \Delta G_q(p_0) \left(1 + \frac{p_1}{p_0} \right) (p_1 - p_T \tan \omega \cos \varphi) \delta \left(\frac{p_0 + p_1}{M} - x \right) \frac{dp_1 d^2 \mathbf{p}_T}{p_0}. \quad (18)$$

The δ -function is modified as

$$\delta \left(\frac{p_0 + p_1}{M} - x \right) dp_1 = \frac{\delta(p_1 - \tilde{p}_1) dp_1}{x/\tilde{p}_0}, \quad (19)$$

where

$$\tilde{p}_1 = \frac{Mx}{2} \left(1 - \left(\frac{p_T}{Mx} \right)^2 \right), \quad \tilde{p}_0 = \frac{Mx}{2} \left(1 + \left(\frac{p_T}{Mx} \right)^2 \right). \quad (20)$$

Modified δ -function allows to simplify the integral

$$g_1^q(x) = \frac{1}{2} \int \Delta G_q(\tilde{p}_0) (M(2x - \xi) - 2p_T \tan \omega \cos \varphi) \frac{d^2 \mathbf{p}_T}{\xi}, \quad (21)$$

where

$$\xi = x \left(1 + \left(\frac{p_T}{Mx} \right)^2 \right). \quad (22)$$

Now we define

$$\Delta q(x, \mathbf{p}_T) = \frac{1}{2} \Delta G_q \left(\frac{M\xi}{2} \right) (M(2x - \xi) - 2p_T \tan \omega \cos \varphi) \frac{1}{\xi}. \quad (23)$$

According to Eq. (40) in [6] we have

$$\Delta G_q \left(\frac{M\xi}{2} \right) = \frac{2}{\pi M^3 \xi^2} \left(3g_1^q(\xi) + 2 \int_{\xi}^1 \frac{g_1^q(y)}{y} dy - \xi \frac{d}{d\xi} g_1^q(\xi) \right). \quad (24)$$

After inserting to Eq. (23) one gets:

$$\begin{aligned} \Delta q(x, \mathbf{p}_T) &= \frac{1}{\pi M^2 \xi^3} \left(3g_1^q(\xi) + 2 \int_{\xi}^1 \frac{g_1^q(y)}{y} dy - \xi \frac{d}{d\xi} g_1^q(\xi) \right) \\ &\quad \times \left(2x - \xi - 2 \frac{p_T}{M} \tan \omega \cos \varphi \right). \end{aligned} \quad (25)$$

This relation allows us to calculate the distribution $\Delta q(x, \mathbf{p}_T)$ from a known input on $g_1^q(x)$. Further, it can be shown, that using the notation defined in Eqs. (9),(10), our result reads

$$-\cos \omega \cdot \Delta q(x, \mathbf{p}_T) = S_L g_1^q(x, p_T) + \frac{\mathbf{p}_T \mathbf{S}_T}{M} g_{1T}^{\perp q}(x, p_T), \quad (26)$$

where

$$g_1^q(x, p_T) = \frac{2x - \xi}{\pi M^2 \xi^3} \left(3g_1^q(\xi) + 2 \int_{\xi}^1 \frac{g_1^q(y)}{y} dy - \xi \frac{d}{d\xi} g_1^q(\xi) \right), \quad (27)$$

$$g_{1T}^{\perp q}(x, p_T) = \frac{2}{\pi M^2 \xi^3} \left(3g_1^q(\xi) + 2 \int_{\xi}^1 \frac{g_1^q(y)}{y} dy - \xi \frac{d}{d\xi} g_1^q(\xi) \right). \quad (28)$$

Apparently, both functions are related in our approach:

$$\frac{g_1^q(x, p_T)}{g_{1T}^{\perp q}(x, p_T)} = \frac{x}{2} \left(1 - \left(\frac{p_T}{Mx} \right)^2 \right) = \tilde{p}_1/M. \quad (29)$$

Finally, with the use of standard input [11] on $g_1^q(x) = \Delta q(x)/2$ we can obtain the curves $g_1^q(x, p_T)$ displayed in Fig. 2. Let us remark, that the curves change the sign at the point $p_T = Mx$. This change is due to the term

$$2x - \xi = x \left(1 - \left(\frac{p_T}{Mx} \right)^2 \right) = 2\tilde{p}_1/M \quad (30)$$

in relation (27). This term is proportional to the quark longitudinal momentum $\tilde{p}_1$ in the proton rest frame, which is defined by given x and p_T . It means, that sign of the $g_1^q(x, p_T)$ is controlled by sign of $\tilde{p}_1$. On the other hand, the function $g_{1T}^{\perp q}(x, p_T)$ does not involve term, which changes the sign. The shape of both functions should be checked by experiment.

To conclude, we presented our recent results on relations between TMDs and PDFs. The study is in progress, further results will be published later.

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EXPERIMENTAL RESULTS

SEARCH FOR NARROW PION-PROTON STATES IN S-CHANNEL AT EPECUR: EXPERIMENT STATUS

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Abstract

An experiment EPECUR, aimed at the search of the cryptoexotic non-strange member of the pentaquark antidecuplet, started its operation at a pion beam line of the ITEP 10 GeV proton synchrotron. The invariant mass range of the interest (1610-1770) MeV will be scanned for a narrow state in the pion-proton and kaon-lambda systems in the information-type experiment. The scan in the s-channel is supposed to be done by the variation of the incident π^- -momentum and its measurement with the accuracy of up to 0.1% with a set of 1 mm pitch proportional chambers located in the first focus of the beam line. The reactions under the study will be identified by a magnetless spectrometer based on wire drift chambers with a hexagonal structure. Because the background suppression in this experiment depends on the angular resolution, the amount of matter in the chambers and setup is minimized to reduce multiple scattering. The differential cross section of the elastic π^-p -scattering on a liquid hydrogen target in the region of the diffractive minimum will be measured with statistical accuracy 0.5% in 1 MeV steps in terms of the invariant mass. For $K_S^0\Lambda^0$ -production the total cross section will be measured with 1% statistical accuracy in the same steps. An important byproduct of this experiment will be a very accurate study of Λ polarization. The setup was assembled and tested in December 2008 and in April 2009 we had the very first physics run. About $0.5 \cdot 10^9$ triggers were written to disk covering pion beam momentum range 940-1135 MeV/c. The talk covers the experimental setup and the current status.

An interest to this experiment originated with the discovery in 2003 by the two experiments LEPS [1] and DIANA [2] a new baryonic state θ^+ with positive strangeness and very small width. Later appeared several strong results where the state was not seen [3] but recent results from LEPS [4] and DIANA [5] still insist on the evidence for this resonance. Quantum numbers of θ^+ are not measured but it is believed that it belongs to pentaquark antidecuplet predicted in 1997 by D. Diakonov, V. Petrov and M. Polyakov [6]. In this case there should also exist a non-strange neutral resonance P11 with mass near 1700 MeV. Certain hints in favour of its presence were found in the modified PWA of GWU group [7] at masses 1680 and 1730 MeV [8]. Recently an indication for this narrow

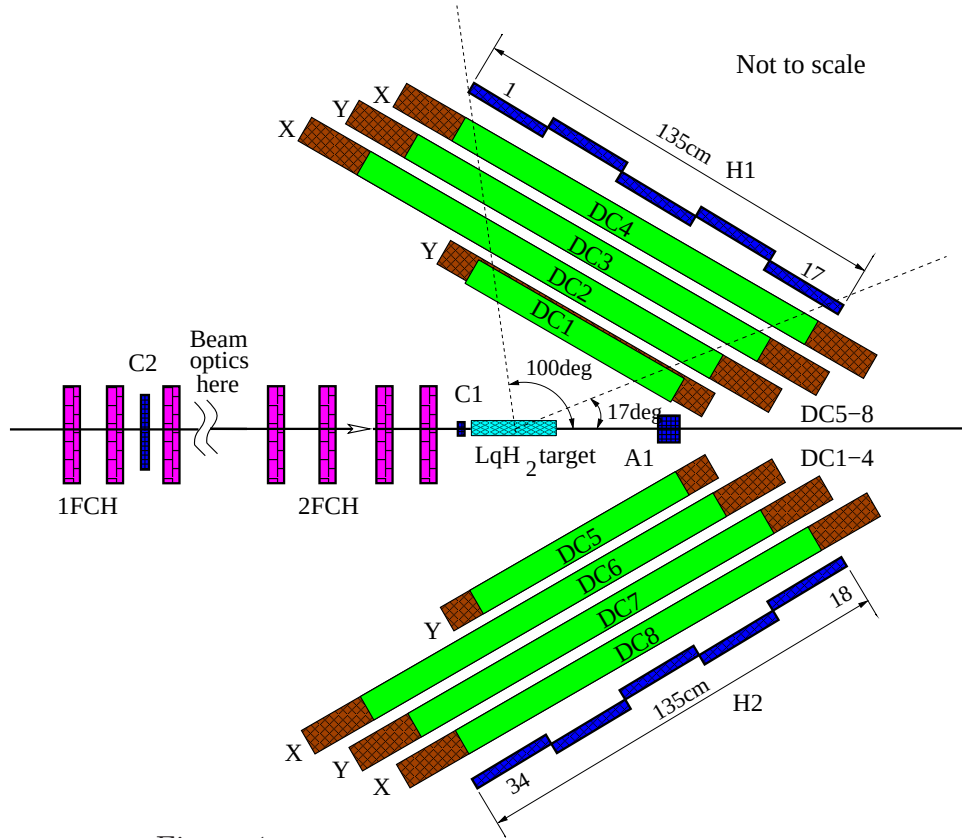


Figure 1: Experimental layout for π^-p elastic scattering.

state was found in η -photoproduction on deuteron in GRAAL [9] and some other experiments. The structure observed has mass 1685 MeV and width < 30 MeV, which was determined by the detector resolution.

Our idea is to search for P11(1700) in formation-type experiment on a pion beam [10]. Precise measurement of the beam momentum and fair statistics will allow us to do a scan with unprecedented invariant mass resolution. We plan to measure differential cross sections of the reactions $\pi^-p \rightarrow \pi^-p$ and $\pi^-p \rightarrow K_S^0 \Lambda^0$ with high statistics and better than a MeV invariant mass resolution. If the resonance does exist our experiment will provide statistically significant result and we will measure its width with the precision better than 0.7 MeV.

The layout dedicated to the elastic scattering measurement is shown in fig. 1. The main parts are: proportional chambers (**1FCH** and **2FCH**), drift chambers (**DC1-8**), liquid hydrogen target (**LqH₂**), scintillation counters (**C1**, **C2** and **A1**) and two scintillation hodoscopes (**H1** and **H2**).

Proportional chambers are placed in the 1st and 2nd focuses of the beam. The chambers are two-coordinate, have square sensitive region of 200×200 mm², 1 mm signal wires pitch, 40 μ m aluminum foil cathode electrodes and 6 mm between the foils. We use "magic" gas mixture (argon-isobutane-freon) to feed the proportional chambers. Beam tests showed efficiency better than 99%.

Main task of the chambers in the 1st focus is to measure the momentum of each pion going to the target. Strong dipole magnets between the internal target and the 1st focus provide horizontal distribution of the particles with different momentum with dispersion 57 mm/%. A distribution over horizontal coordinate in the 1st focus of the events of scattering of the internal beam protons with momentum 1.0 GeV/c over a beryllium

internal target is shown in fig. 2. The peaks observed in the picture correspond to (right to left) the elastic scattering, the first excitation of beryllium nucleus and the second and the third excitations seen as one peak.

The liquid hydrogen target has a mylar cylinder container with diameter 40 mm and the length about 250 mm placed in high vacuum inside beryllium outer shell 1 mm thick. It is connected by two pipes to the liquefier system. One is used for liquid hydrogen inflow and through the other the evaporated gas gets back to liquefier. This design provides minimum of matter for the particles. The refrigeration is provided by liquid helium, which flow is controlled by the feedback supporting constant pressure of the hydrogen in the closed volume. This pressure corresponds to proper ratio between liquid and gas fractions of the hydrogen and thus ensuring that the liquid occupies whole target working volume and that the hydrogen is not frozen. Pressures and temperatures in the target system are monitored and logged.

There are 8 one coordinate drift chambers in the elastic setup. 6 chambers have sensitive region $1200 \times 800 \text{ mm}^2$ and for 2 chambers closest to the target it is $600 \times 400 \text{ mm}^2$. The chambers have double sensitive layers hexagonal structure shown in fig. 3. Comparing to the conventional drift tubes this structure has much more complex fields, but provides significantly less amount of matter on the particle path. Potential wires form nearly regular hexagon with a side of 10 mm. Drift chambers are fed with 70% Ar and 30% CO_2 gas mixture. Beam tests showed better than 99% single layer efficiency and about 0.2 mm resolution.

A unique distributed DAQ system based on the commercial 480 Mbit/s USB 2.0 interface was designed for the experiment [11]. It consists of 100-channel boards for proportional chambers and 24-channel boards for drift chambers, placed on the chambers frames. Each board is connected by two cables (USB 2.0 and power) to the communication box, placed near the chamber. Then the data is transferred to the main DAQ computer by the standard TCP/IP connection. Trigger logic is capable of processing of several trigger conditions firing different sets of detectors.

During the engineering run last December and the first physics run in April this year the main trigger was set as:

$$T = C_1 \cdot C_2 \cdot M_{1FCH} \cdot M_{2FCH} \cdot \overline{A_1}$$

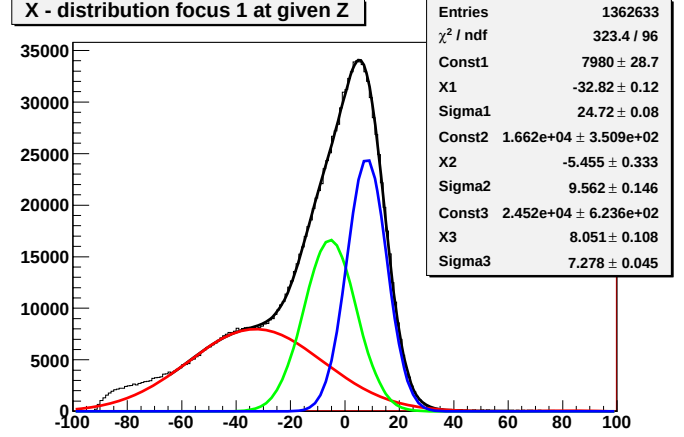


Figure 2: Horizontal distribution in the first focus of internal the accelerator beam protons scattered over internal beryllium target.

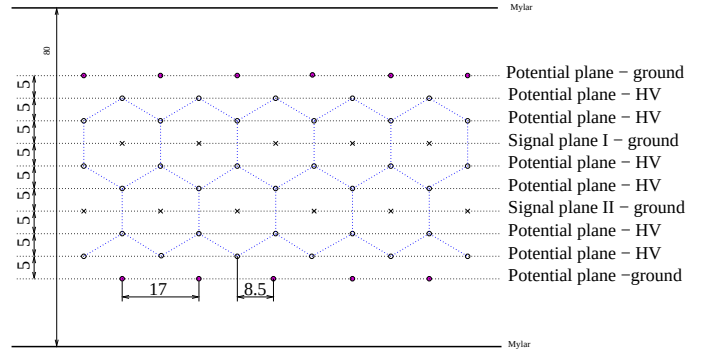


Figure 3: Drift chamber cross section. View along the wires.

where C_1 , C_2 and A_1 - signals from corresponding scintillator counters and M_{1FCH} and M_{2FCH} - majority logic of the proportional chamber planes in the 1st and the 2nd focuses. Other trigger conditions were used to provide beam position and luminosity monitoring. To ensure stable beam momentum an NMR monitoring of the magnetic field of the last dipole was used. We collected over $5 \cdot 10^8$ events in the April run, of which we expect about 5% to be elastic events, in the invariant mass range 1640–1745 MeV. Processing of the April run data had started. An example of the elastic events selection for $50^\circ < \theta_{\text{lab}}^\pi < 60^\circ$ is presented in fig. 4. A distribution of the polar angles difference between pion and proton in the center of mass system is shown. A narrow peak corresponds to the correct assumption which of the secondary particles is pion and which is proton. A wide peak correspond to the wrong assignment. This comes from the fact that in the elastic setup we can't distinguish between pion and proton and can only try some assignment and correct events for which it turned out to be wrong.

As a conclusion:

- Construction of the first stage (elastic) of EPECUR has finished and the setup was successfully commissioned.
- Data taking has started this year and the first $5 \cdot 10^8$ events were collected.
- Data processing is under way.
- We are going to take more runs.
- We are going to proceed to the construction of the second stage (K Λ -production) simultaneously with data taking.

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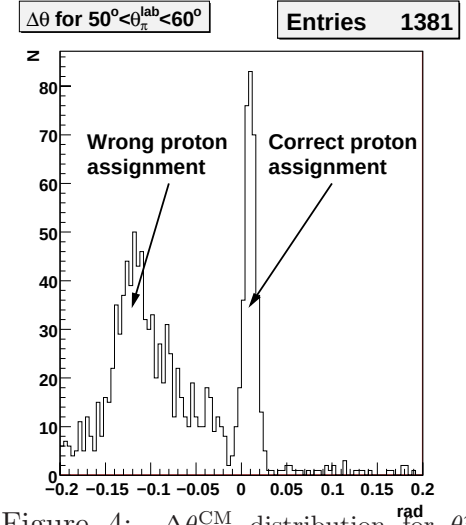


Figure 4: $\Delta\theta^{\text{CM}}$ distribution for θ_{lab}^π range $50^\circ - 60^\circ$.

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MEASUREMENT OF TENSOR POLARIZATION OF DEUTERON BEAM PASSING THROUGH MATTER

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Abstract

The results of measurements and the procedure for handling the data on the tensor polarization of the deuteron beam arising as the beam passes through matter obtained at the Nuclotron during the June 2008 run using an extracted unpolarized 5-GeV/c deuteron beam are described. The observed effect is compared with the calculations made within the framework of the Glauber multiple scattering theory.

In the last years it has become more and more obvious, that many modern problems of nuclear physics and physics of elementary particles cannot be solved without the profound research of polarization effects which are connected with the search for new methodological approaches. For example, it is offered to use the effect of spin filtering, occurring at spin transfer from the polarized electrons in an internal hydrogen gas target to the orbiting antiprotons to produce their polarization in storage rings [1]. This makes it possible to measure the transversity distribution of the valence quarks in the proton with the purpose of overcoming the so-called spin crisis.

The phenomenon of spin filtering can also take place in interactions of others hadrons, in particular, deuterons. In the case of deuterons the cause to this effect is presence of the quadrupole momentum in the deuteron.

Baryshevsky has shown [2] that passage of high-energy particles of spin ≥ 1 through matter is accompanied by new effects: spin oscillations and spin dichroism. These effects can result in polarization of the beam traveling through the target. Tensor polarization of the originally unpolarized beam after its passage through the unpolarized target was first observed for deuterons with the energy up to 20 MeV traveling through carbon foils [3].

The first measurement of the tensor polarization of a beam arising after its passing through a carbon target, was made in Dubna in the unpolarized deuteron beam with the momentum of 5 GeV/c extracted from the Nuclotron [4]. In the present report the results of the recent studies carried out at the Nuclotron with the purpose of confirming the earlier observed effect are described.

The layout of the experimental equipment is shown in Fig. 1. The deuteron beam with an intensity of $5 \times 10^8 - 3 \times 10^9$ particles/spill extracted from the Nuclotron was directed by magnetic elements to the 54, 83 and 137 g/cm²-thick carbon targets *T1* serially placed inside the concrete shield near the focus *F3* of the magnetic-optical channel. The polarization of the deuterons which passed through the targets *T1* was analyzed using the

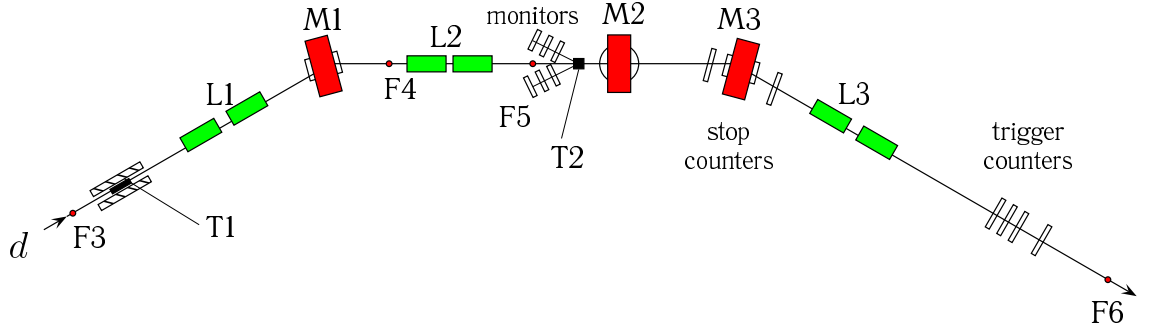


Figure 1: Layout of experiment

target-analyzer $T2$ placed near the focus $F5$. Magnetic lenses and dipoles are designated as $L1, L2, L3$ and $M1, M2, M3$. The part of the magnetic-optical channel behind the target $T1$ up to $F5$ was tuned to the momentum of ~ 5 GeV/ c , and the part behind $F5$ was tuned to 3.3 GeV/ c .

The momenta of the deuterons extracted from the accelerator were adjusted to be exactly 5.0 GeV/ c at the exit from $T1$ with allowance for ionization losses in the target $T1$. The measurements without the target $T1$ were also periodically made. The intensity of the beam was monitored by the ionization chambers placed near to the foci $F3, F4$ and $F5$. The intensity of the secondary beam between $F4$ and $F5$ was $5 \times 10^6 - 3 \times 10^7$ particles per beam spill.

Before the June 2008 run of the Nuclotron the layout of experiment was modernized as follows: 1) an the additional five-electrode ionization chamber was placed behind the target $T1$ to control the position and intensity of the beam at the entrance to the doublet of lenses $L1$; 2) the intensity of the beam incident on the target $T2$ was monitored by two telescopes of the scintillation counters located in the back hemisphere relative to the target $T2$ at the angles of 150° ($M1$) and 210° ($M2$).

The tensor polarization of the deuteron beam that passed through the target $T1$ at 0° was determined by means of the deuteron stripping reaction on the 8-cm thick beryllium target $T2$ placed near $F5$ [5]. The fact that the reaction $d + A \rightarrow p + X$ for proton emission at the zero angle with the momentum $p_p \sim \frac{2}{3}p_d$ has the known tensor analyzing power $T_{20} = -0.82 \pm 0.04$ [6] was used. If the differential cross sections of this reaction in the case of the unpolarized and polarized deuteron beams are designated as σ_0 and σ' respectively, the following relation for the tensor polarization p_{ZZ} holds [4, 7]:

$$p_{ZZ} = \frac{\sqrt{2}}{T_{20}} \left(\frac{\sigma'}{\sigma_0} - 1 \right), \quad (1)$$

where the axis of quantization coincides with the direction of protons emitted forward.

The secondary particles emitted from the target $T2$ at 0° were directed to the focus $F6$ by bending magnets and magnetic lenses. The momentum and polar angle acceptances of the setup determined by the Monte Carlo simulation were $\Delta p/p \sim \pm 2\%$ and ± 8 mrad, respectively.

Coincidences of signals from the scintillation counters located near the focus $F6$ were used as a trigger. Along with the secondary protons, the equipment detected the deuterons that experienced inelastic scattering. The detected particles were identified off line during the processing of the saved data on the basis of the information on their time of flight

over the distance of ~ 28 m between the start and three stop counters. The time-of-flight resolution ~ 0.2 ns allowed one to separate completely secondary protons and deuterons.

The numbers of protons detected at the focus $F6$ in exposures with carbon targets of different thickness and normalized to the monitor counts are shown in Fig. 2.

Here dark circles, stars and crosses refer to the 123-, 83- and 40-g/cm²-thick carbon targets, respectively, and the light circles correspond to the measurements without the target $T1$. The values of these ratios averaged for all the exposures are shown by broken lines. It is seen that the points corresponding to different target thickness are grouped in different regions of the picture. The scatter of the points exceeds statistical errors that are less than the size of the points. The main causes of this scatter are 1) non-stabilities of currents in the magnetic elements of the magnetic-optical channel; 2) nonuniform distribution of the intensity of the extracted deuteron beam within the limits of the spill.

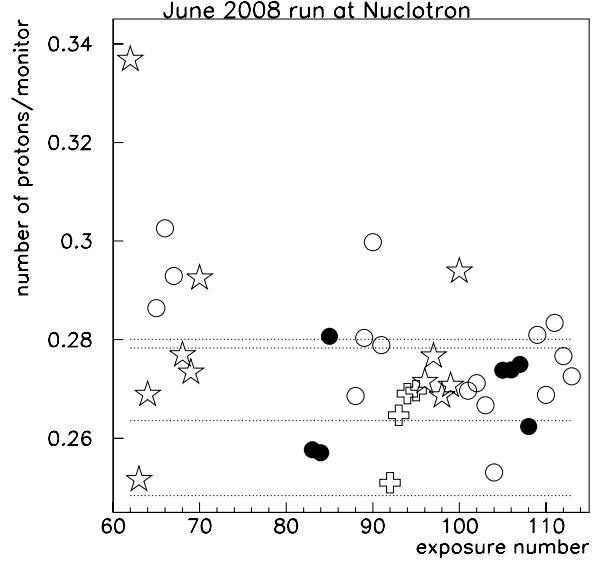


Figure 2: Ratios of proton counts to the monitor ($M1b + M2b$) in the focus $F5$ for targets $T1$ of different thickness: black circles - 137 g/cm², stars - 83 g/cm², crosses - 54 g/cm², light circles - 0 g/cm².

The values of the tensor polarization p_{ZZ} of the deuterons that passed through the target $T1$ were calculated according to expression (1) separately for each channel of registration, and then they were averaged over the channels; the counts without $T1$ were taken as σ_0 . The dependence of the values of the tensor polarization on the thickness of the target $T1$ found in the described run is shown in Fig. 3 by black circles. We note that the given values are found by averaging the results of two independent data processing procedures. Light circles show the results of the previous run [4]. The corridor of errors appropriate to both series of measurements is shown with broken lines. The calculations of the spin alignment of the deuteron beam after passage through matter within the framework of the Glauber multiple scattering model were made in the work [8]. The calculation results for the carbon target are shown in Fig. 3 with the continuous curve.

The measurements performed in June 2008 basically confirm the results of the experimental observation of the spin filtering of deuterons at passage of the beam through a layer of matter obtained in March 2007. In spite of the certain distinctions in the dependences of the tensor polarization of deuterons p_{ZZ} on thickness of the carbon filter ΔC measured in these two experiments, the fact that the deuteron spin alignment increases with increasing ΔC proves to be true. The distinctions mentioned are caused by imperfection of the monitoring system of the magnetic optics of the channel of the slow extraction of particles from the accelerator.

According to requirements of the described experiment the allowable level of deviations of the values of currents in the magnetic elements should not exceed 1 %.

The dependence of the tensor polarization of the deuteron beam on the thickness of the carbon filter obtained in two series of measurements looks like

$$p_{ZZ} = (1.226 \pm 0.209) \Delta C (kg/cm^2),$$

and within the limits of the corridor of errors it agrees with the results of calculations [8].

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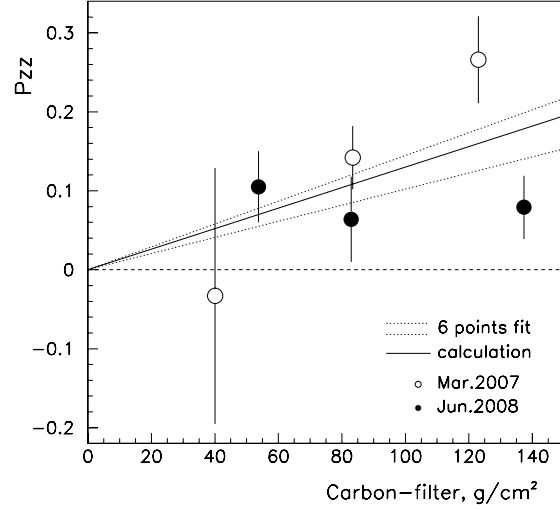


Figure 3: Tensor polarization of deuterons vs thickness of the carbon target *T1*. Light points are the results of work [4], black points are the results of the present work. The broken lines limit the corridor of errors. The continuous curve is the result of the calculation of p_{ZZ} [8] within the framework of the Glauber multiple scattering model.

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PROSPECTS OF MEASURING ZZ AND WZ POLARIZATION WITH ATLAS

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Abstract

The measurement of angular distributions of the decay leptons in di-boson final states allows the reconstruction of the spin density element ρ_{00} in processes where ZZ and WZ final states are produced via quark anti-quark annihilation in proton-proton collisions at the LHC. This note presents the expected sensitivity of such measurements with the ATLAS experiment, using electrons and muons as final state leptons, based on an integrated luminosity of 100 fb^{-1} . Besides the statistical accuracy on the fractions of longitudinally polarized Z or W bosons, various contributions to the systematic uncertainty have been studied.

The W gauge boson is mainly produced singly in $q\bar{q}$ annihilation at hadron-hadron colliders. The same is true for the Z gauge boson, which has also been produced in electron-positron annihilations at LEP. In these processes, the gauge boson has always transverse polarization due to the helicity conservation in the annihilation process. However, some fraction of longitudinal polarization is allowed and expected in the Standard Model (SM) when the annihilation process yields two gauge bosons. Longitudinally produced W bosons have been observed for the first time at LEP [1], in W^+W^- production, and later also at the Tevatron [2] in top-quark decays. The longitudinal polarization degree of freedom for the W and the Z bosons is a direct result of the Higgs mechanism, and it would be a valuable test of the SM to measure the fraction of longitudinally produced bosons in di-boson events at hadron colliders and compare with theory. Any deviation from the SM predictions would be evidence for new physics. Furthermore, di-boson events can be produced as decay products of the Higgs boson, and their polarization state might be of help in the determination of the Higgs spin [3] and its separation from the di-boson continuum.

This note describes feasibility studies using Monte Carlo (MC) events of the ZZ [4] and WZ [5] production channels at LHC, to be measured by the ATLAS detector [6]. LHC is a proton-proton collider at CERN which is expected to start operation at the end of 2009, and eventually will reach a center-of-mass (c.m.s) energy of 14 TeV. This energy value was assumed in our study with an integrated luminosity of 100 fb^{-1} , which is the expected luminosity value, before a possible upgrade of the machine into Super LHC (SLHC), expected to achieve a luminosity increase by an order of magnitude. Background considerations lead us to use only the leptonic decays, $W \rightarrow \ell\nu$ and $Z \rightarrow \ell^+\ell^-$, where $\ell = e, \mu$.

Electron energy is measured at ATLAS by the liquid argon electromagnetic calorimeter with a resolution of 1.2% for electron energy around 100 GeV [7]. The electron direction is obtained from the inner detector consisting of pixel detectors, silicon strip layers and a straw-tube transition radiation detector which also assists electron identification. The electromagnetic calorimeter is inside the hadron calorimeter, which is surrounded by the muon spectrometer measuring the muon momenta with a resolution of (4 - 5)% for muons with transverse momentum, p_T , below 150 GeV [7]. The pseudo-rapidity range for triggering and measuring electrons and muons is $|\eta| < 2.5$. The calorimeters cover a larger range, up to $|\eta| = 5$. This hermetic coverage allows a precise measurement of the missing transverse energy, E_T , which is used to identify the neutrino from the W decay and obtain its transverse momentum.

Monte Carlo event samples have been generated using Pythia [9] for the ZZ and MC@NLO [9] for the WZ channel. These events have been passed through the full ATLAS detector simulation and then through the ATLAS reconstruction program to be used also for real data events. Similar event samples have been used also for the expected background processes (see below).

ZZ candidates are required to have two lepton anti-lepton pairs, each pair with invariant mass within 12 GeV of the Z mass. All leptons are required to have $|\eta| < 2.5$. Two of the leptons must satisfy $p_T > 20$ GeV and the other two must have $p_T > 7$ GeV. Following these cuts, 2194 events are left in our 100 fb⁻¹ pseudo-data sample. Background sources considered were $t\bar{t}$ and $Zb\bar{b}$ events, and the total background level was found to be less than 1% and was neglected.

The WZ channel is more difficult to separate from background and more stringent cuts are needed. WZ candidates are required to have three leptons with $p_T > 25$ GeV and $|\eta| < 2.5$. Out of these three leptons, there must be one lepton anti-lepton pair with invariant mass within 12 GeV of the Z mass. In addition the event must have missing E_T above 25 GeV, and the transverse mass constructed from the missing E_T and the third lepton must be between 50 and 90 GeV. The angle in the transverse plane between the missing E_T vector and the third lepton must exceed 40°. Finally, to further suppress background from $t\bar{t}$ events, no more than one jet is allowed in the event, and its transverse momentum should not exceed 30 GeV. In our 100 fb⁻¹ pseudo-data sample, 2873 WZ candidates are left after these cuts with a background level of 1%, where the background sources considered were $t\bar{t}$, W^+W^- , ZZ, Z+jet and $Z\gamma$ events. This background was neglected in the further analysis.

The kinematic observable which is sensitive to the gauge boson polarization is the boson decay angle, θ_ℓ^* , which is the angle between the lepton (anti-lepton) produced in the W^- , Z (W^+) decay in the rest-frame of the boson with respect to the boson momentum in the di-boson rest-frame. Its distribution is given by,

$$\frac{1}{\sigma} \frac{d\sigma}{d \cos \theta_\ell^*} = \rho_{--} \frac{3}{8} (1 + \cos \theta_\ell^*)^2 + \rho_{++} \frac{3}{8} (1 - \cos \theta_\ell^*)^2 + \rho_{00} \frac{3}{4} \sin^2 \theta_\ell^*,$$

for the W boson, and by,

$$\frac{1}{\sigma} \frac{d\sigma}{d \cos \theta_\ell^*} = \rho_{--} \frac{3}{8} (1 + 2A \cos \theta_\ell^* + \cos^2 \theta_\ell^*) + \rho_{++} \frac{3}{8} (1 - 2A \cos \theta_\ell^* + \cos^2 \theta_\ell^*) + \rho_{00} \frac{3}{4} \sin^2 \theta_\ell^*,$$

for the Z boson. Here, $A = 2\mathbf{l}_\perp a_\ell / (\mathbf{l}_\perp^2 + a_\ell^2)$, where v_ℓ (a_ℓ) is the vector (axial-vector) coupling of the Z boson to leptons. The spin density matrix elements, ρ_{--} , ρ_{++} and

ρ_{00} correspond to transverse left-handed, transverse right-handed and longitudinal polarization of the boson. They satisfy $\rho_{--} + \rho_{++} + \rho_{00} = 1$ and each one of them can be interpreted as the fraction of bosons produced in the corresponding polarization state.

For ZZ events, it is rather straightforward to obtain θ_ℓ^* from the measured kinematic observables of the four leptons. For the WZ events, the longitudinal momentum of the neutrino, p_ℓ^ν , is not measured, but it can be calculated if the neutrino and the corresponding lepton are constrained to the W mass. A quadratic equation is obtained yielding two solutions for p_ℓ^ν . The value taken for p_ℓ^ν is a weighted average of the two solutions, where the weights are the corresponding SM cross-section values [8]. The resulting resolutions in the di-boson invariant mass, $\sqrt{\hat{s}}$, and in $\cos\theta_\ell^*$ are shown in Fig. 1.

Since the gauge boson polarization is expected to depend on $\sqrt{\hat{s}}$, the analysis is done separately for 4 (3) bins in $\sqrt{\hat{s}}$, for the ZZ (WZ) channels. In addition, there is a separation between the W^+Z and W^-Z channels. To correct for detector and analysis effects, the MC events are further divided into 10 bins in $\cos\theta_\ell^*$. For each bin i , the ratio between the number of events at the generator level and the number after the detector and selection is the correction factor, C_i , which accounts for both overall efficiency and resolution effects.

For the ZZ channel, we extract ρ_{00} using the event-by-event projection operator method,

$$\rho_{00} = \frac{1}{N_{\text{events}}} \sum_{\text{events}} C_i \cdot \Lambda_{00}(\cos\theta_\ell^*)$$

where C_i corresponds to the $(\sqrt{\hat{s}}, \cos\theta_\ell^*)$ bin containing the event and Λ_{00} is the projection operator for the longitudinal polarization state, given by, $\Lambda_{00} = 2 - 5\cos^2\theta_\ell^*$, and calculated for the measured $\cos\theta_\ell^*$ of that particular event.

This projection operator method could not be used for the WZ channels, since some of the bins in the MC distribution after the detector were sparse, yielding very large, uncertain and sometimes even infinite correction factors. Therefore, an event-by-event maximum likelihood method was adopted, where the inverse correction factors are used, and they multiply the theoretical $\cos\theta_\ell^*$ distribution for each event.

The systematic effect from the uncertainty in the Parton Distribution Functions (PDF) was studied in the ZZ channel by using different sets of PDFs (CTEQ, MRST, EHLQ2), and in the WZ channels by using the 40 CTEQ6M [9] error sets. Another source of systematic errors considered is MC statistics. For the WZ channels, the dominating systematic effect comes from the replacement of the SM cross-section values used for the weighted average of the two solutions for p_ℓ^ν by theoretical cross-section values [8] calculated with anomalous couplings within the Tevatron limits. The overall systematic errors are typically (20 - 50)% of the statistical ones.

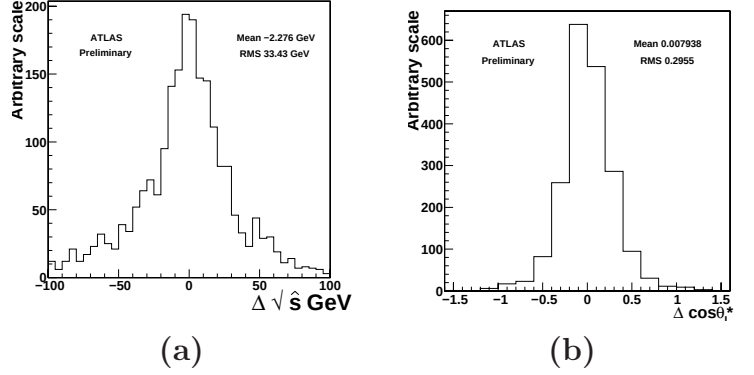


Figure 1: Resolution for WZ events in (a) $\sqrt{\hat{s}}$, (b) $\cos\theta_\ell^*$

Fig. 2 shows the results of the longitudinal polarization fraction, ρ_{00} , for the different channels as function of $\sqrt{\hat{s}}$. The curves correspond to the theoretically expected $\rho_{00}(\sqrt{\hat{s}})$, as obtained from large statistics, generator level MC samples.

In summary, it has been shown in this note that the measurement of the Z and W polarizations in ZZ and WZ events is feasible with an integrated luminosity of 100 fb^{-1} , which will be available only after several years of successful LHC running. The errors are dominated by statistics, and more accurate results are expected with the higher luminosities expected in SLHC.

I would like to thank the organizers of this workshop for the opportunity to present this work, and for the good organization and the nice atmosphere in the workshop. This work was partially supported by the Israel Science Foundation (grant No. 183/07).

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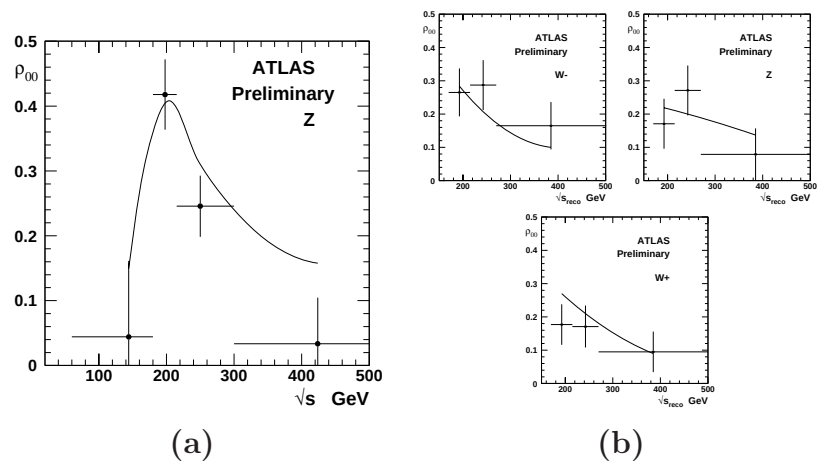


Figure 2: Longitudinal polarization fraction, ρ_{00} , as function of the di-boson mass, $\sqrt{\hat{s}}$, obtained from our 100 fb^{-1} pseudo-data samples, for (a) Z in ZZ events, (b) W^- (left), Z (right), W^+ (bottom)

LATEST RESULTS ON DEEPLY VIRTUAL COMPTON SCATTERING AT HERMES

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Abstract

The HERMES experiment at DESY collected a rich data set for the analysis of Deeply Virtual Compton Scattering (DVCS) utilizing the HERA longitudinally polarized electron or positron beams with an energy of 27.6 GeV and longitudinally and transversely polarized or unpolarized gas targets (H, D or heavier nuclei). The azimuthal asymmetries measured in the exclusive DVCS production allow the access to the real and imaginary parts of certain combinations of Generalized Parton Distributions. Latest results on combined analyzes of beam-spin and beam-charge asymmetries are presented.

1 Introduction

Hard exclusive electroproduction of real photons on nucleons, Deeply Virtual Compton Scattering (DVCS) is one of the theoretically cleanest ways to access Generalized Parton Distributions (GPDs). The theoretical framework of GPDs incorporates knowledge about form factors and parton distribution functions. GPDs depend on four kinematic variables: the squared four-momentum transfer t to the nucleon, the squared four-momentum transferred by the virtual photon $-Q^2$, x and ξ , which represent respectively the average and half the difference of the longitudinal momentum fractions carried by the probed parton in initial and final states. For the proton, there are four twist-2 GPDs per quark flavor: H_q , E_q , $\tilde{H}_q$, and $\tilde{E}_q$.

Data from DVCS are indistinguishable from the electromagnetic Bethe-Heitler (BH) process because of the same final state, see Fig. 1a. The real photon is radiated from the struck quark in DVCS or from the initial or scattered lepton in BH. The cross section of the exclusive photoproduction process can be written in the following form [1]:

$$\frac{d\sigma}{dQ^2 dx_B dt |d\phi} = \frac{x_B e^6}{32(2\pi)^4 Q^4 \sqrt{1+\epsilon^2}} \left[|\tau_{DVCS}|^2 + |\tau_{BH}|^2 + I \right], \quad (1)$$

where $\tau_{DVCS}(\tau_{BH})$ is the DVCS (BH) production amplitude, $I = \tau_{DVCS}^* \tau_{BH} + \tau_{BH}^* \tau_{DVCS}$ is the interference term, x_B is the Bjorken scaling variable. The amplitude of the BH process dominates at HERMES kinematics. However, the kinematic dependence of the cross section terms generate a set of azimuthal asymmetries which depend on the azimuthal angle ϕ between the real-photon production plane and the lepton scattering plane.

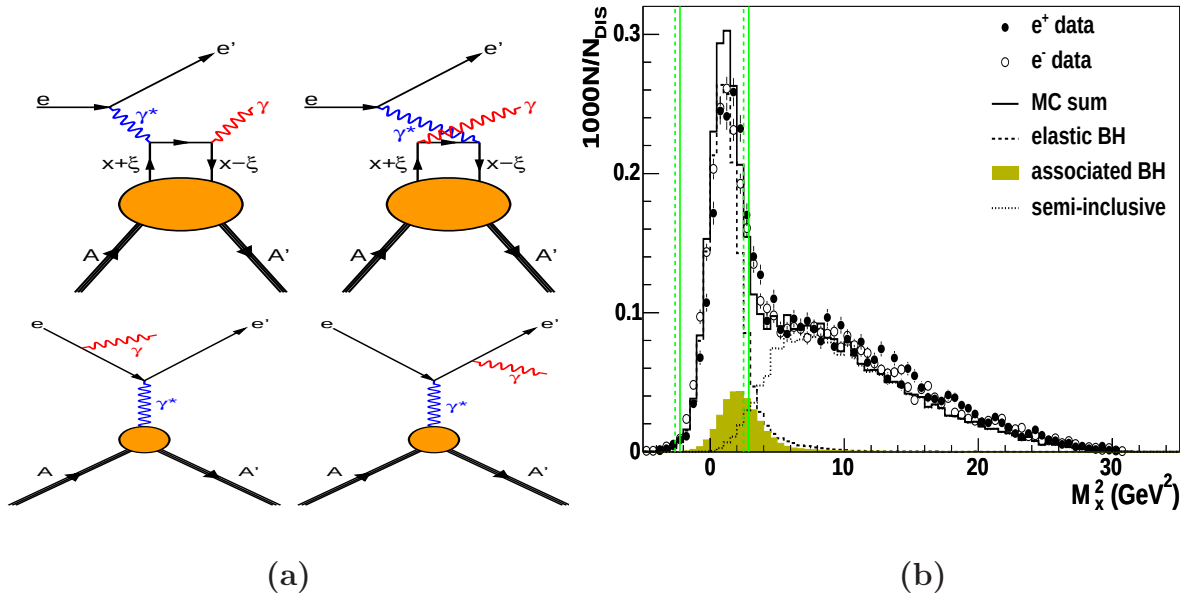


Figure 1: (a) Leading order diagrams for DVCS (top) and Bethe-Heitler process (bottom). (b) The measured distributions of electroproduced real-photon events versus the squared missing mass M_X^2 from positron (open circles) and electron beam (closed circles). The dashed line represents a Monte Carlo simulation including coherent DVCS and BH processes. The “associated” BH process with the excitation of resonant final states is presented in shaded area. The semi-inclusive background is presented as dotted line. Solid line represents the sum of all processes. The region between the two vertical lines indicates the selected “exclusive region”.

2 Azimuthal Asymmetries for DVCS

The cross section for a longitudinally polarized lepton beam scattered off an unpolarized proton target σ_{LU} can be related to the unpolarized cross section σ_{UU} by

$$\sigma_{LU}(\phi, P_B, C_B) = \sigma_{UU} \cdot \left[1 + P_B A_{LU}^{DVCS}(\phi) + C_B P_B A_{LU}^I(\phi) + C_B A_C(\phi) \right], \quad (2)$$

where A_{LU}^I (A_{LU}^{DVCS}) is the charge (in)dependent beam-helicity asymmetry (BSA) and A_C is the beam charge asymmetry (BCA). $C_B(P_B)$ denotes the beam charge (polarization). In the analysis, effective asymmetry amplitudes are extracted which include ϕ dependences from the BH propagators and the unpolarized cross section. Each asymmetry can be expanded in a Fourier harmonics of ϕ angular distributions:

$$A_{LU}^I(\phi) = \sum_{n=0}^2 A_{LU}^{\sin(n\phi)} \sin(n\phi) + \sum_{n=0}^1 A_{LU,I}^{\cos(n\phi)} \cos(n\phi), \quad (3)$$

$$A_{DVCS}^I(\phi) = \sum_{n=0}^2 A_{LU,DVCS}^{\sin(n\phi)} \sin(n\phi) + \sum_{n=0}^1 A_{LU,DVCS}^{\cos(n\phi)} \cos(n\phi), \quad (4)$$

$$A_C(\phi) = \sum_{n=0}^3 A_C^{\cos(n\phi)} \cos(n\phi) + A_C^{\sin(\phi)} \sin(\phi). \quad (5)$$

By combining the data taken with different beam charges and helicities, the amplitudes were fit simultaneously using a Maximum Likelihood method described in detail in Ref. [2].

3 Exclusive DVCS Events

The HERMES experiment [3] exploited longitudinally polarized 27.6 GeV electron or positron beam at HERA storage ring at DESY together with longitudinally or transversely polarized or unpolarized gas targets (H, D or heavier nuclei). Exclusive events were selected requiring the detection of the scattered lepton and one photon. In addition, as the recoiling proton has not been detected, the calculated missing mass was required to match the proton mass within the resolution of the spectrometer, which defines the “exclusive region”, see Fig. 1b.

Without the detection of the recoil proton it is not possible to separate the elastic DVCS and BH events from the “associated” process, where the nucleon in the final state is excited to a resonant state. Within the exclusive region, its contribution is estimated from Monte Carlo simulation to be about 12 %, which is taken as part of signal. The main background contribution of about 3% is originating from semi-inclusive π^0 production and is corrected for. The contribution from exclusive π^0 production is estimated to be less than 0.5%. The systematic uncertainties are obtained from a Monte Carlo simulation estimating the effects of limited acceptance, smearing, finite bin width and alignment of the detectors with respect to the beam. Other sources are background contributions and a shift of the position of the exclusive missing mass peak between the data taken with different beam charges.

4 Beam-Charge and Beam-Spin Asymmetries

In Figs. 1 and 2 results obtained on the hydrogen target are shown [4]. The first four rows of Fig. 1 represent different cosine amplitudes of the BCA, whereas the last row shows the fractional contributions of the associated BH process. In the first column the integrated result is shown, in the other columns the amplitudes are binned in $-t$, x_B and Q^2 variables. The error bars show the statistical and the bands represent the systematic uncertainties. The magnitudes of the first two cosine moments $A_C^{\cos 0\phi}$ and $A_C^{\cos \phi}$ increase with increasing $-t$, while having opposite signs, they are in agreement with theoretical calculations. At HERMES kinematic conditions, both mainly relate to the Compton Form Factor (CFF) which is a convolution of GPD function $\mathcal{H}$ with the hard scattering amplitude [4]. The constant term there is suppressed relative to the first moment. The second cosine moment appears in twist-3 approximation and is found to be compatible with zero like the third cosine moment, which is related to gluonic GPDs. Fraction of “associated” production is presented in the fifth row for each kinematic point.

The first sine moment $A_{LU,I}^{\sin \phi}$ is large and negative in HERMES kinematic conditions, see Fig. 2. This amplitude relates to the imaginary part of the CFF. The results of the GPD model calculations [5,6] based on the framework of double distributions [7] are also shown there. The model includes a Regge-inspired t -ansatz and a factorized t -ansatz with the calculation of the contribution of D-term [6]. We note, that the calculations performed without the contribution of D-term describe well HERMES results on BCA, see Fig. 2. Both BSA model calculations fail to describe the data except for small $-t$.

Results obtained on the deuteron target (not shown) are compatible with the data on a proton for almost all amplitudes in the same kinematic bins [8]. The nuclear-mass dependence of beam-helicity azimuthal asymmetries has been measured for targets rang-

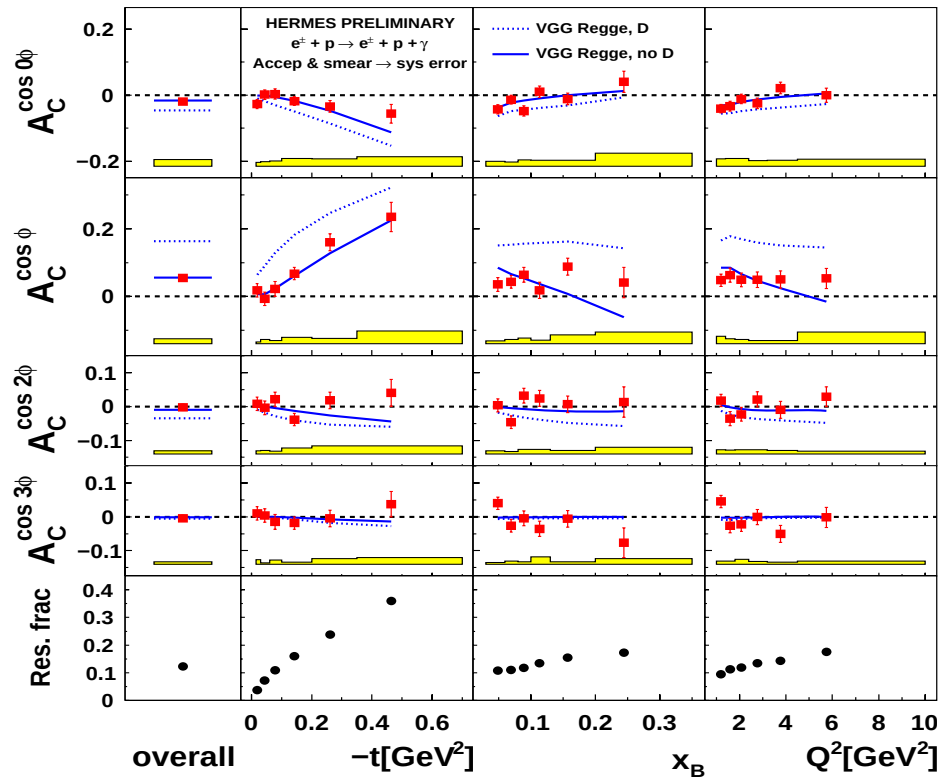


Figure 2: The amplitudes of the beam-charge asymmetry extracted from hydrogen data (red squares) [4]. The error bars (bands) represent the statistical (systematic) uncertainties. The curves are predictions of a double-distribution GPD model [5, 6].

ing from hydrogen to xenon [9]. For hydrogen, krypton and xenon, data were taken with both beam charges. For both DVCS and BH, coherent scattering occurs at small values of $-t$ and rapidly diminished with increasing $|t|$. Coherent and incoherent-enriched samples are selected according to a $-t$ threshold that is chosen to vary with the target such that for each sample approximately the same kinematic conditions are obtained for all target gases. The nuclear-mass dependence of the BCA and BSA is presented separately for the coherent and incoherent-enriched samples in Fig. 4 on the left and right panels, respectively. The $\cos \phi$ amplitude of the BCA is consistent with zero for the coherent-enriched samples for all three targets, while it is about 0.1 for the incoherent-enriched samples. The $\sin \phi$ amplitude of the beam-helicity asymmetry shown in Fig. 4 (right) has values of about -0.2 for both the coherent and incoherent-enriched samples. No nuclear-mass dependence of the BCA and BSA asymmetries is observed within experimental uncertainties. This is in agreement with models that approximate nuclear GPDs by nucleon GPDs neglecting bound state effects. The data do not support the enhancement of nuclear asymmetries compared to the free proton asymmetries for coherent scattering in spin-0 and spin- $\frac{1}{2}$ nuclei as anticipated by various models [10–12]. They also contradict the predicted strong A dependence of the beam-charge asymmetry resulting from a contribution of meson exchange between nucleons to the scattering amplitude [12].

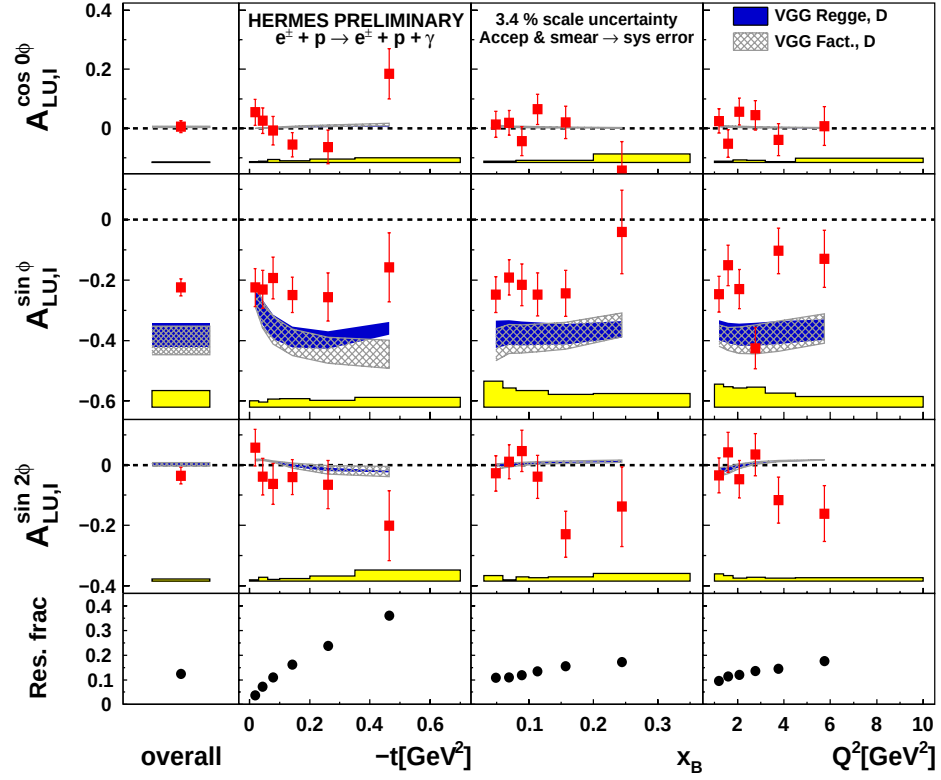


Figure 3: The amplitudes of the beam-helicity asymmetry from the interference term on the unpolarized hydrogen target [4]. The error bars (bands) represent the statistical (systematic) uncertainties. The curves are predictions of a double-distribution GPD model [5, 6].

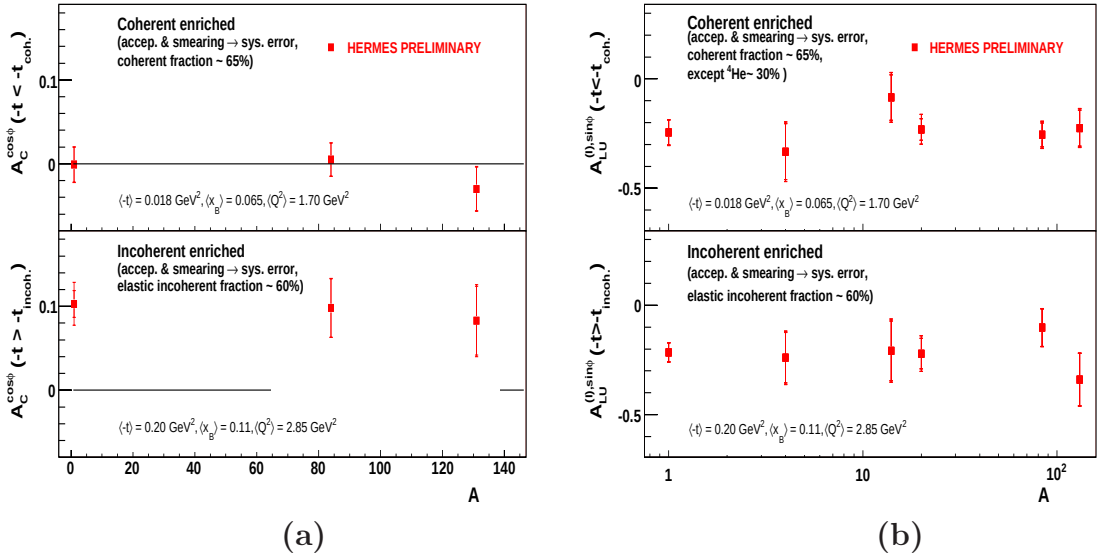


Figure 4: Nuclear-mass dependence of the $\cos \phi$ amplitude of the beam-charge asymmetry (a) and of the $\sin \phi$ amplitude of the beam-helicity asymmetry (b) for the coherent-enriched (upper panels) and incoherent-enriched (lower panels) data samples. The inner (full) errors bar represent the statistical (total) uncertainties.

5 Summary

HERMES has measured significant cosine moment of the beam-charge asymmetry and sine moment of the beam-charge dependent beam-helicity asymmetry in DVCS on hydrogen and deuterium targets. The statistical precision of the data allows us to provide constraints on theoretical calculations. The unknown contribution from the associated process can be understood only from the data taken with the recoil detector.

Beam-charge and beam-helicity asymmetries have been measured for targets ranging from hydrogen to xenon. No nuclear-mass dependence of the asymmetry amplitudes is observed within experimental uncertainties. The obtained results provide constraints on nuclear GPD models.

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POLARIZATION TRANSFER MECHANISM AS A POSSIBLE SOURCE OF THE POLARIZED ANTIPROTONS

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Abstract

We suggest to experimental study the polarization transfer mechanism in the inclusive reaction of the antiproton production by the polarized proton beam.

Introduction

The problem of producing the intense polarized antiproton source with the appropriate polarization becomes very acute at present. Such interest is instigated by the great potential of the physics with the polarized antiprotons [1]. Approximately a dozen suggestions were made about how to get the polarized antiprotons [2], but none of them (exclusion is the filtering technique [3]) was experimentally demonstrated as really working tool. Therefore we suggest the experiment at AGS for measuring the spin transfer tensor from the polarized protons to the secondary antiprotons produced at zero degree. We recall the standard technique of obtaining the unpolarized antiprotons by the unpolarized primary proton beam according to PS CERN scheme. Then we modify this scheme to produce the polarized antiprotons by the polarized proton beam. We also propose the schemes of measuring the spin transfer tensor at $\bar{p}$ production at zero angle as a function of the antiproton momentum.

Experimental evidences for the polarization transfer mechanism.

There is no any theoretical or experimental paper devoted to the direct prediction or proof of the existence of the polarization transfer mechanism in the inclusive reaction $p \uparrow + N \rightarrow \bar{p} \uparrow + X$, where the arrows indicate the polarization of the corresponding particle. Let's remind some other examples of the polarization transfer mechanism.

Neutron polarization by the spin transfer mechanism

Let's consider the production of polarized neutrons from the inclusive reaction $C(p, n)X$ with 590 MeV protons [4]. If polarized protons are used a polarized neutron beam may obtain. *It was found that the transfer of polarization from longitudinally polarized protons to the longitudinally polarized neutrons was most effective.* The absolute neutron beam intensity was about 10^8 n/s for integrated neutron flux over energy range from 100 to 590 MeV. Reaction $p(\rightarrow) + A = n(\rightarrow) + X$ was used for production of the polarized neutrons at zero degree. The initial proton beam momentum, intensity and polarization were

(1.205 ± 0.0012) GeV/c, 10^{10} p/s, 75% respectively. It was shown that the neutron longitudinal polarization reaches about 40% at the kinetic energies $T \geq 450$ MeV. This beam was extensively used to fulfill the complete set of the np -elastic scattering experiments.

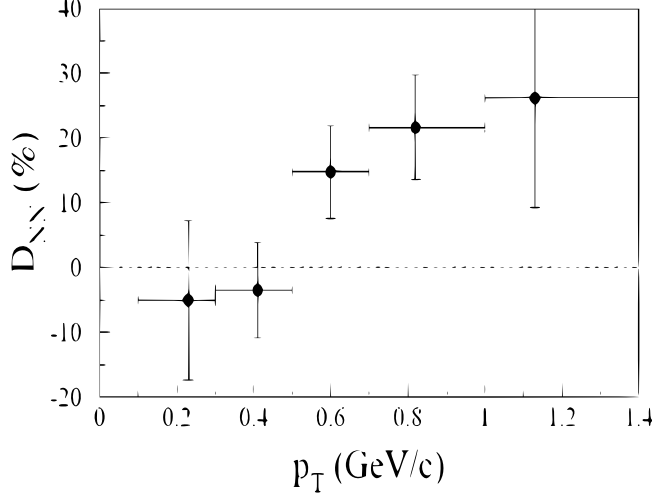


Figure 1: Depolarization D_{NN} data as a function of p_T .

D_{NN} reaches positive values of about 30% at high x_F and $p_T \sim 1.0$ GeV/c (see Fig 2). This result proves the existence of the essential spin transfer mechanism at sufficiently high energy.

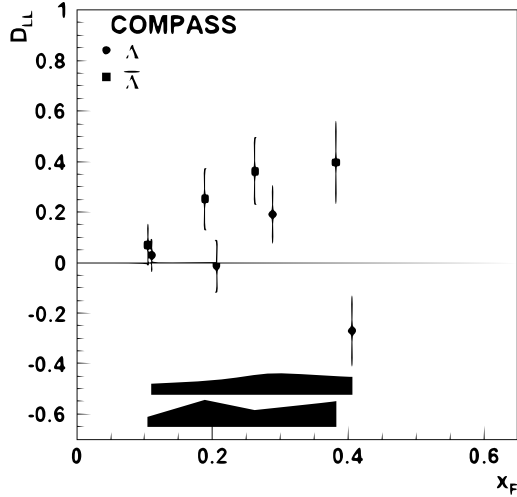


Figure 2: Longitudinal spin transfer of Λ and $\bar{\Lambda}$ in μp collisions

Obviously the spin transfer in this chain should be small as it is seen from this slide. 2) It is surprising that the spin transfer in the case of Λ -hyperon is zero. This is also in variance with the above general statement, since lambda and initial proton have common quarks. On the basis of these data we may expect that the spin transfer from the polarized proton to the final antiproton may be not zero also.

Spin Transfer tensor in Inclusive Λ^0 Production by Transversely Polarized Protons at 200 GeV/c

This is the first and only example of non-zero Spin Transfer Mechanism, experimentally established at 200 GeV/c [5]. The depolarization tensor $D_{NN} = K_{N0;N0}$ (Wolfstein's parameter) in the reaction $p \uparrow + p = \Lambda_0 \uparrow + X$ has been measured with the transversely polarized protons at 200 GeV/c over a wide x_F range and a moderate p_T .

Longitudinal spin transfer of Λ and $\bar{\Lambda}$ in DIS at COMPASS The Fig. 3 presents the results of longitudinal spin transfer of Λ and $\bar{\Lambda}$ in DIS at COMPASS [6]. We would like to emphasize two features of these datas: 1) The spin transfer tensor D_{LL} is not zero for $\bar{\Lambda}$. This is in contrast to the generally accepted statement that the spin transfer is absent in the case when the final and initial baryons do not have any common quark. It means that there is another channel of transferring the polarization. As an example we may refer to professor H.Artru, who indicated us that the polarized proton may emit the polarized gluon which fragments into polarized quark-antiquark pair. The polarized antiquark may fragment into polarized antibaryon. Obviously the spin transfer in this chain should be small as it is seen from this slide. 2) It is surprising that the spin transfer in the case of Λ -hyperon is zero. This is also in variance with the above general statement, since lambda and initial proton have common quarks. On the basis of these data we may expect that the spin transfer from the polarized proton to the final antiproton may be not zero also.

Suggestion of the experiments for measuring the polarization transfer from protons to antiprotons at AGS.

On the scheme (see Fig. 2) the initial slow extracted polarized proton beam strikes on the BTL-1 (Beam transfer line-1) the external production target. Since the polarization is normal to the orbital plane we should rotate it to the longitudinal direction by the Spin Rotator -1 (SR-1). It was found that the transfer of polarization from longitudinally polarized protons to the longitudinally polarized neutrons was most effective [4] so we expect that this method will be applicable to the production of polarized antiprotons. The polarized antiprotons will be transported by the BTL-2 to the polarimeter. Before reaching the polarimeter the longitudinal antiproton polarization should be rotated to transverse polarization. This is done because in strong interaction we can measure only transverse components of the polarization. The preferable direction of the antiproton polarization is the normal to the horizontal plane. So SR-2 makes this function. We have estimated yields of $\bar{p}$ for three different momenta. This is 6.5 GeV/c, where the maximum of yield of $\bar{p}$ is expected. And two other momenta, 1.5 GeV/c and 15.5 GeV/c, where the yields of $\bar{p}$ is smaller by one order of magnitude. So we have two possibilities to work with polarized antiprotons: A and B. A includes the measurements of the antiproton polarization in the momentum range 1.5 GeV/c to 3.1 GeV/c. This is because we would like to use the (g-2) storage ring at AGS. This is useful because we can essentially suppress the pion and kaon backgrounds to the antiproton beam. B includes the measurements of the antiproton polarization in the momentum range 3.1 GeV/c to 15.5 GeV/c. In this case we deal with antiproton beam containing huge backgrounds of pions and kaons. In this case we have to identify the particles in beam by using the Cherenkov counters. For this case the counting rates will be limited by the background particles (mostly by pions). In both cases we use the polarimeter based on the polarized proton target (PPT). So we will use as the analyzing reaction the antiproton-proton elastic scattering. We will use the PPT in polarimeter for two reasons: 1) There are no in literature the data with sufficient precisions on the antiprotons analyzing power in $\bar{p} - p$ and $\bar{p} - A$ -elastic scattering in the whole momentum range of our interest. So using PPT we can measure the analyzing power of the pbar-p-elastic scattering with needed accuracy at any energy. 2) Using the polarized target we can simultaneously measure the analyzing power and the polarization of antiprotons. We suggest to use the propane-diol polarized target [7]. Table presents the counting rate per cycle and the time needed for accumulation of the scattered events for each energy. In calculations we used [9]. To achieve the 1% precision in asymmetry measurement, one needs 10^4 scattering events. For accumulating this statistics we should use the net beam time of order indicated in the table. The detection of scattered particles

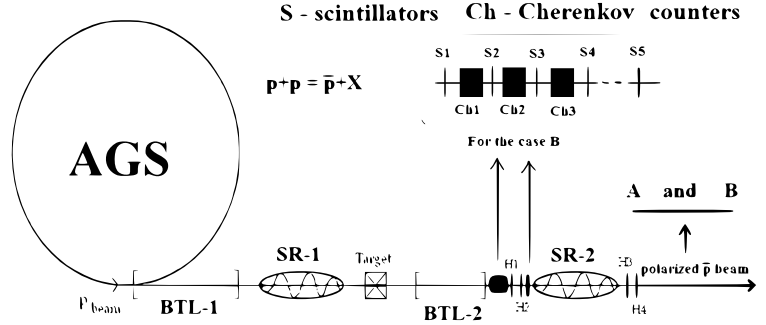


Figure 3: Suggestion of the experiment (a)

Figure 3: Suggestion of the experiment (a)

was done in the interval $-t$ from 0.2 to 1.0 $(GeV/c)^2$ and $\Delta\varphi/2\pi = 0.1$. So these times seem us reasonable for suggested experiment.

	1.5 GeV	3.1 GeV	6.5 GeV	15.5 GeV
$N_{\bar{p}}$	0.2	3	10	1
t, hours	70	5	1.5	14

Conclusions.

We are not aware of any paper devoted to the theoretical or experimental study of the spin transfer mechanism from the polarized protons to the inclusively produced at zero degree antiprotons. Basing on the several experimental data on the spin-transfer effect in elastic and inclusive reactions for protons and lambda hyperons we hope that similar effect may exist also for antiprotons. We would like to emphasize the importance to fulfill the special experiment for measuring the polarization transfer tensor K_{LL} in reaction $p(\rightarrow) + A = \bar{p}(\rightarrow) + X$ at AGS. Initial proton beam with momentum around 22 GeV/c, longitudinal polarization 70%, intensity 210^{11} p/cycle, may produces the longitudinally polarized $\bar{p}$ beam at zero degree. Our estimates show that though such an experiment is difficult but it is realizable.

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PROBING THE HADRON STRUCTURE WITH POLARISED DRELL-YAN REACTIONS IN COMPASS

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Abstract

The study of Drell-Yan (DY) processes involving the collision of an (un)polarised hadron beam on an (un)polarised (proton) target can result in a fundamental improvement of our knowledge on the transverse momentum dependent (TMDs) parton distribution functions (PDFs) of hadrons. The production mechanism of J/ψ and J/ψ - DY duality can also be addressed. The future polarised COMPASS DY experiment is discussed in this context, the most important features briefly reviewed.

1 Transverse spin dependent structure of the nucleon

Spin and transverse momentum dependent semi-inclusive hadronic processes have attracted much interest both, experimentally and theoretically in recent years. These processes provide us more opportunities to study the Quantum Chromodynamics (QCD) and internal structure of the hadrons, as compared to the inclusive hadronic processes or spin averaged processes. Extensive experimental studies been made in different reactions. In particular, the single transverse spin asymmetry (SSA) phenomena observed in various processes [1–6] (such as $H_a + H_b \rightarrow H + X$, $l + H \rightarrow H' + X$, $H_a + H_b \rightarrow l + l' + X$ and $l^+ + l^- \rightarrow H_a + H_b + X$,) have stimulated remarkable theoretical developments. Among them, two approaches in the QCD framework have been most explored: the transverse momentum dependent (TMD) approach see for example [7–9, 11, 10] and the higher twist collinear factorization approach [13, 12, 15, 14]. The first one deals with the TMD distributions in the QCD TMD factorization approach which is valid at small transverse momentum $k_T \ll Q^2$. Such functions generalize the original Feynman parton picture, where the partons only carry longitudinal momentum fraction of the parent hadron. They will certainly provide more information on hadron structure.

In the second approach the spin-dependent differential cross sections can be calculated in terms of the collinear twist three quark-gluon correlation functions in the collinear factorization formalism. This approach is applicable for large transverse momentum $k_T \gg \Lambda_{QCD}$. Recently it was shown that the above two approaches are consistent in the intermediate transverse momentum region $\Lambda_{QCD} \ll k_T \ll Q^2$ where both apply [16, 18, 19].

Since at COMPASS the Drell-Yan data are expected mainly at small transverse momentum in the following we will discuss only the first, namely QCD TMD factorization based approach. More over we limit ourselves by the TMD generated by non-zero transverse momentum $\mathbf{k}_T$.

If we admit a non-zero quark transverse momentum $\mathbf{k}_T$ with respect to the hadron momentum, the nucleon structure function is described, at leading twist, by eight PDF¹: $f_1(x, \mathbf{k}_T^2)$, $g_{1L}(x, \mathbf{k}_T^2)$, $h_1(x, \mathbf{k}_T^2)$, $g_{1T}(x, \mathbf{k}_T^2)$, $h_{1T}^\perp(x, \mathbf{k}_T^2)$, $h_{1L}^\perp(x, \mathbf{k}_T^2)$, $h_1^\perp(x, \mathbf{k}_T^2)$ and $f_{1T}^\perp(x, \mathbf{k}_T^2)$. The first three functions, integrated over $\mathbf{k}_T^2$, give $f_1(x)$, $g_1(x)$ and $h_1(x)$, respectively. The last two ones are T-odd PDFs. The former, known also as Boer-Mulders function describes the unbalance of number densities of quarks with opposite transverse polarization with respect to the unpolarized hadron momentum. The latter, known as Sivers function describes how the distribution of unpolarized quarks is distorted by the transverse polarization of the parent hadron. The correlation between $\mathbf{k}_T$ and parton/hadron transverse polarization is intuitively possible only for a non-vanishing orbital angular momentum of the quarks themselves.

Hence, extraction of $h_1^\perp$ and $f_{1T}^\perp$, as well as of the poorly known h_1 , is of great interest in revealing the partonic (spin) structure of hadrons (see ref. [17] for review). Here, we concentrate mainly on the leading twist transversity h_1 , T-odd Boer-Mulders ($h_1^\perp$) and Sivers ($f_{1T}^\perp$) functions.

In the DY process, as shown in Fig.1, quark and antiquark annihilate with the production of a lepton pair. The fragmentation process is absent in DY, but here we have to deal with the convolution of quark and antiquark PDFs.

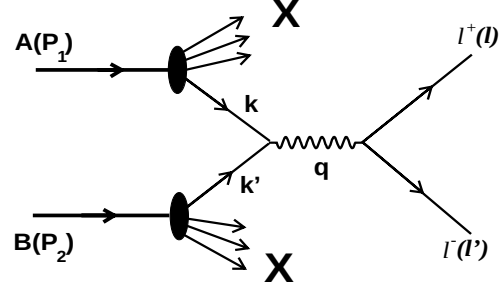


Figure 1: Feynman diagram of the Drell-Yan process; annihilation of a quark-antiquark pair with the production of a lepton pair.

2 Drell-Yan measurements at COMPASS

The proposed measurements of the polarised and unpolarised Drell-Yan process at COMPASS concentrate on the study of transversity (h_1) and of the Boer-Mulders ($h_1^\perp$) and Sivers ($f_{1T}^\perp$) functions. In the next sections we discuss the angular distributions of the dilepton pair as well as some important aspects of the J/ψ formation mechanism in pion-proton interactions.

The COMPASS DY program focuses on valence quark-antiquark ($x_B > 0.1$) annihilation to access the spin-dependent PDFs in the energy scale $Q^2 \gg 1 \text{ (GeV/c)}^2$. Valence antiquarks can be provided by using intense pion beams.

Because of the low cross-section of the Drell-Yan process (fractions of nanobarn) high luminosity and a large acceptance set-ups are required, as well as large value of the polarisation in order to access spin structure information.

All these features are provided by the multi-purpose large acceptance COMPASS spectrometer, in combination with the SPS M2 secondary beams and the COMPASS Polarised Target which combines a 1.2 m length of the target material with a 90% polarization reached for the ammonia target and a large acceptance of the target superconducting magnet.

The possibility to study the DY process with a pion beam at COMPASS was first dis-

¹We are using the so called Amsterdam notations [20]

cussed at the SPSC Meeting at Villars (September 2004) [31]. Compass is also considering the use of an antiproton beam in a second phase of the DY experiment.

3 Description of Drell–Yan processes

In the following we have adopted the formalism of Schlegel [11]. Since COMPASS will not have the possibility to use a polarized hadron beam only reaction of an unpolarised hadron beam (H_a) with a polarized target (H_b) have been considered:

$$H_a(P_a) + H_b(P_b, S) \rightarrow \gamma^*(q) + X \rightarrow l^-(l) + l^+(l') + X \quad (1)$$

where $P_{a(b)}$ is the momentum of the beam (target) hadron, $q = (l + l')$, l and l' are the momenta of the virtual photon, lepton and anti-lepton respectively and S is the four-vector of the target polarization. Among the 12 structure functions which are present in the cross-section for unpolarized, longitudinally or transversely polarized targets, six have a simple interpretation within the LO QCD parton model and 5 of them give rise to azimuthal asymmetries in the dilepton production:

- $A_U^{\cos 2\phi}$ gives access to Boer-Mulders functions of incoming hadrons,
- $A_L^{\sin 2\phi}$ – to Boer-Mulders functions of beam hadron and $h_{1L}^\perp$ function of the target nucleon,
- $A_T^{\sin \phi_S}$ – to Sivers function of the target nucleon,
- $A_T^{\sin(2\phi+\phi_S)}$ – to Boer-Mulders functions of beam hadron and $h_{1T}^\perp$ (pretzelosity) function of the target nucleon,
- $A_T^{\sin(2\phi-\phi_S)}$ – to Boer-Mulders functions of beam hadron and h_1 (transversity) function of the target nucleon.

where θ is polar lepton angle and ϕ and ϕ_S azimuthal lepton angles. Within the QCD TMD PDFs approach the remaining asymmetries can be interpreted as higher order in q_T/q kinematic corrections and as contribution of non-leading twist PDFs.

4 Unpolarised Drell–Yan processes

The angular distribution for the unpolarised case is known since a long time [22, 21] and it is commonly parametrised as

$$\frac{1}{N} \frac{dN}{d\Omega} \equiv \frac{d\sigma}{d^4q d\Omega} \Big/ \frac{d\sigma}{d^4q} \frac{3}{4\pi} \frac{1}{(\lambda + 3)} [1 + \lambda \cos^2 \theta + \mu \sin 2\theta \cos \phi + \frac{\nu}{2} \sin^2 \theta \cos 2\phi], \quad (2)$$

where leptons polar and azimuthal angles are defined in the Collins–Soper frame. In the QCD collinear parton model the coefficients λ , μ and ν are not independent and the Lam–Tung sum rule [23] holds

$$1 - \lambda = 2\nu, \quad (3)$$

which is trivial in collinear LO approximation:

$$\lambda^{LO} = 1, \quad \mu^{LO} = 0, \quad \nu^{LO} = 0. \quad (4)$$

QCD corrections to the Born cross-section allow $\lambda \neq 1$ and $\nu \neq 0$, nevertheless the Lam–Tung sum rule remains valid up to $\mathcal{O}(\alpha_s)$ corrections. However, large azimuthal asymmetries in the distribution of the final leptons observed in high-energetic collisions of pions and anti-protons, with nuclei [24–27], imply a strong violation of the Lam–Tung sum rule. In particular, an unexpectedly large $\cos 2\phi$ modulation of the cross section was observed. As it was mentioned in the previous subsection this modulation appears at twist-2 level for not vanishing Boer Mulders functions of pions and nucleons.

5 Transversely polarised Drell–Yan processes

The spin dependent part of the single transversely polarised Drell–Yan process in general contains 5 non vanishing over azimuthal angle integration asymmetries (Sec. 3). Within the QCD parton model with TMD DFs at twist-2 level only 3 of them survive and give access to the Boer–Mulders function of the pion and to the Siverts ($f_{1T}^\perp$), pretzelosity ($h_{1T}^\perp$) and transversity (h_1) DFs of the nucleon, see previous subsection.

Let us stress that the Siverts and the Boer–Mulders TMD PDFs are T-odd objects. Their field theoretical definition involves a non-local quark–quark correlator which contains the so-called gauge-link operator. While ensuring the colour-gauge invariance of the correlator, this gauge-link operator makes the Siverts and the Boer–Mulders functions process dependent. In fact, on general grounds it is possible to show that the $f_{1T}^\perp$ and the $h_1^\perp$ functions extracted from Drell–Yan processes and those obtained from semi-inclusive DIS should have opposite signs [30], i.e.

$$f_{1T}^\perp \Big|_{DY} = -f_{1T}^\perp \Big|_{DIS} \quad \text{and} \quad h_1^\perp \Big|_{DY} = -h_1^\perp \Big|_{DIS}. \quad (5)$$

An experimental verification of the sign-reversal property of the Siverts function would be a crucial test of QCD in the non-perturbative regime.

5.1 COMPASS DY apparatus acceptance

The acceptance of the COMPASS spectrometer for Drell–Yan events with $\mu^+\mu^-$ pairs in the final state was evaluated using a Monte-Carlo chain starting from PYTHIA 6.2 [32] as event generator following with Comgeant, based on Geant 3.21 [33] program which simulates the particles interaction with the COMPASS apparatus. Three main creation were used to optimize the energy of the incoming pion beam: the total DY event rate, the acceptance of the DY events by the COMPASS apparatus, and the covered range un x_1 and x_2 . Also studies of the combinatorial background and open-charm decays was performed.

As a result a pion beam momentum of 190 GeV/c was chosen.

As one can see from the Figure 2, the COMPASS is sensitive (has a high acceptance in x_p) exactly in the range where the maximal asymmetry is expected. Also visible is that the COMPASS kinematics acceptance is large in the valence region of both q and $\bar{q}$ what corresponds in practice to the pure u-dominance in the quark annihilation and simplify all analysis.

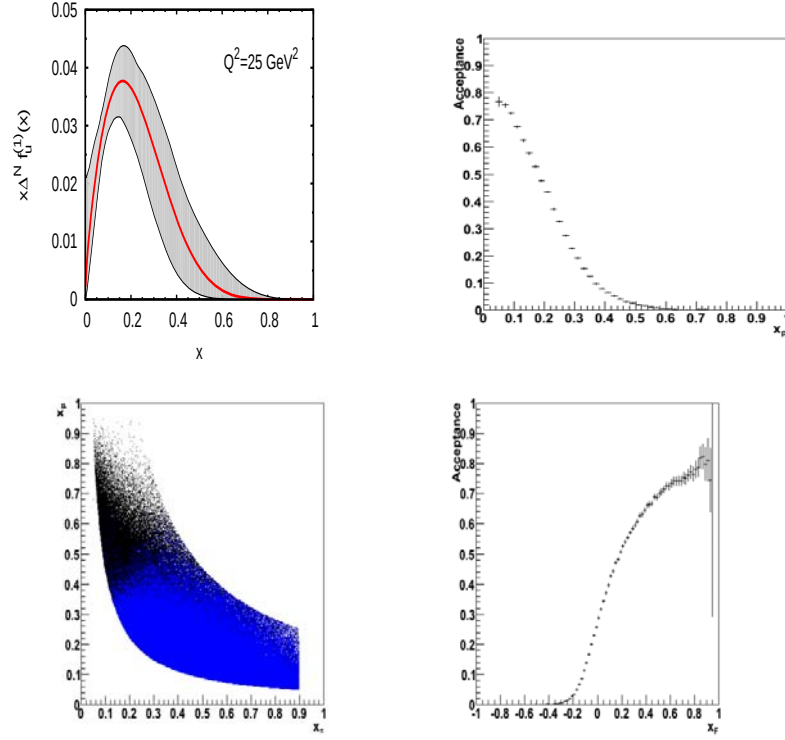


Figure 2: The left upper panel shows the first moment of the Siverts function for the u quark calculated at $Q^2 = 25 \text{ GeV}^2$ from [37]. The left lower panel shows the COMPASS covered kinematic region in x_p versus x_π (in blue). In the right upper and lower panels the COMPASS acceptance as a function of x_p and x_F respectively are shown.

5.2 Expected statistical errors and theory predictions on asymmetries

In our estimates we assume two years of running with the luminosity of $\approx 1.18 \times 10^{32} \text{ s}^{-1} \text{ cm}^{-2}$ and a 120 cm long NH_3 polarized target. Table 1 presents the expected statistical errors for the Siverts asymmetry $\delta A_T^{\sin \phi_S}(x_F)$, taking into account single bin, as a function of the beam energy. In the first column the simulated pion beam momentum is presented as it was used in Monte-Carlo, second column corresponds to the $\mu\bar{\mu}$ mass range $2. - 2.5 \text{ GeV}$, where some contribution from the combinatorial background is expected, and third one – to the most favourable (background free) $\mu\bar{\mu}$ mass range $4. - 9. \text{ GeV}$.

$\delta A_T^{\sin \phi_S}$	$2 < M_{\mu\mu} < 2.5$	$4 < M_{\mu\mu} < 9.$
$p_{\pi^-} = 106 \text{ GeV}/c$	$x^*0.0087$	0.0236
$p_{\pi^-} = 160 \text{ GeV}/c$	$x^*0.0081$	0.0188
$p_{\pi^-} = 190 \text{ GeV}/c$	$x^*0.0079$	0.0174
$p_{\pi^-} = 213 \text{ GeV}/c$	$x^*0.0078$	0.0166

Table 1: Expected statistical errors in the Siverts asymmetry, assuming 2 years of data-taking for pion beam momenta used in the Monte-Carlo and two $\mu\bar{\mu}$ invariant mass ranges.

In figure 3 the expected errors on the Sivers asymmetry are presented, as an example, with different number of x_F bins, for the DY high mass region; together with different theory predictions. As can be seen, depending on the number of bins, a statistical error of 0.01 to 0.02 is reachable.

The asymmetries are evaluated at $\langle M \rangle \simeq 5 \text{ GeV}$ (estimated using PYTHIA). The three lower curves (solid, dashed and dot-dashed lines) represent the estimate of the Sivers single spin asymmetry for COMPASS for 3 different parametrization of the Sivers functions.

Solid and dashed lines correspond respectively to fits I (Eq.6) and II (Eq.7) for Sivers function from Ref. [34] and dot-dashed line corresponds to fit from Ref. [35] (fit III – Eq.8).

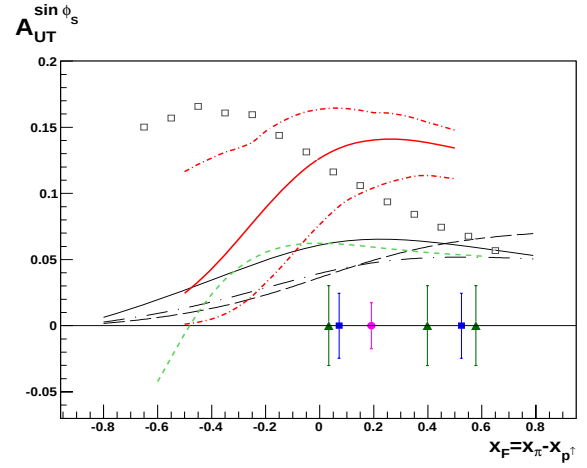


Figure 3: Theoretical predictions and expected statistical errors on Sivers asymmetry in the DY process $\pi^- p \rightarrow \mu^+ \mu^- X$ in the high dimuon mass region $4. < M_{\mu\mu} < 9. \text{ GeV}/c^2$. The three lower curves (black solid, dashed and dot-dashed lines) Ref. [34] and [35], the solid red line Ref. [36] and the predictions obtained in Ref. [29] and Ref. [38] are shown by squares and short-dashed line correspondingly.

$$x f_{1T}^{\perp(1)u} = -x f_{1T}^{\perp(1)d} = 0.4x(1-x)^5, \quad (6)$$

$$x f_{1T}^{\perp(1)u} = -x f_{1T}^{\perp(1)d} = 0.1x^{0.3}(1-x)^5, \quad (7)$$

$$x f_{1T}^{\perp(1)u} = -x f_{1T}^{\perp(1)d} = (0.17...0.18)x^{0.66}(1-x)^5 \quad (8)$$

$$(9)$$

The solid red line shows the Sivers asymmetry for the same mass range, integrated in q_T up to $1 \text{ GeV}/c$ Ref. [36], while the dashed lines shows the 1 sigma error band of the prediction. The predictions obtained in Ref. [29] and Ref. [38] are shown by squares and short-dashed line correspondingly.

6 Competition and complementarity

There are plans for future polarised Drell-Yan experiments at BNL, CERN, Fermilab, GSI, J-PARC and JINR. Some of them are presented in table 2.

Only PAX at FAIR (GSI-Darmstadt) and NICA at JINR (Dubna) plan to measure transverse double polarised Drell-Yan processes. In Dubna it is proposed to study Drell-Yan in proton-proton or deuteron-deuteron polarised beams collisions (access only to interactions between valence quarks and sea anti-quarks). The PAX collaboration plans to polarise anti-protons to study the interactions between valence quarks and valence antiquarks. However the possibility to get a beam of polarised anti-protons still has to be demonstrated. Both these collaborations plan to study e^+e^- final states.

The Drell-Yan programs at RHIC and J-PARC both foresee, like COMPASS, to measure single-spin asymmetries in the Drell-Yan process, but unlike COMPASS, they have only access to valence-sea quarks interactions in $p - \bar{p}$ collisions. The E906 project is

Facility		Type		s (GeV^2)	Timeline
RHIC (STAR)	[39]	collider,	$p^\uparrow p$	200^2	> 2013
E906 (Fermilab)	[40]	fixed target,	pp ,	250	> 2011
J-PARC	[41]	fixed target,	$pp^\uparrow, \pi p^\uparrow$	$60 \div 100$	> 2015
GSI (PAX)	[42]	collider,	$\bar{p}^\uparrow p^\uparrow$	200	> 2017
GSI (Panda)	[43]	fixed target,	$\bar{p}p$	30	> 2016
NICA	[44]	collider,	$p^\uparrow p^\uparrow, d^\uparrow d^\uparrow$	676	> 2014
COMPASS	(this letter)	fixed target,	$\pi^\mp p^\uparrow$	$300 \div 400$	> 2010

Table 2: Future Drell-Yan experiments.

oriented to the study of the sea quark distribution in the proton and can be considered as a good complementary measurement with respect to COMPASS DY.

The Panda experiment is rather designed for the J/ψ formation mechanism study than for the Drell-Yan physics, because of the very small anti-proton beam energy (15 GeV).

7 Conclusion

Single polarised Drell-Yan is a powerful tool to study hadron spin structure. Because of its nature it can provides us with new and complementary information with respect to what we can learn from polarised SIDIS.

If the COMPASS DY program will be approved, COMPASS will have the chance to be the first ever experiment to perform SSA measurements in Drell-Yan processes to access the spin-dependent PDFs in the valence quark region and the statistical significance of the result will be high.

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POLARIZATION OF VALENCE, NON-STRANGE AND STRANGE QUARKS IN THE NUCLEON DETERMINED BY COMPASS

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Abstract

The results of quark helicity distributions are presented which were determined by the COMPASS experiment at CERN. Different approaches have been accomplished in order to access various sets of quark distributions. The analysis of the spin structure function g_1 gives information about the contribution to the spin of the nucleon from the sum of quarks. From the method of difference asymmetry for hadrons with opposite charges one can learn about polarized valence quark distribution. Finally, a recent analysis of the full flavour separation was presented in Quark Parton Model (QPM) and LO QCD approximation. All analyses have been done in perturbative regime $Q^2 > 1 \text{ (GeV/c)}^2$.

1 Introduction

Angular momentum conservation requires that the spin of the nucleon can be represented by the sum rule:

$$\frac{1}{2} = \frac{1}{2}\Delta\Sigma + \Delta G + L_q + L_G, \quad (1)$$

where $\Delta\Sigma$ is a contribution from the sum of different flavours of quark helicity distributions, ΔG the contribution from helicities of gluons, and $L_{q,G}$ stands for quarks and gluons angular orbital momenta. The first experimental measurement of the spin structure of the nucleon started at CERN with the EMC collaboration [1]. At that time it was widely believed that the Ellis-Jaffe sum rule [2] holds and thus quarks carry roughly 60% of the nucleon helicity. Unexpectedly the interpretation of the EMC results indicated very small $\Delta\Sigma$ [1, 3] and that is how the so-called “spin crisis” began. The EMC results have been later confirmed by many other experiments [4–6, 8, 9] and brought the conclusion that quarks carry only about 30% of the nucleon spin.

2 Experimental setup

COMPASS is a fixed target experiment at CERN. One of its purposes is to study the spin structure of the nucleon. Until 2009 two separate programs have been realized: with muon and hadron beam respectively. The muon beam program was divided into sub-projects where the target was polarized either longitudinally or perpendicularly to the beam direction and the target material contained either polarized deuterons (^6LiD) or protons (NH_3). The results presented in this paper are based on the data collected in the

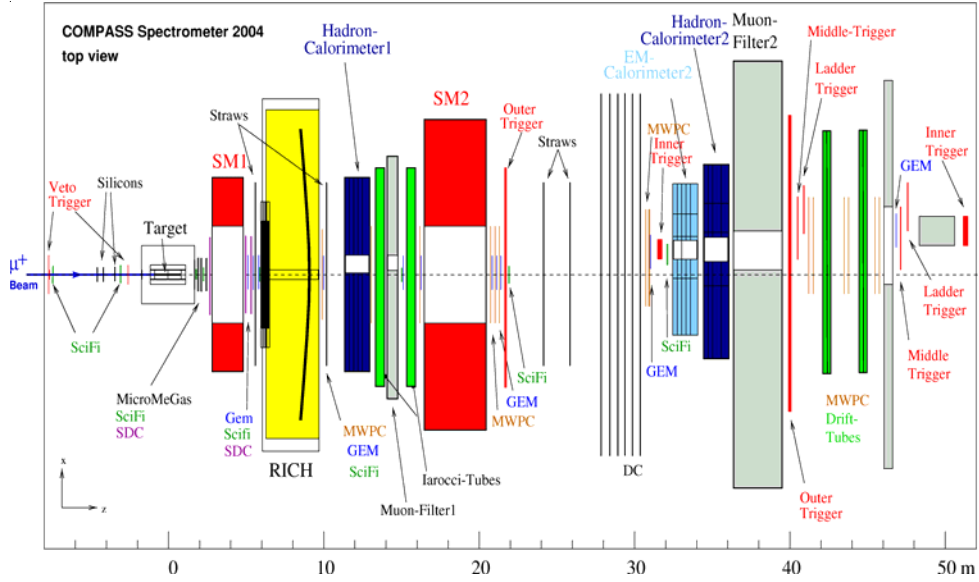


Figure 1: The COMPASS spectrometer.

years 2002-2007 with muons scattered on longitudinally polarized deuteron and proton targets.

The COMPASS muon beam is produced in several steps. Protons of 400 GeV energy coming in a cycles from the Super Proton Synchrotron collide with a solid beryllium target and produce among others charged pions and kaons. The momentum selected pions and kaons decay in a 500 m long tunnel. One of the main decay products are naturally polarized muons. The average polarization is 76% at the energy of 160 GeV.

In the years 2002-2004 the target was composed of two oppositely polarized 60 cm long cells, separated by a 10 cm gap. In 2006 and 2007 three target cells were used instead of two, in order to minimize systematic errors. The cells are located inside a superconducting solenoid which provides a homogenous 2.5 T magnetic field and are exposed to the same beam flux. The target is polarized using the dynamic nuclear polarization technique up to 50% in case of ${}^6\text{LiD}$ and 90% in case of NH_3 .

The COMPASS spectrometer (Fig. 1) is made of two parts which detect tracks in different kinematical regions. Each part contains its own bending magnet, hadron and electromagnetic calorimeter, and a set of tracking detectors of various types. Additionally the first spectrometer is equipped with a large RICH detector, which provides hadron identification. The RICH is designed to separate pions, kaons and protons in a wide range of momenta from 2.5 GeV, up to 50 GeV. A detailed description of COMPASS spectrometer can be found in [10].

3 The spin-dependent structure function g_1^N

The longitudinal spin-dependent structure function of the nucleon is given by:

$$g_1 \approx \frac{F_2}{2x(1+R)} A_1, \quad (2)$$

where F_2 and R are the spin-independent structure functions and A_1 is the cross section longitudinal spin asymmetry, which is related in the quark parton model to the quark

helicity distributions Δq in the following way:

$$A_1(x) = \frac{\sigma^{\uparrow\downarrow} - \sigma^{\uparrow\uparrow}}{\sigma^{\uparrow\downarrow} + \sigma^{\uparrow\uparrow}} \approx \frac{\sum e_q^2 \Delta q(x)}{\sum e_q^2 q(x)}. \quad (3)$$

In the above equation arrows correspond to relative helicity orientation of the incoming muon and the nucleon in the target.

In order to stay in the DIS domain the analysis has been performed for events with $Q^2 > 1 \text{ (GeV/c)}^2$ and $0.1 < y < 0.9$. This region allows COMPASS to collect 135 million events on deuteron target and 93 million events on proton target. The results on g_1^d and g_1^p as a function of x Bjorken are presented in Fig. 2. For comparison the values of the several other experiments have been shown at the same plot. The COMPASS deuteron data, evolved to $Q^2 = 3 \text{ (GeV/c)}^2$, have been used to extract $\Delta\Sigma$ from the first moment of g_1^N [11, 12]:

$$\Gamma_1^N = 0.05 \pm 0.003(\text{stat.}) \pm 0.003(\text{evol.}) \pm 0.005(\text{syst.}). \quad (4)$$

The correction to the first moment from the unmeasured region of $x < 0.004$ and $x > 0.7$ is about 2%. This value can be translated to the flavour-singlet axial current matrix element a_0 using:

$$\Gamma_1^N(Q^2) = \frac{1}{9} \left(1 - \frac{\alpha_s(Q^2)}{\pi} + \mathcal{O}(\alpha_s^2) \left(a_0(Q^2) + \frac{1}{4}a_8 \right) \right), \quad (5)$$

which in $\overline{\text{MS}}$ scheme, and the $Q^2 \rightarrow \infty$ is equal to the fraction of nucleon spin carried by quarks. The value of $a_8 = 0.585 \pm 0.025$ has been taken from hyperon β decays [13], assuming $\text{SU}(3)_f$ flavour symmetry. Finally, the light and strange quark polarization was found to be:

$$\Delta\Sigma = a_{0(Q^2 \rightarrow \infty)} = 0.33 \pm 0.03(\text{stat.}) \pm 0.05(\text{syst.}), \quad (6)$$

$$(\Delta s + \Delta \bar{s}) = \frac{1}{3} (\hat{a}_0 - a_8) = -0.08 \pm 0.01(\text{stat}) \pm 0.02(\text{syst}). \quad (7)$$

4 Valence quark distribution

For the semi-inclusive reaction $\mu N \rightarrow \mu' h X$, the asymmetry in LO QCD can be written in the form

$$A_1^h = \frac{\sum e_q^2 \Delta q(x) D_q^h(z)}{\sum e_q^2 q(x) D_q^h(z)}. \quad (8)$$

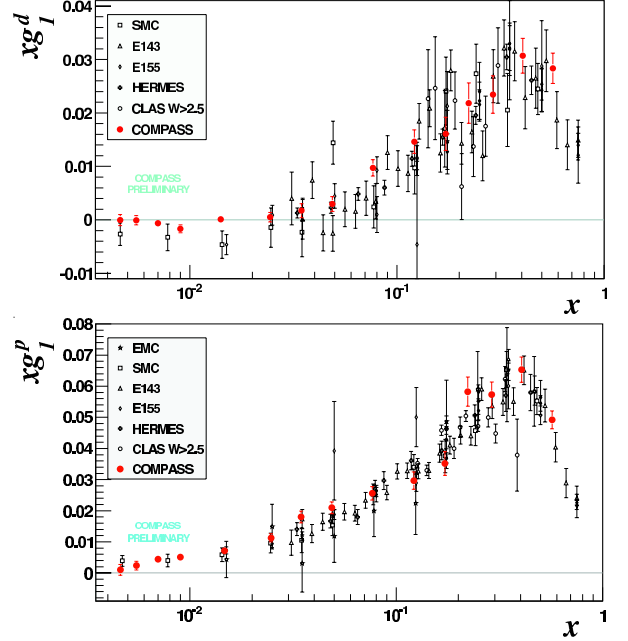


Figure 2: Spin dependent structure function xg_1^d (top) and xg_1^p (bottom) as a function of x at measured Q^2 . The agreement with the other experiments is obvious.

Comparing this formula to the inclusive case (Eq. 3) one can see that additionally fragmentation functions D_q^h are required. Due to this fact semi-inclusive asymmetries evaluated for different hadron types give a possibility to determine the polarization separately for different quark flavours. Such analysis will be presented in the next section. Now let us consider the so-called difference-charge asymmetry defined as follows [14, 15]:

$$A^{h^+-h^-} = \frac{(\sigma_{\uparrow\downarrow}^{h^+} - \sigma_{\uparrow\downarrow}^{h^-}) - (\sigma_{\uparrow\uparrow}^{h^+} - \sigma_{\uparrow\uparrow}^{h^-})}{(\sigma_{\uparrow\downarrow}^{h^+} - \sigma_{\uparrow\downarrow}^{h^-}) + (\sigma_{\uparrow\uparrow}^{h^+} - \sigma_{\uparrow\uparrow}^{h^-})}. \quad (9)$$

In the leading order QCD approximation the fragmentation functions cancel out in the $A^{(h^+-h^-)}$, giving direct access to valence quark polarization. Moreover for the deuteron target hadron identification is not needed, since

$$A^{h^+-h^-} = A^{\pi^+-\pi^-} = A^{K^+-K^-} = \frac{\Delta u_v(x) + \Delta d_v(x)}{u_v(x) + d_v(x)}. \quad (10)$$

The difference asymmetry can be obtained from single hadron asymmetries as $A^{h^+-h^-} = \frac{A^{h^+} - r A^{h^-}}{1-r}$, where $r = \sigma^{h^-}/\sigma^{h^+}$. This ratio of the cross-sections has been determined with the help of Monte Carlo simulation. Figure 3 shows difference asymmetry as a function of x Bjorken. In addition to standard selections as in the inclusive case the cut on the fraction of photon energy carried by hadron $0.2 < z < 0.85$ was implemented. The valence quark polarization was evaluated with formula (10) after correction for R and the deuteron D-state contribution:

$$\Delta u_v + \Delta d_v = \frac{u_v + d_v}{(1+R)(1-1.5\omega_D)} A^{h^+-h^-}.$$

As an unpolarized valence distribution MRST04 at LO has been used [16]. It is worth to notice that contribution from sea can be neglected in last three bins ($x > 0.3$). This gives opportunity to use an alternative approach in this region:

$$\Delta u_v + \Delta d_v = \frac{36}{5} \frac{g_1^d(x, Q^2)}{(1-1.5\omega_D)} - \left[2(\Delta \bar{u} + \Delta \bar{d}) + \frac{2}{5}(\Delta s + \Delta \bar{s}) \right]. \quad (11)$$

One can judge from Fig. 4 (left) that this method is much more precise. The integral over the measured x range is equal $0.40 \pm 0.07(stat.) \pm 0.05(syst.)$ supporting models with asymmetric sea polarization $\Delta \bar{u} = -\Delta \bar{d}$.

5 Flavour separation of helicity distributions

Recently COMPASS published results of helicity densities of valence, non-strange sea and strange quarks obtained in LO QCD approximation from deuteron data [17]. Just now

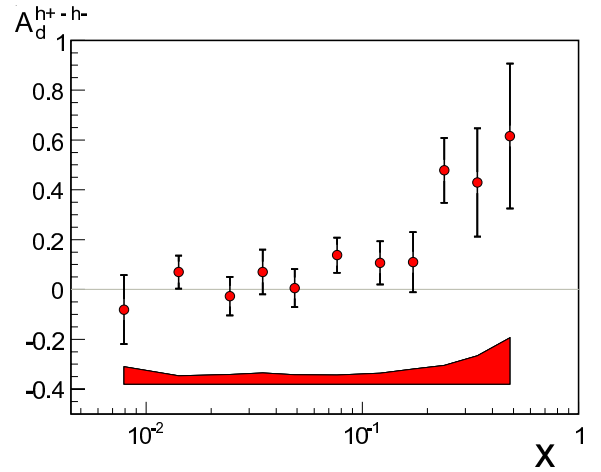


Figure 3: Difference-charge asymmetry of unidentified hadrons $A_d^{h^+-h^-}$ determined by COMPASS. The error bars are statistical ones and the band at the bottom represents systematical uncertainties.

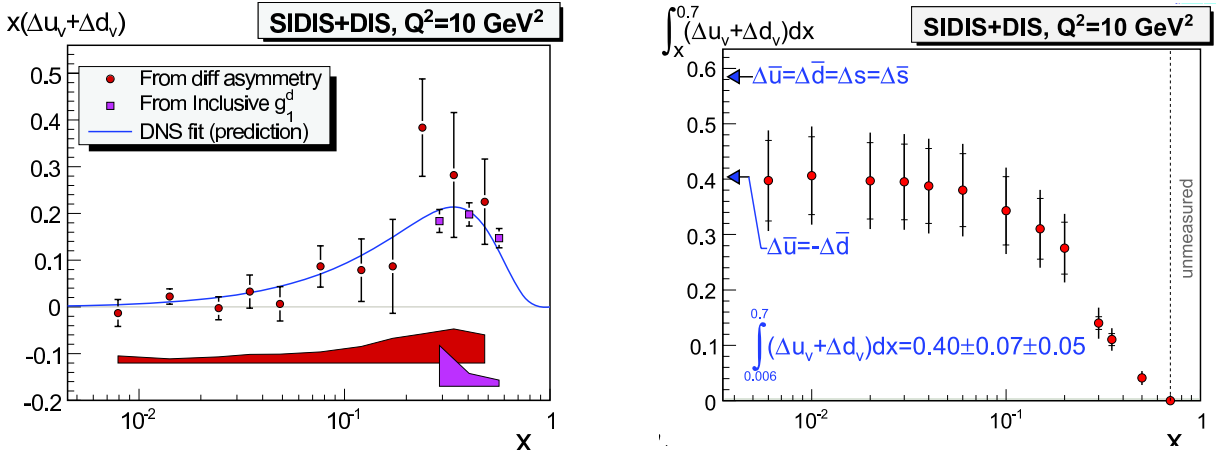


Figure 4: Left: Polarized valence quark distribution obtained from difference asymmetry. The line shows DNS fit in which presented data wasn't used. The additional points at high x come from g_1^d . Right: Integral of $\Delta u_v + \Delta d_v$ as a function of low x limit of integration.

the analysis of data collected in 2007 on the proton target has been finished, allowing to separate helicity distributions for different quark flavours. In order to do that asymmetries from identified hadrons have been measured. The identification was provided by the RICH detector. In total COMPASS has collected 43 (8) million of events with identified charged pions (kaons) on deuteron target and 25 (6.5) million of events with identified charged pions (kaons) on proton target. The same kinematical cuts were used for the proton and the deuteron data, mainly $Q^2 > 1 \text{ GeV}^2$ to select DIS events. Hadrons were required to have momenta in the range $10 < p < 50 \text{ GeV}/c$ where the RICH identification is most accurate.

Figure 1 presents the asymmetries as a function of x together with the HERMES results [18]. One can observe a good agreement between two experiments. It is also worth to notice that the proton asymmetries are significantly larger than deuteron ones. In LO QCD approximation we can use Eq. 8 to access quark helicities. Writing explicitly

$$\begin{aligned}
 A_{1,d} &= \frac{5(\Delta u + \Delta d) + 5(\Delta \bar{u} + \Delta \bar{d}) + 4\Delta s}{5(u + d) + 5(\bar{u} + \bar{d}) + 2(s + \bar{s})} \\
 A_{1,d}^h &= \frac{(4D_u^h + D_d^h)(\Delta u + \Delta d) + (4D_{\bar{u}}^h + D_{\bar{d}}^h)(\Delta \bar{u} + \Delta \bar{d}) + 2(D_s^h + D_{\bar{s}}^h)\Delta s}{(4D_u^h + D_d^h)(u + d) + (4D_{\bar{u}}^h + D_{\bar{d}}^h)(\bar{u} + \bar{d}) + 2(D_s^h s + D_{\bar{s}}^h \bar{s})} \\
 A_{1,p} &= \frac{4(\Delta u + \Delta \bar{u}) + (\Delta d + \Delta \bar{d}) + 2\Delta s}{4(u + \bar{u}) + (d + \bar{d}) + (s + \bar{s})} \\
 A_{1,p}^h &= \frac{4(D_u^h \Delta u + D_{\bar{u}}^h \Delta \bar{u}) + (D_d^h \Delta d + D_{\bar{d}}^h \Delta \bar{d}) + (D_s^h + D_{\bar{s}}^h)\Delta s}{4(D_u^h u + D_{\bar{u}}^h \bar{u}) + (D_d^h d + D_{\bar{d}}^h \bar{d}) + (D_s^h s + D_{\bar{s}}^h \bar{s})}
 \end{aligned} \tag{12}$$

one can see that quark distributions are multiplied by different weights for various types of asymmetries. These weights are made from the combinations of fragmentation functions D_q^h .

We assumed that polarized strange sea is symmetric ($\Delta s = \Delta \bar{s}$), so in total we have only 5 parameters ($\Delta u, \Delta d, \Delta \bar{u}, \Delta \bar{d}, \Delta s$). Thus with 10 different equations of asymmetries ($A_{1,d}, A_{1,d}^{\pi^+}, A_{1,d}^{\pi^-}, A_{1,d}^{K^+}, A_{1,d}^{K^-}, A_{1,p}, A_{1,p}^{\pi^+}, A_{1,p}^{\pi^-}, A_{1,p}^{K^+}, A_{1,p}^{K^-}$) it is possible to determine a quark helicities using standard least-square method. The final results as a function of x are

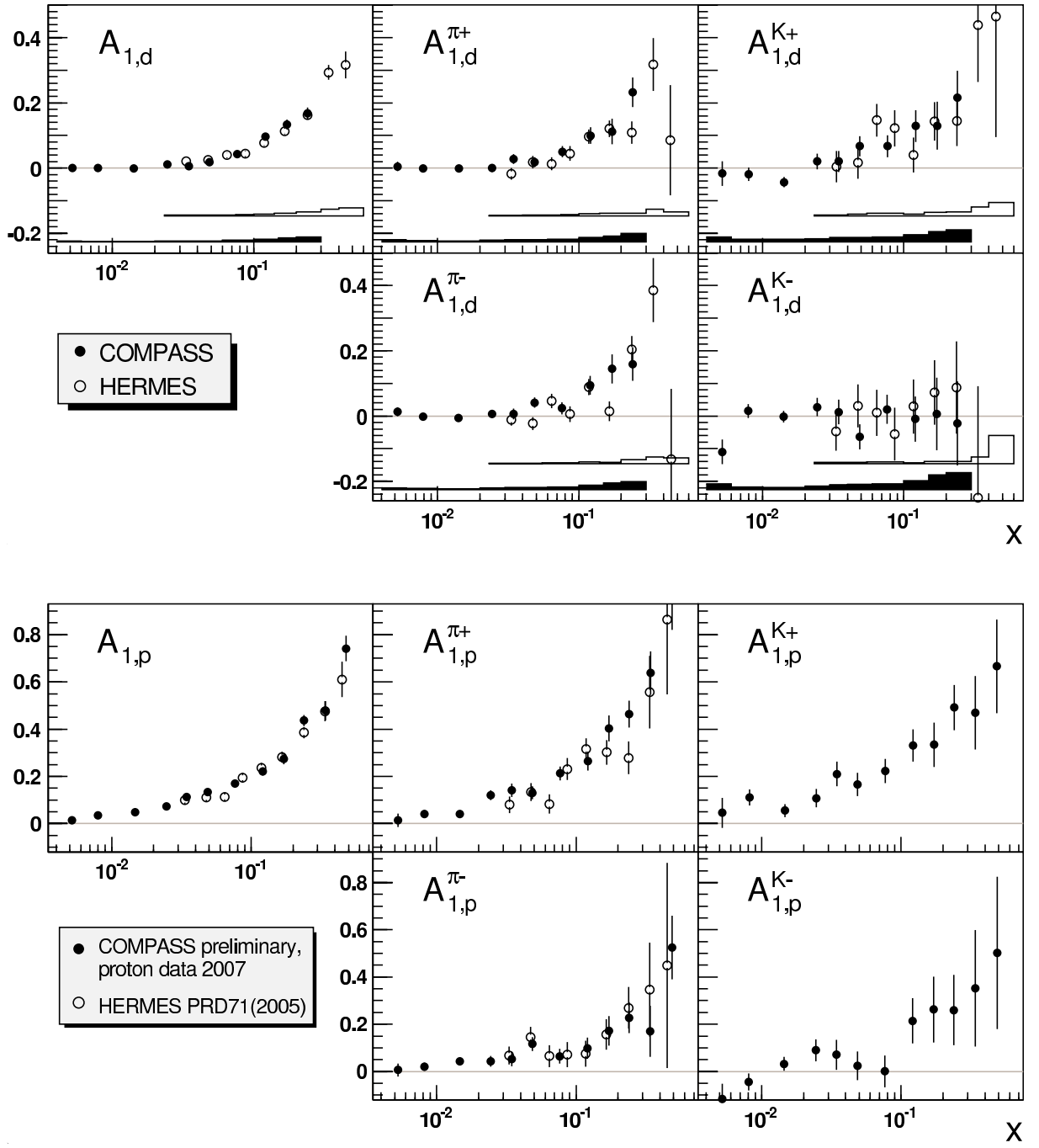


Figure 5: Comparison of COMPASS inclusive and hadron asymmetries of deuteron (top) and proton (bottom) data with HERMES results [18].

shown in Fig. 6.

The parametrization of quark helicity distributions estimated in LO by DNS [19] is also shown. It fits very well to the COMPASS data for light quarks while the disagreement for strange quarks is visible. As an unpolarized parton distributions MRST04 LO QCD [16] was used. For fragmentation functions DSS parametrizations at LO QCD was chosen [20], which are the most recent published ones. The first moments truncated to measured x -range are:

$$\begin{aligned}\Delta u &= 0.45 \pm 0.02(\text{stat.}) \pm 0.03(\text{syst.}) \\ \Delta d &= -0.25 \pm 0.03(\text{stat.}) \pm 0.02(\text{syst.}) \\ \Delta \bar{u} &= 0.01 \pm 0.02(\text{stat.}) \pm 0.004(\text{syst.}) \\ \Delta \bar{d} &= -0.04 \pm 0.03(\text{stat.}) \pm 0.01(\text{syst.}) \\ \Delta s &= \Delta \bar{s} = -0.01 \pm 0.01(\text{stat.}) \pm 0.02(\text{syst.}).\end{aligned}$$

6 Conclusions

A review of the COMPASS results on helicity quark distributions has been given. COMPASS analyzed all available longitudinal data from the years 2002-2007. The 2002-2006 data were collected on a deuteron target while in 2007 the proton target was used. The g_1 structure function has been measured and provided information about $\Delta\Sigma$ and polarization of strange quarks. The valence helicity distribution has been precisely determined using the so-called difference asymmetry. Flavour separation was possible using full set of semi-inclusive data. Last result depends on the chosen parametrization of the fragmentation functions.

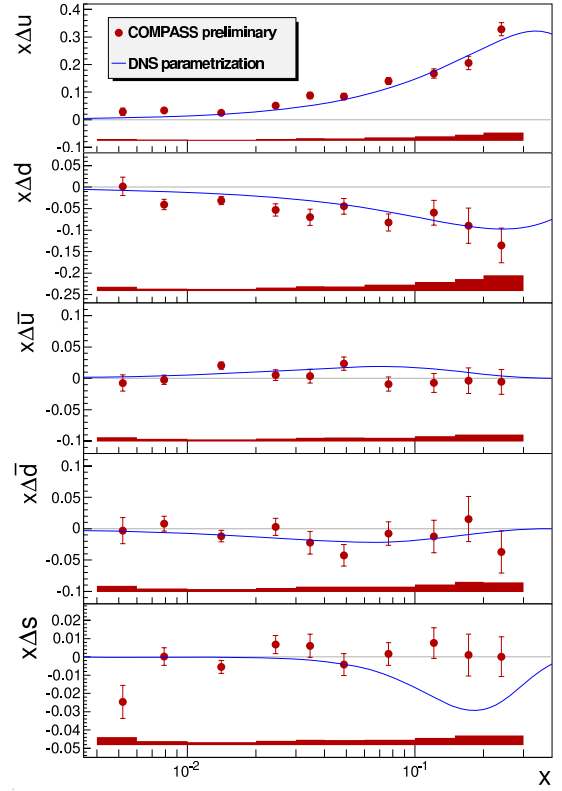


Figure 6: The quark helicity distributions evaluated at $Q^2 = 3 \text{ (GeV/c)}^2$ as a function of x . For strange quarks $\Delta s = \Delta \bar{s}$ was assumed. Bands at the bottom of each plot represent systematic errors. Solid line is a LO DNS parametrization.

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NEW MONTE-CARLO GENERATOR OF POLARIZED DRELL-YAN PROCESSES.

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Abstract

The new Monte-Carlo generator of unpolarized and single-polarized Drell-Yan events is presented. Its performance and physical applications are discussed.

The MC generator of polarized Drell-Yan events is necessary for *first*, estimation of the single-spin asymmetry (SSA) feasibility on the preliminary (theoretical) stage (without details of experimental setup). *Second*, as an input for detector simulation software (for example, GEANT [1] based code) on both planning of experimental setup and the data analysis stages. Until recently there was no in the free access any generator of Drell-Yan events except for the only PYTHIA generator [2]. However, regretfully, in PYTHIA there are only unpolarized Drell-Yan processes and, besides, they are implemented in PYTHIA without correct q_T and $\cos 2\phi$ dependence, which is absolutely necessary to study Boer-Mulders effect. Recently did appear the first generator of polarized DY events [3], [4] where q_T and angle dependencies are properly taken into account. However, this generator have some strong disadvantages (see below). That is why we wrote the new generator of polarized DY events. The scheme of generator is quite simple and very similar to the event generator GMC_TRANS [5] which was successfully used by HERMES collaboration for simulation of the Sivers effect in semi-inclusive DIS processes [6]. Briefly, the scheme of DY event generation look as follows. First, the generator performs the choice of flavor q of annihilating $q\bar{q}$ pair and choice does the given hadron (for example, polarized) contains annihilating quark or, alternatively, antiquark of given (chosen) flavor. It is done in accordance with the total unpolarized DY cross-sections for each flavor and each alternative choice for given hadron in initial state (annihilating quark or antiquark inside). Then, the variables x_F and Q^2 are selected according to the part of unpolarized cross-section (see, for example, Eq. (1) in Ref. [7]) which does not contain the angle dependencies. At the next step the polar angle θ is selected from $\sin\theta(1 + \cos^2\theta)$ distribution in that cross-section. Then, the Gaussian model for $f_{1q}(x, k_T)$ is applied and after that the transverse momentum of lepton pair q_T is selected from exponential distribution $\exp(-q_T^2/2)/2\pi$. At the next step ϕ and ϕ_S angles are selected in accordance with $\cos 2\phi$, $\sin(\phi - \phi_S)$ and $\sin(\phi + \phi_S)$ dependencies of the single-polarized DY cross-section (see, for example, Eq. (2) in Ref. [7]). The k_T dependencies of Boer-Mulders $h_{1q}^\perp(x, k_T)$ and Sivers $f_{1T}^q(x, k_T)$ PDFs are fixed by the Boer model [8] and Gaussian ansatz [9], [10], respectively. At this stage of ϕ and ϕ_S selection x_F , Q^2 , θ and q_T variables are already fixed¹ that essentially increases the rate of ϕ and ϕ_S selection. All variables are generated using the standard von Neumann acceptance-rejection technique (see, for example [2]).

¹Certainly, one can select all variables simultaneously. However, such scheme essentially decreases the rate of events generation.

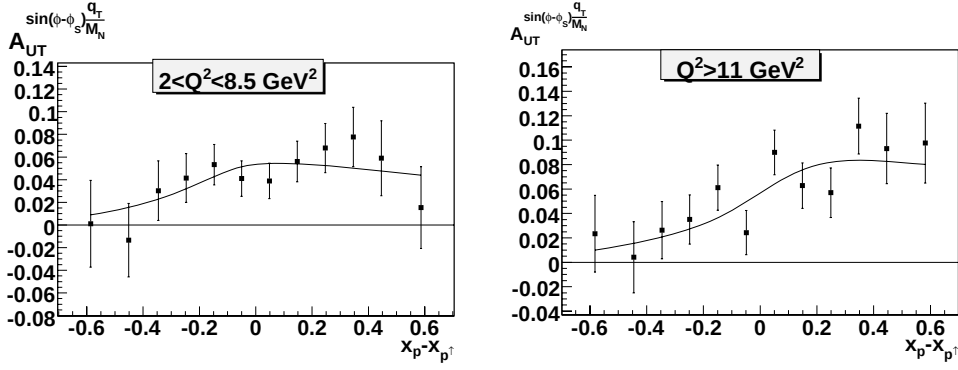


Figure 1: Estimation of asymmetries $A_{UT}^{\sin(\phi-\phi_S)\frac{q_T}{M_N}} \Big|_{pp\uparrow}$ for NICA, $s = 400\text{GeV}^2$. Here fit from Ref. [10] is used. The points with error bars are obtained by using simulations with event generator at the applied statistics 100K pure Drell-Yan events. $\langle Q^2 \rangle \simeq 3.5\text{GeV}^2$ and $\langle Q^2 \rangle = 15\text{GeV}^2$ for the left and right plots, respectively.

Let us stress once again that the generator elaborated in [3], [4] is the first generator of polarized DY events and in many respects it helped us to construct our generator. Now we briefly consider the advantages of the new generator.

As it was discussed above, one of the main requirement to any generator of polarized DY events is to properly include the nontrivial q_T dependence of DY cross-sections. This is of especial importance for the q_T weighted objects we deal with in the paper. In the earlier generator [3], [4] the approximations are applied for the convolution calculations in the $h_{1q}^\perp$ containing parts of DY cross-section just as it was done in the original paper [8] (see the discussion around Eqs. (47) and (53) in Ref. [8]). However, the direct calculations show that implementation of these approximation in the generator leads to essential distortion of the q_T weighted objects, such as SSA $A_{UT}^{\sin(\phi+\phi_S)\frac{q_T}{M_N}}$: the values of SSA obtained from simulated data significantly differ from the respective values calculated directly from the input parametrizations/models for f_{1q} , $h_{1q}^{\perp(1)}$ and h_{1q} . That is why we avoid the such kind of approximations. Instead, we calculate the convolutions numerically for the large discrete set of q_T values and then we perform the standard spline interpolation procedure to reconstruct the calculated convolution as the continuous function of q_T . As a result, the values of SSA $A_{UT}^{\sin(\phi+\phi_S)\frac{q_T}{M_N}}$ reconstructed from simulated data are in a good agreement with the respective values calculated directly from the parametrizations/models entering the generator as an input.

Another advantage of the new generator is that it can be much more easy combined with the PYTHIA generator (where almost all possible processes are included) which is necessary for the analysis of background processes that can produce false DY events (misidentification of lepton pair). The point is that, on the contrary to generator [3], [4], in our generator, just as in PYTHIA, the DY processes are generated for each flavor in separation. Besides, the constructed generator have some technical advantages: due to the chosen generation scheme the rate of event generation is much higher than for generator [3], [4] (about 1000 events per second on standard PC for our generator and about 10 for generator developed in Refs. [3], [4]), where all kinematical variables are thrown simultaneously; the search for cross-sections maximums for the von Neumann algorithm is performed automatically as well as the calculation of the total cross-section

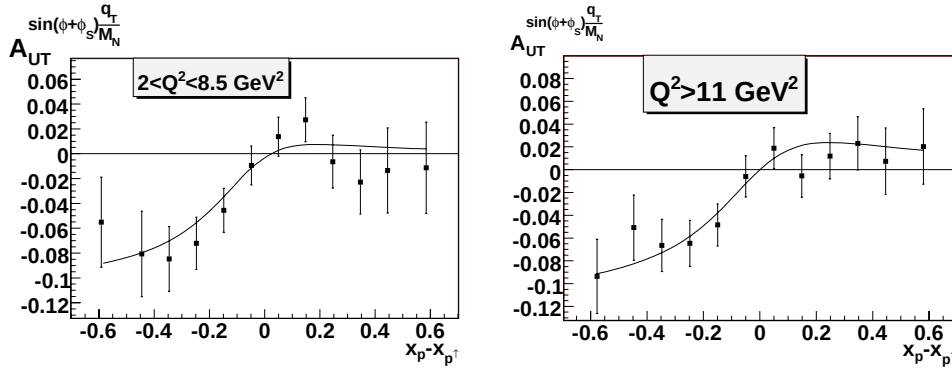


Figure 2: Estimation of asymmetries $A_{UT}^{\sin(\phi+\phi_S)\frac{q_T}{M_N}}|_{pp^\dagger}$ for NICA, $s = 400\text{GeV}^2$. Here the evolution model with the input ansatz $h_{1q,\bar{q}} = \Delta q, \bar{q}$ at $Q_0^2 = 0.23\text{GeV}^2$ is used. GRV94 [12] parametrization for $q(x)$ and GRSV95 [13] parametrization for $\Delta q(x)$ are used. The points with error bars are obtained with the developed generator of polarized DY events at the applied statistics 100K (top) and 50K (bottom) pure Drell-Yan events. $\langle Q^2 \rangle \simeq 3.5\text{GeV}^2$ and $\langle Q^2 \rangle = 15\text{GeV}^2$ for the left and right plots, respectively.

at the end of each generation run.

To demonstrate the output of generator we prepared the sample with applied statistics 100K of pure Drell-Yan events for each of two Q^2 ranges: $2 < Q^2 < 8.5\text{GeV}^2$ and $Q^2 > 11\text{GeV}^2$. Cut $2 < Q^2 < 8.5\text{GeV}^2$ is applied to avoid misidentification of lepton pairs due to numerous background processes below $Q^2 = 2\text{GeV}^2$ and to exclude lepton pairs coming from J/ψ region. Cut $Q^2 > 11\text{GeV}^2$ is also applied to avoid the lepton pairs coming from J/ψ region. We choose NICA experimental conditions (proton-proton collider with $s = 400\text{GeV}^2$).

The results are presented in Fig. 1 and 2. From these figures one can see that the reconstructed from the simulated data asymmetry is in very good agreement with the respective values calculated directly from the parametrization (solid line) for both types of single-spin asymmetries. We performed also a large number of another statistical tests which proved the generator to be working perfectly.

In summary, new generator of polarized DY events was developed which allow to estimate the feasibility of both weighted with $\sin(\phi - \phi_S)$ and $\sin(\phi + \phi_S)$ single-spin asymmetries is constructed. The performance of the generator essentially increased in comparison with the preceding generator [3]. The generator was successfully tested and now it applied in a lot of physical applications. For example, the estimations performed for proton-proton collisions demonstrate that the single-spin asymmetries are presumably measurable even at the applied statistics 50K pure Drell-Yan events. While $A_{UT}^{\sin(\phi-\phi_S)\frac{q_T}{M_N}}$ is presumably measurable in both kinematical regions $x_p > x_{p^\dagger}$ and $x_p < x_{p^\dagger}$, the asymmetry $A_{UT}^{\sin(\phi+\phi_S)\frac{q_T}{M_N}}$ could be measured only in the region $x_p < x_{p^\dagger}$, where it takes quite considerable values. It is of importance that while for the symmetric collider mode (RHIC, NICA) there is no any difficulties to reach both $x_p > x_{p^\dagger}$ and $x_p < x_{p^\dagger}$ kinematical regions, for the fixed target mode (COMPASS, J-PARC), there is the problem to reach the region $x_p < x_{p^\dagger}$. This is hardly possible for the running COMPASS experiment, where only unpolarized proton beam and forward geometry spectrometer is available. At the same time it still could be achieved by J-PARC, where at present the option with polarized

proton beam is planned.

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HIGH ENERGY SPIN PHYSICS WITH THE PHENIX DETECTOR AT RHIC

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Abstract

The Relativistic Heavy Ion Collider (RHIC) at Brookhaven National Laboratory has demonstrated the unique ability to collide beams of polarized protons at center of mass energies from $\sqrt{s}=62.4$ to 500 GeV. Such collisions have been analyzed by members of the PHENIX collaboration to shed light on the size and shape of the spin-dependent distribution function of the gluon $\Delta g(x)$, and on the origin of large asymmetries in inclusive hadron production from collisions involving transversely polarized protons. After a brief review of RHIC and the PHENIX detector, our recent results in spin physics will be presented, along with our prospects for determining quark flavor-separated spin-dependent distribution functions extracted using parity violation in the production and decay of W bosons at RHIC.

1 Introduction

From polarized deep-inelastic scattering experiments at SLAC, CERN, and DESY, we have learned that quark and antiquarks carry perhaps 25% of the spin of the proton [1]. This fraction was unexpectedly small, and understanding the remainder, which must come from the intrinsic spin of the gluons $\Delta g(x)$, and orbital angular momentum of the quarks L_q , and gluons L_g , presents an outstanding challenge to theorists and experimentalists.

The Relativistic Heavy Ion Collider (RHIC) at Brookhaven National Laboratory has the unique ability to collide beams of highly polarized protons. This has provided experimentalists with a new tool to measure, with leading order sensitivity, the gluon contribution to the spin of the proton [2]. In addition, RHIC has reached pp center of mass energies of 500 GeV, where measurements of Δu , $\Delta \bar{u}$, Δd , and $\Delta \bar{d}$ are possible by exploiting parity-violation in the process $pp \rightarrow W^\pm + X$ [2].

2 RHIC : The Polarized Proton Collider

RHIC has two rings, each of which can contain 120 bunches of polarized protons. In the interaction regions, a bunch from one ring collides with a bunch from the other ring, then these same bunches circulate in opposite directions 3833 m around the RHIC rings and collide again about 13 μ s later [3].

The spin is preserved during acceleration and store through the use of a pair of Siberian snakes. The spins can be rotated from transverse to longitudinal at the interaction region by spin rotators. In the most recent run in 2009, RHIC was run at $\sqrt{s}=500$ GeV for 30

Run	Year	$\sqrt{s}$ (GeV)	Polarization	Recorded Luminosity
1,2	2001,2002	200	14%, Transverse	0.15 pb ⁻¹
3	2003	200	34%, Longitudinal	0.35 pb ⁻¹
4	2004	200	46%, Longitudinal	0.12 pb ⁻¹
5	2005	200	47%, Longitudinal 47%, Transverse	3.4 pb ⁻¹ 0.4 pb ⁻¹
6	2006	200 62.4	55%, Longitudinal 50%, Longitudinal	7.5 pb ⁻¹ 0.08pb ⁻¹
8	2008	200	44%, Transverse	5.2 pb ⁻¹
9	2009	500 200	34%, Longitudinal 56%, Longitudinal	11 pb ⁻¹ 16 pb ⁻¹

Table 1: Summary of polarized proton runs at RHIC and integrated luminosity recorded by PHENIX.

days, and demonstrated 35% polarization and a peak luminosity of $8.5 \times 10^{31} \text{ cm}^{-2}\text{s}^{-1}$, and store average of $5.5 \times 10^{31} \text{ cm}^{-2}\text{s}^{-1}$. RHIC also ran at 200 GeV, at roughly 55% polarization. The polarization of the beams was measured every few hours using pC CNI polarimeters, which are calibrated using a polarized hydrogen gas jet target to better than 5% accuracy. A summary of the polarized proton runs at RHIC is given in Table 1.

3 PHENIX Detectors

The PHENIX detector (see Fig. 1) surrounds the interaction region at 8 o'clock in the RHIC ring, and is composed of two forward detectors and four spectrometer arms [4]. The forward detectors (BBC and ZDC) are used to determine the collision vertex position and time, proton beam luminosity and polarization, and can form a minimum bias trigger. The four spectrometers have specialized functions. The east and west central arms have tracking, particle identification, and an electromagnetic calorimeter at central rapidity. The muon spectrometer arms are specialized for the detection of muons.

The north and south muon arms cover a pseudorapidity interval $1.2 < |\eta| < 2.4(2.2)$ respectively, with full azimuthal coverage. They employ a hadron absorber, radial magnetic field, and cathode strip tracking chambers. Trigger electronics identify muons using layers of proportional tubes interspersed with hadron absorber downstream of the magnet arms. Of particular interest to the spin program is the ability to trigger on muons from J/ψ decay, and high energy muons from W decay. The latter channel will allow the flavor separated measurement of the u and d polarized quark and antiquark distributions when a trigger upgrade is complete.

The east and west central arms are in a solenoidal field, and can identify electrons, photons, and charged hadrons in a pseudorapidity range $-0.35 < \eta < 0.35$ over $|\Delta\phi| = 90^\circ$ in azimuth. Tracking is provided by drift chambers (DC), time expansion chamber (TEC), and pad chambers (PC), while particle identification is furnished by a time-of-flight system, (TOF) and ring-imaging Cerenkov (RICH). Of particular importance is the electro-

magnetic calorimeter (EMCal). Individual calorimeter towers are made of lead scintillator (PbSc) or lead glass (PbGl) and subtend $\Delta\eta \times \Delta\phi \approx 0.01 \times 0.01$. This fine segmentation allows the two photons from π^0 decay to be resolved up to p_T of 12 GeV/c. Shower shape analysis extends this range beyond 20 GeV/c, and allows the two photon invariant mass peak from π^0 decay to be used for energy calibration. Fine segmentation is also important for distinguishing rare direct photon events from the background from photons originating from neutral meson decay. This allows the measurement of the theoretically clean direct-photon asymmetry [2].

In addition to having well-understood detectors, PHENIX has selective triggers for final states of interest to the spin program, including high p_T pions and photons. Such events can be recorded at rates above 5 kHz with the PHENIX data acquisition system. Data analysis is done in parallel at the RHIC Computing Facility and the CCJ facility at RIKEN in Wako, Japan.

4 Recent Physics Results

PHENIX acquires sensitivity to Δg by measuring double helicity asymmetries in the production of particular final states, such as π^0 . The asymmetries are defined :

$$A_{LL}^{\pi^0} = \frac{d\sigma_{++}/dp_T - d\sigma_{+-}/dp_T}{d\sigma_{++}/dp_T + d\sigma_{+-}/dp_T}$$

where $\sigma_{++}(\sigma_{+-})$ is the production cross section for π^0 from pp collisions with like (unlike) helicities. In leading order, this can be factorized as the sum of all partonic subprocesses $ab \rightarrow cX$ where parton c fragments into the detected π^0 . In this framework, A_{LL} can be understood as the convolution :

$$A_{LL}^{\pi^0} = \frac{\sum_{abc} \Delta f_a \Delta f_b \hat{\sigma}(ab \rightarrow cX) \hat{a}_{LL}(ab \rightarrow cX) D_c^{\pi^0}}{\sum_{abc} f_a f_b \hat{\sigma}(ab \rightarrow cX) D_c^{\pi^0}},$$

where $f_a(\Delta f_a)$ are the unpolarized (polarized) parton distribution functions, and $D_c^{\pi^0}$ is the fragmentation function of c into π . The spin-averaged partonic scattering cross-section for $ab \rightarrow cX$ is denoted by $\hat{\sigma}$, and $\hat{a}_{LL}$ is the analyzing power; both of which are calculable at next-to-leading order [2].

The $A_{LL}(pp \rightarrow \pi^0 X)$ analysis have been PHENIX' most sensitive probes of Δg since we can trigger on and reconstruct $\pi^0 \rightarrow 2\gamma$ with high efficiency. Also, at $\sqrt{s} = 200$ GeV, the cross-section is dominated by qg scattering for $p_{T\pi} > 5$ GeV, so the process is directly sensitive to Δg .

In the analysis, the invariant mass spectrum is formed from all pairs of photons, and the asymmetry is formed around the π^0 mass peak in a window from 112-162 MeV. The

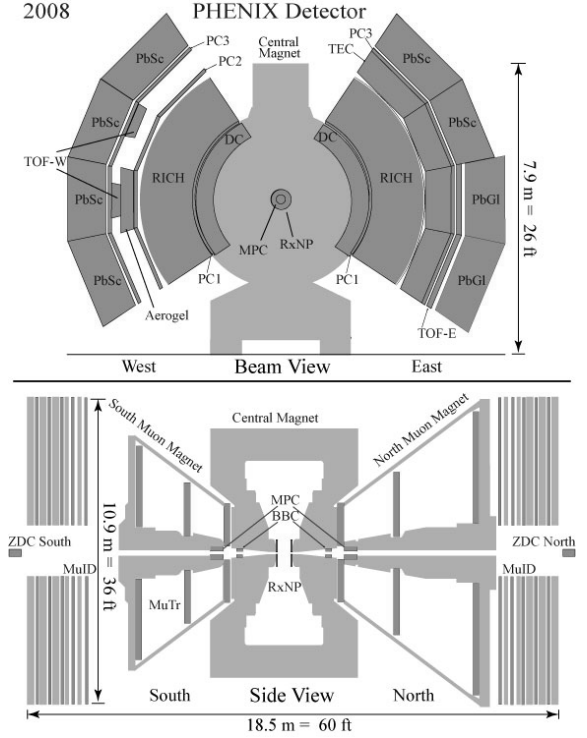


Figure 1: Beam and side views of the PHENIX detector

background fraction, and asymmetry in the background, are estimated from sidebands around the mass peak. Background fractions vary from 16% in the π^0 p_T range of 2-3 GeV, to 6% in the 9-12 GeV range at $\sqrt{s} = 200$ GeV.

In 2006, PHENIX recorded 0.08 pb $^{-1}$ of pp collisions at $\sqrt{s} = 62.4$ GeV. This run was intended primarily to improve upon ISR results on inclusive pion cross-sections. At these low center of mass energies, next-to-leading order (NLO) perturbative QCD (pQCD) predictions of inclusive pion cross-sections tend to under-predict the observed yields. Including threshold resummation at next-to-leading-logarithm (NLL) accuracy improves the agreement with fixed target [5] and collider data [6]. The PHENIX measurement of the inclusive π^0 cross-section confirmed the necessity of such corrections [7]. In addition, the measurement of $A_{LL}(pp \rightarrow \pi^0 X)$ at $\sqrt{s} = 62.4$ GeV [7], made with one week

of data, excluded the GRSV scenarios $\Delta g(x) = \pm g(x)$ [8] (see Fig. 2). The NLL predictions for $A_{LL}^{\pi^0}$ are smaller than for NLO pQCD. The asymmetry measurement also extended our sensitivity to $\Delta g(x)$ to higher x in comparison to the measurements at $\sqrt{s} = 200$ GeV.

From the 2005 and 2006 runs at $\sqrt{s} = 200$ GeV, PHENIX published results on A_{LL} of π^0 [9] (see Fig. 3). These asymmetries are consistent with zero, and clearly lie below the GRSV standard model (which reflected the best fit to polarized deep inelastic scattering (DIS) data), thus favoring a smaller $\Delta g(x)$, and smaller integral Δg . An estimate of the integral of $\Delta g(x)$ can be made from these data in the framework of the GRSV model. We use the notation $\Delta g^{[a,b]}$ to denote the integral of $\Delta g(x)$ from $a < x < b$. The polarized DIS data were refit using a constraint on the value of $\Delta g^{[0,1]}$, at a scale $\mu^2 = 0.4$ GeV ^{2} , to extract a set of polarized parton distribution functions (PDFs) using the GRSV parametrization. From the resulting polarized PDFs, a prediction for $A_{LL}^{\pi^0}$ is made and compared with the measurement. The resulting χ^2 is plotted as a function of $\Delta g_{GRSV}^{[0.02,0.3]}$, evolved to $\mu^2 = 4$ GeV ^{2} , in Fig. 3. The curve is not parabolic since π^0 can be produced from gg , qg , and qq' scattering. The former has quadratic sensitivity to Δg , leading to two minima in the χ^2 curve. The qg contribution is linear in Δg , breaking the symmetry. Including the statistical uncertainty, the results suggest $\Delta g_{GRSV}^{[0.02,0.3]} = 0.2 \pm 0.1(1\sigma)$ and $0.2_{-0.8}^{+0.2}$ (3σ) with an additional systematic uncertainty of ± 0.1 from the experimental uncertainty in beam polarization and relative luminosity. Using other parametrizations yields similar results, suggesting the asymmetry data are sensitive primarily to the integral $\Delta g^{[0.02,0.3]}$ and not to the specific shape of $\Delta g(x)$. A similar analysis showed that changes of a factor of 2 in the renormalization, factorization, and fragmentation scales used in the NLO pQCD prediction, lead to changes in the value of $\Delta g_{GRSV}^{[0.02,0.3]}$ extracted of $\pm 0.1_{-0.4}^{+0.1}$ at $\Delta\chi^2 = 1(9)$.

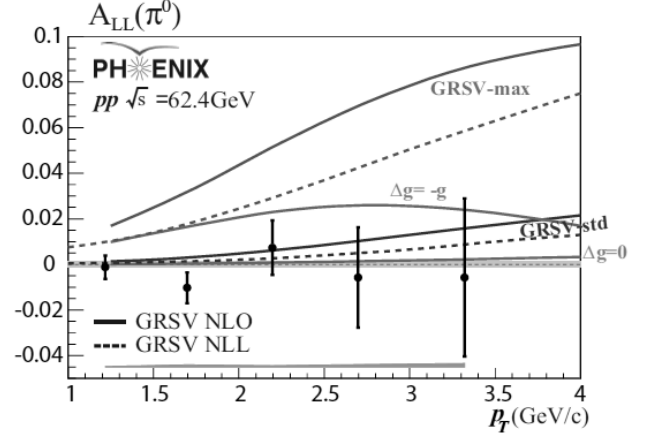


Figure 2: A_{LL} for π^0 at $\sqrt{s} = 62.4$ GeV as a function of p_T , showing statistical uncertainties. (14% uncertainty in polarization not shown.) Four GRSV NLO pQCD (solid curves) and NLL pQCD (dashed curves) are shown for comparison. See text for details.

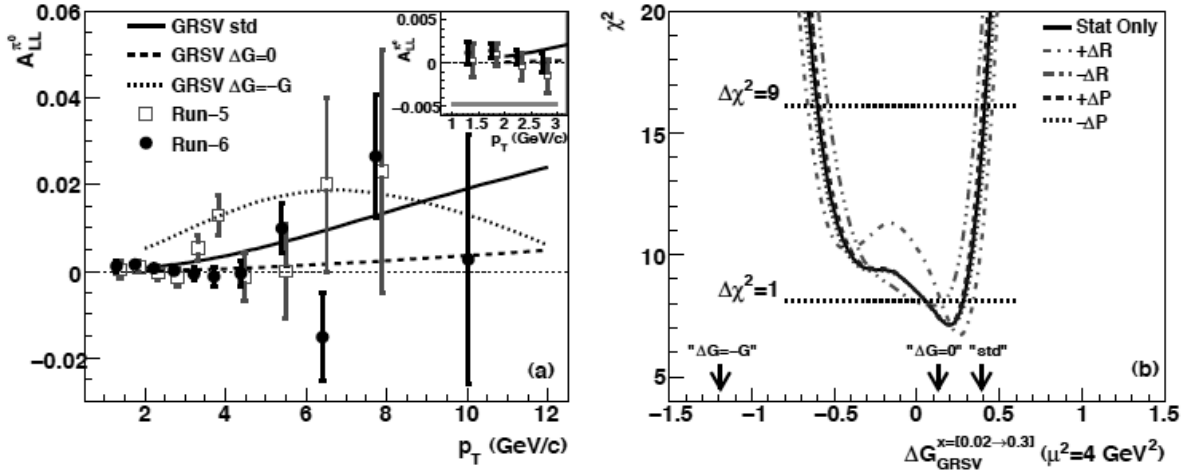


Figure 3: Left : A_{LL} for π^0 at $\sqrt{s} = 200$ GeV as a function of p_T , showing statistical uncertainties (8.3% uncertainty in beam polarization not shown.) Right : χ^2 profile as function of $\Delta g_{GRSV}^{[0.02,0.3]}$ using 2005 and 2006 $A_{LL}^{\pi^0}$ data.

The above results are consistent with a comprehensive global analysis of almost all polarized DIS and RHIC data [1, 10]. The RHIC data now pose the tightest constraint on $\Delta g(x)$ in the range $0.05 < x < 0.2$. The global analysis suggests a truncated first moment $\Delta g^{[0.001,1]} = 0.013^{+0.702}_{-0.314}$ at $Q^2 = 10$ GeV², and full first moment $\Delta g^{[0,1]} = -0.084$. These results are very interesting. First, they suggest that large, anomaly-inspired values for Δg are unlikely. Second, the difference between the truncated and full first moment highlights the importance of extending the measurements at RHIC to lower x . This may be possible in the near future as PHENIX' barrel and forward silicon vertex detectors are installed, and a proposed forward calorimeter (covering $1 < \eta < 3$) is developed. These will increase the acceptance of the PHENIX detector substantially, and will allow new, correlated asymmetry measurements of γ -jet, hadron-jet, and jet-jet. These new observables will allow us to constrain the kinematics of the hard-scattered partons, and improve our sensitivity to $\Delta g(x)$ at lower momentum fraction x , perhaps to a few $\times 10^{-3}$.

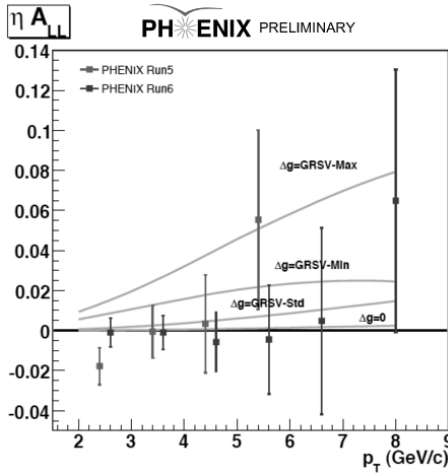
The PHENIX collaboration also has preliminary asymmetry results on several other inclusive channels. The A_{LL}^{η} measurement has a slightly different sensitivity to Δg than π^0 since the flavor content of the η is different. The asymmetry measured is consistent with 0, and is shown in Fig. 4 along with GRSV model predictions [8]. PHENIX measurement of A_{LL} of direct photons is shown also in Fig. 4. This asymmetry too is consistent with 0, and is heavily statistics limited. The virtues of this measurement are that it is dominated by qg scattering, and so is linear in Δg , and that the photon carries information about the hard scattering kinematics and is not the result of a fragmentation process.

In Fig. 5 is shown the $A_{LL}(pp \rightarrow h^+ X)$ measured at $\sqrt{s} = 62.4$ GeV (where h^+ denotes a non-identified charged hadron (π, K, p)). These data were taken with the minimum bias trigger, with the momentum of the hadron measured using the drift chamber (DC). The PHENIX RICH has a threshold for $\pi^\pm$ of 4.7 GeV, so a RICH veto ensures the events are $h^\pm$ with $p_T < 4.7$ GeV. The GRSV model curves are shown on the plot. The GRSV max curve, corresponding to $\Delta g(x) = g(x)$ is not shown as it is clearly excluded by this data. Similar data exist for h^- .

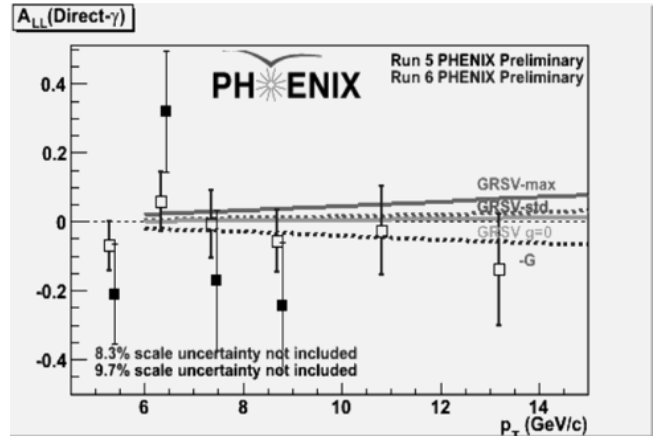
At $\sqrt{s} = 200$ GeV, $A_{LL}(pp \rightarrow \pi^\pm X)$ has been measured at $p_T > 5$ GeV. Here we trigger on the EMCal, and look for associated tracks, measuring the particle momentum in the DC. The hadrons can be identified as pions as now they are above the RICH threshold. The charged pions are not as well measured as π^0 since PHENIX lacks a dedicated charged hadron trigger. Still, the channel is important because it has high analyzing power for Δg . This is apparent from two observations. First, at $p_T > 5$ GeV, π production is dominated by qg scattering, giving leading order sensitivity to Δg . Second, we can write $A_{LL}^{\pi^+} \propto \Delta g(x_1) \otimes (\Delta u(x_2)D_u^{\pi^+} + \Delta \bar{d}(x_2)D_d^{\pi^+})$. Since Δu is large and positive, if Δg is positive, we expect $A_{LL}^{\pi^+} > 0$. Similarly, $A_{LL}^{\pi^-} \propto \Delta g(x_1) \otimes (\Delta d(x_2)D_d^{\pi^-} + \Delta \bar{u}(x_2)D_u^{\pi^-})$. Since Δd is negative, this yields smaller asymmetries than for π^+ . The $A_{LL}^{\pi^0}$ will lie in between. Thus we predict an ordering of the asymmetries $A_{LL}^{\pi^+} > A_{LL}^{\pi^0} > A_{LL}^{\pi^-}$ if $\Delta g > 0$. We also observe that $A_{LL}^{\pi^+}$ has the largest analyzing power for Δg of these channels.

Finally we note with great excitement that the W programs at PHENIX and STAR have their first data from the RHIC run in 2009 at $\sqrt{s} = 500$ GeV. PHENIX accumulated roughly 11 pb^{-1} of data during this 4 week run. The goals of the program (see [2,11] for details), are to extract the flavor-separated spin-dependent quark distribution functions, in particular $\Delta \bar{u}$ and $\Delta \bar{d}$. These have been measured in semi-inclusive DIS experiments (SMC, HERMES, and COMPASS) but with limited accuracy. The virtue of the RHIC program is that these quantities will be measured at very high momentum scales, without uncertainties due to fragmentation functions.

With the recorded luminosity, PHENIX expects $\approx 200 e^+$ from $pp \rightarrow W^+ \rightarrow e^+ \nu_e$ with $p_{Te} > 25$ GeV in the central arms and $\approx 35 e^-$ with $p_{Te} > 25$ GeV from $pp \rightarrow W^- \rightarrow e^- \bar{\nu}_e$. In the analysis, we look for > 25 GeV in the calorimeter with an isolated, high-momentum charged track passing a time-of-flight cut. The resolution of the DC is sufficient to distinguish e^+ from e^- even at $p_T > 40$ GeV, allowing us to identify the parent as W^+ or W^- . The main background is from energetic charged hadrons showering in the EMCal. Other backgrounds include $e^\pm$ from photon conversion from π^0 decay, $e^\pm$ from



(a)



(b)

Figure 4: (a) A_{LL} for η at $\sqrt{s} = 200$ GeV as a function of p_T , showing statistical uncertainties and GRSV model predictions (8.3% uncertainty in beam polarization not shown). (b) A_{LL} for direct photons at $\sqrt{s} = 200$ GeV as a function of p_T , showing statistical uncertainties and GRSV model predictions.

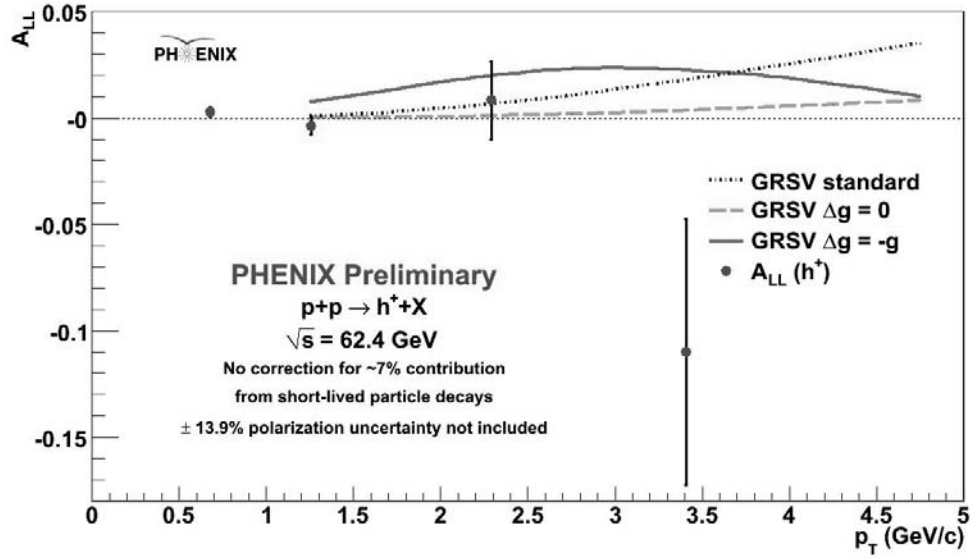


Figure 5: A_{LL} for h^+ versus p_T measured at $\sqrt{s} = 62.4$ GeV, showing statistical uncertainties, and GRSV model predictions.

heavy quark decay, cosmics, and accidentals. The process $pp \rightarrow Z \rightarrow e^+e^-$ constitutes roughly 6% background for W^+ , and 30% for W^- . We have clear evidence for a W signal above these backgrounds, and after all cuts, and considering detectors inefficiencies and dead regions, we expect an uncertainty on the single spin asymmetry $\delta A_L^{W^+} \approx 0.3$. With a few more runs, this uncertainty should reach a few percent.

5 Summary

Results from the RHIC spin experiments, PHENIX and STAR, now pose the tightest constraints on $\Delta g(x)$, which is smaller than anticipated from polarized DIS experiments. There is considerable uncertainty in $\Delta g(x)$ in the low- x ($x < 0.02$) region, which will be addressed by the experiments in the future. The process $pp \rightarrow W \rightarrow e\nu$ has been observed at PHENIX, and efforts are underway to estimate the cross-section and single-spin asymmetry A_L^W . These are very exciting and productive times for the spin program.

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MEASUREMENTS OF THE A_{yy} , A_{xx} , A_{xz} AND A_y ANALYZING POWERS IN THE $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ REACTION AT THE ENERGY $T_d=270$ MEV.

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Abstract

The experimental results on the analyzing power T_{20} in the $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reaction with excitation of levels of a nucleus ^{13}C at the energy $T_d = 140, 200$ and 270 MeV for at emission angle $\Theta_{cm}=0^\circ$ are presented in this work. The data on the tensor A_{yy} , A_{xx} , A_{xz} and vector A_y analyzing powers for the $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reaction at energy $T_d = 270$ MeV in the angular range from 4° to 18° in laboratory system are also obtained.

Introduction

The $(\vec{d}, p)$ stripping reaction has proved a useful tool for probing single particle aspects of nuclear structure. With the advent of beams of exotic nuclei there has been renewed interest in charged particle spectroscopy and $(\vec{d}, p)$ stripping in particular as a means of investigating the structure of neutron-rich nuclei via reactions in inverse kinematics [1].

Direct measurement of the capture reactions at energies of astrophysical interest is, in some cases, nearly impossible due to the low reaction yield, especially, if the capture involves exotic nuclei. Alternative indirect methods, such as the asymptotic normalization coefficient (ANC) method, based on the analysis of breakup or transfer reactions, have been used as a tool to obtain astrophysical S-factors. The advantage of indirect approaches comes from the fact that transfer and breakup reactions can be measured at

higher energies, where the cross sections are much larger[2].

Single-nucleon transfer reactions that probe the degrees of freedom of single particles have been extensively used to study the structure of stable nuclei. The analysis of such reactions provides the angular momentum transfer, which gives information on the spin (j) and parity (π) of the final state. The sensitivity of the cross sections to the single-nucleon components allows for the extraction of spectroscopic factors. The recent indications of reduced occupancies of single-particle states reveal that reliable measurements of spectroscopic factors in exotic nuclei are highly desirable [3].

The $(\vec{d}, p)$ stripping reaction has long been used as a means of probing the single particle structure of nuclei. In particular, through distorted wave Born approximation (DWBA) analysis it has been used to determine the orbital angular momentum and spectroscopic factors of specific states in the recoil nucleus [1].

In this report, the data on the tensor A_{yy} , A_{xx} , A_{xz} and vector A_y analyzing powers for the $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reaction at energy $T_d = 270$ MeV in the angular range from 4° to 18° in laboratory system obtained at RIKEN are presented. The experimental results on analyzing powers at energy $T_d = 140, 200$ and 270 MeV for $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reaction with excitation of levels of a nucleus ^{13}C reaction at emission angle $\Theta_{cm}=0^\circ$ are also presented.

Experiment

The experiment was performed at RIKEN Accelerator Research Facility (RARF). Details of the experiment can be found in ref.[4].

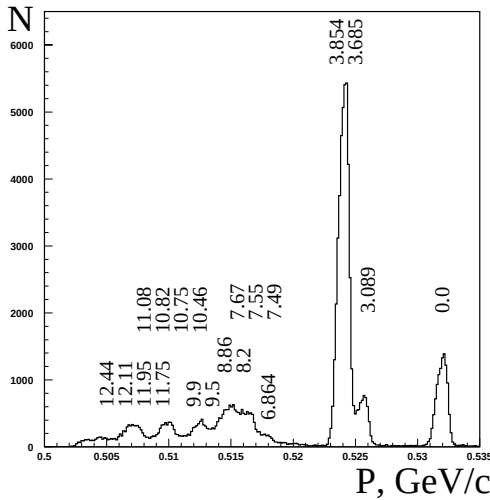


Figure 1: Typical momentum spectra for $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reaction and $\Theta_{cm}=0^\circ$ ($T_d=140$ MeV). Peaks corresponding to the ^{13}C states are labeled by their excitation energies in MeV.

with MWDC followed by the three plastic scintillators at the second focal plane. Criteria used for the identification of the scattered protons from the $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reaction are the following: particle must be registered in the all three scintillation detectors and it was selected by the correlation of the energy losses in the 1st and the 2nd and the 1st and the 3rd scintillation detectors; radio frequency signal of the cyclotron was used as a reference for the time-of-flight measurement.

In this experiment, four spin modes were used: the mode 0 - unpolarized mode, mode 1 - pure tensor mode, mode 2 - pure vector mode and mode 3 is mixed mode. The obtained polarization values were $\sim 75\%$ of the ideal values. The direction of symmetric axis of the beam polarization was controlled with a Wien filter located at the exit of polarized ion source. The polarized deuteron beam was accelerated up to 140, 200 and 270 MeV by the combination of the AVF cyclotron and Ring cyclotron. The beam polarizations were measured with D-room polarimeter located at D-room and Swinger polarimeter placed in the front of the target.

Scattered particles (^3H , ^3He or p) were momentum analyzed with quadrupole and dipole magnets (Q-Q-D-Q-D) and detected

Typical momentum spectrum for the $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reaction is shown in figure 1. Peaks corresponding to the ^{13}C states are labeled by their excitation energies in MeV.

Results and discussion

The results on the tensor analyzing power T_{20} for the $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reaction are presented in figure 2. The circles, squares and triangles represent the data on T_{20} for $^{12}\text{C}(\vec{d}, p)^{13}\text{C}(g.s.)$, $^{12}\text{C}(\vec{d}, p)^{13}\text{C}(3.089\text{MeV})$ and $^{12}\text{C}(\vec{d}, p)^{13}\text{C}(3.6845 + 3.854\text{MeV})$ reactions, respectively. The data on T_{20} for this reaction at the energy $T_d = 270$ MeV agree each other. The difference in the analyzing power for the $^{12}\text{C}(\vec{d}, p)^{13}\text{C}(g.s.)$ and for the reaction with the excitation of ^{13}C levels at the energies $T_d = 140$ and 200 MeV is observed.

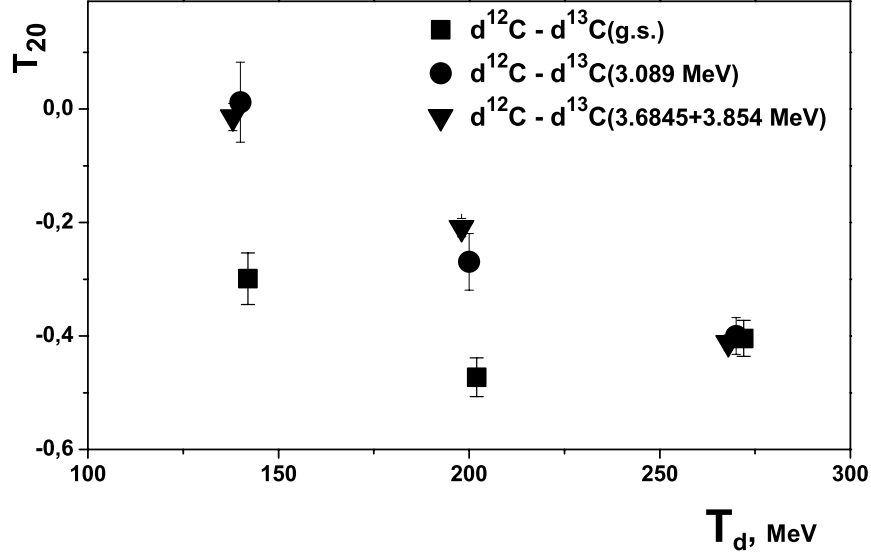


Figure 2: The experimental results on T_{20} in $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reactions at the energy $T_d = 140, 200$ and 270 MeV and emission angle $\Theta_{cm}=0^\circ$.

The experimental results on the vector A_y and tensor A_{yy} , A_{xx} and A_{xz} analyzing powers of the $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ and $^{12}\text{C}(\vec{d}, p)^{13}\text{C}(g.s.)$ reactions at $T_d = 270$ MeV are presented by the circles and triangles respectively in figure 3. The sign of A_y data for $^{12}\text{C}(\vec{d}, p)^{13}\text{C}(g.s.)$ and for $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reactions is different. Values of the tensor A_{yy} and A_{xz} analyzing powers for these reactions are positive and in accordance within experimental accuracy. The sign of A_{xx} data is negative.

Summary

The results on the tensor analyzing power T_{20} for the $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reaction at the energies $T_d = 140, 200$ and 270 MeV and emission angle $\Theta_{cm}=0^\circ$ have been obtained.

The tensor analyzing power for $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reactions is decreasing with increasing of the kinetic energy of an incident deuteron.

The data on T_{20} for $^{12}\text{C}(\vec{d}, p)^{13}\text{C}(3.089\text{MeV})$ and $^{12}\text{C}(\vec{d}, p)^{13}\text{C}(3.6845 + 3.854\text{MeV})$ reactions agree each other.

The experimental results on the vector A_y and tensor A_{yy} , A_{xx} , A_{xz} analyzing powers of the $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reaction at the energy $T_d = 270$ MeV in the angular range from 4°

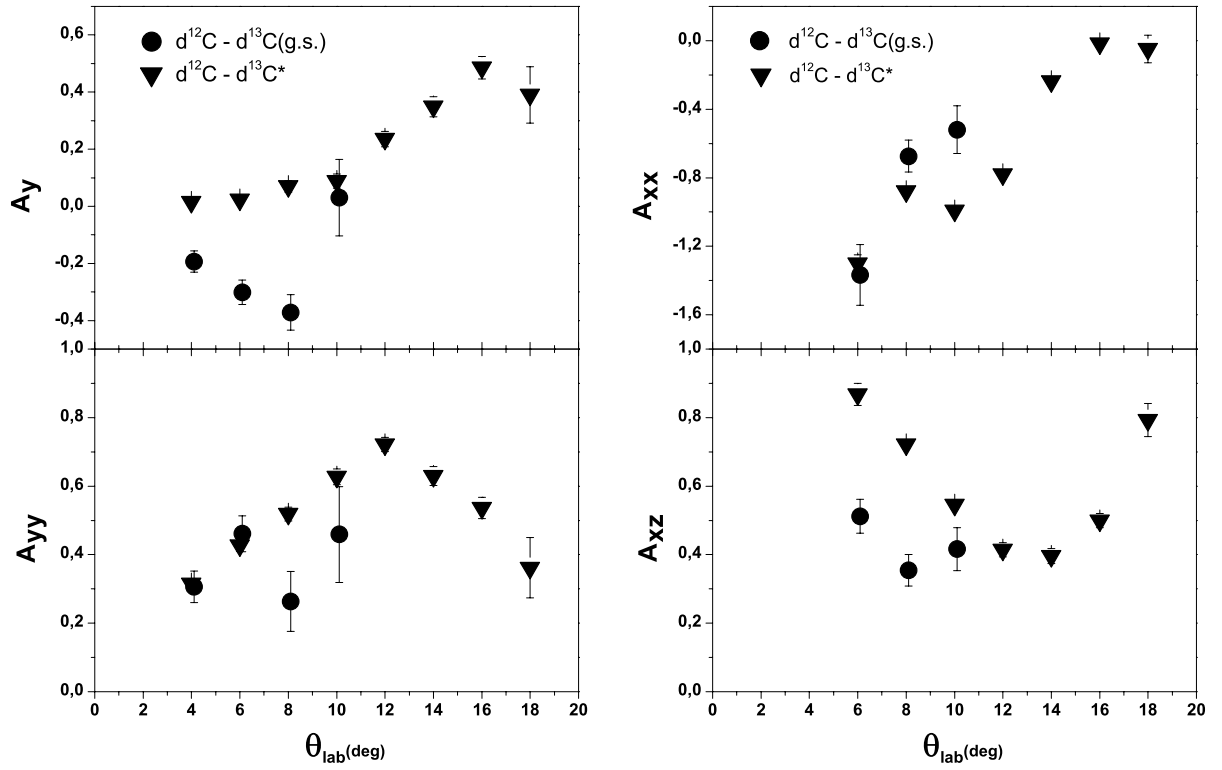


Figure 3: The experimental results on A_y , A_{yy} , A_{xx} and A_{xz} analyzing powers of the $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reaction at energy $T_d = 270$ MeV.

to 18° in laboratory system have been also obtained.

The A_{yy} data for $^{12}\text{C}(\vec{d}, p)^{13}\text{C}(\text{g.s.})$ and for $^{12}\text{C}(\vec{d}, p)^{13}\text{C}^*$ reactions are in accordance within achieved precision. The sign of tensor analyzing power for these reactions is positive, while the sign of the vector A_y analyzing power is different.

Acknowledgments

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OVERVIEW OF RECENT HERMES RESULTS

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Abstract

An overview of recent HERMES results is presented. The review topics include: the nucleon structure function $F_2(x)$, the polarized structure function $g_2(x)$, the search for a two-photon exchange effects in polarized DIS, the Collins and Sivers asymmetries, the azimuthal asymmetries in hadron electro-production off unpolarized target, and strange quark distributions, $S(x)$ and $\Delta S(x)$.

1 Introduction

The HERMES experiment at DESY was designed to investigate the spin structure of the nucleon with a clear motivation to study/resolve the “Proton Spin Puzzle”. The “Puzzle” originated due to a measurement by the European Muon Collaboration of the quark spin contribution to the proton spin, which appeared to be surprisingly small [1]. Twelve years of running by the HERMES experiment, from 1995 to 2007, has led to many exciting, unexpected results. New fields of investigations gained importance, such as the study of the third twist-2 quark distribution function of the nucleon, so-called transversity, the study of Transverse Momentum Dependent distribution and fragmentation functions, the study of Generalized Parton Distributions. This paper presents part of the results obtained by the HERMES Collaboration in the past two years based on inclusive and semi-inclusive DIS measurements only. Measurements of exclusive reactions are presented in other contributions of the HERMES Collaboration at this Workshop.

The HERMES experiment studies the scattering of a longitudinally polarized lepton beam (e^+ or e^-) of 27.6 GeV off a transversely/longitudinally polarized or unpolarized gas target internal to the HERA storage ring. Scattered leptons and coincident hadrons were detected by the HERMES spectrometer [2]. Leptons were identified with an efficiency exceeding 98% and a hadron contamination of less than 1%. The HERMES dual-radiator ring-imaging Čerenkov detector allows full hadron identification in the momentum range $2 \div 15$ GeV.

2 Inclusive DIS

Structure function $F_2(x)$. During its data-taking, HERMES collected about 58 million DIS events with (un)polarized hydrogen and deuterium targets. This statistics is highly competitive with respect to other fixed target DIS experiments, and motivates a new extraction of structure function $F_2(x)$ performed by HERMES in the same kinematic domain explored by its polarized measurements. Extraction of such a function requires the experimental data are corrected for detector smearing and radiative effects with an unfolding procedure [3], and a precise knowledge of the luminosity [2]. To extend the

kinematic range to low values of Bjorken x , a low 0.1 GeV^2 threshold is used for Q^2 . The structure function $F_2(x, Q^2)$ of the proton is presented in Fig. 1. The overall normalization uncertainty, which is mainly due to the luminosity normalization, is estimated to be 6.4%. Results of other experiments performed at DESY, CERN, and SLAC are presented as well for comparison. A previously unexplored region in the plane x - Q^2 is covered by this new measurement. At the same time there is a good agreement with world data in the overlap region. The structure function $F_2(x)$ of the deuteron is measured as well but not shown.

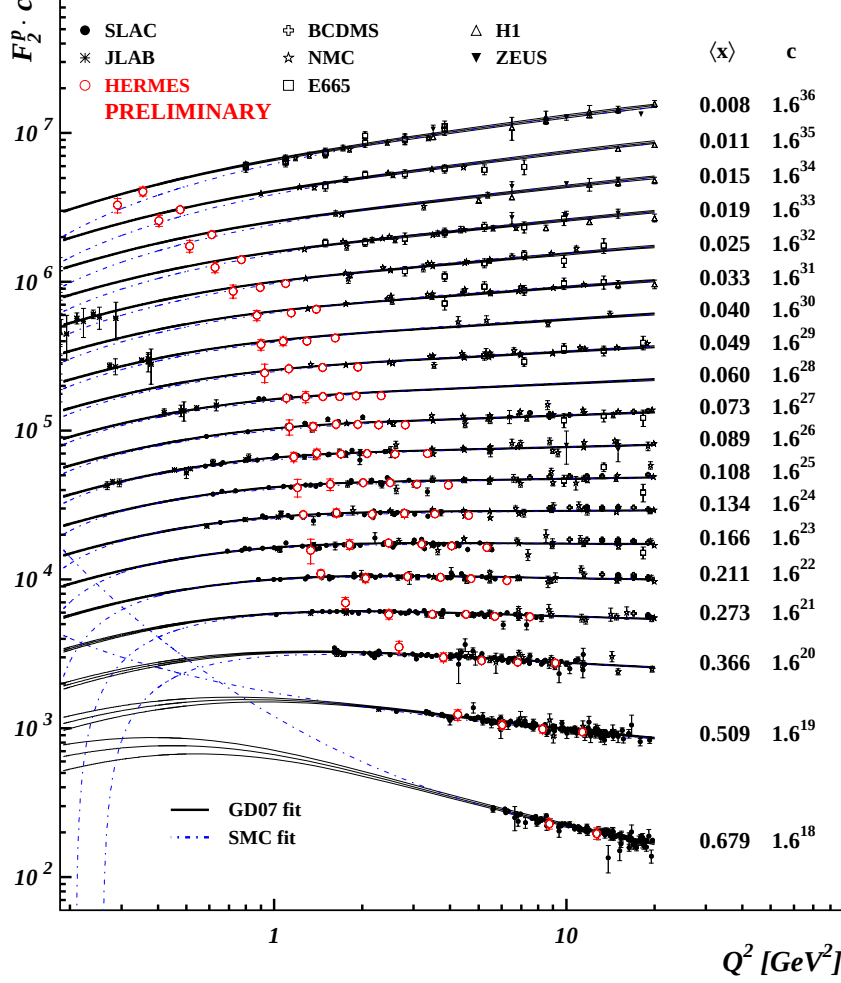


Figure 1: Structure function $F_2(x, Q^2)$ of the proton as a function of Q^2 , for different bins in x . Inner error bars are total minus the normalization uncertainty, shown as an outer error bar. The HERMES data are compared to measurements from NMC, BCDMS, JLab, HERA and SLAC. The phenomenological parameterizations of world data by the SMC [4] and GD07 [5] are shown by curves. Both parameterizations do not include the HERMES data points.

Structure function $g_2(x)$. The spin-dependent nucleon structure function $g_2(x)$ has no probabilistic interpretation in simple quark parton model. It has contributions from quark-gluon correlations and can be written as a sum of two terms, $g_2(x, Q^2) = g_2^{WW}(x, Q^2) + \bar{g}_2(x, Q^2)$, where $g_2^{WW}(x, Q^2) = -g_1(x, Q^2) + \int_x^1 g_1(y, Q^2) dy/y$ is the twist-2 Wandzura-Wilczek part, while $\bar{g}_2(x, Q^2)$ is the unknown twist-3 part. A measurement of the structure function $g_2(x)$ requires a longitudinally polarized beam and a transversely polarized target. The polarization-dependent part of the cross-section is given by:

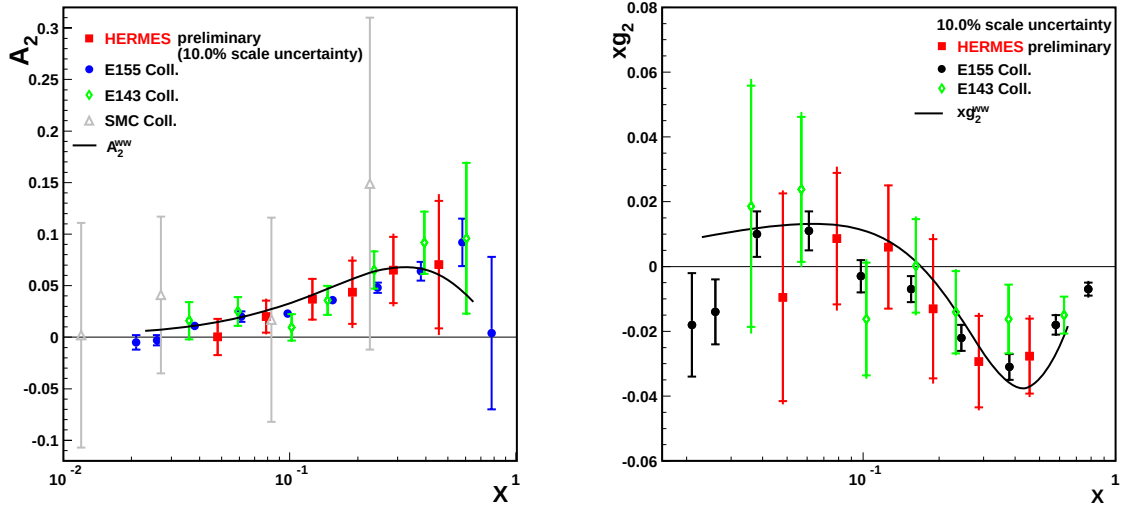


Figure 2: Left panel: the virtual photon asymmetry A_2^p as a function of x . Right panel: the structure function of the proton xg_2^p as a function of x . HERMES data are shown together with data from E155, E143 and SMC experiments. The lines correspond to predictions following Wandzura-Wilczek relation.

$$\frac{d^3\Delta\sigma}{dxdy d\phi_S} = -P_B P_T \cos\phi_S \frac{e^4}{4\pi^2 Q^2} \gamma \sqrt{1-y} \left(\frac{y}{2} g_1(x, Q^2) + g_2(x, Q^2) \right), \quad (1)$$

where P_B and P_T are the beam and target polarizations, while ϕ_S is the azimuthal angle about the beam direction between the lepton scattering plane and the “upwards” target spin direction.

During the 2003–2005 running period, an average polarization value $\langle P_B \cdot P_T \rangle$ of 0.24 was achieved. The experimental data are unfolded for detector smearing and radiative effects analogously to the studies of the structure function $g_1(x, Q^2)$ [3]. The HERMES results are presented in Fig. 2, where the virtual photon asymmetry A_2^p (left panel) and the structure function xg_2^p (right panel) are shown as functions of x . Data from experiments at CERN and SLAC are presented as well for comparison. HERMES results are in good agreement with the most accurate measurement made by the E155 experiment, and consistent with expectations based on the Wandzura-Wilczek term $g_2^{WW}(x)$.

Two-photon exchange contribution. The interpretation of spin asymmetries in DIS is based on the single-photon exchange assumption. Interference between one- and two-photon exchange amplitudes would lead to a $\sin\phi_S$ azimuthal asymmetry in inclusive DIS off transversely polarized target. Here, ϕ_S is the same angle defined in previous section. The asymmetry is expected to be proportional to the charge of the incident lepton and to $M/\sqrt{Q^2}$, where M is the nucleon mass. The Q^2 range was divided into a “DIS region”, with $Q^2 > 1 \text{ GeV}^2$, and a “low- Q^2 region”, with $Q^2 < 1 \text{ GeV}^2$, to test for a possible enhancement of the asymmetry due to the factor $M/\sqrt{Q^2}$. The amplitudes $A_{UT}^{\sin\phi_S}$ were extracted [6] separately for electrons and positrons. The asymmetries are consistent with zero for both the “DIS region” as well as for the “low- Q^2 region”. As conclusion of the study, no signal of the two-photon exchange was found within the uncertainty, which is of order 10^{-3} .

3 Semi-Inclusive DIS

Collins and Sivers asymmetries. HERMES has published its first results on azimuthal single spin asymmetries for charged pions produced in semi-inclusive DIS [7] off transversely polarized target using about 10% of the collected statistics. Preliminary results for pions and charged kaons obtained with full statistics have been presented at many conferences. The main contributions to the asymmetries were identified as due to the Collins [8] and Sivers [9] mechanisms. The Collins mechanism produces a correlation in the fragmentation process between the transverse spin of the struck quark and the transverse momentum $P_{h\perp}$ of the produced hadron. This correlation is described by the Collins fragmentation function $H_1^\perp(z)$. The Sivers mechanism produces a correlation between the transverse polarization of the target nucleon and the transverse momentum p_T of the struck quark. This correlation is described by the Sivers distribution function $f_{1T}^\perp(x)$. The contribution of the Collins mechanism to the asymmetry is proportional

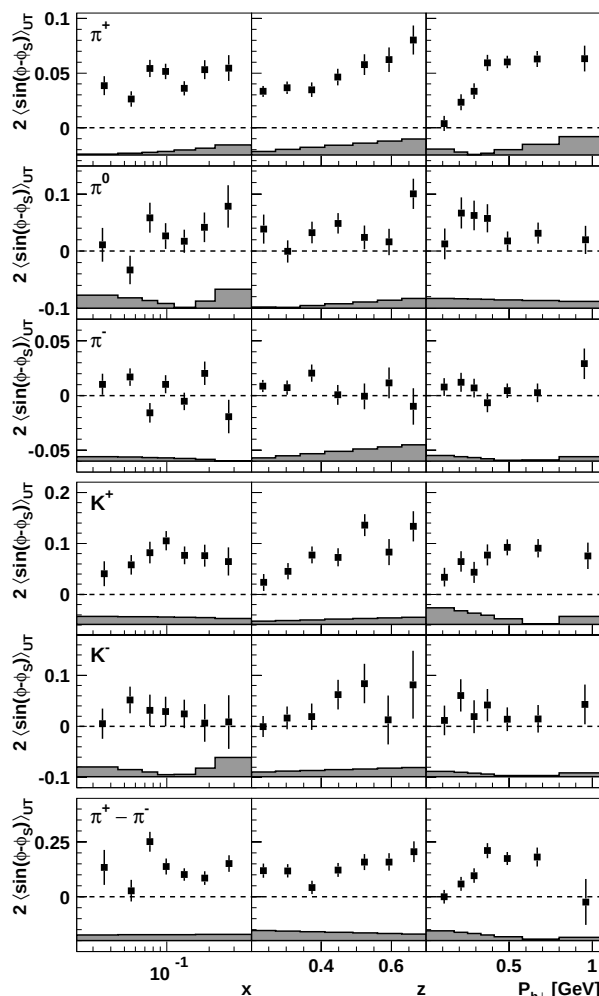


Figure 3: Sivers amplitudes for pions, charged kaons, and the pion-difference asymmetry (as denoted in the panels) as functions of x , z , or $P_{h\perp}$. The systematic uncertainty is shown as a band at the bottom of each panel. In addition there is a 7.3% scale uncertainty from the target-polarization measurement.

to the convolution of the transversity with the Collins function, while the contribution of the Sivers mechanism is proportional to the convolution of the Sivers function with the usual spin-independent fragmentation function. Fortunately, the contributions can be easily disentangled as they produce different azimuthal modulation of the asymmetry, $\sin(\phi + \phi_S)$ and $\sin(\phi - \phi_S)$ for the Collins and Sivers cases, respectively. Here, azimuthal angles ϕ and ϕ_S are the angles of the hadron momentum $\vec{P}_h$ and of the transverse component $\vec{S}_T$ of the target-proton spin, respectively, about the virtual-photon direction.

Final results for the Sivers amplitudes were obtained recently [10]. The resulting Sivers amplitudes for pions, charged kaons, and for the pion-difference asymmetry are presented in Fig. 3 as functions of x , z , or $P_{h\perp}$. The amplitudes are positive and increase with increasing z , except for π^- , for which they are consistent with zero. In order to test for higher-twist effects, the x dependence of the Sivers amplitudes for π^+ and K^+ is presented in Fig. 4, where each x -bin is divided into two Q^2 regions below and above the corresponding average Q^2 ($\langle Q^2(x_i) \rangle$) for that x bin. While the averages of the kinematics integrated over in those x bins do not differ significantly, the $\langle Q^2 \rangle$ values for the two Q^2 ranges change by a factor of about 1.7. The π^+ asymme-

tries in the two Q^2 regions are fully consistent, while there is a hint of systematically smaller K^+ asymmetries in the large Q^2 region.

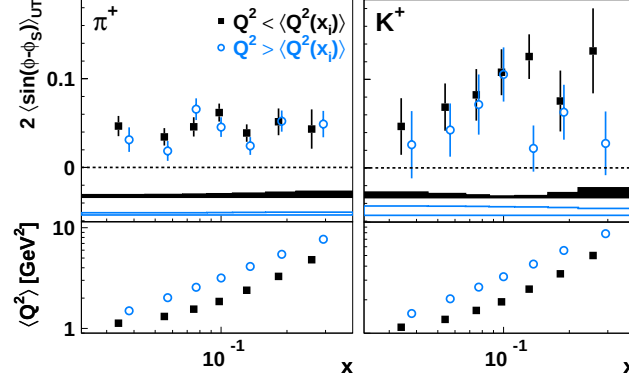


Figure 4: Sivers amplitudes for π^+ (left panel) and K^+ (right panel) as functions of x . The Q^2 range for each bin was divided into the two regions above and below $\langle Q^2(x_i) \rangle$ of that bin. In the bottom panel the average Q^2 values are given for the two Q^2 ranges.

A two-dimensional extraction of the Collins and Sivers amplitudes may be of great interest as it provides detailed information for the phenomenological study of specific distribution and fragmentation functions. A measurement of such two-dimensional amplitudes was done for charged pions (the statistics of charged kaons and neutral pions is too low for such a study) in variables x , z , and $P_{h\perp}$. As an example, in Fig. 5 the two-dimensional Collins amplitude $2 \langle \sin(\phi + \phi_S) \rangle_{UT}^\pi$ is presented as a function of z in bins of x (top panels) and as a function of x in bins of z (bottom panels).

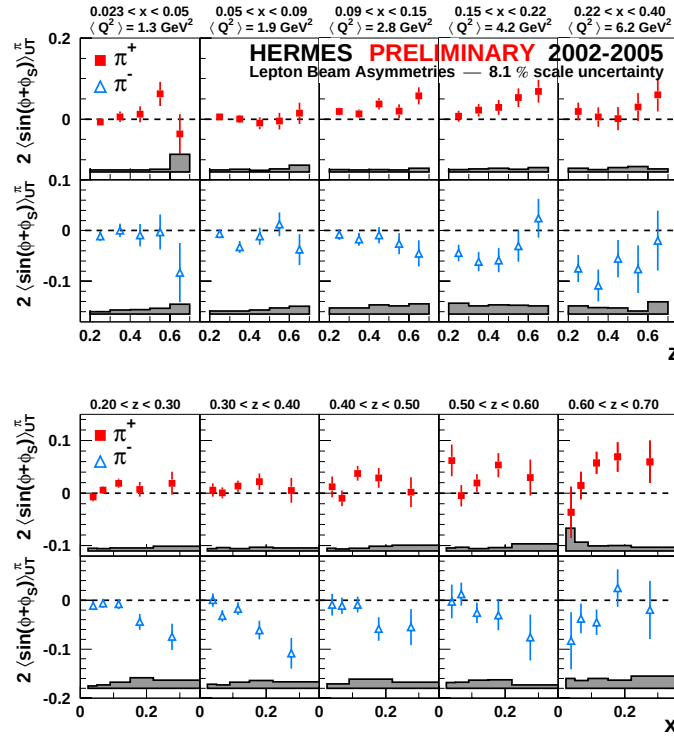


Figure 5: Collins amplitudes for π^+ (full squares) and π^- (open triangles). Upper panels: z dependence in five bins of Bjorken x . Lower panels: x dependence in five bins of z .

Unpolarized target asymmetries. The azimuthal asymmetry for hadron production in lepton DIS off unpolarized target was predicted to be non-zero many years ago due to the fact that the kinematics is noncollinear when the quark intrinsic transverse momentum is taken into account. Other possible sources of the asymmetry are perturbative gluon radiation and Boer-Mulders mechanism, which is due to the correlation of the quark intrinsic transverse momentum and intrinsic transverse spin. This correlation is described by the Boer-Mulders distribution function $h_1^\perp(x, k_T)$, which represents the transverse-polarization distribution of quarks inside an unpolarized nucleon.

The cross-section for hadron production in lepton-nucleon DIS $lN \rightarrow l'hX$ for the case of unpolarized target and unpolarized beam can be written in the following form [11]:

$$\frac{d^5\sigma}{dxdydzdP_{h\perp}^2d\phi_h} \propto A(y)F_{UU,T} + B(y)F_{UU,L} + C(y)\cos\phi_h F_{UU}^{\cos\phi_h} + B(y)\cos 2\phi_h F_{UU}^{\cos 2\phi_h} \quad (2)$$

Here, ϕ_h is the angle between the lepton scattering plane and the hadron production plane, $A(y)$, $B(y)$, and $C(y)$ are kinematic factors, $F_{UU,T(L)}$, $F_{UU}^{\cos\phi_h}$, and $F_{UU}^{\cos 2\phi_h}$ are structure functions responsible for 1, $\cos\phi_h$, $\cos 2\phi_h$ azimuthal modulations, respectively.

At HERMES, the extraction of the unpolarized modulations was performed using a multi-dimensional unfolding procedure to correct for radiative and acceptance effects on hydrogen and deuterium data, separately for positive and negative hadrons. The $\cos\phi_h$ amplitudes as function of x , y , z , and $P_{h\perp}$ for positive and negative hadrons produced off hydrogen target are presented in top panel of Fig. 6. The $\cos 2\phi_h$ amplitudes are presented in bottom panel of Fig. 6. An important feature shown by HERMES data is the different behaviour of the amplitudes for positive and negative hadrons. Such difference can be considered as an evidence of a non-zero Boer-Mulders function.

Analogous results were obtained for the hadron asymmetries with deuterium target.

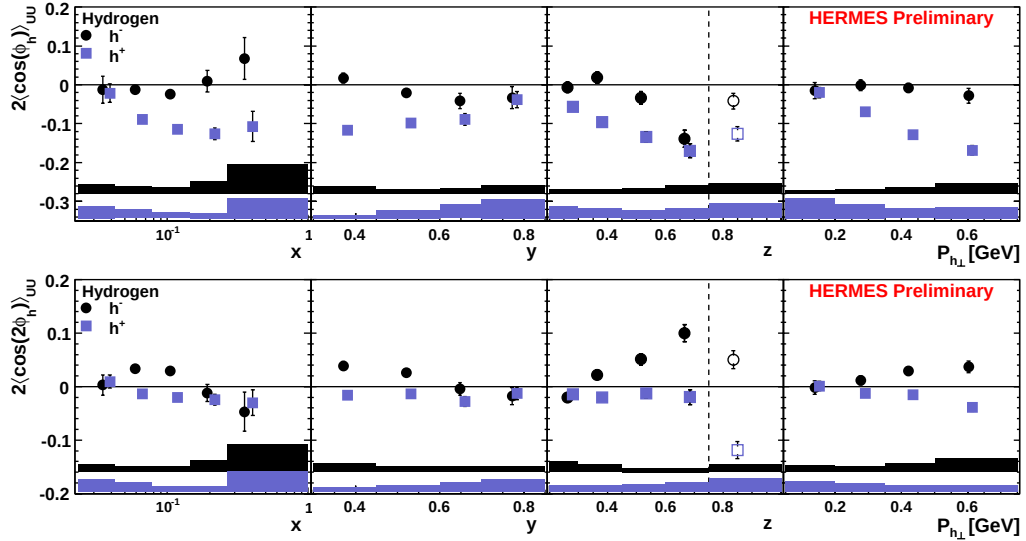


Figure 6: $2\langle\cos(\phi_h)\rangle_{UU}$ amplitude (top panel) and $2\langle\cos(2\phi_h)\rangle_{UU}$ amplitude (bottom panel) for negative (circles) and positive (squares) hadrons. Data are from hydrogen target. The error bars represent the statistical errors while the band represents the systematic errors. The open points at high z are not included in the projections over the other variables.

Strange quark distributions. The strange quark momentum, $S(x) = s(x) + \bar{s}(x)$, and helicity, $\Delta S(x) = \Delta s(x) + \Delta \bar{s}(x)$, distributions were measured [12] using data of semi-inclusive production of charged kaons off deuterium target, an isoscalar target. The analysis is based on the assumption of charge conjugation invariance in fragmentation and isospin symmetry between proton and neutron.

The multiplicity of charged kaons produced off deuterium target can be written at leading order in the following form:

$$\frac{dN^K(x)}{dN^{DIS}(x)} = \frac{Q(x) \int \mathcal{D}_Q^K(z) dz + S(x) \int \mathcal{D}_S^K(z) dz}{5Q(x) + 2S(x)}, \quad (3)$$

where $Q(x) = u(x) + \bar{u}(x) + d(x) + \bar{d}(x)$ is the non-strange parton distribution, $\mathcal{D}_Q^K(z)$ and $\mathcal{D}_S^K(z)$ are the fragmentation functions of non-strange and strange quarks into kaons, respectively. Using parameterization of $Q(x)$ from CTEQ6L one may extract the momentum distribution $S(x)$. The results are presented in Fig. 7 (left panel), together with parameterizations of $xS(x)$ and $x(\bar{u}(x) + \bar{d}(x))$ from CTEQ6L. The shape of $xS(x)$ measured by HERMES is incompatible with $xS(x)$ from CTEQ6L as well as with the assumption it corresponds to the average distribution of the non-strange sea.

In the extraction of the helicity distribution $\Delta S(x)$, only the double-spin asymmetry $A_{\parallel,d}^K(x, Q^2)$ for all charged kaons, irrespective of charge, and the inclusive asymmetry $A_{\parallel,d}(x, Q^2)$ are used. They can be expressed in terms of the non-strange quark helicity distribution $\Delta Q(x) = \Delta u(x) + \Delta \bar{u}(x) + \Delta d(x) + \Delta \bar{d}(x)$ and the strange quark helicity distribution $\Delta S(x)$:

$$\begin{aligned} A_{\parallel,d}(x) \frac{d^2 N^{DIS}(x)}{dx dQ^2} &= \mathcal{K}_{LL}(x, Q^2) [5\Delta Q(x) + 2\Delta S(x)], \\ A_{\parallel,d}^K(x) \frac{d^2 N^K(x)}{dx dQ^2} &= \mathcal{K}_{LL}(x, Q^2) \left[\Delta Q(x) \int \mathcal{D}_Q^K(z) dz + \Delta S(x) \int \mathcal{D}_S^K(z) dz \right]. \end{aligned} \quad (4)$$

Here, $\mathcal{K}_{LL}$ is a kinematic factor. These equations permit the simultaneous extraction of the helicity distribution $\Delta Q(x)$ and the strange helicity distribution $\Delta S(x)$. The results are presented in Fig. 7 (right panel). The strange helicity distribution is consistent with zero over the measured range.

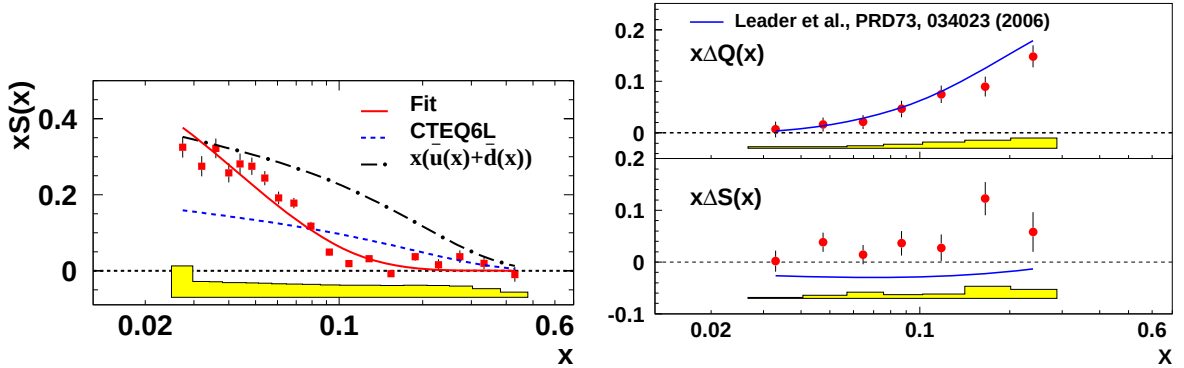


Figure 7: Left panel: The strange quark distribution $xS(x)$ at $Q^2 = 2.5 \text{ GeV}^2$. The solid curve is a fit of the data, the dashed curve shows the result from CTEQ6L, and the dot-dash curve is the average of light antiquark distributions from CTEQ6L. Right panel: nonstrange and strange quark helicity distributions at $Q^2 = 2.5 \text{ GeV}^2$. The error bars are statistical, and the bands represent the systematic uncertainties. The curves are the LO results of world data analysis.

4 Summary

Since the end of data-taking in 2007, the HERMES Collaboration has continued in refining the quality of and analyzing data.

The large statistics of inclusive DIS events collected by HERMES with hydrogen and deuterium targets allowed to perform a new measurement of the structure function $F_2(x)$ for proton and deuteron. The data cover a new still unexplored region and are in agreement with world data in the overlap region. A measurement of the spin-dependent structure function of the proton $g_2(x)$ is performed. This is the first measurement of HERMES where both, longitudinal polarization of the beam and transverse polarization of the target are important. The results are in agreement with the most accurate data obtained by the E155 experiment and do not show any evidence, within the statistical accuracy, of a contribution of the twist-3 part $\bar{g}_2(x)$. A search of the two-photon exchange effects using the transversely polarized hydrogen target gave no evidence within the uncertainty, at the level of 10^{-3} .

The final result on the extraction of the Sivers amplitudes from semi-inclusive production of pions and charged kaons using unpolarized lepton beam and transversely polarized hydrogen target was obtained. The results obtained based on a small part of the collected statistics are confirmed on the full data sample and new results on the amplitudes for production of neutral pions and charged kaons are presented. The amplitudes for positive kaons is larger than for positive pions. Azimuthal asymmetries for semi-inclusive production of positive and negative hadrons with unpolarized hydrogen and deuterium targets were studied. The different behaviour of positive and negative hadrons can be considered as an evidence of a non-zero Boer-Mulders distribution function. Strange quark distributions, $S(x)$ and $\Delta S(x)$, were measured using production of charged kaons off deuterium target. The measured shape of $xS(x)$ is in a contradiction with known parameterizations. The strange helicity distribution is consistent with zero.

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NLO QCD PREDICTIONS FOR GLUON POLARIZATION FROM OPEN-CHARM ASYMMETRIES MEASURED AT COMPASS

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Abstract

Recently published by the COMPASS Collaboration result on gluon polarization, $\Delta G/G$, from the open-charm channel has been obtained in the LO QCD approximation. The NLO QCD corrections to the analyzing power for the Photon-Gluon Fusion process are calculated in the COMPASS kinematical domain. The method based on the LO QCD Monte-Carlo generator with Parton Shower is proposed to simulate Phase Space for the NLO QCD processes. This approach can be easily implemented in the weighted method for estimation of gluon polarization, used in the COMPASS analysis. The new result for the gluon polarization in the NLO QCD approximation, based on published asymmetries for open-charm channel from COMPASS, is shown. The proposed method is compared with the approach where events are generated from uniformly distributed kinematical variables and then weighted by calculated NLO QCD cross section for Photon-Gluon Fusion process. This allows to control the potential bias related to the non-properly treated Phase Space for NLO QCD processes in the LO QCD Monte Carlo with Parton Shower concept.

1 Introduction

The COMPASS is a fixed target experiment at CERN laboratory. One of its goals is the direct measurement of the gluon polarization, important for understanding the spin structure of the nucleon. The experiment is using a 160 GeV polarized muon beam from SPS at CERN scattered off a polarized ^6LiD target [1].

In LO QCD approximation the only subprocess which probes gluons inside nucleon is Photon-Gluon Fusion (PGF). There are two ways allowing direct access to gluon polarization via the PGF subprocess available in the COMPASS experiment: the open-charm channel where the events with reconstructed D^0 mesons are used and the production of two hadrons with relatively high- p_T in the final state. The estimation of the gluon polarisation in the open-charm channel is much less Monte-Carlo (MC) dependent than in the two high- p_T hadrons method, where the complicated background requires very good MC description of the data. On the other hand the statistical precision in high- p_T hadrons method is much higher than in the open charm channel. To increase statistical precision the weighting method has been used in the open-charm analysis, recently published by COMPASS Collaboration [2]. The analysis has been performed in LO QCD approximation. The resolved photon contribution has been checked to be unimportant in the kinematical domain covered by COMPASS experiment. The estimation of the gluon

polarization as well as a construction of the statistical weight used in the analysis requires the knowledge of the so-called analyzing power, a_{LL} , (the ratio of polarized over unpolarized partonic cross sections) and the signal strength event-by-event basis. The analyzing power is calculated in the LO QCD approximation and as a signal identified D^0 meson (reconstructed from its decay) is used. In contrast to the resolved photon contribution NLO QCD corrections to unpolarized and polarized cross sections are supposed to be large in the COMPASS kinematical domain. To allow to use the COMPASS data in the independent analysis the open-charm asymmetries in bins in p_T and energy of D^0 meson was also published [2]. In this paper I would like to present the method of computing analyzing power event-by-event basis in the NLO QCD approximation. The method is based on LO MC with Parton Shower (PS) and can be easily allied in the COMPASS analysis scheme, also in the weighting method. The calculations are done for NLO QCD corrected PGF process and the new gluon polarization result based on published asymmetries are also presented. There is another part of NLO QCD corrections to muoproduction of open-charm: light quark originated processes where emitted gluon produces charm-anticharm quark pair. These processes contribute to the background because they don't probe gluons inside the nucleon. It introduces the complication into the analysis because signal defined as an observed D^0 meson is polluted by these higher order processes¹. They are not considered in this paper but the extension of the proposed method is rather straightforward however technically more complicated. It requires modification of the asymmetry decomposition equation which is a basis of the COMPASS analysis (see. eqs 5 and 6 in [2]). The paper is organized as follows: in the next section the NLO QCD corrections to the PGF process are discussed. The Monte-Carlo method and the Parton Shower concept is presented in section 3. The new results for the gluon polarization obtained using published asymmetries are presented in section 4. The discussion of the approximation used in the computation of the analyzing power is presented in section 4. Conclusions are presented in section 5.

2 NLO QCD corrections to the open-charm cross sections

In the LO QCD approximation PGF is the only process which contributes to the open-charm production. Moreover for the energy range covered by the COMPASS experiment the QCD evolution is not able to produce significant fraction of charm sea inside nucleon and only hard part of the cross section is responsible for open-charm production. As it was mentioned above the resolved photon contribution is also small and so-called intrinsic charm content inside nucleon is in the considered kinematics suppressed. Therefore observation of the charm signal seems to be an ideal to probe gluons inside nucleon for the COMPASS experiment. It is however known that the unpolarized cross section (averaged over spin states) is not precisely described by the LO QCD approximation and the NLO QCD corrections are important for the photoproduction of the open-charm, [4]. Also spin-dependent (polarized) cross section, calculated recently shows the important dependence of the approximation used [5]. It was also argued that the naive expectation

¹Notice that background taken into account in the COMPASS analysis is a combinatorial background; see also [3]

that NLO QCD corrections can be factorized and canceled out on the level of asymmetry, is not true.

Therefore it is important to estimate the order of the NLO QCD effects in the COMPASS experiment analysis. In this paper corrections to PGF process are only considered. As it was said in the Introduction the place where QCD approximation is used is the analyzing power, a_{LL} . The measured asymmetry is related to the gluon polarization as follows:

$$A^{exp} = P_b P_t f a_{LL} \frac{S}{S+B} \frac{\Delta g}{g} + A^{bgd} \quad (1)$$

where P_b , P_t and f are beam polarization, target polarization and dilution factor, respectively. A combinatorial background asymmetry, A^{bgd} , is extracted together with the signal (D^0 meson reconstructed in the event). $S/(S+B)$ is a signal purity and $\Delta g/g$ is a gluon polarization, integrated over kinematically accessible region by COMPASS measurement. For more details the reader is referred to [2].

There are two types of the NLO QCD corrections to PGF process: virtual and soft ones, where kinematics is the same as for LO PGF process (two into two particle kinematics) and the corrections with extra gluon emission (two into three particle kinematics). To guarantee the proper cancellation of the infrared and soft parts the virtual and real corrections should be added together (after integration over unobserved gluon in the final state) to obtain the so-called reduced cross section, free of any divergences [4]. COMPASS is using polarized muon beam what means that virtuality of the photon, Q^2 never reach the photoproduction limit. The smallest value of Q^2 allowed by kinematics is related to the muon mass. Nevertheless the calculation of the a_{LL} in the photoproduction limit is a very good approximation and can be used in the COMPASS analysis. The collection of the formulae for polarized and unpolarized cross sections for the finite Q^2 for PGF process in LO QCD approximation can be found in [6]. The NLO QCD corrections are partially listed in [4, 5] while the missing, finite parts are available by request [7]. To compute analyzing power a_{LL} event-by-event basis the knowledge of kinematics on of the parton level is needed. The measurement does not allow to reconstruct kinematical variables on the partonic level as the only one produced, charmed D^0 meson is reconstructed in the event at COMPASS. Therefore the exact kinematics of the event have to be simulated with the help of MC technique.² In the next section the MC approach for a_{LL} calculation in LO QCD approximation is discussed and the extension of the method including NLO QCD corrections is proposed.

3 Monte-Carlo approach for the calculation of the analyzing power

As the COMPASS analysis is performed in LO QCD approximation LO MC generator AROMA [8] is used to calculate analyzing power, a_{LL} . The COMPASS apparatus is simulated using GEANT package and produced output is in the form of the real data.

²In the ideal case where two charmed particles produced in the final states are reconstructed the LO QCD PGF process could be calculated using measured kinematics of charmed mesons. If unobserved gluons are radiated (NLO QCD corrections to PGF) the kinematics on the partonic level cannot be reconstructed from reconstructed heavy system.

Finally the COMPASS reconstruction program CORAL is used on the simulated events. This procedure allows to take into account all effects related to real data taking as acceptance, efficiency of the detector, reconstruction procedure efficiency etc, see [1]. The simulated event sample after all selection criteria applied [2] are used for computing a_{LL} event-by-event basis. The calculated analyzing powers, a_{LL} , are then parameterized by set of kinematical quantities like Q^2 or p_T of D^0 meson, accessible in the direct measurement for every event. A multi-dimensional Neural Network (NN) approach for the parameterization is used what allows to treat correctly correlations between kinematical variables [9]. After training on the MC sample the Neural Network is run on the real data computing a_{LL} for every real event. Finally a_{LL} is used to construct the statistical weight.

In the NLO QCD approximation however the problem with Phase Space for gluon emission processes appears. LO QCD approach in MC is not able to reproduce correctly kinematics of the events unless the so-called Parton Shower is switched on in the generation. The PS concept has been developed to improve the real data description by MC and allows to simulate multi-gluon emissions in some approximation [8]. Energy of all gluons emitted in the PS in the event can be considered as a limit of integration over unobserved gluons associated with the NLO QCD real corrections to PGF process. This procedure allows to calculate polarized and unpolarized cross sections in the LO and NLO QCD approximation, where real gluon emissions are integrated out over energy allowed by PS emissions in the event. The calculation of the analyzing power is then straightforward.

Two important problems, however, related to the method using PS concept appears. First, the CMS energy of the simulated event can be calculated using final charm-anticharm pair, or using initial photon-gluon system. In the LO QCD approximation there is no difference due to energy-momentum conservation but it is not longer true in the case of PS switched on. Performing integration over unobserved gluons with CMS energy of the event determined on the partonic level by heavy quark system the gluon distribution should be included into integrand function because of the fact that x_g , the momentum fraction carried by initial gluon is related to the energy of the real gluon emitted in the PS. It required the assumption about gluon distribution what is not convenient in the "direct" gluon polarization measurement. Therefore the CMS energy of the event on the partonic level should be defined by initial photon-gluon system. This case is even more similar to the real data analysis where only one charmed meson is reconstructed. The integrand contains only calculable partonic cross sections while gluon distribution is factorized as in the LO QCD. In the calculations presented in this paper both possibilities have been considered to check the potential effect related to this inconsistency in the treatment of the partonic kinematics and the three different polarized gluon distribution models were used.

The second problem is related to the fact that PS concept is not equivalent to MC in the NLO QCD approximation. There is still a big discussion how to use LO MC with PS to simulate effectively correct NLO processes but the subject is difficult and the satisfactory solutions exist only in some cases [10]. To test the correctness of the proposed method based on the LO MC with PS for PGF process for charm muoproduction the kinematics of the events with real gluon emission has been generated using uniformly distributed kinematical variables and then the events were weighted by the cross section calculated in the NLO QCD approximation (weighted MC method). This method guarantees that

NLO Phase Space is simulated correctly especially for hard gluon emissions which are not properly treated in the PS concept. On the other hand hard emissions of the real gluons should not be important because the probability of such process is rather small. The comparison of these two approaches is discussed in the next section. The weighted MC approach cannot be used in the COMPASS analysis scheme because the acceptance corrections and detector performance cannot be simulated in such a simple MC generator. As it will be seen in the next section the results obtained with AROMA generator and PS and using weighted MC approach are very similar what justify the method based on PS concept.

4 Predictions for the gluon polarization in NLO QCD approximation for PGF process

The gluon polarization predictions presented in this paper are estimated from the published asymmetries by COMPASS Collaboration. To calculate a_{LL} - only pure MC generator without apparatus simulation was used. Hence the statistical error is slightly underestimated. Asymmetries published in [2] are weighted by the weight composed of depolarization factor and signal strength ($S/(S+B)$). In the calculations presented here signal strength is assumed to be one because MC reproduces only signal (open charm events are generated). In the weighted MC approach the Peterson fragmentation function fitted to data from BELLE Collaboration was used [11] while for AROMA generation JETSET was applied. Three different models of gluon polarization, $\Delta G/G$, were used (in the case where partonic CMS energy is defined by final charm quark pair system): $\Delta G/G = \text{const}$ and two solutions from COMPASS QCD fits analysis, published in [12]. For an unpolarized gluons MRST 2004 PDF set was used.

To validate the method the gluon polarization in the LO QCD approximation (AROMA MC generator with PS switched off) and using the published asymmetries was found: $\Delta G/G = -0.47 \pm 0.23$. The good agreement between fully weighted COMPASS published result: $\Delta G/G = -0.49 \pm 0.27$, where acceptance and apparatus simulation have been taken into account justifies also COMPASS binning used for presenting D^0 asymmetries.

The gluon polarization results obtained in the NLO QCD approximation for the gluon polarization model independent case is: $\Delta G/G = 0.032 \pm 0.23$ and should be compared with the weighted MC result $\Delta G/G = 0.005 \pm 0.22$. The good agreement is seen what justifies the approach based on PS concept to simulate NLO QCD real emissions. Finally the gluon polarization results for three considered models of gluon polarizations (for the case where partonic CMS is defined by charm quark system) are: -0.051 ± 0.24 for $\Delta G/G = \text{const}$, -0.036 ± 0.24 for $\Delta G/G > 0$ and -0.057 ± 0.24 for $\Delta G/G < 0$ solutions of COMPASS QCD fits, respectively. In the light of these results the inconsistency related to definition of partonic CMS energy, discussed in the previous section, is not important.

5 Conclusions

The predictions for gluon polarization based on recently measured asymmetries for open-charm muoproduction by COMPASS Collaboration are presented. The results are ob-

tained in the NLO QCD approximation. The method of calculations of the analyzing power in NLO QCD approximation, based on AROMA MC generator with PS concept is discussed in details. The presented approach can be easily applied in the statistically weighted analysis as reported by the COMPASS Collaboration [2].

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STUDY OF LIGHT NUCLEI SPIN STRUCTURE FROM $p(d,p)d$, ${}^3\text{He}(d,p){}^4\text{He}$ AND $d(d,p){}^3\text{H}$ REACTIONS.

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Abstract

The data on the vector A_y and tensor A_{yy} , A_{xx} , A_{xz} analyzing powers for the $dd \rightarrow {}^3\text{H}p$ and $dd \rightarrow {}^3\text{H}en$ reactions at $E_d=200$ and 270 MeV are presented. The tensor analyzing powers for these reactions demonstrate the sensitivity to the ${}^3\text{H}$, ${}^3\text{He}$ and deuteron spin structure. The essential disagreements between the experimental results and the theoretical calculations within the framework of the one-nucleon exchange model are observed. Obtained experimental data are important for the preparation of the experiments at the LHEP-JINR Nuclotron-M at higher energies. It is planned to obtain new experimental data on the polarization observables for the $p(d,p)d$, ${}^3\text{He}(d,p){}^4\text{He}$ and $d(d,p){}^3\text{H}$ reactions at intermediate and high energies. These data will ensure new information on the spin structure of light nuclei at the short internucleonic distances, where the relativistic effects and the three-nucleon forces play an important role.

At the present time, there is a great deal of interest in understanding the short-range spin structure of deuteron and other light nuclei, as ${}^3\text{H}$ and ${}^3\text{He}$. They represent an important testing ground for models of the NN interaction and for studies of the many-body aspects of the strong interaction in nuclei. In addition, ${}^3\text{H}$ and ${}^3\text{He}$ are the simplest systems in which the effects of possible three-body nuclear forces can be explored.

The essential amount of the experimental data sensitive to the structure of light nuclei, have been accumulated at several last decades [1], [2]. Large discrepancies between the Nd data and theoretical predictions based on exact solution of the Faddeev equations

with only modern NN potential are reported. The inclusion 2π -exchange 3NF models into theoretical calculations removes many of them. In contrast, theoretical calculations with 3NFs still have difficulties in the reproducing of some spin observables and cross sections at backward scattering. One of the tools for understanding of the reason of these disagreements is the investigation of the processes sensitive to the deuteron, ${}^3\text{He}({}^3\text{H})$ spin structure and 3-nucleon forces(3NFs). Binary reactions like $dd \rightarrow {}^3\text{Hp}$, $dd \rightarrow {}^3\text{Hen}$, $dp \rightarrow pd$ and $d{}^3\text{He} \rightarrow p{}^4\text{He}$ are sensitive to the structure of deuteron and 3-nucleon forces. In addition, the $dd \rightarrow {}^3\text{Hp}({}^3\text{Hen})$ reaction can be used as an effective tool to investigate the structure ${}^3\text{H}$ and ${}^3\text{He}$ at short distances.

In this report the data on the angular distributions of the vector A_y and tensor A_{yy} , A_{xx} , A_{xz} analyzing powers in the $\vec{d}d \rightarrow {}^3\text{Hp}$ and $\vec{d}d \rightarrow {}^3\text{Hen}$ reactions at 200 and 270 MeV of the deuteron kinetic energy are presented. These data are important for the preparation of the experiments at the LHEP-JINR Nuclotron-M at higher energies. It is planned to obtain new experimental data on the polarization observables for the $p(d,p)d$, ${}^3\text{He}(d,p){}^4\text{He}$ and $d(d,p){}^3\text{H}$ reactions. These data will ensure new information on the spin structure of light nuclei at the short internucleonic distances.

Experiment on the measurement of the analyzing powers in the $\vec{d}d \rightarrow {}^3\text{Hp}$ and $\vec{d}d \rightarrow {}^3\text{Hen}$ reactions was performed at RARF(RIKEN, Japan). The details on the experiment are described elsewhere [3]. The experimental results on the vector and tensor analyzing powers of the $\vec{d}d \rightarrow {}^3\text{Hp}({}^3\text{Hen})$ reactions are presented in Fig.1. The errors of the experimental values include both the statistical and the systematic errors. The systematic errors were derived from the errors of the beam polarizations measurements. The solid and long-dashed curves are the results of ONE calculations [4] using CD-Bonn and Paris deuteron and ${}^3\text{He}$ wave

functions [5], respectively. One can see that ONE calculations are in the qualitative agreement with the data on the T_{20} . But they don't reproduce the behavior of the tensor A_{yy} , A_{xx} and A_{xz} analyzing powers in the full angular range. The vector analyzing power A_y equals zero in the ONE model. Some structures in the experimental results on A_y indicate possibility other than ONE mechanisms in these reactions. The reason of the discrepancy can be in the non-adequateness of the 3N-bound state spin structure and/or more complicated reaction mechanism. The multiple scattering calculations are in progress now.

We plan to measure the cross section, vector A_y and tensor A_{yy} and A_{xx} analyzing powers in dp -elastic scattering [6] at large angles in c.m.s. in the energy range 0.3-2.0 GeV at Internal Target Station at Nuclotron-M. At deuteron energies below 300 MeV, Faddeev calculations can provide good description of cross section data, by introducing a three nucleon force, over the whole angular range. However, their reproduction of the

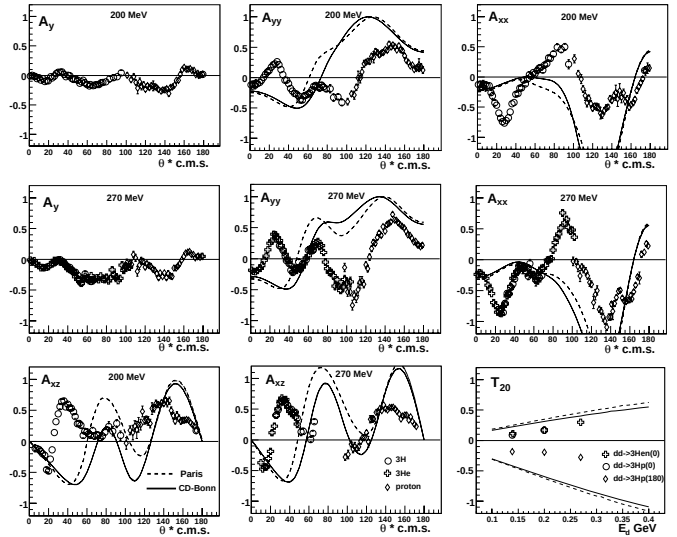


Figure 1: The analyzing powers of the $\vec{d}d \rightarrow {}^3\text{Hp}({}^3\text{Hen})$ reactions. The solid and long-dashed curves are the results of ONE calculations.

polarization observables are not so good as that for cross section data, which may be regarded as an insufficiency of our knowledge on spin dependence of three nucleon force. Measurement of energy dependence of polarization observables for dp reaction in the region of cross section minimum can give an irreplaceable clue to the problem. The Fig.2 show the dependence of the vector A_y and tensor A_{yy} analyzing powers in dp - elastic scattering at the fixed angles in the c.m.s. as a function of transverse momentum p_T . The open and solid symbols represent the data obtained at RIKEN, Saclay, ANL and at Nuclotron (Dubna), respectively. The change of the sign in the both A_y and A_{yy} analyzing powers values at p_T 600-700 MeV is observed. Further precise measurements are required to understand the reason of such behavior.

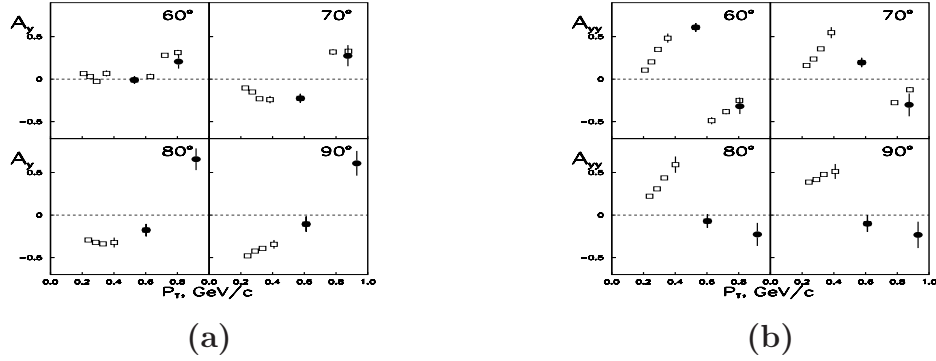


Figure 2: (a) The dependence of the vector A_y analyzing power in dp - elastic scattering at the fixed angles in the c.m.s. as a function of transverse momentum p_T . (b) The dependence of the tensor A_{yy} analyzing power in dp - elastic scattering at the fixed angles in the c.m.s. as a function of transverse momentum p_T .

The goal of the ${}^3\text{He}(d,p){}^4\text{He}$ reaction study at Nuclotron-M is to understand the reason of the long staying puzzle, namely, the behavior of the tensor analyzing power T_{20} in dp - backward elastic scattering [7]. T_{20} in dp - scattering (see Fig.3(a)) demonstrates unexplained strange structure at the internal momentum $k \sim 0.3\text{--}0.5$ GeV/c (in the vicinity of the D -wave dominance). The experiments performed at RIKEN at the energies below 270 MeV have shown that the polarization correlation coefficient, $C_{//} = 1 - \frac{1}{2\sqrt{2}}T_{20} + \frac{3}{2}C_{y,y}$, for the ${}^3\text{He}(d,p){}^4\text{He}$ reaction may be a unique probe to the D -state admixture in deuteron [8]. Tensor analyzing power T_{20} , spin correlation $C_{y,y}$ and polarization correlation coefficient $C_{//}$ for the ${}^3\text{He}(d,p){}^4\text{He}$ reaction are shown in Fig.3(b). Solid lines in the figures represent calculations based on an impulse approximation proposed in [9]. The full symbols are the data obtained at RIKEN. The open squares show the expected precision for the data at Nuclotron.

The main goal of the experiment is to obtain the data on $C_{//}$ in the energy region of 1.0–1.75 GeV, where the contribution from the deuteron D -state is expected to reach a maximum in ONE approximation, to obtain new information on the strange structure observed in the behavior of T_{20} in the dp - backward elastic scattering and to realize experiment on the full determination of the matrix element of the ${}^3\text{He}(d,p){}^4\text{He}$ reaction in the model independent way. These data will help us also to understand the short-range spin structure of deuteron and effects of non-nucleonic degrees of freedom.

New experimental data will ensure the important information on the spin structure of light nuclei at the short internucleonic distances, where the relativistic effects and the

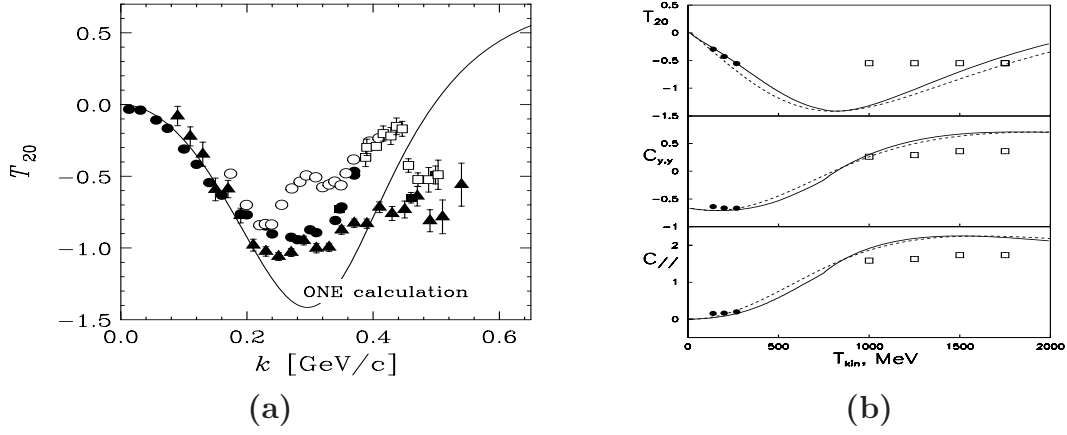


Figure 3: (a) T_{20} data taken at Dubna and Saclay plotted as a function of the internal momentum k . (b) Tensor analyzing power T_{20} , spin correlation $C_{y,y}$ and polarization correlation coefficient $C_{//}$ for the ${}^3\text{He}(d,p){}^4\text{He}$ reaction.

three-nucleon forces play an important role.

Acknowledgments

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EXCLUSIVE ELECTROPRODUCTION OF ρ^0 , ϕ , AND ω MESONS AT HERMES

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Abstract

Spin Density Matrix Elements (SDMEs) in exclusive ρ^0 - and ϕ -meson electroproduction have been measured in the HERMES experiment with a 27.5 GeV longitudinally polarized electron(positron) beam impinging on unpolarized hydrogen and deuterium targets at kinematics of $1 < Q^2 < 7 \text{ GeV}^2$, $3 < W < 6.3 \text{ GeV}$, and $-t < 0.4 \text{ GeV}^2$. The ρ^0 SDMEs deviate from expectations based on s -channel helicity conservation (SCHC) by factors of from 3 to 10 of the total measurement uncertainties σ_{tot} , while no signal of SCHC violation is observed in the ϕ meson SDMEs. Ratios of natural-parity-exchange helicity amplitudes T_{11} ($\gamma_T^* \rightarrow \rho_T$), T_{01} ($\gamma_T^* \rightarrow \rho_L$) to T_{00} ($\gamma_L^* \rightarrow \rho_L$) are extracted from the data. For the unnatural-parity-exchange amplitude U_{11} , the value $|U_{11}/T_{00}|$ is obtained. The Q^2 and t dependences of all the ratios are compared with perturbative QCD predictions. Azimuthal target spin asymmetries for ω -meson production on the transversely polarized proton were also measured.

1 Introduction

Exclusive electroproduction of vector mesons (V) on nucleons (N) provides valuable information both on the strong interaction and the structure of the hadrons involved in the process. Renewed interest in exclusive meson production in deep-inelastic scattering (DIS) of leptons has resulted from the possibility to describe the process in terms of generalized parton distributions (GPD), which contain a wealth of information about the nucleon structure. In the one-photon-exchange approximation, all the observables in electroproduction can be expressed through the helicity amplitudes $F_{\lambda_V \nu_N \lambda_\gamma \lambda_N}$ for the process

$$\gamma^*(\lambda_\gamma) + N(\lambda_N) \rightarrow V(\lambda_V) + N(\nu_N), \quad (1)$$

in which the helicities are given in parentheses. The amplitudes depend on the heavy photon (γ^*) virtuality $Q^2 = -q^2 > 0$, the center-of-mass energy of the γ^*N system W , and the squared four-momentum transfer to the nucleon t . Since the spin-density matrix of the virtual photon is well known from quantum electrodynamics and the SDMEs of vector mesons may be extracted from the angular distributions of the decay particles, the properties of the helicity amplitudes can be studied in detail.

The formalism for the description of the process (1) for unpolarized targets was developed in [1], in which expressions for the SDMEs $r_{\lambda_V \lambda'_V}^\alpha$ ($\alpha = 0, 1, \dots, 8$) in terms of the helicity amplitudes were established. It was first pointed out in [2] that the physics

of electroproduction from polarized targets is richer than that for polarized lepton beams and unpolarized targets, and the formalism was generalized for the case of polarized targets. Recently a new general formalism for the description of vector-meson SDMEs was proposed in [3]. From the symmetry properties [1–3] of the helicity amplitudes it follows that there are 18 independent amplitudes. Since the SDMEs depend on ratios of these amplitudes, there are 34 independent real parameters needed to determine all the SDMEs. However the number of SDMEs for a given measurement can exceed 34 as, for example, in the case of the 45 [3] SDMEs observable from electroproduction from a transversely polarized target with an unpolarized lepton beam. This means that the SDMEs themselves are not independent and amplitude ratios A_i provide a more economic basis for fitting the angular distributions of the decay particles. The hierarchy of amplitudes established at HERMES kinematics [4], namely

$$|T_{00}|^2 \sim |T_{11}|^2 \gg |U_{11}|^2 > |T_{01}|^2 \gg |T_{10}|^2 \sim |T_{1-1}|^2, \quad (2)$$

permits one to reduce the number of essential free amplitude ratios to 5: $A_1 = \text{Re}(T_{11}/T_{00})$, $A_2 = \text{Im}(T_{11}/T_{00})$, $A_3 = \text{Re}(T_{01}/T_{00})$, $A_4 = \text{Im}(T_{01}/T_{00})$, and $A_9 = |U_{11}/T_{00}|$. The ratios $A_5 - A_8$ are found to be very small, where $A_5 = \text{Re}(T_{10}/T_{00})$, $A_6 = \text{Im}(T_{10}/T_{00})$, $A_7 = \text{Re}(T_{1-1}/T_{00})$, $A_8 = \text{Im}(T_{1-1}/T_{00})$. The quantities $T_{\lambda_V \nu_N \lambda_\gamma}$ are a shorthand notation for the amplitudes $T_{\lambda_V \frac{1}{2} \lambda_\gamma \frac{1}{2}}$ for natural-parity exchange (NPE) without nucleon spin flip; the amplitude of unnatural-parity exchange (UPE) is $U_{\lambda_V \nu_N \lambda_\gamma \lambda_N}$ and $|U_{11}|^2 = |U_{1\frac{1}{2}1\frac{1}{2}}|^2 + |U_{1-\frac{1}{2}1\frac{1}{2}}|^2$. These NPE and UPE amplitudes are related to helicity amplitudes $F_{\lambda_V \nu_N \lambda_\gamma \lambda_N}$ by the equations [1–3] $T_{\lambda_V \nu_N \lambda_\gamma \lambda_N} = (F_{\lambda_V \nu_N \lambda_\gamma \lambda_N} + (-1)^{\lambda_V - \lambda_\gamma} F_{-\lambda_V \nu_N - \lambda_\gamma \lambda_N})/2$, $U_{\lambda_V \nu_N \lambda_\gamma \lambda_N} = (F_{\lambda_V \nu_N \lambda_\gamma \lambda_N} - (-1)^{\lambda_V - \lambda_\gamma} F_{-\lambda_V \nu_N - \lambda_\gamma \lambda_N})/2$. In Regge phenomenology, the NPE amplitudes correspond to exchanges of reggeons of natural parity with $P = (-1)^J$ (Pomeron, ρ , f_2 , a_2 , ...), whereas exchanges of reggeons of unnatural parity with $P = -(-1)^J$ (π , a_1 , b_1 , ...) contribute to the UPE amplitudes.

2 Selection of Exclusive Vector Meson Events

The ρ^0 and ϕ events used for the SDME analysis were produced in DIS of longitudinally polarized electrons and positrons with an energy 27.57 GeV from unpolarized hydrogen and deuterium targets; the ω events were collected from DIS with a transversely polarized proton and unpolarized beam. The scattered electron (positron) and particles from the decays $\rho^0 \rightarrow \pi^+ \pi^-$, $\phi \rightarrow K^+ K^-$, $\omega \rightarrow \pi^+ \pi^- \pi^0$ ($\pi^0 \rightarrow 2\gamma$) were detected and identified by the HERMES spectrometer [5]. In the period 1996–2000 particles were identified using a threshold Čerenkov detector, lead glass electromagnetic calorimeter and preshower detector; since 2000 a dual-radiator RICH detector [6] replaced the threshold detector. The constraints $0.6 < M_{\pi\pi} < 1.0$ GeV, $0.99 < M_{KK} < 1.04$ GeV, and $0.71 < M_{\pi\pi\pi} < 0.87$ GeV ($0.1 < M_{\gamma\gamma} < 0.16$ GeV) were applied to the invariant mass of the final particles to identify ρ^0 , ϕ , and ω mesons. Since the recoil nucleon was not detected exclusive events were selected by the requirement $\Delta E = (M_X^2 - M^2)/(2M) < 0.6$ GeV where M denotes the nucleon mass and M_X is the missing mass of the recoiling hadronic system. Diffractive events for ρ^0 and ϕ were selected by the constraint $-t' < 0.4$ GeV² while for ω production $-t' < 0.6$ GeV² was used. Here $t' = t - t_0$ with $-t_0$ being the smallest value of $-t$ at fixed W , Q^2 , and vector-meson mass M_V . The semi-inclusive DIS background was

subtracted using simulated events generated by PYTHIA. The fraction of semi-inclusive DIS events for the entire kinematic region was found to be 8% for ρ^0 , 1.6% for ϕ , and 29% for ω mesons. The SDMEs were considered as free parameters in fitting the angular distributions of scattered leptons and decay hadrons using the formalism of Ref. [1]. The maximum likelihood method was used to determine the best fit with uncertainty calculations being performed by MINUIT (see details in [4]).

3 Spin Density Matrix Elements and Azimuthal Target Spin Asymmetry

The SDMEs for ρ^0 and ϕ mesons extracted for the entire kinematic region $1 < Q^2 < 7 \text{ GeV}^2$, $3 < W < 6.3 \text{ GeV}$, and $0 < -t' < 0.4 \text{ GeV}^2$ are presented in Fig. 1. As this data was collected with longitudinally polarized beam of polarization ± 0.53 , it was possible to make the first determination of 8 LU SDMEs (LU: Longitudinally polarized beam, Unpolarized target) of the matrices r_{ij}^3 , r_{ij}^7 , and r_{ij}^8 (in the Schilling-Wolf notation). They are shown in Fig. 1 in yellow bands. The other 15 UU SDMEs may be measured in experiments with both an Unpolarized beam and Unpolarized target; those from HERMES are also shown in Fig. 1. According to the hierarchy (2) discussed above, the SDMEs are categorized into five classes. Class A comprises SDMEs with dominant terms proportional to $|T_{00}|^2/\mathcal{N}$ or $|T_{11}|^2/\mathcal{N}$ where $\mathcal{N}$ is the normalization factor (see definition in [4]) common to all SDMEs. Class B contains elements whose main terms are proportional to $\text{Re}(T_{11}T_{00}^*)/\mathcal{N}$ for UU SDMEs and $\text{Im}(T_{11}T_{00}^*)/\mathcal{N}$ for LU matrix elements, where T_{00}^* is the complex conjugate amplitude. This corresponds to the general rule: the dominant term of UU (LU) SDMEs is proportional to the real (imaginary) part of the product of two amplitudes. Class C contains SDMEs whose dominant terms are proportional to the product of the amplitude T_{01} (transition of the transverse photon to the longitudinal vector meson, $\gamma_T^* \rightarrow V_L$) and T_{00}^* ($\gamma_L^* \rightarrow V_L$) or T_{11}^* ($\gamma_T^* \rightarrow V_T$). For r_{00}^1 the T_{01} contribution is quadratic. Classes D and E are composed of SDMEs in which the main terms contain products of T_{10} ($\gamma_T^* \rightarrow V_L$) and T_{1-1} ($\gamma_{-T}^* \rightarrow V_T$), respectively, by T_{11}^* . The SDMEs in Fig. 1 are multiplied by numerical factors to make the dominant terms equal to $\text{Re}(T_{ij}T_{mm}^*)/\mathcal{N}$ or $\text{Im}(T_{ij}T_{mm}^*)/\mathcal{N}$ ($i, j = 0, \pm 1, m = 0, 1$).

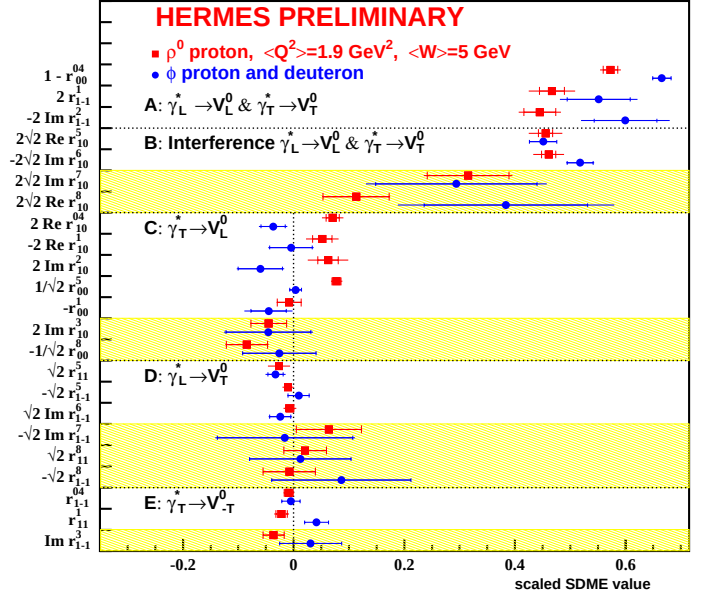


Figure 1: The 23 SDMEs for ρ^0 production on the proton (squares) and for ϕ production on proton and deuteron (circles). The inner error bars show statistical uncertainty while the outer ones indicate statistical and systematic uncertainties added in quadratures. The unshaded (shaded) areas show unpolarized UU (polarized LU) SDMEs divided into five classes (see text for more explanations).

Class A comprises SDMEs with dominant terms proportional to $|T_{00}|^2/\mathcal{N}$ or $|T_{11}|^2/\mathcal{N}$ where $\mathcal{N}$ is the normalization factor (see definition in [4]) common to all SDMEs. Class B contains elements whose main terms are proportional to $\text{Re}(T_{11}T_{00}^*)/\mathcal{N}$ for UU SDMEs and $\text{Im}(T_{11}T_{00}^*)/\mathcal{N}$ for LU matrix elements, where T_{00}^* is the complex conjugate amplitude. This corresponds to the general rule: the dominant term of UU (LU) SDMEs is proportional to the real (imaginary) part of the product of two amplitudes. Class C contains SDMEs whose dominant terms are proportional to the product of the amplitude T_{01} (transition of the transverse photon to the longitudinal vector meson, $\gamma_T^* \rightarrow V_L$) and T_{00}^* ($\gamma_L^* \rightarrow V_L$) or T_{11}^* ($\gamma_T^* \rightarrow V_T$). For r_{00}^1 the T_{01} contribution is quadratic. Classes D and E are composed of SDMEs in which the main terms contain products of T_{10} ($\gamma_T^* \rightarrow V_L$) and T_{1-1} ($\gamma_{-T}^* \rightarrow V_T$), respectively, by T_{11}^* . The SDMEs in Fig. 1 are multiplied by numerical factors to make the dominant terms equal to $\text{Re}(T_{ij}T_{mm}^*)/\mathcal{N}$ or $\text{Im}(T_{ij}T_{mm}^*)/\mathcal{N}$ ($i, j = 0, \pm 1, m = 0, 1$).

As is seen in Fig. 1 values of the elements of classes A and B from the ρ^0 are close to those from the ϕ meson. If the SCHC approximation is valid, only elements of classes A and B may be nonzero. There is no statistically significant violation of SCHC observed in the ϕ meson elements, while elements of class C from the ρ^0 meson deviate from zero with factors of from 3 to 10 of the total uncertainties σ_{tot} . Since $T_{01} \equiv 0$ in the absence of longitudinal quark motion [7], this could mean that the quark motion effects in the ρ^0 are stronger than in the ϕ meson due to the heavier valence quarks ($s\bar{s}$) compared to those of the ρ^0 meson ($u\bar{u}$ and $d\bar{d}$). Note however that the H1 Collaboration has observed a small violation of SCHC [8] for the ϕ meson.

As was shown in [9], contributions to the azimuthal target spin asymmetry $A_{UT}^{\sin(\phi-\phi_s)}$ of u and d quarks cancel for the ρ^0 and have the same sign for the ω meson. This makes the ω meson an excellent candidate for studying the GPD E . HERMES accumulated 429 ω production events from the transversely polarized proton. The results show that the sign of $A_{UT}^{\sin(\phi-\phi_s)}$ is negative, in agreement with the prediction of Ref. [9], and deviates from zero at $-t' = 0.12 \text{ GeV}^2$ by twice the total uncertainty σ_{tot} .

4 Direct Extraction of Helicity Amplitude Ratios

Fitting the angular distribution of decay pions in correlation with the scattered lepton provides the parameters of the amplitude ratios $A_1 - A_9$ for ρ^0 production if the SDMEs are expressed in terms of these amplitudes. The Q^2 dependence of $\text{Re}\{T_{11}/T_{00}\}$ ($\text{Im}\{T_{11}/T_{00}\}$) is shown in left (right) panel of Fig. 2 both for the hydrogen and deuterium targets.

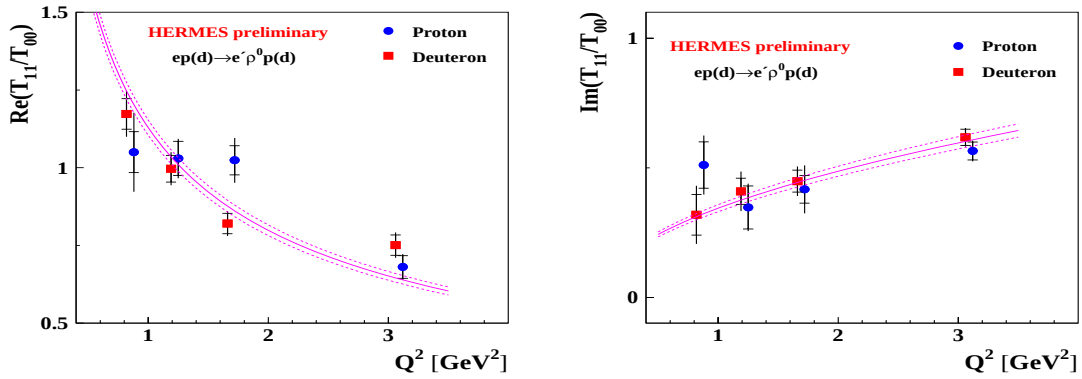


Figure 2: The Q^2 dependence of $A_1 = \text{Re}(T_{11}/T_{00})$ and $A_2 = \text{Im}(T_{11}/T_{00})$ for hydrogen and deuterium targets. The inner error bars show the statistical uncertainties while the outer ones indicate statistical and systematic uncertainties added in quadrature. The parameterization of the curves is given in text. The solid lines are calculated with mean values of the fit parameters; dashed lines correspond to a one standard deviation change in the curve parameters.

The parameterization $A_1 = a/Q$, $A_2 = bQ$ describes the Q^2 dependence well in the HERMES kinematic region for both targets. Since the angular distributions for production on the proton and deuteron are compatible within the experimental accuracy, a fit to the combined data was performed. The results $a = (1.129 \pm 0.024) \text{ GeV}$ with χ^2 per degree of freedom $\chi^2/N_{df} = 1.02$, $b = (0.344 \pm 0.014) \text{ GeV}^{-1}$, $\chi^2/N_{df} = 0.87$ correspond to the curves presented in Fig. 2; the solid lines are calculated using the mean values of the

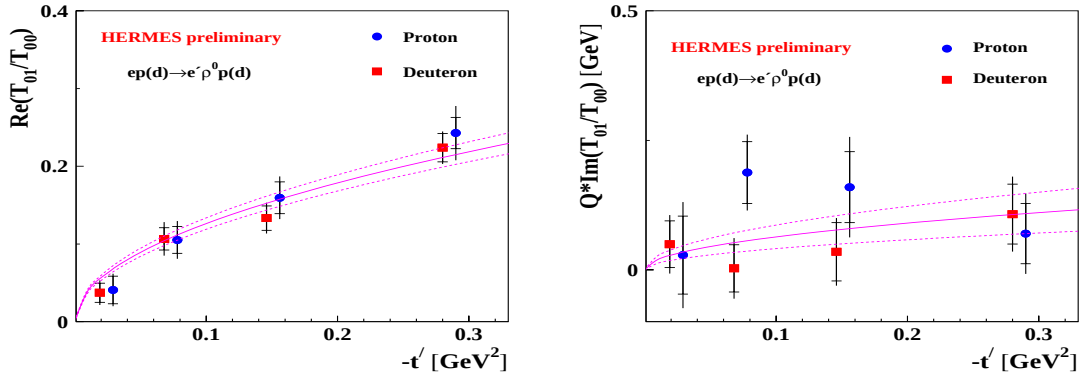


Figure 3: The t' dependence of $A_3 = \text{Re}(T_{01}/T_{00})$ and $A_4 = \text{Im}(T_{01}/T_{00})$ for hydrogen and deuterium targets. Inner (outer) error bars show the statistical (total) uncertainties, as in Fig. 2. The different curves represent the fit parameterizations and their uncertainties, also as in Fig. 2.

parameters a and b while the dashed curves correspond to $\pm$ one standard deviation of the total uncertainty. No t' dependence of A_1 and A_2 was observed in the region $-t' < 0.4$ GeV^2 . The result for A_1 is in agreement with the prediction [10, 7] valid at high Q^2 , $T_{11}/T_{00} \propto M_V/Q$, obtained within a perturbative QCD (pQCD) framework. The linear increase of A_2 with Q disagrees with the asymptotic behaviour of T_{11}/T_{00} predicted in pQCD. The phase difference δ_{11} between the amplitudes T_{11} and T_{00} increases with Q^2 (since $\tan \delta_{11} = A_2/A_1 = bQ^2/a$) and its mean value is about 30° . This is in a sharp disagreement with the prediction of Ref [9] based on the GPD approach in pQCD.

According to pQCD calculations [10, 7], the asymptotic behaviour of T_{01}/T_{00} at high Q^2 is $\sqrt{-t'}/Q$. The best fit of the combined $p + d$ data however, is to $A_3 = c\sqrt{-t'}$, giving $c = 0.40 \pm 0.02$ GeV^{-1} with $\chi^2/N_{df} = 0.72$; fitting A_4 to $d\sqrt{-t'}/Q$ yields $d = 0.20 \pm 0.07$ with $\chi^2/N_{df} = 1.09$. The result of the fit for A_3 does not support the $1/Q$ dependence predicted by the pQCD asymptotic behaviour while the behaviour of A_4 is in accordance with it. The phase difference δ_{01} between T_{01} and T_{00} is given by $\tan \delta_{01} = d/(cQ)$ for these parameterizations; it decreases with Q and is equal to $(29^\circ \pm 9^\circ)$ at $Q^2 = 0.8$ GeV^2 . A comparison of the curves calculated with the parameters c and d with values of A_3 and A_4 in four $-t'$ bins is presented in Fig. 3.

Finally, we note that the result of the fit of the combined $p + d$ data to $|U_{11}/T_{00}| = g$, yields $g = 0.391 \pm 0.013$ with $\chi^2/N_{df} = 0.44$. This results contradicts the pQCD asymptotic behaviour $U_{11}/T_{00} \propto M_V/Q$. The absence of any t dependence means that U_{11} cannot correspond solely to single-pion-exchange.

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SEMI-INCLUSIVE DIS AND TRANSVERSE MOMENTUM DEPENDENT DISTRIBUTION STUDIES AT CLAS

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Abstract

Transverse Momentum Dependent parton distribution functions were introduced to describe both longitudinal and transverse momentum distributions of partons inside a nucleon. Great progress has been made in recent years in understanding these distributions measuring different spin asymmetries in semi-inclusive processes. Here we present an overview of the ongoing studies at CLAS and programs planned for the 6 and 12 GeV activity.

1 Introduction

The spin structure of the nucleon has been of particular interest since the EMC [1] measurements, subsequently confirmed by a number of other experiments [2–7], showed that the helicity of the constituent quarks account for only a fraction of the nucleon spin. Possible interpretations of this result include significant polarization of either the strange sea (negatively polarized) or gluons (positively polarized) and the contribution of the orbital momentum of quarks. The semi-inclusive deep inelastic scattering (SIDIS) experiments, when a hadron is detected in coincidence with the scattered lepton provide access to spin-orbit correlations. Observables are spin azimuthal asymmetries, and in particular single spin azimuthal asymmetries (SSAs), of the detected hadron, which are due to the correlation between the quark transverse momentum and the spin of the quark/nucleon.

2 Transverse Momentum Distributions

Significant progress has been made recently in understanding the role of partonic initial and final state interactions [8–10, 2]. The interaction between the active parton in the hadron and the spectators leads to gauge-invariant transverse momentum dependent (TMD) parton distributions [8, 10–13]. Furthermore, QCD factorization for semi-inclusive deep inelastic scattering at low transverse momentum in the current-fragmentation region has been established in Refs. [14, 15]. This new framework provides a rigorous basis to study the TMD parton distributions from SIDIS data using different spin-dependent and independent observables. TMDs are probability densities for finding a (polarized) parton with a longitudinal momentum fraction x and transverse momentum $\vec{k}_T$ in a (polarized) nucleon (see Tab. 1). The diagonal elements of the table are the partonic momentum, longitudinal and transverse spin distribution functions, and are the only ones surviving after

Table 1: Leading-twist transverse momentum-dependent distribution functions. U , L , and T stand for transitions of unpolarized, longitudinally polarized, and transversely polarized nucleons (rows) to corresponding quarks (columns).

N/q	U	L	T
U	$\mathbf{f}_1$		$h_1^\perp$
L		$\mathbf{g}_1$	$h_{1L}^\perp$
T	$f_{1T}^\perp$	g_{1T}	$\mathbf{h}_1$ $h_{1T}^\perp$

$\vec{k}_T$ -integration. Off-diagonal elements require non-zero orbital angular momentum and are related to the wave function overlap of $L=0$ and $L=1$ Fock states of the nucleon [16].

Similar quantities arise in the hadronization process, where a set of fragmentation functions can be introduced, describing the probability that a (polarized) parton fragment into a (polarized) hadron. For unpolarized hadrons, like for example pions and kaons, only two leading twist fragmentation functions are accessible: the usual unpolarized fragmentation function D_1 and the Collins T -odd fragmentation function $H_1^\perp$ [9] describing fragmentation of transversely polarized quarks into unpolarized hadrons.

3 TMDs measurements with CLAS at 6 GeV

Measurements of SSAs in SIDIS kinematics have been performed with CLAS at the Jefferson Laboratory using a 5.7 GeV energy electron beam. Scattering of unpolarized or longitudinally polarized electron beams off an unpolarized H_2 or longitudinally polarized NH_3 target has been studied in a wide range of kinematics. The average beam polarization was $\approx 75\%$ and the average target polarization was $\approx 70\%$. The scattered electron and charged or neutral pions were detected in the CLAS with DIS cuts $Q^2 > 1 \text{ GeV}^2$ and $W > 2 \text{ GeV}$.

3.1 Single Spin Asymmetry with longitudinally polarized target

As was first discussed by Kotzinian and Mulders in 1996 [17], the spin-orbit correlation in a longitudinally polarized nucleon gives rise to a Single Spin Asymmetry which involves the leading twist distribution function $h_{1L}^\perp$ and the Collins fragmentation function $H_1^\perp$

$$\sigma_{UL}^{\sin 2\phi} \propto S_L(1-y) \sin 2\phi \sum_{q,\bar{q}} e_q^2 h_{1L}^\perp(x) H_1^\perp(z) \quad (1)$$

where ϕ is the azimuthal angle of the hadron with respect to the lepton plane and x, y, z are the fraction of the nucleon momentum carried by the struck quark, the fraction of the electron energy carried by the virtual photon and the fraction of the virtual photon momentum carried by the detected hadron, respectively. A recent measurement of the $\sin 2\phi$ asymmetry for charged pions has been performed by Hermes [18] and is consistent with zero, as shown by the empty squares in Fig. 1. Non-zero asymmetries are predicted at large x ($x > 0.2$), a region well covered by CLAS. Indeed, the preliminary CLAS results at 6 GeV for charged pions, reported with full triangles in Fig. 1, show significant negative SSAs, while the results for neutral pions are consistent with zero.

The yellow band in Fig. 1 is the result of a calculation by Efremov *et al.* [19], using the Kotzinian-Mulders distribution function $h_{1L}^\perp$ from the chiral soliton model evolved at $Q^2 = 1.5 \text{ GeV}^2$. With the current statistical errors on the π^+ and π^0 measurements, the experimental data seem to be in agreement with the calculation, while for the π^- the model predicts positive asymmetry while the measured SSA is negative.

Higher twist contributions are expected to generate $\sin \phi$ asymmetries in the cross section. Preliminary CLAS results of the $\sin \phi$ contributions to σ_{UL} are shown in Fig. 2 as a function of the transverse momentum P_T with full squares, together with the leading twist $\sin 2\phi$ contributions (full circles).

These results indicate non-negligible higher twist contributions, of the same sign and comparable in size for π^+ and π^0 and with opposite sign for π^- .

New 6 GeV data have been taken in 2009, with the expected statistical accuracy on the measured SSA shown by the projected data points in the top part of Fig. 1. They will allow to draw much more statistically significant conclusions on the size of the asymmetry and on the possible agreement with theoretical calculations.

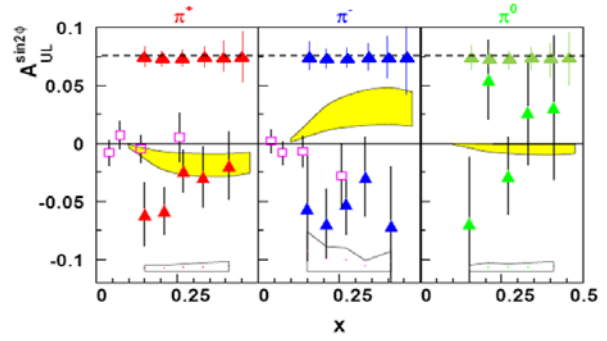


Figure 1: The target SSA measured by CLAS for pions (full triangles) compared with Hermes results (empty squares). The yellow band is the result of the calculation from [19]. The triangles on top of each plot show the CLAS projected results for the new 2009 data.

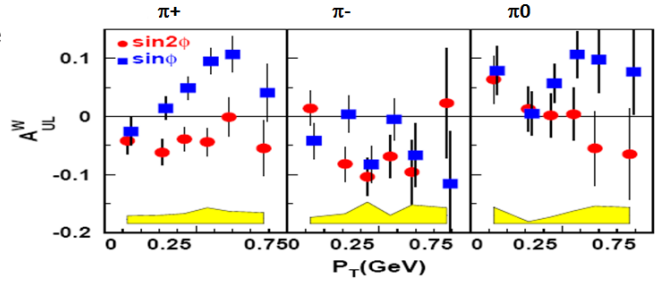


Figure 2: The $\sin 2\phi$ (full circles) and $\sin \phi$ (full squares) contributions to σ_{UL} measured by CLAS for pions.

3.2 Beam Single Spin Asymmetry with unpolarized target

Beam single spin asymmetries with an unpolarized target are zero at leading twist. Higher order terms involve either the leading twist distribution functions f_1 and $h_1^\perp$ convoluted with higher twist fragmentation functions or the leading twist unpolarized D_1 and Collins H_1 distribution functions convoluted with higher twist fragmentation functions. The resulting asymmetry contains the $\sin \phi$ modulation

$$\sigma_{LU}^{\sin 2\phi} \propto \lambda \frac{M_p}{Q} \sin \phi F_{LU} \quad (2)$$

where M_p is the proton mass and the structure function F_{LU} encodes all the relevant distribution and fragmentation functions. The preliminary CLAS results obtained at 6 GeV in three different run periods are shown in Fig. 3 as a function of x . As can be seen, positive and neutral pions have positive and comparable asymmetry. Since the π^0 Collins distribution function is expected to be small, these results seem to indicate that the Collins-type contribution to the beam SSA could be small.

3.3 Single spin asymmetry with a transversely polarized target

With unpolarized beam and a transversely polarized target, at leading twist the cross section contains several single spin asymmetry terms. The Sivers asymmetry

$$\sigma_{UT}^{\sin(\phi-\phi_S)} \propto S_T \sin(\phi - \phi_S) \sum_{q,\bar{q}} e_q^2 f_{1T}^\perp(x) D_1^\perp(z) \quad (3)$$

where ϕ_S is the azimuthal angle of the target spin with respect to the lepton plane and S_T is the target polarization, contains the Sivers function $f_{1T}^\perp$, describing unpolarized quarks in a transversely polarized target. It arises from the correlations between the transverse momentum of quarks and the transverse spin of the target and requires both orbital angular momentum, as well as non-trivial phases from the final state interaction.

The Collins asymmetry

$$\sigma_{UT}^{\sin(\phi+\phi_S)} \propto S_T(1-y) \sin(\phi + \phi_S) \sum_{q,\bar{q}} e_q^2 h_1(x) H_1^\perp(z) \quad (4)$$

contains the transversity distribution function h_1 . This distribution function, describing transversely polarized quarks in a transversely polarized nucleon, is as fundamental as f_1 and g_1 , but it is presently much less known than the other two. In fact, being charge conjugation-odd, the transversity can appear in the cross section only convoluted with another charge conjugation-odd distribution function, like the Collins function H_1 , making its extraction impossible in DIS. It does not mix

with gluons, and in the non relativistic limit it is equal to g_1 . Thus it probes the relativistic nature of quarks and it has very different Q^2 evolution than g_1 .

First measurements of the Sivers and Collins effects for pions have been obtained by Hermes [20]. The results for charged and neutral pions are shown in Fig. 4 (for Sivers) and in Fig. 5 (for Collins) by the empty squares. The measured asymmetries are in general large, however the limited range in x explored by Hermes does not allow to fully discriminate between the models. In fact, the various models ([21], [22], [23] for the Sivers effect and [21], [24] for the Collins effect) tend to give comparable results at small x and tend to diverge only for $x > 0.2$. This region can be well explored by the new CLAS experiment [25] planned for 2011 to run with a polarized HD -ice target at 6 GeV. The data points on top of each plot of Fig. 4 and 5 show the expected statistical accuracy after about two months of data taking.

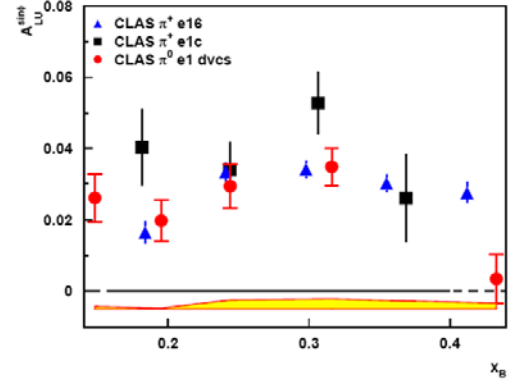


Figure 3: Beam single spin asymmetry for π^+ (squares and triangles) and π^0 (circles) as a function of x measured by CLAS.

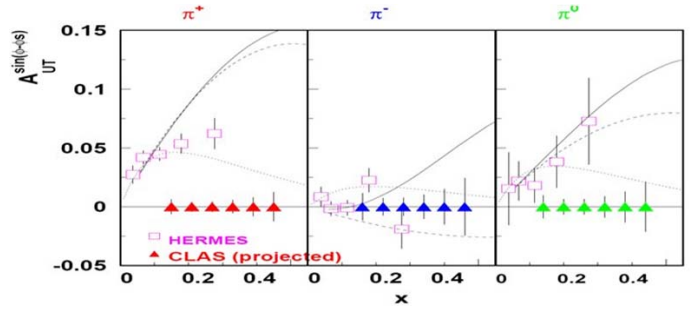


Figure 4: Sivers effect measured by Hermes as a function of x for pions (empty squares). The curves are calculated from [21] (solid and dashed) and [22], [23] (dotted). The full triangles on top of each plot shows the expected accuracy of the planned CLAS measurements at 6 GeV.

At leading twist, a third asymmetry

$$\sigma_{UT}^{\sin(3\phi-\phi_S)} \propto S_T \sin(3\phi - \phi_S) \sum_{q,\bar{q}} e_q^2 h_{1T}^\perp(x) H_1^\perp(z) \quad (5)$$

contains the distribution function $h_{1T}^\perp$, which is responsible for the non-spherical shape of the nucleon [26]. In certain class of models, like the bag model and the spectator model [27], this distribution function is related to g_1 and h_1 through the relation

$$g_1(x) - h_1(x) = h_{1T}^\perp(x) \quad (6)$$

thus giving information on the relativistic nature of the quarks in the nucleon. A first look at this new function could be obtained by the new CLAS measurements at 6 GeV, as well as for the double spin asymmetry A_{LT} with longitudinally polarized beam. The latter contains the distribution function g_{1T} , describing longitudinally polarized quarks in a transversely polarized nucleon.

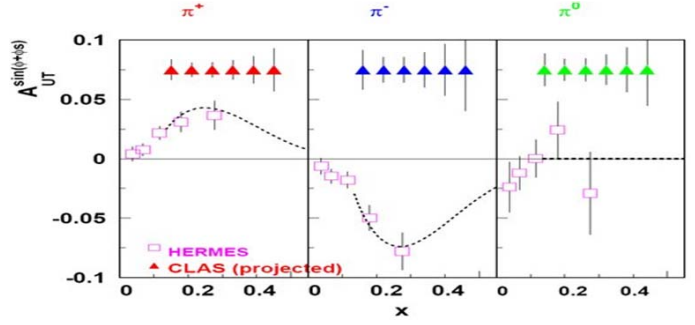


Figure 5: Collins effect measured by Hermes as a function of x for pions (empty squares). The curve is calculated for the Hermes kinematics in [21] with a k_T gaussian parameterization from [24]. The full triangles on top of each plot shows the expected accuracy of the planned CLAS measurements at 6 GeV.

4 TMDs measurements with CLAS at 12 GeV

The measurement of the transverse momentum distributions is one of the main scientific goals driving the 12 GeV upgrade of the Jefferson Laboratory [28], which is expected to start to be operational in 2015. The extended x and Q^2 range of the accessible kinematic plane, together with the expected luminosity of $10^{35} \text{ cm}^{-2} \text{ s}^{-1}$, and a new large angle detector (CLAS12) will allow the measurement of single and double spin asymmetries with at least an order of magnitude better accuracy than the present, or near to come, experimental data.

One of the key point of the new detector should be the possibility to identify semi-inclusive kaons with high rejection power from proton and pion background. In fact, kaon measurements are of fundamental importance in order to achieve flavor separation in the extraction of the distribution functions from the measured asymmetries. For this, the construction of a RICH detector has been proposed [29]. Such kind of detector would greatly improve the CLAS12 capability in terms of particle identification, allowing in particular kaon to pion separation with a rejection factor better than 1:100 in the whole range of kaon momenta above 1 GeV/c.

The CLAS12 experimental program on the TMDs includes measurements of all the eight leading twist parton distribution functions in Tab.1 with the proper choice of beam and/or target polarizations. A full list of the relevant observables and the expected accuracy can be found in the several proposals already approved by the Jefferson Laboratory PAC [29, 30].

5 Conclusions

Transverse Momentum Dependent distribution and fragmentation functions have been found as one of the main tool to access spin-orbit correlations of the partons in the nucleon. Measurements of Single (and Double) Spin Asymmetry have shown that the transverse degrees of freedom are crucial in understanding of the nucleon structure. In the last years, a significant amount of experimental data from several Laboratories has begun to come out. In this scenario, CLAS play an important role, being the only experiment having access to the large- x region, where model predictions have the biggest differences among each other.

With the foreseen large step forward in terms of luminosity and detection capability, the measurements planned to start in 2015 at the Jefferson Laboratory with the CLAS12 detector will provide information on the internal structure of the nucleon with unprecedented quality.

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THE COMPLETION OF SINGLE-SPIN ASYMMETRY MEASUREMENTS AT THE PROZA SETUP

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Abstract

Single spin asymmetry in inclusive π^0 -production was measured in the polarized target fragmentation region using 50 GeV proton beam. The asymmetry is in agreement with asymmetry measurements in the polarized beam fragmentation region carried out at higher energies. The measurement completed 30-years history of polarized measurements at the PROZA setup.

Introduction

This report is sad in some sense, because it concludes the 30 year's history of PROZA experiment. Nevertheless the new experimental program is planning at IHEP (see S. Nurushev's talk [1]). We present the recent results in inclusive π^0 production as well as the highlights of the previous polarization study at PROZA.

Asymmetry in the unpolarized beam fragmentation region

Single (left-right) spin asymmetry was measured in the reaction $\pi^- + d_{\uparrow} \rightarrow \pi^0 + X$ in the beam fragmentation region. The experimental Setup is presented in Fig. 1. γ -quanta were measured using a lead-glass electromagnetic calorimeter placed at 8 m downstream the target.

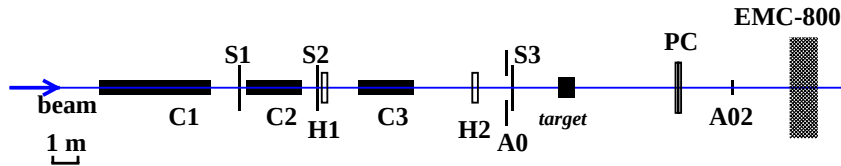


Figure 1: The Experimental Setup PROZA, $S1$ - $S3$ – trigger scintillator counters, $A0$, $A02$ –beam anti-coincidence counters, $H1$ - $H2$ – beam hodoscopes, PC –Proportional chamber, EMC -800 – Electromagnetic calorimeter

The asymmetry $A_N^{meas}(\phi)$ was calculated for each angle ϕ as the difference of the normalized numbers of pions $n_{\downarrow\uparrow}$ (equivalent to the differential cross-section) for opposite signs of target polarization:

$$A_N^{meas}(\phi) = \frac{D}{P_{targ}} \cdot A_N^{raw} = \frac{D}{P_{targ}} \cdot \frac{n_{\uparrow} - n_{\downarrow}}{n_{\uparrow} + n_{\downarrow}} \quad (1)$$

An average polarization value (P_{targ}) of fully deuterized propane-diol target was 35%, dilution factor $D=2.5$ to 5 decreasing with x_F increases. The procedure to measure carefully the dilution factor is described in detail elsewhere [2]. The asymmetry sign was selected to be consistent with all polarized beam experiments.

Final asymmetry A_N was calculated by fitting the measured asymmetry by linear function

$$A_N^{meas}(\phi) = A_0 + A_N \cdot \cos(\phi). \quad (2)$$

The last procedure allowed to eliminate systematic errors caused by beam monitor instability. Asymmetry was measured in the range of $0.6 < x_f < 1.0$ and $1.0 < p_T < 2.0$ GeV/c.

The results are presented in Fig. 2 and in Table 1.

Table 1: A_N in the reaction $\pi^- + d_{\uparrow} \rightarrow \pi^0 + X$

$x_F / p_T, \text{ GeV/c}$	1.0-1.2	1.2-1.4	1.4-1.6	1.6-2.0	1.0-2.0
0.6-0.7	-8.6 ± 8.0	2 ± 10.0	-23.0 ± 16.0	2.0 ± 25.0	-6.5 ± 5.7
0.7-0.8	17.0 ± 8.0	30.0 ± 9.5	24.0 ± 11.0	6.0 ± 16.0	21.1 ± 5.1
0.8-0.9	7.0 ± 6.6	19.0 ± 8.0	8.9 ± 8.7	13.0 ± 15.0	11.3 ± 4.2
0.9-1.0	12.4 ± 6.9	8.0 ± 7.0	10.7 ± 10.8		10.3 ± 4.5
0.7-1.0	11.5 ± 4.1	16.8 ± 4.6	13.6 ± 5.8	9.7 ± 10.9	13.6 ± 2.6

Surprisingly asymmetry is notable near the edge of phase space: $A_N = (13.6 \pm 2.6(stat) \pm 2.0(syst))\%$ at $p_T > 1$ (GeV/c) and $0.7 < x_F < 1.0$.

Non-zero asymmetry in the unpolarized beam fragmentation region can be explained in some models [4,5]. The asymmetry value in unpolarized π^- -beam fragmentation region close to the one found in the same reaction in the polarized target fragmentation region [6], and also to the neutron polarization measured in the reaction $\pi^- + p_{\uparrow} \rightarrow \pi^0 + n$ at the same transferred momentum [7].

Asymmetry was also measured in the reaction $K^- + d_{\uparrow} \rightarrow \pi^0 + X$ at $x_F > 0.7$. $A_N = (-0.4 \pm 6.1)\%$ at $p_T(\text{GeV/c}) \leq 1.2$ and $A_N = (11 \pm 6.1)\%$ at $p_T > 1.2$ (GeV/c).

Asymmetry in the polarized target fragmentation region

A_N in the polarized target fragmentation region was measured earlier at PROZA at 40 GeV [6] and 70 GeV [8] (see Fig. 3). We present an asymmetry measurement in the

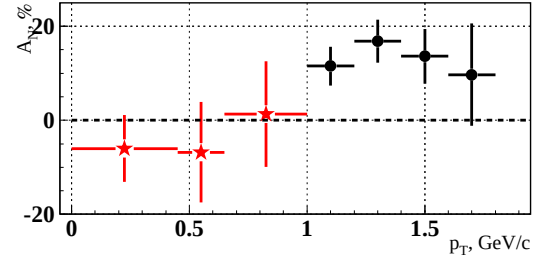


Figure 2: A_N in the reaction $\pi^- + d_{\uparrow} \rightarrow \pi^0 + X$ in the beam fragmentation region. Circles – currents measurements, stars – previous data [2,3]

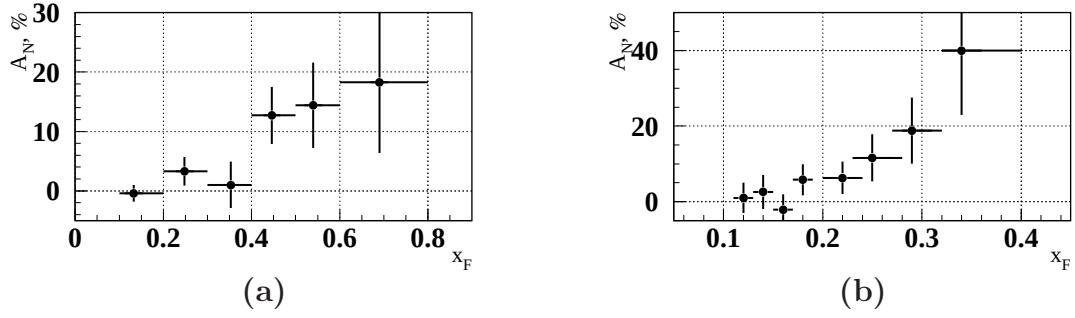


Figure 3: A_N in the polarized target fragmentation region: (a) in the reaction $\pi^- p_{\uparrow} \rightarrow \pi^0 X$ at 40 GeV [6]; (b) in the reaction $pp_{\uparrow} \rightarrow \pi^0 X$ at 70 GeV [8]. Results are presented in the polarized proton fragmentation region to be consistent with the polarized beam data.

reaction $p + p_{\uparrow} \rightarrow \pi^0 + X$ in the polarized target fragmentation region at 50 GeV. Two sets of data (2005 and 2007) are being used for analysis. The experiment was carried out at the upgraded PROZA-2M setup. The electromagnetic calorimeter was placed at 2.3 m downstream the target at 30° respect to the beam direction. The geometry was selected to detect neutral pions in the backward hemisphere. A special trigger on transverse momentum p_T allowed to enrich data with negative values of x_F . A special algorithm was developed to reconstruct γ 's which hit the detector at large angles (up to 20 degrees) [9]. The π^0 mass spectrum is presented in Fig. 4a.

A single-arm experimental setup was used. A special procedure was used to eliminate systematic errors [6]. The asymmetry was measured at $-0.6 < x_F < -0.2$. The result is presented in Fig. 4b and in Table. 2. x_F and asymmetry values are inverted to be consisted with the existing polarized beam data.

Table 2: A_N in the reaction $p + p_{\uparrow} \rightarrow \pi^0 + X$ at 50 GeV

x_F	0.11-0.13	0.13-0.17	0.17-0.21	0.21-0.25	0.25-0.35	0.35-0.45	0.45-0.60
$A_N, \%$	0.8 ± 2.2	2.0 ± 1.8	1.5 ± 1.9	2.3 ± 2.1	5.8 ± 1.8	6.8 ± 4.1	11.0 ± 5.8

The asymmetry measured at $-0.6 < x_F < -0.25$ ($6.2 \pm 1.5\%$) is in a very good

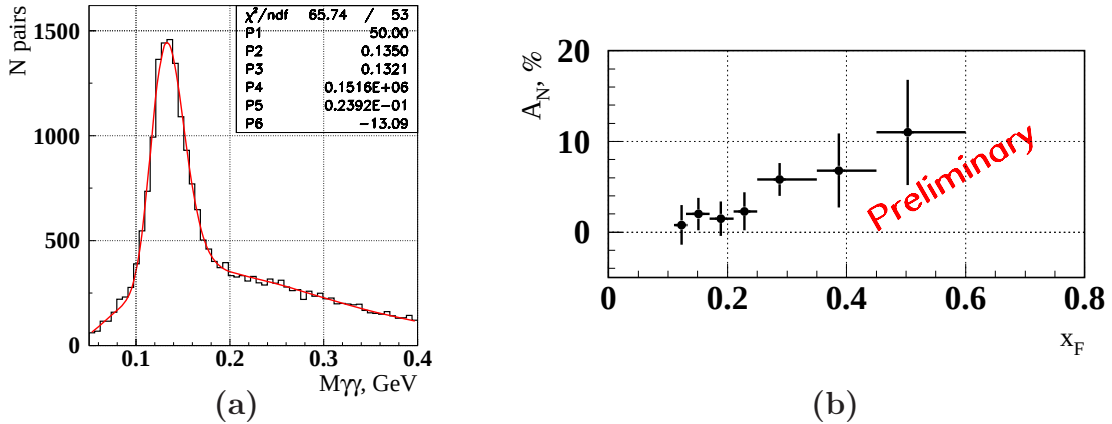


Figure 4: (a) $\gamma\gamma$ -mass spectrum; (b) A_N in the reaction $p + p_{\uparrow} \rightarrow \pi^0 + X$ at 50 GeV.

agreement with the previous PROZA data at 40 GeV ($6.9 \pm 2.8\%$) [6], the E704 data ($6.3 \pm 0.7\%$) [10] and with the STAR data [11]. Single-spin asymmetry does not depend on energy in a very wide range. Intermediate energies give us the possibility to measure asymmetry of a variety types of particles with excellent accuracy.

Highlights of the previous PROZA results

Asymmetry was measured in different exclusive charge-exchange reactions: $\pi^- p_{\uparrow} \rightarrow \pi^0(\eta, \eta'(958), \omega(783), f_2(1270))n$ at 40 GeV [12]- [15]. These results are presented in another talk [1]. Only two out of several ones results are presented in Fig. 5

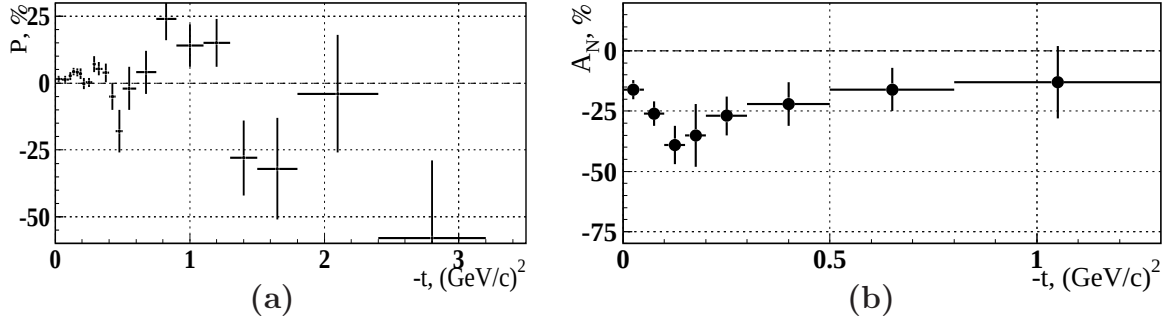


Figure 5: (a) Polarization in the reaction $\pi^- + p_{\uparrow} \rightarrow \pi^0 + n$ [12]. (b) A_N in the reaction $\pi^- + p_{\uparrow} \rightarrow f_2(1270) + n$ at 70 GeV [13].

Let me remind here the most interesting features of the observed polarization (asymmetry) behavior:

- polarization has a minimum when differential cross-section changes it's slope;
- there are oscillations in polarization behavior;
- there is an indication that asymmetry is bigger for heavier particles.

Lessons from these data will be discussed later. All these asymmetries were measured more than 25 years ago, nevertheless there are no theoretical models describing all the data together.

PROZA experiment was one of the first to measure single spin asymmetry in inclusive production. A_N in inclusive π^0 -production in the central region is presented in Fig. 6.

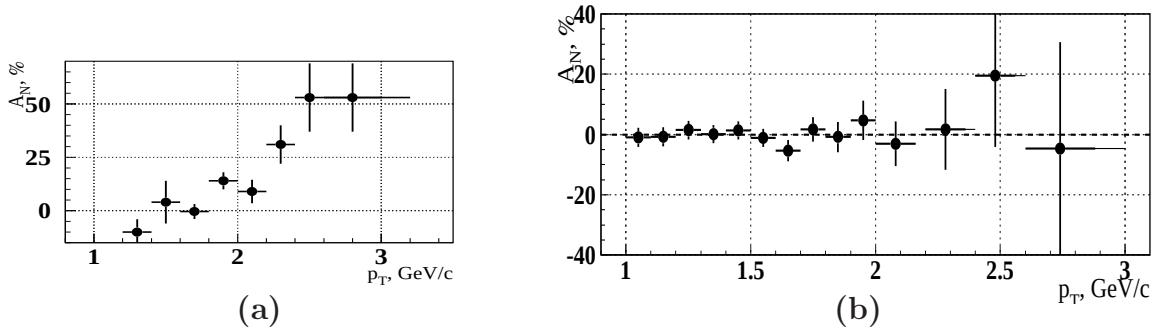


Figure 6: A_N in the central region in the reactions: (a) $\pi^- + N(p, d)_{\uparrow} \rightarrow \pi^0 + X$ at 40 GeV [16]. (b) $p + p_{\uparrow} \rightarrow \pi^0 + X$ at 70 GeV [17]. Asymmetry sign is inverted to be consistent with all polarized beam data.

Unexpectedly large single spin asymmetry in π^0 -production at π^- -beam in the central region [16] pointed out that polarization effects are beam quark flavour dependent, since A_N in $pp_{\uparrow}$ interactions is zero at the same energy.

Answers and *questions* instead of Conclusion

PROZA experiment found many interesting effects both in exclusive and inclusive reactions. Nevertheless even *more questions have to be discussed*.

Let's first summarize what we know and what we *can not explain* in exclusive reactions.

- A significant polarization (asymmetry) was found in the all exclusive reactions [12]-[15]. *Does the asymmetry magnitude increase with meson mass?*
- There is an indication on asymmetry oscillations. *Is it real effect for all particles?* Better accuracy is required.
- Polarization changes it's sign in the dip region on the $\pi^- + p_{\uparrow} \rightarrow \pi^0 + n$ differential cross-section. *Is it valid for other reactions? What is the theoretical explanation of this effect?*
- Simple Regge model can not describe polarization. Modification was required. One of the possible solution is Odderon pole in addition to ρ -pole. *There is no predictions for the most of the reactions except $\pi^- + p_{\uparrow} \rightarrow \pi^0 + n$ (see [18] for example). Another interesting prediction is that $P(\pi^0) + 2P(\eta) = P(\eta')$. It is very interesting also to measure asymmetry in $a_0(980)$ production [19]).*

Similar complicated situation is for inclusive reactions.

- Asymmetry mainly does not depend on energy (see also E704, BNL, RHIC data). *We have very good possibility to measure asymmetry in different channels at intermediate energies with good accuracy at the SPASCHARM experiment.*
- A significant asymmetry was found for u - and d - quark particles. The asymmetry is quark flavor dependent (at least for pion and proton beams). The asymmetry in the η -production is bigger than in the π^0 production (see also STAR data). *What is the asymmetry for $ss - \bar{b}$ and heavier states (ϕ and others)?.*
- Asymmetry increases with p_T at the central region in the reaction $\pi^- + p_{\uparrow} \rightarrow \pi^0 + X$ *Most of the models can not predict non-zero asymmetry in the central region and describe p_T behavior.*
- A threshold effect and a scaling was observed. Asymmetry in the non-polarized beam and in the polarized target (beam) regions close to the edge of phase space are equal each other in the reaction $\pi^- + p_{\uparrow} \rightarrow \pi^0 + X$ *It is very important to measure asymmetry in a wide kinematic region in different channels to discriminate between different models.*

We may conclude that we have found a lot of interesting spin effects. Nevertheless we all have desires, possibilities and duties trying to find much more inviting and unpredictable.

Acknowledgement

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THE GENERALIZED PARTON DISTRIBUTION PROGRAM AT JEFFERSON LAB

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Abstract

Recent results on the Generalized Parton Distribution (GPD) program at Jefferson Lab (JLab) will be presented. The emphasis will be in the Hall A program aiming at measuring Q^2 –dependences of different terms of the Deeply Virtual Compton Scattering (DVCS) cross section. This is a fundamental step before one can extract GPD information from JLab DVCS data. The upcoming program in Hall A, using both a 6 GeV beam (2010) and a 11 GeV beam ($\approx$ 2015) will also be described.

1 Introduction

Understanding of the fundamental structure of matter requires an understanding of how quarks and gluons are assembled to form the hadrons.

The asymptotic freedom of Quantum Chromodynamics (QCD) makes it possible to use perturbation theory to treat interactions of quarks and gluons at short distances. In high-energy scattering, short-distance and long-distance physics may be separated to leading power in momentum transfer—an approach known as “factorization”. For example, the cross sections for hard electron-proton or proton-proton collisions that transfer large momentum can be expressed as a product of a short-distance partonic (quark or gluon) cross section, calculable in perturbative QCD, and parton distribution functions that encode the long-distance information on the structure of the proton.

Traditionally, the spatial view of nucleon structure provided by lattice QCD and the momentum-based view provided by the parton picture stood in stark contrast. Recent theoretical developments have clarified the connection between these two views, that of momentum and that of spatial coordinates, such that ultimately it should be possible to provide a complete “space-momentum” map of the proton’s internal landscape. These maps, referred to as Generalized Parton Distributions (GPDs), describe how the spatial shape of a nucleon changes when probing different ranges of quark momentum. Projected along one dimension, GPDs reproduce the form factors; along another dimension they provide a momentum distribution. With enough information about the correlation between space and momentum distributions, one can construct a full “tomographic” image of the proton. The weighted integrals, or moments, of the GPDs contain information about the forces acting on the quarks bound inside the nucleon. These moments can now be computed using lattice QCD methods. Not only do they have an attractive physical interpretation, but they can also be compared directly to unambiguous predictions

from the underlying theory. Higher-order QCD corrections for many observables sensitive to GPDs have now also been calculated, strengthening the theoretical underpinning for extraction of GPDs.

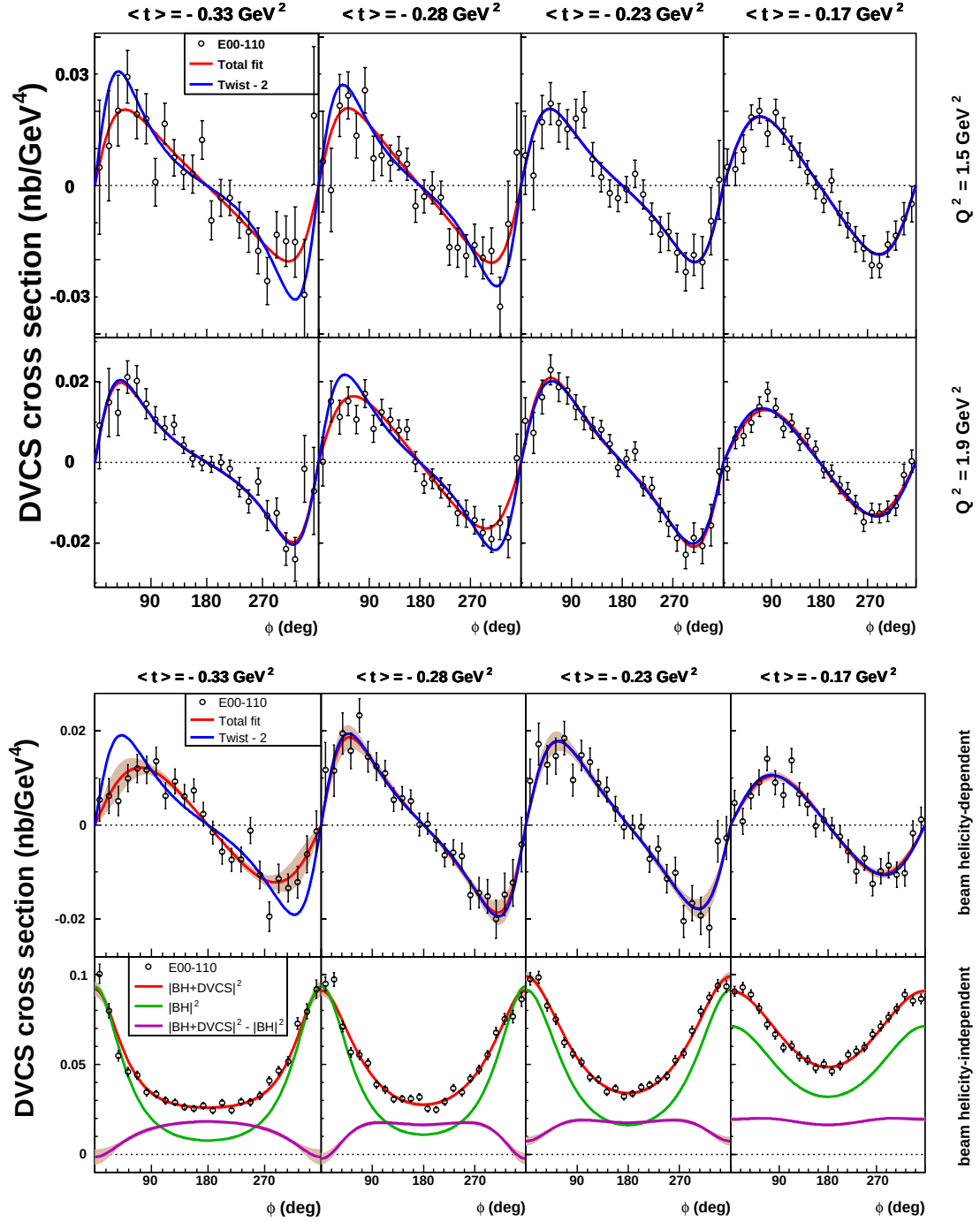


Figure 1: Data and fit to the helicity-dependent $d^4\Sigma/dQ^2 dx_B dt d\phi$, and the helicity-independent $d^4\sigma/[dQ^2 dx_B dt d\phi]$ cross sections, as a function of the azimuthal angle ϕ between the leptonic and hadronic planes. All bins are at $\langle x_B \rangle = 0.36$. Error bars show statistical uncertainties. Solid lines show total fits with one- σ statistical error bands. The green line is the $|\text{BH}|^2$ contribution to $d^4\sigma$. The first three rows show the helicity-dependent cross section for values of $Q^2 = 1.5, 1.9$ and 2.3 GeV^2 , respectively. The last row shows the helicity-independent cross section for $Q^2 = 2.3 \text{ GeV}^2$.

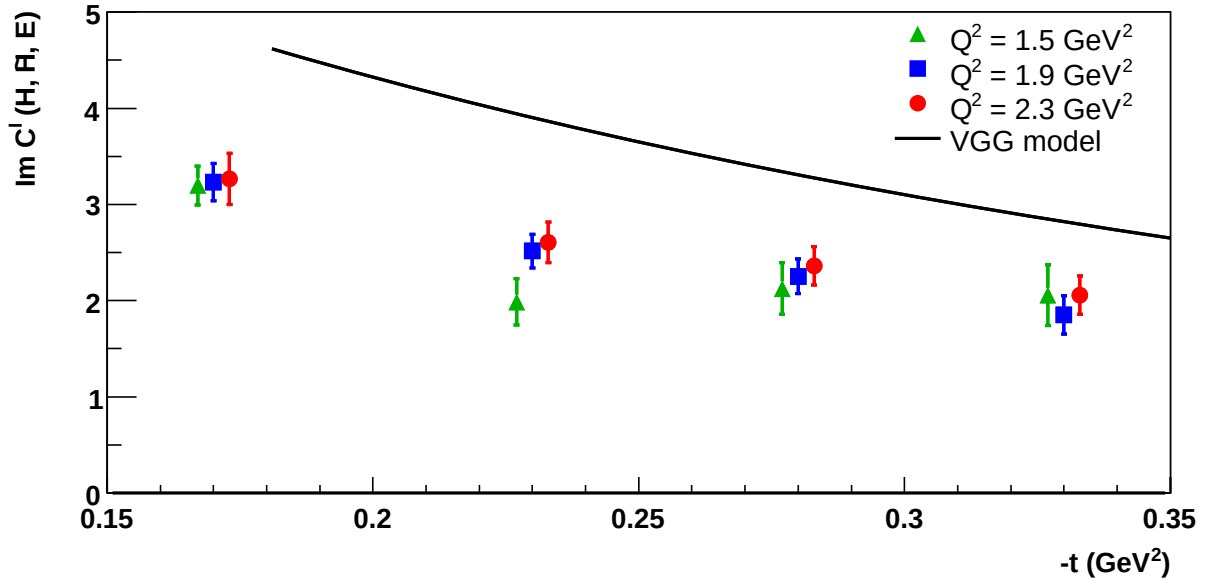


Figure 2: Extracted imaginary part of the twist-2 angular harmonic as function of t . Superposed points are offset for visual clarity. Their error bars show statistical uncertainties.

The simplest and cleanest experimental process that gives access to GPDs is Deeply Virtual Compton Scattering (or DVCS), where a virtual photon of high 4-momentum Q^2 scatters off a nucleon producing a real photon on the final state ($\gamma^* p \rightarrow \gamma p$). Experimentally, DVCS occurs in the electroproduction of photons off the nucleon: $ep \rightarrow ep\gamma$, where there is a competing channel, the Bethe-Heitler (BH) process, where the photon of the final state is radiated by the electron and not the proton. The BH process is, however, calculable in Quantum Electrodynamics given our knowledge of the proton form factors.

Extracting GPDs from DVCS requires the fundamental demonstration that DVCS is well described by the leading twist mechanism at finite Q^2 . This leading twist mechanism consists on the scattering on a single quark (or gluon) inside the nucleon.

2 Current experimental situation

The H1 [1, 2] and ZEUS [3] collaborations measured the cross section for $x_B \approx 10^{-3}$. The HERMES collaboration measured relative beam-helicity [4] and beam-charge asymmetries [5, 6]. Relative beam-helicity [7] and longitudinal target [8] asymmetries were measured at the Thomas Jefferson National Accelerator Facility (JLab) by the CLAS collaboration. The first dedicated DVCS experiment ran in Hall A at JLab [9]. Helicity-correlated cross section of electroproduction of photons were measured with high statistical accuracy as a function of Q^2 (Fig. 1).

The combination of Compton form factors (CFF) extracted from these data (Fig. 2) was found to be independent of Q^2 , within uncertainties. This exciting result strongly supports the twist-2 dominance of the imaginary part of the DVCS amplitude. GPDs can then be accessed experimentally at the kinematics of JLab.

A dedicated DVCS experiment on the neutron (E03-106) followed E00-110. Using a deuterium target, E03-106 provided the first measurements of DVCS on the neutron,

particularly sensitive to the least known GPD E , and DVCS on the deuteron [10]. A first constraint on the orbital angular momentum of quarks inside the nucleon (Fig. 3) was obtained relying on the VGG model of GPDs [11].

3 Future DVCS program with CEBAF at 6 GeV

A new experiment E07-007 was recently approved with the highest scientific rating (A) by the JLab Program Advisory Committee (PAC-31) and will run next year. E07-007 will use the important lessons we learned in E00-110 to further test and explore the potential of DVCS to measure GPDs.

In E00-110 we managed to determine the photon electroproduction helicity-independent (unpolarized) cross section. This additional information was not in the initial goal of the E00-110 experiment, and we were only able to measure it at one single value of $Q^2 = 2.3 \text{ GeV}^2$ (Fig. 1). We found that the total cross section was much larger than the BH contribution. The excess in the cross section is then coming from both the interference term (BH–DVCS) and the DVCS² contribution. These contributions contain different kinds of GPD combinations. The interference term is proportional to a *linear* combination of GPD integrals, whereas the DVCS² is related to a *bilinear* combination of GPD integrals.

The harmonic ϕ structure of the DVCS cross section does *not* allow the independent determination of the BH–DVCS interference and the DVCS² contributions. Indeed, the interference $\mathcal{I}$ and DVCS² terms have the following harmonic structure (with $Q = \sqrt{Q^2}$):

$$\begin{aligned} \mathcal{I} &= \frac{i_0/Q^2 + i_1 \cos \varphi / Q + i_2 \cos 2\varphi / Q^2 + i_3 \cos 3\varphi / Q}{\mathcal{P}_1 \mathcal{P}_2} \\ \text{DVCS}^2 &= d_0/Q^2 + d_1 \cos \varphi / Q^3 + d_2 \cos 2\varphi / Q^4. \end{aligned} \quad (1)$$

The product of the BH propagators reads:

$$\mathcal{P}_1 \mathcal{P}_2 = 1 + \frac{p_1}{Q} \cos \varphi + \frac{p_2}{Q^2} \cos 2\varphi. \quad (2)$$

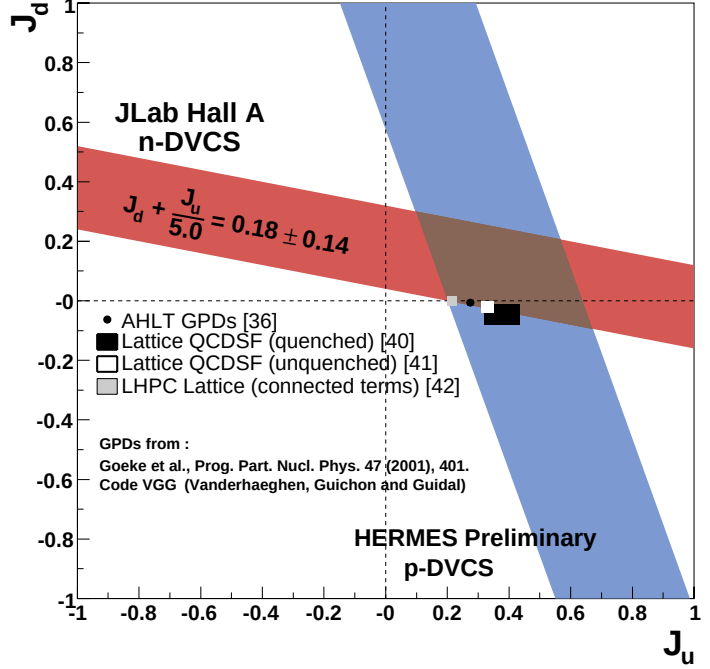


Figure 3: Experimental constraint on J_u and J_d quark angular momenta from the present n-DVCS results [10]. A similar constraint from the p-DVCS target spin asymmetry measured by HERMES, and different lattice QCD based calculations are also shown.

Reducing to a common denominator ($\times \mathcal{P}_1 \mathcal{P}_2$), one obtains:

$$\begin{aligned}
\mathcal{P}_1 \mathcal{P}_2 \mathcal{I} + \mathcal{P}_1 \mathcal{P}_2 \text{DVCS}^2 = & \boxed{(i_0 + d_0)/Q^2} + d_1 p_1/2/Q^4 + p_2 d_2/2/Q^6 \\
& + [i_1/Q + (p_1 d_0 + d_1)/Q^3 + (p_1 d_2 + p_2 d_1)/2/Q^5] \cos \varphi \\
& + [i_2/Q^2 + (p_2 d_0 + p_1 d_1/2 + d_2)/Q^4] \cos 2\varphi \\
& + [i_3/Q + (p_1 d_2 + p_2 d_1)/2/Q^5] \cos 3\varphi \\
& + [p_2 d_2/4/Q^6] \cos 4\varphi.
\end{aligned} \tag{3}$$

One sees in Eq. (3) that the interference $\mathcal{I}$ and the DVCS^2 terms *mix at leading order in $1/Q$* in the azimuthal expansion.

3.1 Interference and DVCS^2 separation

E07-007 was approved to measure the photon electroproduction helicity-independent cross section as a function of Q^2 and at different incident beam energies. The different beam energy dependence of the DVCS^2 and $\text{BH} \cdot \text{DVCS}$ interference term complement the azimuthal analysis of the cross section in order to provide a complete separation of these two contribution.

3.2 Separation of the longitudinal π^0 electroproduction cross section

In addition to these fundamental measurements on DVCS, E07-007 will concurrently measure the π^0 electroproduction cross section for different beam energies. We successfully measured the π^0 electroproduction cross section in E00-110. However, since we only ran at one incident beam energy, the longitudinal and transverse separation of the cross section could not be done. E07-007 will allow the first separation of the longitudinal cross section, using the Rosenbluth technique. A factorization theorem holds for the longitudinal π^0 electroproduction cross section that relates it, at sufficiently large Q^2 , to a different flavor combination of GPDs. This channel is an essential element to make a flavor decomposition of GPDs, which is impossible to obtain from DVCS on the proton alone.

E07-007 will test the Q^2 -dependence of the longitudinal cross section and check if the factorization is applicable at the Q^2 values currently accessible experimentally (around $Q^2 \sim 2 \text{ GeV}^2$). If this test turns out positive, very interesting and complementary information on GPDs will be available from this channel.

3.3 E08-025: Rosenbluth-like separation on the neutron

An equivalent DVCS program at 6 GeV using a deuterium target was approved earlier this year and will run concurrently with E07-007. It will allow to separate the interference and DVCS^2 terms on neutron observables. This will complement E07-007 on the proton.

4 12 GeV program

The experimental GPD program is at the heart of the scientific motivation of the major Jefferson Lab upgrade to 12 GeV.

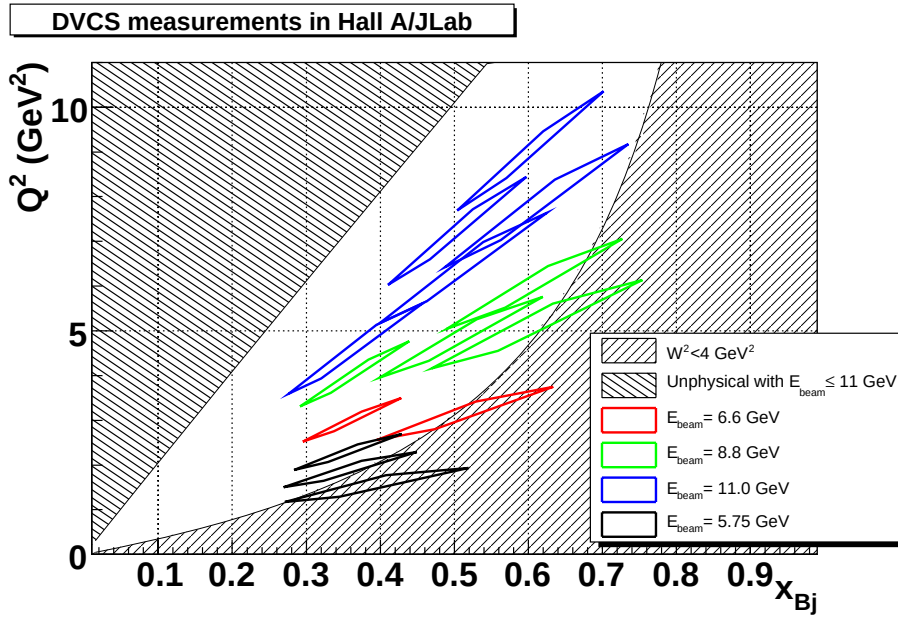


Figure 4: Approved DVCS kinematics for $H(e, e'\gamma)p$ measurements in Hall A with 3, 4, and 5 pass beams of CEBAF at 12 GeV. The diamond shapes trace the approximate acceptance of the Hall A High Resolution Spectrometer in each setting. The boundary of the unphysical region corresponds to the maximum possible Q^2 at a given x_B for 11 GeV. This corresponds to 180° electron scattering, equivalent to $\theta_q = 0^\circ$. The points at $E_{Beam} = 5.75$ GeV were obtained in E00-110.

E12-06-114 [12] is a wide experimental program of DVCS high precision cross-section measurements in Hall A with CEBAF at 12 GeV, enthusiastically approved by JLab PAC-30. With polarized 6.6, 8.8, and 11 GeV beams incident on the liquid hydrogen target, we will provide a major survey of proton DVCS in almost the entire kinematic range accessible with CEBAF at 12 GeV (Fig. 4).

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LAMBDA PHYSICS AT HERMES

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Abstract

Results for Λ and $\bar{\Lambda}$ hyperon polarization measured by the HERMES experiment using the 27.6 GeV longitudinally polarized positron beam incident on polarized or unpolarized gas targets are reported. The longitudinal spin transfer from the beam to the hyperon D_{LL}^{Λ} has been measured at $Q^2 > 0.8 \text{ GeV}^2$ predominantly at positive x_F and values of $D_{LL}^{\Lambda} = 0.102 \pm 0.056(stat) \pm 0.020(syst)$ for Λ and $D_{LL}^{\Lambda} = 0.152 \pm 0.139(stat) \pm 0.030(syst)$ for $\bar{\Lambda}$ have been obtained. The spin transfer K_{LL}^{Λ} from the longitudinally polarized target to the hyperon has been measured in the quasi-real photoproduction regime ($Q^2 \approx 0 \text{ GeV}^2$) with an average photon energy of $\simeq 15 \text{ GeV}$, resulting in $K_{LL}^{\Lambda} = 0.024 \pm 0.008(stat) \pm 0.003(syst)$ for Λ and $K_{LL}^{\Lambda} = 0.002 \pm 0.019(stat) \pm 0.008(syst)$ for $\bar{\Lambda}$, respectively. The spontaneous transverse Λ and $\bar{\Lambda}$ polarization (not related to beam or target polarization) has also been studied in a wide range of nuclear targets.

The HERMES experiment offers a good opportunity to study Λ polarization both in semi-inclusive, $eN \rightarrow e'\bar{\Lambda}X$, and inclusive, $eN \rightarrow \bar{\Lambda}X$ reactions, using the 27.6 GeV polarized positron beam of the HERA accelerator incident on internal polarized or unpolarized gas targets. The HERMES detector, described in detail in [1], is a forward magnetic spectrometer with angular acceptance $\pm(40\text{--}140)$ mrad in the vertical direction and ± 170 mrad in the horizontal direction. During the data taking period the typical beam polarization was 0.4-0.5. The beam helicity was reversed about once per month.

HERMES used polarized hydrogen and deuterium targets with typical polarization of 85% and unpolarized targets ^1H , ^2D , ^3He , ^4He , ^{14}N , ^{20}Ne , ^{84}Kr and ^{132}Xe in a wide range of atomic numbers.

The Λ hyperon polarization P^{Λ} was measured through its $p\pi^-$ decay channel. In this parity-violating decay, the angular distribution of the proton has a form

$$\frac{dN}{d\Omega_p} = \frac{dN_0}{d\Omega_p} (1 + \alpha P^{\Lambda} \cdot \cos\theta_p), \quad (1)$$

where θ_p is the angle between the proton momentum and the direction of the Λ polarization in the Λ rest frame, and $\alpha = 0.642 \pm 0.013$ is the analyzing power of the weak decay. The symbols $dN/d\Omega_p$ and $dN_0/d\Omega_p$ denote the distributions for the decay of polarized and unpolarized Λ samples, respectively.

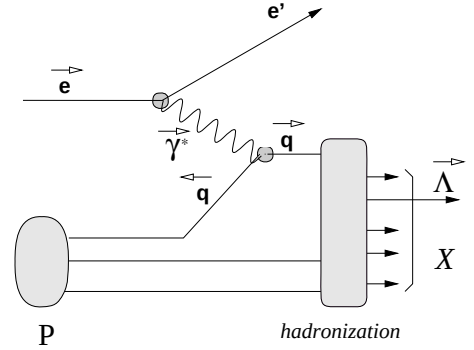


Figure 1: The single-quark scattering mechanism leading to Λ production in the current fragmentation region of polarized deep inelastic electron scattering.

According to Eq. 1 the protons are preferentially emitted along the spin direction of their parent Λ which provides an opportunity to measure Λ polarization by measuring the "forward-backward" asymmetry of the decay products in the Λ rest frame. Λ events were selected by requiring the presence of at least two hadron candidates of opposite charge. In the event selection the fact has been used that at HERMES kinematics the detected proton (or anti-proton) is always the leading particle. The combinatorial background has been mainly suppressed using this kinematical criterion and a restriction imposed on the distance between the primary (Λ production) and secondary (Λ decay) vertices.

The spin transfer from the longitudinally polarized beam to the Λ hyperon was studied in semi-inclusive deep-inelastic scattering (DIS) $eN \rightarrow e'\Lambda X$ (Fig.1), where the scattered positrons were detected in coincidence with Λ events. The requirements $Q^2 > 0.8 \text{ GeV}^2$ and $W > 2 \text{ GeV}$, where $-Q^2 = q^2$ is the four-momentum transfer squared of the exchanged virtual photon and W is the invariant mass of the photon-nucleon system, were imposed on the positron kinematics to ensure that the events originated from the DIS domain. In addition, the requirement $y = 1 - E'/E < 0.85$ was imposed to exclude the large contribution of radiative corrections, where $E(E')$ is the energy of the positron before (after) the scattering process.

The component of the polarization transferred along the L' direction from the virtual photon to the produced Λ hyperon is given by $P_{L'}^\Lambda = P_b D(y) D_{LL'}^\Lambda$, where P_b is the longitudinal beam polarization, $D(y) \simeq [1 - (1 - y)^2]/[1 + (1 - y)^2]$ is the depolarization factor and L is the primary quantization axis, directed along the virtual photon momentum, and it is assumed that the spin of the virtual photon is directed along its momentum.

The spin transfer coefficient $D_{LL'}^\Lambda$ describes the probability that the polarization of the virtual photon is transferred via the struck quark to the Λ hyperon along a secondary quantization axis L' . In this analysis, the quantization axis L' is chosen along the virtual photon direction L , i.e. $L' = L$. It is important to remember that the direction L is taken in the Λ rest frame.

The extraction of D_{LL}^Λ from the data is based on the moments method applied to a beam helicity balanced data set [2]. The data have also been analyzed with the maximum likelihood method. The results has been found to be very close to those obtained by the moments method.

In order to estimate the systematic uncertainty of the obtained results an identical

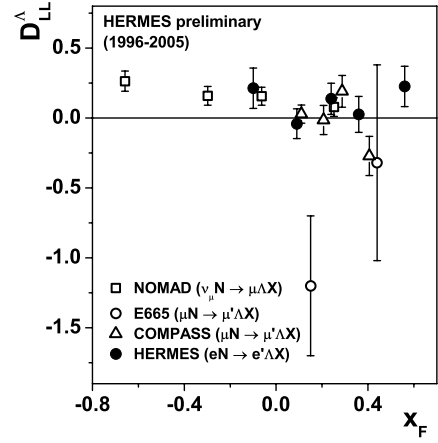


Figure 2: Compilation of the world data on the longitudinal spin-transfer coefficient D_{LL}^Λ on x_F .

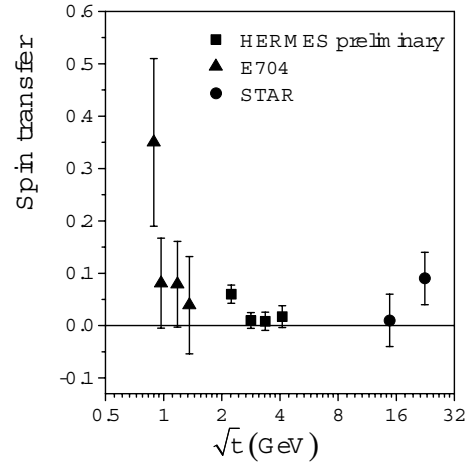


Figure 3: Compilation of the world data on spin transfer from a polarized nucleon to the Λ (for HERMES K_{LL}^Λ) vs $\sqrt{t}$ variable.

analysis was carried out reconstructing h^+h^- hadron pairs produced in invariant-mass windows below and above the Λ ($\bar{\Lambda}$) mass peak. As an additional check of a possible false polarization, the decay $K_s^0 \rightarrow \pi^+\pi^-$ was studied. These analyzes resulted in $D_{LL}^{h^+h^-} = -0.002 \pm 0.011$ and $D_{LL}^{K_s^0} = -0.016 \pm 0.035$.

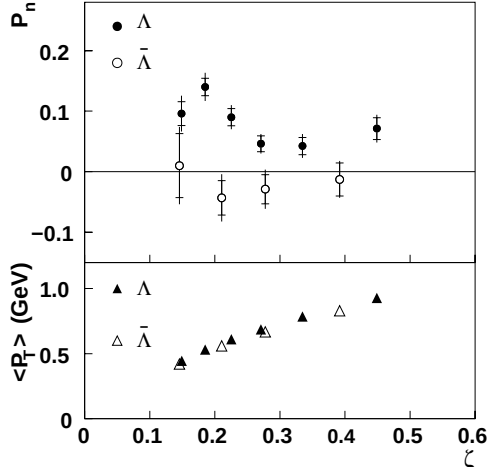


Figure 4: Transverse polarizations P_Λ and $P_{\bar{\Lambda}}$ (upper panel) and mean $\langle p_T \rangle$ (lower panel) as functions of $\zeta = (E_\Lambda + p_{z\Lambda})/(E_e + p_e)$.

similar to the D_{LL}^Λ analysis [7]. The kinematical dependence of K_{LL}^Λ as a function of $\sqrt{t} = \sqrt{-(p_\Lambda - p_p)^2}$, where p_Λ (p_p) is the four-momentum of the Λ hyperon (target nucleon), is presented in Fig. 3 together with results obtained by the E704 ($p^\dagger p \rightarrow \bar{\Lambda} X$, $p_{beam} = 200$ GeV/c) [8] and STAR ($\bar{p}p \rightarrow \bar{\Lambda} X$, $\sqrt{s} = 200$ GeV) [9] collaborations. The HERMES data confirm the trend of K_{LL}^Λ increasing with decreasing $\sqrt{t}$ previously observed by the E704 collaboration. Averaged over kinematics the HERMES result for the spin transfer is $K_{LL}^\Lambda = 0.024 \pm 0.008(stat) \pm 0.003(syst)$ for the Λ hyperon, while for the $\bar{\Lambda}$ hyperon it is compatible with zero: $K_{LL}^{\bar{\Lambda}} = 0.002 \pm 0.019(stat) \pm 0.008(syst)$.

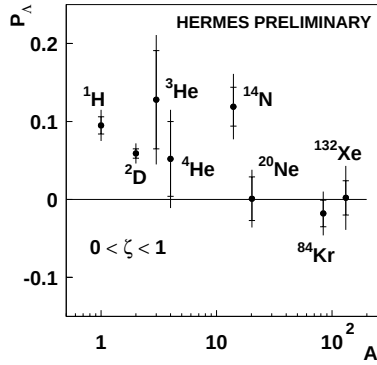


Figure 5: Transverse Λ polarization vs atomic number of nuclei.

The formalism of the extraction of transverse Λ polarization is based upon the up/down symmetry of the HERMES detector and the moments method [2, 18]. Averaged over the experimental kinematics, the net Λ polarization extracted from the data collected in the years 1996-2005 with various targets was found to be significantly positive while the net

The HERMES data for D_{LL}^Λ [3] are presented in Fig. 2 as a function of x_F together with data of the NOMAD experiment at CERN [4], the E665 experiment at Fermilab [5] and the COMPASS experiment at CERN [6]. As seen from Fig. 2, the NOMAD, COMPASS and HERMES results are compatible in the kinematic region of overlap $-0.1 < x_F < 0.3$. The HERMES results, averaged over kinematics are $D_{LL}^\Lambda = 0.102 \pm 0.056(stat) \pm 0.020(syst)$ for Λ and $D_{LL}^{\bar{\Lambda}} = 0.152 \pm 0.139(stat) \pm 0.030(syst)$ for $\bar{\Lambda}$, respectively.

The spin transfer from the longitudinally polarized ^1H and ^2D targets to the hyperon, K_{LL}^Λ , was measured in the quasi-real photoproduction regime. The K_{LL}^Λ extraction procedure is similar to the D_{LL}^Λ analysis [7].

The kinematical dependence of K_{LL}^Λ as a function of $\sqrt{t} = \sqrt{-(p_\Lambda - p_p)^2}$, where p_Λ (p_p) is the four-momentum of the Λ hyperon (target nucleon), is presented in Fig. 3 together with results obtained by the E704 ($p^\dagger p \rightarrow \bar{\Lambda} X$, $p_{beam} = 200$ GeV/c) [8] and STAR ($\bar{p}p \rightarrow \bar{\Lambda} X$, $\sqrt{s} = 200$ GeV) [9] collaborations. The HERMES data confirm the trend of K_{LL}^Λ increasing with decreasing $\sqrt{t}$ previously observed by the E704 collaboration. Averaged over kinematics the HERMES result for the spin transfer is $K_{LL}^\Lambda = 0.024 \pm 0.008(stat) \pm 0.003(syst)$ for the Λ hyperon, while for the $\bar{\Lambda}$ hyperon it is compatible with zero: $K_{LL}^{\bar{\Lambda}} = 0.002 \pm 0.019(stat) \pm 0.008(syst)$.

Transverse Λ polarization was observed and investigated in many high-energy scattering experiments, with a wide variety of hadron beams and kinematic settings [10, 11]. Because of the parity-conserving nature of the electromagnetic and the strong interaction, any final-state hadron (hyperon) polarization in an inclusive reaction with unpolarized beam and target can only point along the normal $\hat{n} = \frac{\vec{p}_e \times \vec{p}_\Lambda}{|\vec{p}_e \times \vec{p}_\Lambda|}$ of the production plane, where $\vec{p}_e$ is the (positron) beam momentum and $\vec{p}_\Lambda$ is the Λ momentum, respectively. Possible mechanisms for the origin of this polarization were discussed for example in Refs. [12], [13], [14] [15] [16], [17].

$\bar{\Lambda}$ polarization was consistent with zero: $P_{\Lambda} = 0.078 \pm 0.006(stat) \pm 0.012(syst)$ and $P_{\bar{\Lambda}} = -0.025 \pm 0.015(stat) \pm 0.018(syst)$ [18].

The wide variety of target nuclei used until the end of data taking in 2007 allowed to study also the A-dependence of the transverse Λ polarization. The resulting A-dependence is presented in Fig. 5. There is a clear indication of an A-dependence of the spontaneous Λ polarization: the polarization for light nuclei ($^1H, ^2D$) is statistically significant positive, while for heavier nuclei ($^{84}Kr, ^{132}Xe$) the polarization is compatible with zero within the statistical uncertainties of the data.

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SPIN PHYSICS AT NICA

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Abstract

Recently JINR launched the project on constructing the new collider NICA which will operate both in heavy ion collisions mode and in polarized proton-proton collisions mode. The aspects of the physical program of the latter option (SPD/NICA) are considered. The preliminary estimations on rates and possible layout of the spectrometer are presented.

Since the famous “spin crisis“ in 1987, the problem of the nucleon spin structure remains one of the most intriguing puzzle of high energy physics. The central component of this problem attracting for many years enormous both theoretical and experimental efforts, is a search for answering the questions, how the spin of the proton is build up from spins and orbital momenta of its constituents. The searches brought up a concept of the parton distribution functions in nucleon, at the beginning two of them, one, f_1 , for unpolarized and second, g_1 , for polarized nucleons. Now we know that must be about 50 different parton distributions functions for a complete description of the nucleon structure. While today a part of the polarized distributions can be considered as sufficiently well known, there is a number of PDFs which either are absolutely unknown, or poorly known, especially the spin dependent ones, These are longitudinally polarized distributions of valence light sea, strange quarks and gluons, both sea and valence transversely polarized distributions of all flavours. This new class of PDFs, is characterized by its the non-trivial dependence on a transverse quark momentum. The most significant among them are Sivers and Boer-Mulders PDFs. The studies of these open questions of the nucleon structure is the first priority task for the scientific program of the second IP the NICA facility. The parameters of the NICA collider allows to perform the important spin and polarization effects studies of:

- *Drell-Yan (DY) processes with longitudinally and transversely polarized p and d beams. Extraction of unknown (poor known) parton distribution functions (PDFs);*
- *PDFs from J/ψ production processes;*
- *Spin effects in baryon, meson and photon productions; Effects in various exclusive reactions;*
- *Diffraction processes;*
- *Cross sections, helicity amplitudes and double spin asymmetries (Krisch effect) in elastic reactions;*
- *Spectroscopy of quarkoniums.*

To determine Boer-Mulders and Sivers PDFs at NICA the following measurements must be performed:

- **Unpolarized and single polarized Drell-Yan (DY) processes with pp and pd collisions;**

Table 1: Status of the DY experiments

Experiment	Status	Remarks
CERN: E615	Finished	Fixed target mode. Unpolarized target and beam
CERN: NA10,38,50	Finished	Fixed target mode. Unpolarized target and beam
FERMILAB: E886,906	Running	Fixed target mode. Unpolarized target and beam
RHIC: STAR,PHENIX	Running up to ?	Collider mode. Detector upgrade for DY studies polarized beams
GSI: PANDA	Plan > 2016	Fixed target mode. Unpolarized
GSI: PAX	Plan > 2016	Collider mode. Polarized beams.
		Problem with anti-proton beam polarization
CERN: COMPASS	Plan > 2011	Fixed target mode. Polarized beam. Valence PDFs
J-PARC	Plan > 2011	Fixed target mode. Unpolarized beam and target, $s \sim 60 - 100 GeV^2$
SPASCHARM	Plan ?	Fixed target mode. Unpolarized beam and target, $s \sim 140 GeV^2$
SPD-NICA	Plan > 2014	Collider mode. $s \sim 670 GeV^2$, high luminosity, polarized proton and deuteron beams, energy scan.

Table 2: The estimation of the cross-sections and number of DY events for NICA and PAX kinematics

	σ_{DY} total, nb	L, $cm^{-2}s^{-1}$	Thousands events
PAX, $\sqrt{s} = 14.6 GeV$	~ 2	$\sim 10^30$	~ 10
NICA, $\sqrt{s} = 20 GeV$	~ 1	$\sim 10^30$	~ 5
PAX, $\sqrt{s} = 26 GeV$	~ 1.3	$\sim 10^30$	~ 7

– J/ψ production processes with unpolarized and single polarized pp and pd collisions, which can not be completely duplicated by other experiments (COMPASS [1], RHIC [2], PAX [3] and J-PARC [4], as it is seen from Table 1.

First, at COMPASS and J-PARC the fixed target mode with unpolarized beam and polarized target is only possible. As it will be shown later, for this option the transversity and Boer-Mulders PDFs are hardly accessible for pp and pd collisions. COMPASS can reach the transversity via π^-p collisions. However, in this case the unknown pion PDFs will be involved, introducing the additional theoretical uncertainty [5].

Second, PAX plans to access transversity via measurements of double polarized DY processes, which will very difficult because of the antiproton polarization problems [3].

Third, the RHIC experiments cover the kinematical region different from the respective region covered by NICA.

Fourth, the advantage of NICA in comparison with all above mentioned experiments is that only NICA can provide an access to transversity, Boer-Mulders and Sivers PDFs for both u and d flavors.

The Table 2 shows the estimation of the cross-sections and possible statistics for DY events, which can be collected with SPD at NICA during one month in comparison with those from the PAX proposal.

To estimate the possible precision of measurements of the asymmetries the set of

original software packages (MC simulation, generator etc.) was developed. The SSA asymmetries, which can be measured with SPD NICA detector are estimated to be about 5-10% and the simulations shows that the such asymmetries are be visible within the errors at the statistics about 100K events, which corresponds to two years of data taking.

The measurements of J/ψ production processes are also very important for tests of duality model and due to unique possibility to contribute the data in unpolarized J/ψ production. The data on both unpolarized and polarized J/ψ production obtained with SPD NICA can essentially improve the theoretical models of unpolarized J/ψ production. The possibility of NICA facility to change the energy beams allows to scan different kinematical regions and to obtain the unique information.

The preliminary design of (SPD) detector for spin effects studies is based on the requirements imposed by the DY and J/ψ productions studies. These requirements are the following: almost 4π geometry for secondary particles; precise vertex detector; precise tracking system ; precise momentum measurement of secondary particles; good particle identification capabilities (μ , π , p , e , etc.) high trigger rate capabilities. The most of these requirements are also good for other studies mentioned above. Basing on these requirements several possible scheme of SPD are considered, one of them is similar to the detector of PAX experiment (close to NICA in kinematics) [3] at FAIR GSI, the second one is the SPD of limited possibilities, providing the muon pair detection only and the third is the scheme of SPD based on so-called Muon Range System, which is considered as detector for PANDA muon system. This scheme is shown in Fig. 1 The main parts of this scheme is described below.

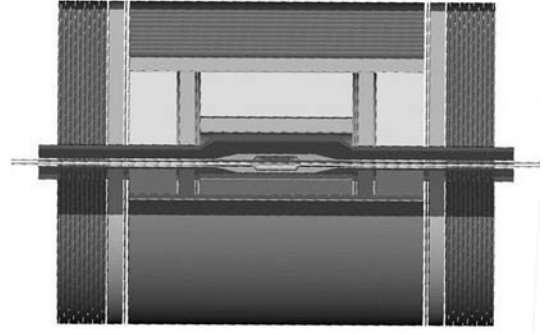


Figure 1: The design of SPD with Range System shown at the edge of the detector. The magnet system is shown in red, EC- in yellow.

The toroid magnet of the spectrometer provides a field free region around the interaction point and does not disturb the trajectories. The toroid magnet can consist of 8 superconducting coils symmetrically placed around the beam axis. A support ring upstream of the target hosts the supply lines for electric power and for liquid helium. At the downstream end, an hexagonal plate compensates the magnetic forces to hold the coils in place. The field lines of ideal toroid magnet are always perpendicular to the path of the particles originating from the beam line. Since the field intensity increases inversely proportional to the radial distance: greater bending power is available for particles scattered at smaller angles, which have higher momenta. These properties help to design a compact spectrometer that keeps the investment costs for the detector tolerable. The production of such a field requires the insertion of the coils into the tracking volume shadowing part of the azimuthal acceptance. Preliminary studies show that the use of superconducting coils, made by a Nb3Sn-Copper core surrounded by a winding of aluminum for support and cooling, allows one to reach an azimuthal detector acceptance in excess of 85%, while the radius of the inner magnet volume can be about 0.3 m and outer - about 0.7m, with $\int B dL \sim 0.8 - 1 Tm$. Several layers of double-sided Silicon strips can provide a precise vertex reconstruction and tracking of the particles before they reach the magnet. The design should use a small number of silicon layers to minimize the radiation length of the

tracking material. With a pitch of 50-100 μm it is possible to reach an spatial resolution of 20-30 μm . Such a spatial resolution would be provide 50-80 μm for precision of the vertex reconstruction, and permits to reject the secondary decays of mesons into leptons.

The coordinate resolution of 150-200 μ can be achieved with conventional drift chambers. The chambers can be assembled as modules consisting of several pairs of tracking planes with wires at -30 ; 0 ; 0 ; +30 deg. with respect to the direction parallel to the magnetic field lines. This can provide the momentum resolution of the order of 1-3 % over the kinematic range of the detector. The calorimeter can consists of “shashlyk” modules with the application of new readout technics based on AMPD technology working in the magnetic field. The modules can have an area of 4x4cm² and a length of 30-40 cm. The expected energy resolution can be $\sigma(E)/E = (5 - 8)\%/\sqrt{E}$. The calorimeter also can be used for the triggering of DY electrons. Sets of hodoscope planes are used for triggering. To improve the lepton identification, the passive Pb radiator (about 2 radiation lengths) can be placed in front of the external hodoscope to initiate the electromagnetic showers. The system of mini-drift layers with Fe layers called by Range System (RS). It can provide the clean ($\sim 99\%$) muon identification for muon momenta grater than 1 GeV. The combination of responses from EM calorimeter and RS can be used for the identification of pions and protons in the wide energy range.

The final version of the SPD will be defined after detailed Monte-Carlo simulations and consideration of requirements for other spin effects studies. The purpose is to have the simple universal detector.

SPD-NICA project is under preparation at 2nd interaction point of NICA collider. The purpose of this experiment is the study of the nucleon spin structure with high intensity polarized light nuclear beams using the following planned features of NICA collider: very high collision proton (deuteron) energy up to $\sqrt{s} \sim 26(12)GeV$; the average luminosity up to $10^{30}(10^{29})cm^2/s$ both proton and deuteron beams can be effectively polarized, with the polarization degree not less than 50%.

We welcome to the collaboration on SPD-NICA project.

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MEASUREMENTS OF TRANSVERSE SPIN EFFECTS IN THE FORWARD REGION WITH THE STAR DETECTOR

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Abstract

Measurements by the STAR collaboration of the cross section and transverse single spin asymmetry (SSA) of neutral pion production at large Feynman x (x_F) in pp-collisions at $\sqrt{s} = 200$ GeV were reported previously. The x_F dependence of the asymmetry can be described by phenomenological models that include the Sivers effect, Collins effect or higher twist contributions in the initial and final states. Discriminating between the Sivers and Collins effects requires one to go beyond inclusive π^0 measurements. For the 2008 run, forward calorimetry at STAR was significantly extended. The large acceptance of the Forward Meson Spectrometer (FMS) allows us to look at heavier meson states and $\pi^0 - \pi^0$ correlations. Recent results, the status of current analyzes and near-term plans will be discussed.

Contrary to simple perturbative QCD (pQCD) predictions, measurements of inclusive pion production in polarized proton-proton collisions at center-of-mass energies ($\sqrt{s}$) up to 20 GeV found large transverse single spin asymmetries (A_N) [1, 2]. Measurements with the Solenoidal Tracker at RHIC (STAR) detector confirmed that the π^0 asymmetry survives at $\sqrt{s} = 200$ GeV and grows with increasing Feynman x ($x_F = 2p_L/\sqrt{s}$) [3, 4]. Similar large spin effects have recently been found in electron-positron and semi-inclusive deep inelastic scattering [7, 8]. Significant developments in theory in the past few years suggest common origins for these effects, but the large A_N in inclusive pion production is not yet fully understood.

A number of phenomenological models extend pQCD by introducing parton intrinsic transverse momentum (k_T) and considering correlations between parton k_T and proton spin in the initial state (Sivers effect [5]) or final state interactions of a transversely polarized quark fragmenting into a pion (Collins effect [6]). These models can explain the observed x_F dependence of A_N , but their expectations of decreasing asymmetry with increasing pion transverse momentum (p_T) is not confirmed by the experimental data [4]. According to the current theoretical understanding both the Sivers and Collins mechanisms contribute to π^0 A_N and discriminating between the two should be possible by measuring the asymmetry in direct photon production or forward jet fragmentation.

Hadron production in the forward region in pp collisions probes large- x quark on low- x gluon interactions and is a natural place to study spin effects. First measurements of that type at RHIC were done with the STAR Forward Pion Detector (FPD), a modular electromagnetic calorimeter placed at pseudorapidity $\eta \sim 3-4$ [3]. Prior to the 2008 run, forward calorimetry at STAR was extended with the FMS which replaced the FPD modules on one side (west) of the STAR interaction region. The FMS is a matrix of 1264 lead glass cells. It provides full azimuthal coverage in the region $2.5 < \eta < 4.0$ and

has ~ 20 times larger acceptance than the FPD. In addition to the study of inclusive π^0 production, this allows us to reconstruct heavier mesons and to look at “jet-like” events and $\pi^0 - \pi^0$ correlations.

The 2008 run at RHIC with transversely polarized proton beams measured an integrated luminosity of $\sim 7.8 \text{ pb}^{-1}$ at average beam polarization $P_{beam} \sim 45\%$. The first step in the analysis of FMS data was to look at the inclusive π^0 asymmetry to make a point of contact with prior FPD measurements. $A_N(x_F)$ was found to be comparable to the previous results [9]. The 2π coverage in azimuthal angle (ϕ) made possible a study of the $\cos\phi$ dependence of the asymmetry, which is well described by a linear function, as expected [9]. The FMS also allowed us to extend the measured $\pi^0 p_T$ region to $\sim 6 \text{ GeV}/c$ and added new data to investigate the p_T dependence of A_N [10].

Another use of the large FMS acceptance was a reconstruction of forward neutral pion pairs from pp collisions. This analysis followed a few steps:

- All photons in the FMS were reconstructed and a list was made of those satisfying the condition that the cluster energy was above 2 GeV and a fiducial volume requirement.
- π^0 reconstruction considered all possible two-photon combinations; a pion was identified if the di-photon invariant mass was between 0.05 and 0.25 GeV/c^2 . The “leading” pion was required to have transverse momentum $p_{T,L} > 2.5 \text{ GeV}/c$. The “subleading” pion was formed from the remaining photons and satisfied the requirement $1.5 \text{ GeV}/c < p_{T,S} < p_{T,L}$.
- The difference in azimuthal angle between the leading and subleading pions ($\Delta\phi$) was calculated and the event distribution $dN/d\Delta\phi$, normalized by the number of leading pions, was plotted.

The invariant mass for both pions is presented in Fig. 1a. Coincidence probability as a function of $\Delta\phi$ is shown in Fig. 1b. The distribution has two clear peaks corresponding

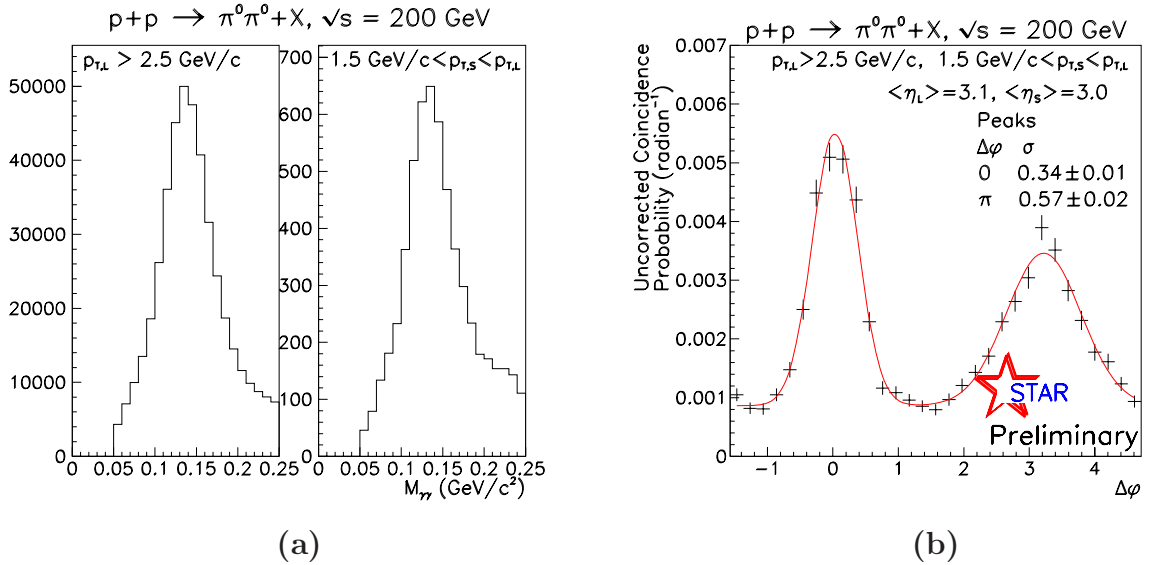


Figure 1: (a) Invariant mass distributions for leading (left) and subleading (right) neutral pions. (b) Azimuthal correlations between two neutral pions in the FMS.

to “near-side” and “away-side” correlations, and is fitted by two Gaussian functions plus a constant background. The correlations are not yet normalized. Simulations to obtain efficiency corrections are under way.

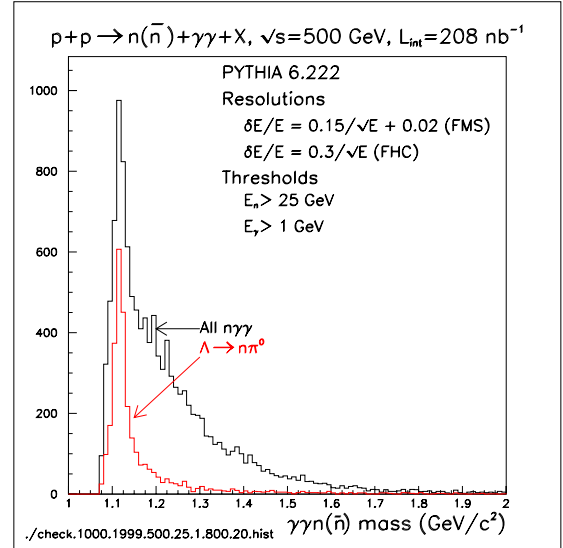
The above analysis can be extended to look at the correlations sorted by the spin state of the colliding protons. Back-to-back ($\Delta\phi \approx \pi$) di-hadron measurements can provide access to the Sivers function, as suggested in [11], assuming that the neutral pions serve as jet surrogates. Near-side hadron correlations can provide sensitivity to the Collins fragmentation function and transversity.

Near-term plans for forward physics at STAR include measurements of the cross section and transverse SSA in inclusive π^0 production in polarized proton-proton collisions at $\sqrt{s} = 500$ GeV and a proposal to add a Forward Hadron Calorimeter (FHC) behind the FMS.

The 2009 run at RHIC was the first physics run at $\sqrt{s} = 500$ GeV. STAR sampled 10 pb^{-1} of data with longitudinally polarized beams at this energy. A first look at the FPD data shows that with the lead-glass matrices alone, π^0 events can be reconstructed up to $x_F \sim 0.25$. Each FPD module also includes a two-plane scintillation shower maximum detector that provides essential data to separate photons from decays of pions at higher x_F . Measurements with transversely polarized beams at $\sqrt{s} = 500$ GeV are tentatively planned for the 2011 run.

The proposed addition of the FHC (two matrices of 9×12 lead-scintillator detectors) to the STAR detector is motivated by the following physics goals:

- to measure the transverse SSA for full jets that should allow to isolate the contribution from the Sivers mechanism to the observed π^0 asymmetry;
- measurements of polarization transfer coefficients through Λ polarization in the $\Lambda \rightarrow n\pi^0$ channel to test pQCD predictions.



PYTHIA simulation of $pp \rightarrow n(\bar{n}) + 2\gamma + X$ events in the FMS+FHC have been done using a fast event generator method. Reconstructed mass for all $n\gamma\gamma$ events and for the $\Lambda \rightarrow n\pi^0$ process is shown in Fig. 2.

Figure 2: Three-cluster mass distributions from the simulation of $pp \rightarrow n(\bar{n}) + 2\gamma + X$ events.

In summary, precision measurements of $\pi^0 A_N$ with the FPD allow for a quantitative comparison with theoretical models. The FMS allows us to go beyond inclusive π^0 production to heavier mesons, “jet-like” events and particle correlation studies. Measurements of A_N for direct photons or jets are needed to disentangle the dynamical origins. Prospects for the more distant future include the development of a RHIC experiment to measure transverse SSA for Drell-Yan production of dilepton pairs.

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THE FIRST STAGE OF POLARIZATION PROGRAM SPASCHARM AT THE ACCELERATOR U-70 OF IHEP

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Abstract

The first stage of the proposed polarization program SPASCHARM includes the measurements of the single-spin asymmetry (SSA) in exclusive and inclusive reactions with production of stable hadrons and the light meson and baryon resonances. In this study we foresee of using the variety of the unpolarized beams (pions, kaons, protons and antiprotons) in the energy range of 30-60 GeV. The polarized proton and deuteron targets will be used for revealing the flavor and isotopic spin dependencies of the polarization phenomena. The neutral and charged particles in the final state will be detected.

Introduction

In shaping the new polarization program at U-70 we were guided by three conditions: by our own experiences. by theoretical status of subject and the reliability of the new program. As concerns of the first condition one may refer on the comparative study of polarizations in the elastic scattering of particles and antiparticles by using the polarized proton target [1], [2], measurements of the spin transfer tensor [3], [4] (HERA Collaboration), study of the SSA in the exclusive and inclusive charge exchange reactions at 40 GeV/c [5] (PROZA Collaboration), study of polarization effects at 200 GeV/c by using the polarized proton and antiproton beams (E581/E704 Collaboration, FNAL) [6]. The fourth example of polarization data came recently from the STAR Collaboration at energy in the center of mass $\sqrt{s}=200$ GeV [7], [8].

The second condition is the status of the relevant theoretical models. Since the energies and transfer momenta with which we are dealing are not large enough, so there is a doubt about the possible application of the perturbative quantum chromodynamics p(QCD). Therefore either the specific models should be used for the interpretation of the experimental data or the general asymptotic predictions might be applied.

The third condition is relevant to the reliability of the experiment, that is, availability of robust equipments, manpower, money and other resources.

Below we shall briefly describe all these conditions.

1 The preceding study of polarization effects at U70

In 1970-1976 at U70 Collaboration of physicists from Saclay (France), Protvino, Dubna and Moscow (Russia) (HERA Collaboration) had performed the measurements of the polarization parameters P and R (spin rotation parameter) in elastic scattering of particles and antiparticles at ~ 40 GeV/c by using the polarized proton target. The polarization data are presented in Figure 1 with one panel for pair of particle and antiparticle.

In papers [9] and [10] it was considered some consequences of the hypothesis of the approximate γ_5 invariance of the strong interactions. According to this hypothesis at high energies and large momentum transfers $s, -t \gg m^2$ (m is the mass of the particles involved in reactions) the polarization in any elastic scattering of particles or antiparticles should be equal to zero. From Figure 1 the following results stem out of:

1. polarizations are not zero in reactions induced by pions, protons and antiprotons. It means that the hypothesis of γ_5 invariance does not work for that reactions yet,
2. the polarizations are zero for reactions induced by kaons. It means that the hypothesis of γ_5 invariance may work in these reactions. But the large error bars in the measured polarizations make this statement doubtful. The future experiments measuring the polarizations in kaon induced elastic scattering with better statistics are needed.

For all of above reactions the next comment follows. Though the asymptotic regime was reached for s it's not fulfilled for t , since for $|t| > 1(\text{GeV}/c)^2$ the experimental errors are large. This is the next item for the future measurements with high statistics. There is a good measurement of the polarization parameter in $\pi^\pm p, K^\pm p, pp, \bar{p}p$ elastic scattering at 6 GeV/c [11]. In this case the γ_5 invariance does not work too for all reactions, exception is $\bar{p}p$, where polarization is compatible with zero in small $-t$ region in frame of the large error bars.

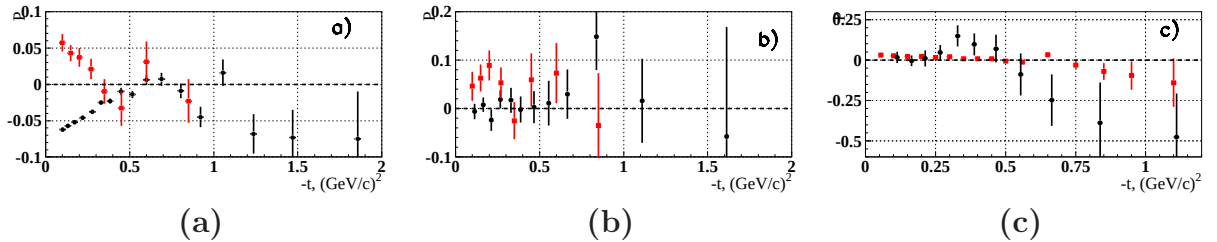


Figure 1: (a) P in elastic scattering: $\pi^- p(\bullet)$ and $\pi^+ p(\blacksquare)$. (b) P in elastic scattering: $K^- p(\bullet)$ and $K^+ p(\blacksquare)$. (c) P in elastic scattering: $\bar{p}p(\bullet)$ and $pp(\blacksquare)$.

In paper [12] the study was made of the asymptotic relations between polarizations in cross channels of a reaction. Using the crossing symmetry and Fragman - Lindeloff theorem they arrived at the following result: polarizations in the elastic scattering processes (see Figure 1) induced by particle and antiparticle at a given energy and a given angle should be equal in magnitude and opposite in sign. If we look at Figure 1 one may note that this statement is approximately correct for the pion induced reactions, not correct for reactions initiated by proton and antiproton and ambiguous for reactions involving kaons (thanks to the small statistics). Therefore the new elastic scattering experiment should clarify this interesting problem by gathering a large statistics, specially at large transverse momenta.

The HERA Collaboration making use of the simple Regge pole model concluded that the elastic scattering polarizations induced by pions and kaons follow the predictions of such model, while polarization in elastic pp scattering reveals the drastic deviation from the prediction of the Regge pole model. Such behavior may be explained by assuming that at 40 GeV/c momentum the dominant contribution to the polarization in elastic pp scattering comes from the pomeron with the spin flip term of the order of 10% with respect to the spin non flip term. Involving in the analysis the data on the spin rotation parameter [3], [4] they strengthened their conclusion. But the statistics are not so large to be unambiguous in such conclusion. One needs more measurements.

The PROZA Collaboration measured the single spin asymmetries in the charge exchange binary and inclusive reactions at the incident beam momentum 40 GeV/c [5]. With the different statistics the results were obtained for the exclusive reactions containing in the final states the mesons of the different mass and quantum numbers Figure 2 , Figure 3.

The polarization data for reactions 1, 2 and 3 (the mesons in the final states are spinless) were extensively analyzed in frame of the different models. For example, in the asymptotic model [12] the polarizations in all of these three reactions should be zero. But such predictions are in contrast to the experimental data (see Figure 2). In the Regge pole model with inclusion of the odderon [13], [14], the best approximation predicting the new dip in polarization around the crossing point at $-t \sim 0.2(\text{GeV}/c)^2$ was obtained in the model [13]. After analyzing the reaction (Figure 2a) the authors of the paper [14] noted:

The surprising results of the recent 40 GeV/c Serpukhov measurement of the polarization in $\pi^- p \rightarrow \pi^0 n$ are shown to support the conjecture that the crossing-odd amplitude may grow asymptotically as fast as is permitted by general principles.

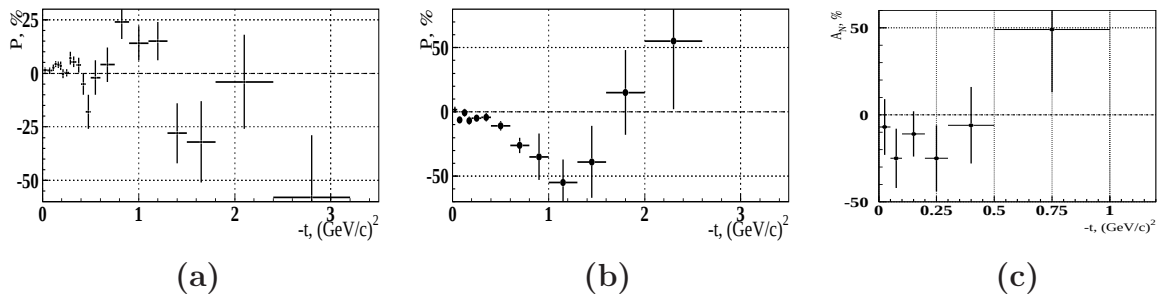


Figure 2: (a) Polarization in reaction $\pi^- + p \rightarrow \pi^0 + n$. (b) Polarization in reaction $\pi^- + p \rightarrow \eta + n$. (c) Polarization in reaction $\pi^- + p \rightarrow \eta' + n$.

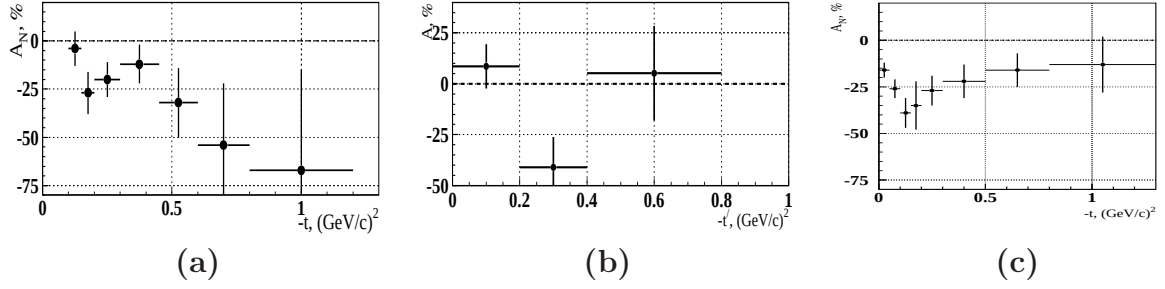


Figure 3: (a) asymmetry in reaction $\pi^- + p \rightarrow \omega + n$. (b) asymmetry in reaction $\pi^- + p \rightarrow a_2 + n$. (c) asymmetry in reaction $\pi^- + p \rightarrow f + n$.

This model in contrast to other ones predicts the shift of the left zero crossing point farther to left and the increase of polarization with growth of the incident momentum. This is an attractive subject for experimental check.

The reactions 2, 3 were also analyzed in frame of the Regge pole model and data are consistent with model prediction. Other reactions 4-6 showing also the significant spin effects (see Figure 3) did not yet attract the attention of theoreticians.

These data are unique in the sense that nobody made (30 years later) the similar measurements at higher energies. This fact may confirm the assumption that the U-70 accelerator occupies a good niche for such studies of exclusive reactions dying rapidly with growth of energy.

By using the experimental set-up PROZA the inclusive asymmetries were measured at 40 GeV/c in the following charge exchange reactions:

$$\pi^- + p \rightarrow \pi^0 + X \quad (1), p + p \rightarrow \pi^0 + X \quad (2)$$

in the central, polarized target and unpolarized beam fragmentation regions [5]. In central region the asymmetry about 30% was found at $p_T > 2$ GeV/c, while in the beam fragmentation region it was around 10-15% at $p_T > 1$ GeV/c in reaction (1). In contrary the analyzing power for reaction (2) is almost zero at the same regions. These are the puzzling results of the (SSA) measurements in the inclusive charge exchange reaction (1) at the incident momentum of 40 GeV/c (PROZA). There is no any independent experimental confirmation of these results.

2 First stage of the SPASCHARM polarization program

In composing the new scientific program we are guided by the recent theoretical and experimental developments in polarization physics. It's obvious also that this program is also strongly influenced by our own experiences, by resources and competitions with other collaborations over the world. Therefore we attempt of using efficiently our proton synchrotron U70, existing experimental equipments and fit to the environmental requirements. So we are going to propose the following first stage polarization program:

1. Asymmetry measurements in charge exchange exclusive reactions at 34 GeV/c with

emphasis to increase the statistics of the most of reactions shown in the Figures 2 and 3 by approximately by factor 10 and move to the larger $\sqrt{s}$ region.

2. Comparative studies of asymmetries induced by particles and antiparticles in binary and inclusive reactions.
3. Study of spin transfer mechanism by using the unstable spin carrying particles like hyperons, vector mesons, etc. We emphasize, that only fixed target experiments, like ours, may measure spin transfer tensors for stable final particles, like antiprotons and protons.
4. Asymmetry measurements in inclusive productions of various stable hadrons containing partons of different flavors (u,d,s,c quarks).
5. The systematic studies of the isospin dependence of single spin asymmetry.
6. The comparative studies of asymmetries in production of particles and antiparticles in final state.
7. Asymmetry studies by using the light ion beams and polarized target.
8. Check more accurately the puzzle caused by differences of the single spin asymmetries induced by pion and proton beams in the central and fragmentation regions at 34 GeV/c.
9. The new upcoming polarized proton beam will lead to the measurements of the inclusive single spin asymmetries with unprecedented precisions.
10. Finally with the availability of the polarized beam and polarized target the way will be opened for the intense studies of the double spin asymmetries in many reactions.

For the inclusive reactions the extensive Monte Carlo simulations were made for beam particles π^- , K^- , p and $\bar{p}$ for momentum of 34 GeV/c. For the sake of brevity we present, as an example, only the results of calculations for the $\bar{p}$ beam. According to the negative beam composition the fraction of the $\bar{p}$ particles is only 0.3% at momentum of 34 GeV/c. Therefore the absolute flux of $\bar{p}$ beam is only $9 * 10^3 \bar{p}/cycle$. The yields of the secondary particles on this beam are very important for comparison to the yields of the same particles in the proton beam. For detection of the secondary resonances with the rare decay modes produced on the propandiole target the request was imposed: the energy deposit in calorimeters should be > 2 GeV. Knowing the $\bar{p}$ interaction cross section with target the yields of the secondary particles with higher cross section for 30 days beam run are presented in the next Table 1.

Two comments are in order to this Table 1. First, the number of events for polarized protons are less than indicated in Table 1 by one order of magnitude. Second on the level of several percents the asymmetries produced by protons and antiprotons may be compared if the statistics are larger than 10^5 . It means that such comparisons may be done for sure for pions, kaons, baryons, antibaryons, but doubtful for Xi^- . The estimates were done for ideal apparatus and not taking into accounts the real backgrounds.

Table 1: The estimated yields N_{EV} of the secondary particles from the propandiole target ($C_3H_8O_2$, 20 cm long) stricken by the $\bar{p}$ beam of 34 GeV/c. One month beam run was assumed (3.610^8 interactions). B/S means the background to signal ratio.

#	particle	N_{EV}	*	#	particle	N_{EV}	B/S
1	π^+	2.1×10^8	*	7	n	1.6×10^7	*
2	π^-	2.6×10^8	*	8	$\bar{n}$	1.4×10^8	*
3	K^+	1.7×10^7	*	9	$\Lambda \rightarrow \bar{p} + \pi^+$	2.1×10^6	0.1
4	K^-	2.2×10^7	*	10	$\Lambda \rightarrow \bar{n} + \pi^0$	1.1×10^6	8.0
5	p	1.6×10^7	*	11	$\Delta^{--} \rightarrow \bar{p} + \pi^-$	4.2×10^7	7.0
6	$\bar{p}$	1.8×10^8	*	12	$\Xi^- \rightarrow \Lambda + \pi^-$	1.0×10^5	0.1

3 The experimental apparatus

The experimental apparatus for the SPASCHARM program consists of the following elements:

1. Beam apparatus consisting of the scintillation and Cherenkov counters, scintillation hodoscopes for detecting and identifying the beam particles (not shown in Figure 4).
2. The polarized proton (deuteron) target (target in Figure 4).
3. The guard system surrounding the polarized target (PT).
4. The polarization building-up and holding magnet (target magnet).
5. GEM1, GEM2
6. The large aperture magnetic spectrometer.
7. Micro drift chambers (MDC).
8. Two large aperture multichannel threshold Cherenkov counters for identifications of the secondary particles.
9. Multiwire proportional chambers (MWPC).
10. Electromagnetic calorimeter (ECAL).
11. Hadron calorimeter (HCAL).
12. Muon system.
13. Scintillation hodoscopes.

The layout of the SPASCHARM detectors is presented in Figure 4. The main elements of the apparatus, their structure, positions and sizes are listed in the Table 2.

Conclusions

The SPASCHARM program presents the natural extension of our previous polarization experiments. Proposed experiment will open new and wider perspectives due to the several reasons. First it contains the magnetic spectrometer with the high resolution tracking detectors allowing to register all secondary charged particles with precise angular and momentum resolutions. Secondly it has the fast particle identification system allowing to reconstruct the resonances with high probability. Third, the electromagnetic and hadronic

Table 2: The parameters of the main elements of the experimental apparatus. L-distance from detector to polarized target (PT), DS-detector structure, WS-wire spacing, GS-gross size, N_{ch} -number of channels

detector	L, m	DS	WS	GS	N_{ch}
GEM1	0.5	X, Y	Strip 0.4	20 x 20	1000
GEM2	0.75	X, Y	Strip 0.4	30 x 30	1500
1.MDC	1.0	X, X', Y, Y', U, V'	6	65 x 54	1200
2.MDC	1.5	X, X', Y, Y', U, V'	6	111 x 81	1920
3.MDC	2.0	X, X', Y, Y', U, V'	6	150 x 111	2610
4.MWPC	3.5	X, Y, U, V	2	150 x 100	2500
5.MWPC	6.5	X, Y, U, V	2	150 x 100	2500

calorimeters practically allow (together with threshold Cherenkov counters) to identify all hadrons in final state having sufficiently large cross sections. The large angular and momentum acceptances will finally allow to increase by a factor of 10 the statistics than in previous PROZA experiments. Using the forward detectors with the guard counters around the polarized target one can select the binary charge exchange reactions with one order better statistics and also detect new reactions. The detections of hyperons and vector mesons permit to study not only polarization but also the spin transfer mechanism in strong interaction. It is assumed that the full apparatus for the first stage of the SPASCHARM polarization program will be ready to 2013 beam run.

The distinct feature of our program will be the comparative studies of the polarization phenomena induced by the particles and antiparticles.

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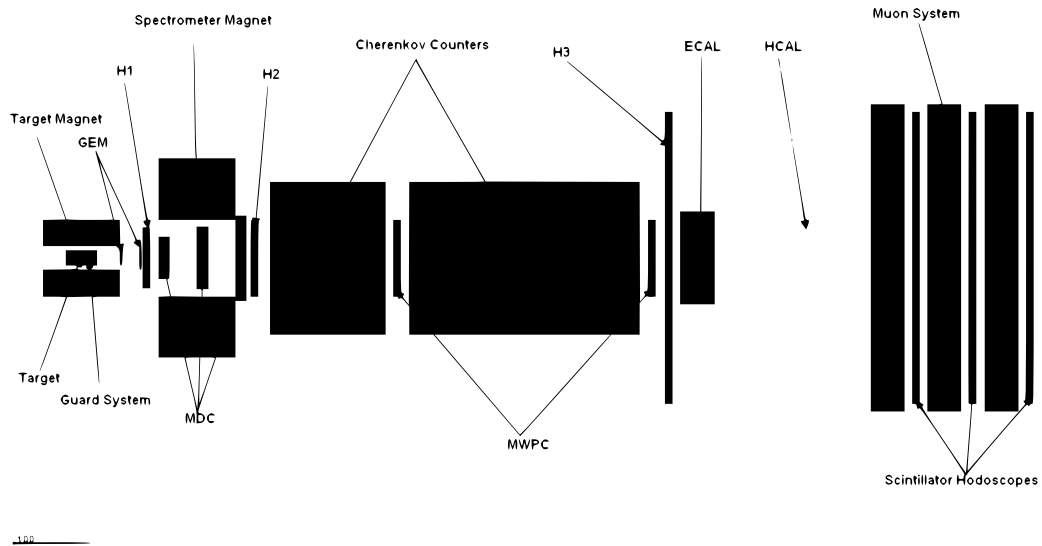


Figure 4: Layout of the SPASCHARM experimental apparatus.

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POLARIZATION MEASUREMENTS IN PHOTOPRODUCTION WITH CEBAF LARGE ACCEPTANCE SPECTROMETER

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Abstract

A significant part of the experimental program in Hall-B of the Jefferson Lab is dedicated to the studies of the structure of baryons. CEBAF Large Acceptance Spectrometer (CLAS), availability of circularly and linearly polarized photon beams and recent addition of polarized targets provides remarkable opportunity for single, double and in some cases triple polarization measurements in photoproduction. An overview of the experiments will be presented.

Among the most exciting and challenging topics in sub-nuclear physics today is the study of the structure of the nucleon and its different modes of excitation, the baryon resonances. Initially, most of the information on these excitations came primarily from partial wave analysis of data from πN scattering. Recently, these data have been supplemented by a large amount of information from pion electro- and photoproduction experiments. Yet, in spite of extensive studies spanning decades, many of the baryon resonances are still not well established and their parameters (i.e., mass, width, and couplings to various decay modes) are poorly known. Much of this is due to the complexity of the nucleon resonance spectrum, with many broad, overlapping resonances. While traditional theoretical approaches have highlighted a semi-empirical approach to understanding the process as proceeding through a multitude of s -channel resonances, t -channel processes, and non-resonant background, more recently attention has turned to approaches based on the underlying constituent quarks. An extensive review of the quark models of baryon masses and decays can be found in Ref. [1]. Most recently lattice QCD is making significant progress in calculations of baryon spectrum. While these quark approaches are more fundamental and hold great promise, all of them predict many more resonances than have been observed, leading to the so-called “missing resonance” problem. One possible reason why they were not observed because they may have small coupling to the πN . At the same time they may have strong coupling to other final states ηN , $\eta' N$, KY , $2\pi N$.

A large part of the experimental program of Jefferson Lab and CLAS in particular is dedicated to baryon spectroscopy. There were several CLAS running periods with circularly and linearly polarized photon beams. High quality data for the cross sections of π^0 [2], π^+ [3], η [4], η' [5] and kaon photoproduction were obtained. In addition to the cross sections [6] for $K^+\Lambda$ and $K^+\Sigma^0$ final states the polarization of hyperons P [7] and polarization transfer C_x/C_z were measured [8]. Artru, Richard, and Soffer [9] pointed out that for a circularly polarized beam, there is a rigorous inequality

$$R^2 \equiv P^2 + C_x^2 + C_z^2 \leq 1. \quad (1)$$

The data from the former two experiments allowed us to test this inequality. The remarkable result is that the Λ hyperon is produced fully polarized at all values of W and scattering angle for a fully circularly polarized beam. This is something which was not expected and yet is to be explained.

The beam asymmetry Σ was measured with linearly polarized beam. Significant amount of data for cross sections and beam asymmetries was also accumulated at ELSA, MAMI, GRAAL and LEPS. However, without double polarization measurement it is still impossible to resolve all ambiguities in the reaction amplitude. Several experiments [10] were proposed to measure double polarization observables in all reaction channels $\pi^0 p$, $\pi^+ n$, ηp , $\eta' p$, KY , $\pi^+ \pi^- p$ with CLAS, circularly and linearly polarized photons and longitudinally and transversely polarized target.

Experimental Hall-B at Jefferson Lab provides a unique set of instruments for these experiments. One instrument is the CLAS [11], a large acceptance spectrometer which allows detection of particles in a wide range of θ and ϕ . The other instrument is a broad-range photon tagging facility [12] with the recent addition of the ability to produce linearly-polarized photon beams through coherent bremsstrahlung. The remaining component essential for the double polarization experiments is a frozen-spin polarized target (FROST) .

The Hall B photon tagger [12] covers a range in photon energies from 20 to 95% of the incident electron beam energy. Unpolarized, circularly polarized and linearly polarized tagged photon beams are presently available.

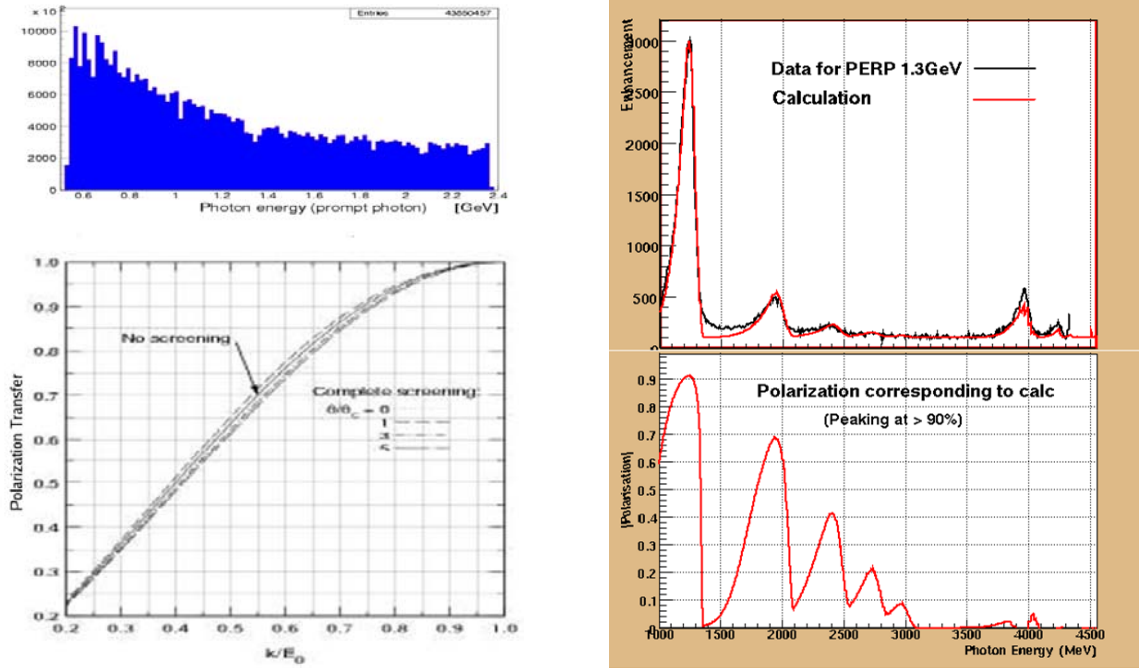


Figure 1: An example of tagged incoherent bremsstrahlung photon spectrum (top left) and polarization transfer from electron to photon (bottom left). An example of coherent bremsstrahlung spectrum (top right) and calculated photon polarization as a function of energy

With a polarized electron beam incident on the bremsstrahlung radiator, a circularly polarized photon beam can be produced. The degree of circular polarization of the photon beam depends on the ratio $k = E_\gamma/E_e$, and is given by [13]

$$P_\odot = P_e \cdot \frac{4k - k^2}{4 - 4k + 3k^2}. \quad (2)$$

The magnitude of $P_\odot$ ranges from 60% to 99% of the incident electron beam polarization P_e for photon energies E_γ between 50% and 95% of the incident electron energy. CEBAF accelerator routinely delivers electron beam with polarization of 85% and higher. An example of tagged circularly polarized photon spectrum and its degree of polarization is shown in Fig. 1 (left).

A linearly polarized photon beam is produced by the coherent bremsstrahlung technique, using an oriented diamond crystal as a radiator. Figure 1 (right) shows an example of collimated linearly polarized tagged photon spectrum obtained in Hall-B. The degree of polarization is a function of the fractional photon beam energy and collimation and can reach 80% to 90%. With linearly polarized photons, over 80% of the photon flux is confined to a 200-MeV wide energy interval.

An essential piece of the hardware for this experiment is a polarized target capable of being polarized transversely and longitudinally with a minimal amount of material in the path of outgoing charged particles. The Hall-B polarized target [14] used in electron beam experiments is a dynamically polarized target. The target is longitudinally polarized with a pair of 5 Tesla Helmholtz coils. These massive coils limit available aperture to 55 degrees in forward direction. For photon beam experiments, a frozen-spin target is a much more attractive choice. Butanol was chosen as target material. In frozen-spin mode, the

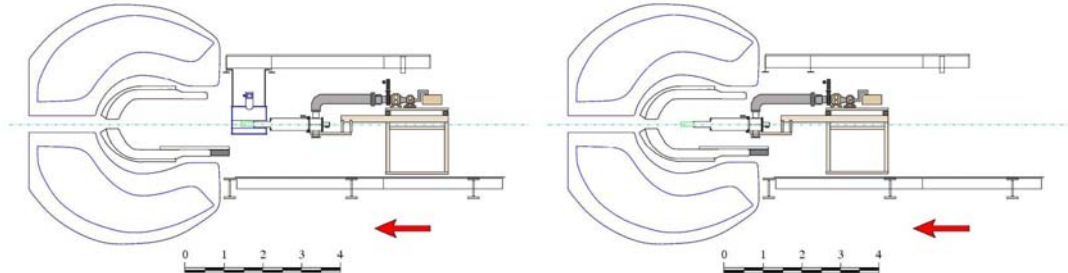


Figure 2: The target is pulled out from the CLAS and inserted in the polarizing magnet (left panel). The target is inside the CLAS in its normal configuration for data taking (right panel).

target material is dynamically polarized in a strong magnetic field of 5 Tesla outside of CLAS. After maximal polarization is reached the cryostat is turned to the “holding” mode with much lower magnetic field of 0.5 Tesla at a temperature of about 30 mK, and then moved in CLAS. A photon beam does not induce noticeable radiation damage and does not produce significant heat load on the target material. Under these conditions the target can hold its polarization with a relaxation time on the order of several days before re-polarization is required. Since the holding field is relatively low, it is possible to design a “transparent” holding magnet with a minimal amount of material for the

charged particles to traverse on their way into CLAS. Figure 2 shows two configurations of the target. On the top panel the target is pulled out from the CLAS and inserted in the polarizing magnet. The bottom panels shows the target inside the CLAS in its normal configuration for data taking.

The frozen spin target has been built by the JLab polarized target group. Table 1 summarizes main parameters of the target. There are two holding magnets available and the target can be configured either for longitudinal or transverse polarization. During the first round of experiments the target demonstrated excellent performance running continuously for three and a half months.

Table 1: Parameters of the FROST.

	Expectation	Result
Base temperature	50 mK	28mK no beam 30 mK with beam
Cooling power	10 μ W (frozen) 20 mW (polarizing)	800 μ W@50 mK 60 mW@300 mK
Polarization	$\pm 85\%$	+82% -85%
Relaxation time	500 hours (5% per day)	2700 hours (+Pol) 1400 hours (-Pol) (<1.5% per day)

The first round of experiment with FROST was conducted in November 2007 – February 2008. In this set of experiments target polarization was longitudinal. We used both, linearly and circularly polarized photons. Photon energy range covered 0.5 – 2.3 GeV. Combination of circularly polarized beam and longitudinally polarized target allows us to measure helicity asymmetry E . With linearly polarized photons incident on longitudinal target we have access to the observable G .

Here we present the first look at the data collected. A preliminary analysis of the helicity asymmetry in the reaction $\gamma p \rightarrow \pi^+ n$ was done. Selected events were binned in photon energy and $\cos \theta_{c.m.}$. In each bin we determined measured asymmetry

$$A = \frac{N_{\Rightarrow} - N_{\Leftarrow}}{N_{\Rightarrow} + N_{\Leftarrow}}, \quad (3)$$

where $N_{\Rightarrow}$ and $N_{\Leftarrow}$ is number of events beam and target polarization is parallel and anti-parallel respectively. Then, to get helicity asymmetry, E , this raw asymmetry was corrected for beam and target polarization and effective dilution factor. At this stage of the analysis we used very rough estimate of these corrections. A very preliminary results for helicity asymmetry for π^+ photoproduction are presented in Figure 3. One can see that asymmetry is large and energy evolution displays various structures in angular dependence. This is just one example of what is coming in near future. The analysis of all other final states is underway.

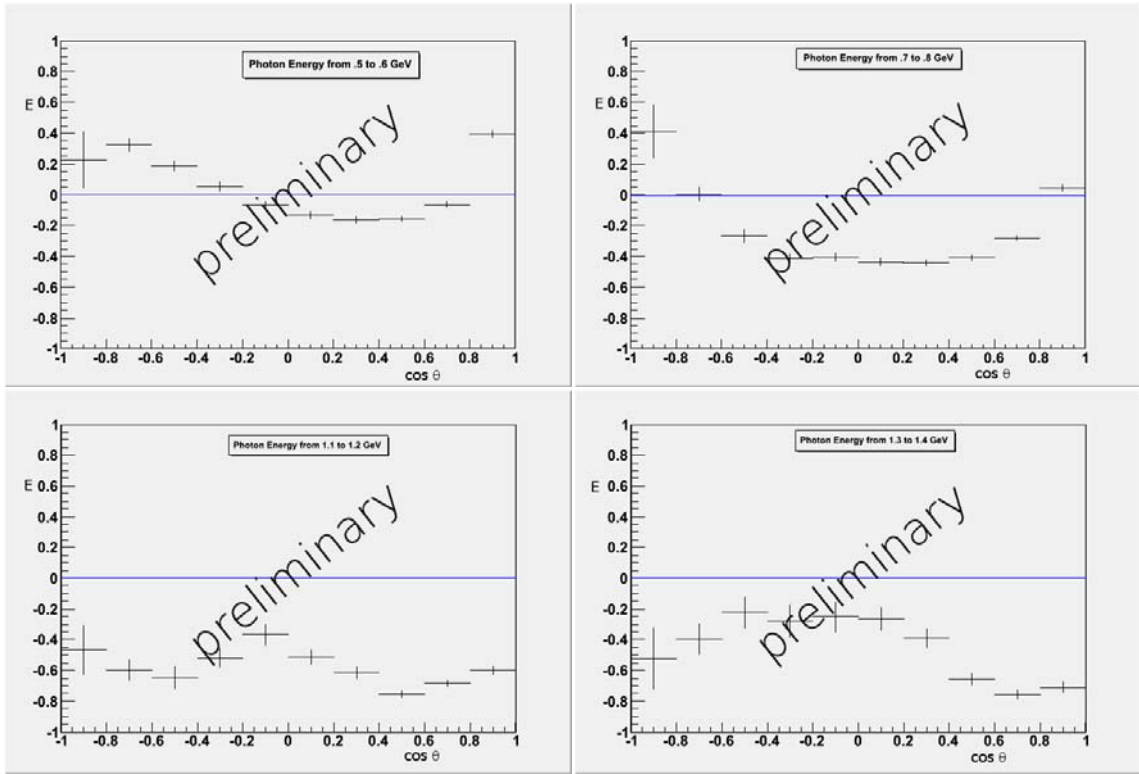


Figure 3: Preliminary helicity asymmetry for $\gamma p \rightarrow \pi^+ n$

The addition of Frozen Spin Target, with both, longitudinal and transverse polarization significantly advances our experimental capabilities. The first round of the double polarization photoproduction experiments with longitudinally polarized target has been complete and experimental data are being analyzed. The second part of the experiment is scheduled to run in spring of 2010 and will use transversely polarized target.

Upon completion of the experiment it will be possible for the first time to perform complete experiment of KY photoproduction and nearly complete for other final states. Entire program is more than just a sum of several experiments, observables for all final states are measured simultaneously under the same experimental conditions and have the same systematic uncertainties.

Another essential part of the program which was not described here involves photoproduction experiments on the deuteron target which allow to study different isospin states of the baryon resonances. Several CLAS experiments with polarized photons and unpolarized deuteron target were complete. Double polarization experiments with HD-Ice polarized target are in preparation.

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INVESTIGATION OF DP-ELASTIC SCATTERING AND DP-BREAKUP AT ITS AT NUCLOTRON

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Abstract

The experiment on the light nuclei structure (*LNS*) studies at Internal Target Station at Nuclotron is discussed. The results on the vector A_y and tensor A_{yy} , A_{xx} analyzing powers for dp - elastic scattering obtained at $T_d = 880$ MeV are shown. The status of the setup upgrade (detectors, high voltage and DAQ systems etc.) is presented.

The purpose of Light Nuclei Structure experimental program is to obtain the information on the spin τAY dependent parts of 2-nucleon and 3-nucleon forces from two processes: dp -elastic scattering in a wide energy range and dp - non-mesonic breakup with two protons detection at energy 300 - 500 MeV.

Nowadays a new generation of the NN potentials (AV-18, CD-Bonn, Nijmegen etc.) was obtained. They reproduce data on the nucleon nucleon scattering up to 350 MeV with very good accuracy. However, these modern 2N forces fail to provide the experimental binding energies of few-nucleon systems. Moreover the data on the dp - elastic scattering and deuteron breakup are not described properly. Incorporation of the 3N forces makes it possible to reproduce the binding energy of the three-nucleon bound systems and also data on unpolarized dp - interaction. Nevertheless, polarization data for the reactions with participation of three and more nucleons are not described even with inclusion of 3NF. Therefore, the obtaining of the additional polarization data in the dp - interaction in the wide energy range more is very desirable for the study of the spin structure of 2N and 3N forces [1].

Another goal is to find a suitable reaction to provide efficient polarimetry of high energy deuteron. dp elastic scattering at large angles ($\theta_{cm} \geq 60^\circ$) has been proposed for

the polarimetry at the energies between 0.88 and 2.0 GeV [1]. These data are necessary to provide good precision of the beam polarization measurement for DSS project [2] at the LHEP-JINR.

The experimental data on the deuteron analyzing powers for large phase space were obtained at 130 MeV at KVI [3]. A_y , A_{yy} and A_{xx} analyzing powers in dp -breakup will be investigated at Internal Target Station at 200-500 MeV. The predictions for tensor analyzing power A_{yy} and differential cross section in the selected dp -breakup configurations at 400 MeV are shown in Fig.1. The light shaded band contains the theoretical predictions based on CD-Bonn, AV18, Nijm I, II and Nijm 93. The darker band represents predictions when these NN forces are combined with the TM 3NF. The solid line is for AV18+Urbana IX and the dashed line for CD Bonn+TM. One can see large sensitivity of these observables to the model of 3NF.

The details of the experiment on the measurements of analyzing powers in dp -elastic scattering at Nuclotron can be found in ref. [4]. The polarization of the deuteron beam was measured at 270 MeV, where high-precision data on the analyzing powers from dp -elastic scattering from RIKEN exist [5]. Two particles from the dp -elastic scattering were clearly distinguished by their time-of-flight differences from the target and their energy losses in the plastic scintillators. Values on the vector and tensor polarizations of the deuteron beam were obtained. These values are in good agreement with the results obtained by low energy polarimeter based on $^3\text{He}(d, p(0^\circ))^4\text{He}$ reactions [6].

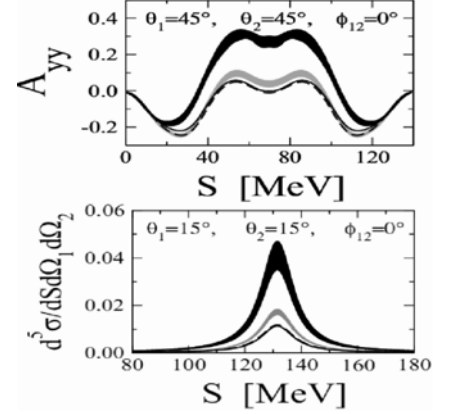


Figure 1: The tensor analyzing power A_{yy} and differential cross section in selected breakup configurations at 400 MeV. The curves are explained in the text.

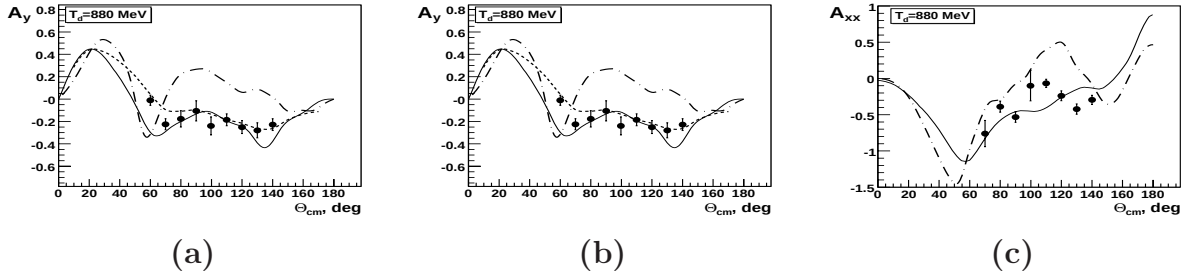


Figure 2: (a) Vector A_y and (b) tensor A_{yy} , (c) A_{xx} analyzing powers for the dp - elastic scattering at 880 MeV. The curves are described in the text.

The true numbers of the dp - elastic events at 880 MeV were obtained by imposing of the cut on two-dimensional plot of ADC signals from deuterons and protons detectors. Additionally, the graphical cuts were applied on the time-of-flight difference between deuterons and protons from the target. The contribution of the carbon dependent on the scattering angle was subtracted.

The experimental data obtained at 880 MeV are compared with the theoretical predictions in Fig.2. Solid curves correspond to Faddeev calculation using the CD-Bonn NN potential. One can see that the Faddeev calculations reproduce the behavior of all the analyzing powers. Dotted curves correspond to the calculations in optical potential

framework with the use of deuteron wave function (DWF), derived from the dressed bag model of Moscow-Tuebingen group [7]. Dash-dotted curves conform to the calculations performed within the multiple scattering expansion formalism with the use of CD-Bonn DWF [8]. The parameterization of the NN t-matrix [9] has been used to take the off-shell effects into account. The deviation of the data from the calculations can indicate the sensitivity to the short-range 3N correlations.

The dp breakup reaction will be investigated using ΔE - E techniques for the detection of protons. 8 detectors of this kind are planned to use in the experiment. Each detector consists of two scintillators ΔE and E .

The first scintillator has the cylindrical shape with the height 10 mm and the diameter 100 mm. Two PMTs-85 view the scintillator and they are located the friend opposite to the friend. Two planes have been made on the scintillator to increase the area of the optical contact between the scintillator and photocatode of the PMT-85. These planes have been polished (see Fig.3). ΔE scintillator it is covered by a white paper. Digital dividers of voltage are used for PMT-85.

E scintillator also has the cylindrical form with the height 200 mm and diameter 100 mm. PMT-63 view through this scintillator. Given PMT has been chosen because of the suitable size of the cathode (100 mm) with both good time and amplitude properties. E scintillator has been wrapped by a white paper. The part which is located to ΔE scintillator has been covered by a black paper. It is made to exclude possibility of the light hit from one plastic scintillator to another.

The beam test of two ΔE - E detectors at ITS has been performed at the initial deuteron kinetic energy of 2.3 GeV in June 2008. During this test a new DAQ-system based on the trigger LT320D module has been used. One of the important advantages of this module is the possibility to control the status of majority coincidence circuit on-line. The trigger used was the coincidence of the signals from 2 ΔE - E detectors located in the horizontal plane on the left and right from the initial beam direction.

First results for the beam test of ΔE - E counters are presented in Fig.4. The results on the correlation of the energy depositions in two E - scintillation counters and their time-of-flight difference are given in Fig.4.

The modernization of ΔE - E detectors consists in the diminution of the ΔE scintillator diameter to 8 cm. It has allowed to reduce the solid angle for detected nucleons. Solid angle reduction allows to exclude nucleons which take off through a E - scintillator lateral side. This detector was tested with cosmic muons. Results of the calibration are presented on Fig.5. Pedestal is disposed in 23-rd bin. Mean size of amplitude distribution from E - scintillator is equal to 243 channels. 10 channels are equal 1.8 MeV for this configuration of the detector.

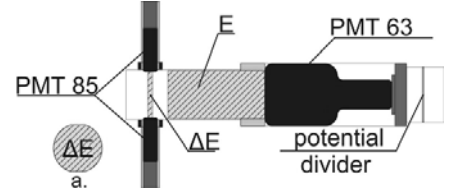


Figure 3: The schematic view of the ΔE - E detector for dp τAY breakup reaction study.

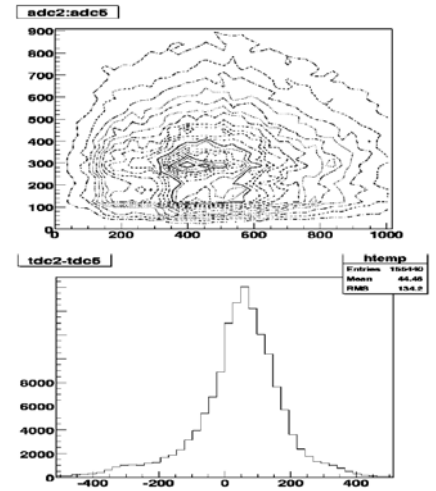


Figure 4: The correlation of amplitudes from two E -detectors and their time-of-flight difference.

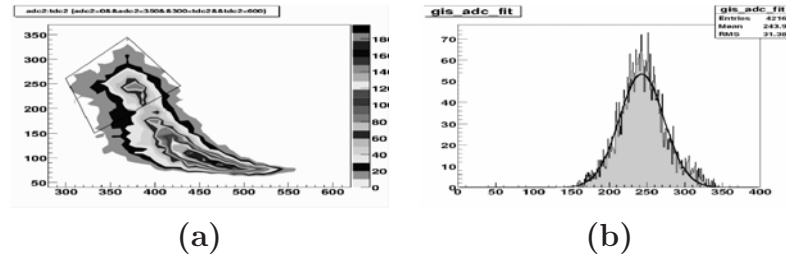


Figure 5: (a) the correlation of amplitude and time from one E detector and (b) the amplitude from PMTs-63 with imposed cut.

The high voltage digital dividers of PMTs-85 were controlled by the module connected with computer through the bus RS-232. This system was designed at LHEP-JINR [10]. The high voltage system for PMTs-63 is based on "Wenzel Elektronik" module, whose voltage is adjusted and checked on-line using DAC and ADC CAMAC modules.

The versatile DAQ system for middle range physics experiments MIDAS [11] was used for on-line control of high voltage provided by "Wenzel Elektronik". The system demonstrated good time stability during beam test.

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SPIN PHYSICS WITH CLAS

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Abstract

Inelastic scattering using polarized nucleon targets and polarized charged lepton beams allows the extraction of double and single spin asymmetries that provide information about the helicity structure of the nucleon. A program designed to study such processes at low and intermediate Q^2 for the proton and deuteron has been pursued by the CLAS Collaboration at Jefferson Lab since 1998. Our inclusive data with high statistical precision and extensive kinematic coverage allow us to better constrain the polarized parton distributions and to accurately determine various moments of spin structure function g_1 as a function of Q^2 . The latest results are shown, illustrating our contribution to the world data, with comparisons of the data with NLO global fits, phenomenological models, chiral perturbation theory and the GDH and Bjorken sum rules. The semi-inclusive measurements of single and double spin asymmetries for charged and neutral pions are also shown, indicating the importance of the orbital motion of quarks in understanding of the spin structure of the nucleon.

1 Spin Structure Functions

One fundamental goal of Nuclear Physics is the description of the structure and properties of hadrons, and especially nucleons, in terms of the underlying degrees of freedom, namely quarks and the color forces between them. Much progress has been made over the last decades towards this goal, both experimentally (e.g., through structure function and form factor measurements) and theoretically (effective theories like the quark model, chiral perturbation theory as well as complete solutions of QCD on the lattice). At the same time, there are many important questions that require further investigation, such as: What is the quark structure of nucleons in the valence region, in particular in the limit of large momentum fraction carried by a single quark, $x \rightarrow 1$? How can we describe the transition from hadronic degrees of freedom to quark degrees of freedom for the nucleon? How can we describe the nucleon in three dimensions and what are the correlations between transverse momentum and spin? How does quark orbital angular momentum contribute to the spin of the nucleon? A program designed to study these questions, and utilizing the CLAS detector, 6 GeV polarized electron beam, and longitudinally polarized solid ammonia targets (NH_3 and ND_3) has been pursued by the CLAS Collaboration at Jefferson Lab since 1998. This program entails both inclusive measurements of inelastic electron scattering as well as coincident detection of leading hadrons (pions etc.) produced in such events. Due to the large acceptance of CLAS, a large kinematical region is accessed simultaneously. Both the scattered electrons and leading hadrons from the hadronization of the struck quark are detected, allowing us to gain information on its flavor.

1.1 Nucleon Helicity Structure at large x

The photon-nucleon asymmetry $A_1(x, Q^2)$ reflects the valence spin structure of the nucleon. Valence quarks are the irreducible kernel of each hadron, responsible for its charge, baryon number and other macroscopic properties. The region $x \rightarrow 1$ is a relatively clean region to study the valence structure of the nucleon since this region is dominated by valence quarks while the small x region is dominated by gluon and sea densities. Due to its relative Q^2 -independence in the DIS region, the virtual photon asymmetry A_1 , which is approximately given by the ratio of spin-dependent to spin averaged structure functions,

$$A_1(x) \approx \frac{g_1(x)}{F_1(x)}, \quad (1)$$

is one of the best physics observables to study the valence spin structure of the nucleon. At leading order,

$$A_1(x, Q^2) = \frac{\sum e_i^2 \Delta q_i(x, Q^2)}{\sum e_i^2 q_i(x, Q^2)}, \quad (2)$$

where $q = q \uparrow + q \downarrow$ and $\Delta q = q \uparrow - q \downarrow$ are the sum and difference between quark distributions with spin aligned and anti-aligned with the spin of the nucleon. The x dependence of the parton distributions provide a wealth of information about the quark-gluon dynamics of the nucleon. In particular spin degrees of freedom allow access to information about the structure of hadrons not available through unpolarized processes. Furthermore, the spin dependent distributions are more sensitive than the spin-averaged ones to the quark-gluon dynamics responsible for spin-flavor symmetry breaking. Several models make specific predictions for the large x behavior of quark distributions.

1.2 Moments and Sum Rules

The spin structure function g_1 is important in understanding the quark and gluon spin components of the nucleon spin, and their relative contributions in different kinematic regions. At high Q^2 , g_1 provides information on how the nucleon spin is composed of the spin of its constituent quarks and gluons. At low Q^2 , hadronic degrees of freedom become more important and dominate the measurements. There is particular interest in the first moment of g_1 , $\Gamma_1(Q^2) = \int_0^{1-} g_1(x, Q^2) dx$, which is constrained at low Q^2 by the Gerasimov-Drell-Hearn sum rule [1] and at high Q^2 by the Bjorken sum rule [2] and previous DIS experiments. In our definition the upper limit of the integral does not include the elastic peak. Ji and Osborne [3] have shown that the GDH sum rule can be generalized to all Q^2 via

$$S_1(\nu = 0, Q^2) = \frac{8}{Q^2} [\Gamma_1(Q^2) + \Gamma_1^{el}(Q^2)], \quad (3)$$

where $S_1(\nu, Q^2)$ is the spin-dependent virtual photon Compton amplitude. S_1 can be calculated in Chiral Perturbation Theory (χ PT) at low Q^2 and with perturbative QCD (pQCD) at high Q^2 . Therefore, Γ_1 represents a calculable observable that spans the entire energy range from hadronic to partonic descriptions of the nucleon. Higher moments are also of interest: generalized spin polarizabilities, γ_0 and δ_{LT} , are linked to higher moments of spin structure functions by sum rules based on similar grounds as the GDH sum rule. Higher moments are less sensitive to the unmeasured low- x part since they are more weighted at high- x .

1.3 Flavor Decomposition of the Helicity Structure

If we want to understand the three-dimensional structure of the nucleon, we have to go beyond inclusive measurements that are only sensitive to the longitudinal momentum fraction x carried by the quarks. The large acceptance of CLAS allows us to collect data on semi-inclusive (SIDIS) reactions simultaneously. In these reactions, a second particle, typically a meson, is detected along with the scattered lepton. By making use of the additional information given by the identification of this meson, one can learn more about the polarized partons inside the nucleon than from DIS alone. The asymmetry measured by DIS experiments is sensitive to combinations of quark and anti-quark polarized parton distribution functions ($\Delta q + \Delta \bar{q}$), as well as (via NLO analysis) the gluon PDF ΔG . SIDIS experiments exploit the statistical correlation between the flavor of the struck quark and the type of hadron produced to extract information on quark and antiquark PDFs of all flavors separately. Combined NLO analysis of DIS and SIDIS data can therefore give a more detailed picture of the contribution of all quark flavors and both valence and sea quarks to the total nucleon helicity. Beyond the determination of the polarized PDFs, SIDIS data can also yield a plethora of new insights into the internal structure of the nucleon as well as the dynamics of quark fragmentation. For instance, looking at the z - and p_T -dependence of the various meson asymmetries (both double spin asymmetries and single spin target or beam asymmetries), one can learn about the intrinsic transverse momentum of quarks and their orbital angular momentum.

2 Measurements and Data Analysis

A_1 and g_1 were extracted from measurements of the double spin asymmetry $A_{\parallel}$ in inclusive ep scattering:

$$g_1 = \frac{F_1}{1 + \gamma^2} [A_{\parallel}/D + (\gamma - \eta)A_2], \quad (4)$$

where F_1 is the unpolarized structure function, A_2 is the virtual photon asymmetry, and γ , D and η are kinematic factors. F_1 and A_2 are calculated using a parametrization of the world data, and $A_{\parallel}$ is measured. The spin asymmetry for ep scattering is given by:

$$A_{\parallel} = \frac{N_- - N_+}{N_- + N_+} \frac{C_N}{f P_b P_t f_{RC}} + A_{RC}, \quad (5)$$

where $N_-(N_+)$ is the number of scattered electrons normalized to the incident charge with negative (positive) beam helicity, f is the dilution factor needed to correct for the electrons scattering off the unpolarized background, f_{RC} and A_{RC} correct for radiative effects, and C_N is the correction factor associated with polarized ^{15}N nuclei in the target. $A_{\parallel}$ was measured by scattering polarized electrons off polarized nucleons using a cryogenic solid polarized target and CLAS in Hall B. The raw asymmetries were corrected for the beam charge asymmetry, the dilution factor and radiative effects. Since the elastic peak is within the acceptance range, the product of beam and target polarization was determined from the known ep elastic asymmetry.

The longitudinally polarized electrons were produced by a strained $GaAs$ electron source with a typical beam polarization of $\sim 70\%$. Two solid polarized targets were used:

$^{15}\text{ND}_3$ for polarized deuterons and $^{15}\text{NH}_3$ for polarized protons. The targets were polarized using the method of Dynamic Nuclear Polarization, with the typical polarization of 70-90% for protons, and 10-35% for deuterons. Besides the polarized targets, three unpolarized targets (^{12}C , ^{15}N , liquid ^4He) were used for background measurements. The scattered electrons were identified using the CLAS package [4], consisting of drift chambers, Cherenkov detector, time-of-flight counters and electromagnetic calorimeters. Data were taken with beam energies of 1.6, 2.4, 4.2 and 5.7 GeV, covering a kinematic range of $0.05 < Q^2 < 4.5 \text{ GeV}^2$ and $0.8 < W < 3.0 \text{ GeV}$. The data include multi-particle final states, making it possible to investigate exclusive and semi-inclusive pion production, deeply virtual Compton scattering and other exclusive channels.

3 Results

3.1 Large x behaviour of $A_1(x, Q^2)$

The photon-nucleon asymmetry A_1^d is shown in Figure 1. Along with the recent CLAS data, the plot shows results from previous experiments, and predictions from several models.

The models [5] include the suppression of transitions to states in the lowest even and odd parity multiplets with combined quark spin $S = \frac{3}{2}$, the suppression of transitions to states with helicity $h = \frac{3}{2}$, and the suppression of transitions to the states which couple only through symmetric components of the spin-flavor wavefunction. Also shown is the prediction of the hyperfine-perturbed quark model, which involves spin-spin interaction between quarks, mediated by one gluon or pion exchange [6]. Our data show a preference for the pQCD limit as $x \rightarrow 1$, and are also consistent with the hyperfine-perturbed quark model.

3.2 Moments of $g_1(x, Q^2)$

The first moments of g_1^p and g_1^d are shown in Figure 2. The parametrization of world data is used to include the unmeasured contribution to the integral down to $x = 0.001$. Only the Q^2 bins in which the measured part constitutes at least 50% of the total integral are included. For the proton, the parametrization at high x ($1.09 < W < 1.14$ (1.15) GeV) is used for the low (high) energy data). For the deuteron, the integration is carried out up to the nucleon pion production threshold at high x , excluding the quasi-elastic and electro-disintegration contributions. The integral is observed to turn over at low Q^2 , consistent with the slope predicted by the GDH sum rule. In general the data are well described by the phenomenological models of Burkert and Ioffe [10] and Soffer and Teryaev [11]. The low Q^2 Γ_1 data are shown in more detail in the right-hand panels of Figs. 2. It is possible to make a quantitative comparison between our results for Γ_1^p and Γ_1^d at low Q^2 and the next-to-leading order χ PT calculation by Ji, Kao and Osborne [12], who find $\Gamma_1^p(Q^2) = -\frac{\kappa_p^2}{8M^2}Q^2 + 3.89Q^4 + \dots$ and $\Gamma_1^d(Q^2) = -\frac{\kappa_n^2}{8M^2}Q^2 + 3.15Q^4 + \dots$

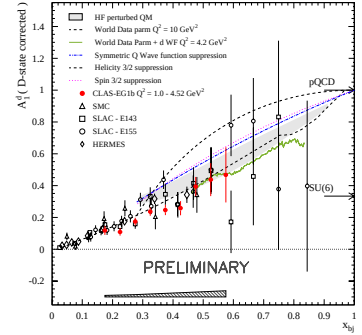


Figure 1: Asymmetry A_1^d plotted vs x could differentiate between the different models of valence spin structure of the nucleons.

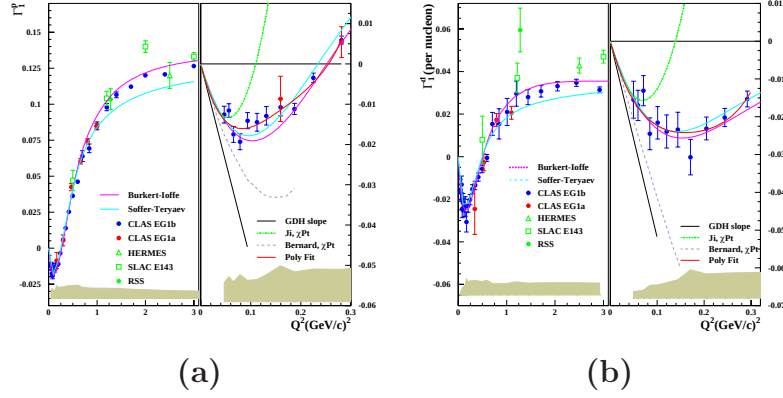


Figure 2: (a) Γ_1^p as a function of Q^2 . (b) Γ_1^d as a function of Q^2 . The EG1a [7], SLAC [8] and Hermes data [9] are shown for comparison. The filled circles represent the present data, including an extrapolation over the unmeasured part of the x spectrum using a model of world data.

Treating the deuteron as the incoherent sum of a proton and a neutron and correcting for the D-state,

$$\Gamma_1^d(Q^2) = \frac{1}{2}(1 - 1.5\omega_D) \{ \Gamma_1^p(Q^2) + \Gamma_1^n(Q^2) \}, \quad (6)$$

one finds that $\Gamma_1^d(Q^2) = -0.451Q^2 + 3.26Q^4$. The low Q^2 results for Γ_1^p and Γ_1^d have been fit to a function of the form $aQ^2 + bQ^4 + cQ^6 + dQ^8$ where a is fixed at -0.455 (proton) and -0.451 (deuteron) by the GDH sum rule.

For the proton, $b = 3.81 \pm 0.31$ (stat) $+0.44 - 0.57$ (syst) is extracted and for the deuteron, $b = 2.91 \pm 0.52$ (stat) ± 0.69 (syst) was obtained, both consistent with the Q^4 term predicted by Ji *et.al.*

Our fit is shown in the right-hand panel of plots in Figs. 2 along with Ji's prediction. We find that the Q^6 term becomes important even below $Q^2 = 0.1$ GeV² and that this term needs to be included in the χ PT calculations in order to extend the range of their validity.

Higher moments of g_1 are interesting as well. In our kinematic domain these moments emphasize the resonance region over DIS kinematics because of extra factors of x in the integrand. The generalized forward spin polarizability of the nucleon is given by [13]

$$\gamma_0(Q^2) = \frac{16\alpha M^2}{Q^6} \int_0^{x_0} x^2 \left\{ g_1(x, Q^2) - \frac{Q^2 x^2}{4M^2} g_2(x, Q^2) \right\} dx, \quad (7)$$

where α is the fine structure constant. First results for the generalized forward spin polarizability of the proton for a range of Q^2 from 0.05 to 4 GeV² are shown in Fig. 3. Our data lie closest to the MAID 2003 [14] model, which is a phenomenological fit to single pion production data and includes only the resonance region. However, since γ_0 is weighted by an additional factor of x^2 compared to Γ_1 , the contribution to the integral from the DIS part of the spectrum is rather small. The MAID model follows the trend of the data but significantly underpredicts them numerically.

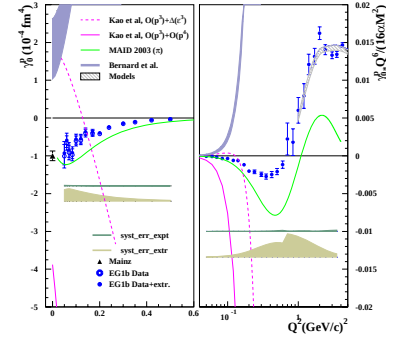


Figure 3: Generalized forward spin polarizability γ_0^p as a function of Q^2 for the full integral (closed circles) and the measured portion of the integral (open circles)

The 4th order Heavy Baryon Chiral Perturbation calculation by Kao, Spitzenberg and Vanderhaeghen [15], shown by the dashed line in Fig. 3, also underpredicts the data. The authors note that the $O(p^4)$ correction term is of opposite sign to the $O(p^3)$ term and shows no sign of convergence. A leading order correction to account for $\Delta(1232)$ degrees of freedom, not shown, is also negative. By contrast the χ PT calculation of Bernard, Hemmert and Meissner [16], indicated by the grey band, including the resonance contribution, overpredicts the data. The $\Delta(1232)$ and vector meson contribution is negative (around $-2 \times 10^{-4} \text{ fm}^4$) but the discrepancy with the data suggests that this has been underestimated.

3.3 Factorization tests

The semi-inclusive measurements of double spin asymmetries allow us to study factorization of $x, z = E_h/\nu$ and p_T dependency for charged and neutral pions. We study the quantity g_1/F_1 as a function of kinematic variables for all three pion flavors [18]. The z -dependence of semi-inclusive g_1/F_1 is examined in Fig. 4. We compare the data with LO pQCD predictions obtained from the GRSV2000 [17] parametrization. The ratio should be approximately independent of z , broken by the different weights given to the polarized u and d quarks by the favored and unfavored fragmentation functions. This is indeed observed in the data, which are in good agreement with the model in both magnitude and z -dependence up to $z = 0.7$. The observed drop-off at high z for π^- is also expected due to increased importance of $d(x)$ with increasing z . The x -dependence of g_1/F_1 for π^+, π^-, π^0 are consistent with each other and follow the same trend as the inclusive result, which is expected if factorization works. The trend is for the ratio to increase with x , due to increasing dominance of polarized $u(x)$ at high x . The ratio g_1/F_1 has also been studied as a function of transverse component of the hadron momentum p_T . The data suggest that at small p_T , g_1/F_1 tends to decrease for π^+ and to increase for π^- . This result indicates that quarks aligned and anti-aligned with the nucleon spin might have different transverse momentum distributions, but more data is needed to study this behavior.

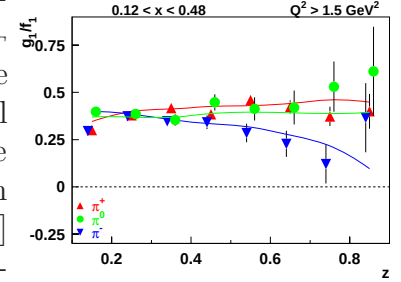


Figure 4: z -dependence of semi-inclusive g_1/F_1 for the proton target

4 Summary and Outlook

In conclusion, the spin asymmetries for the proton and the deuteron have been measured over a vast kinematic region, allowing systematic studies of various aspects of the nucleon spin structure at low and intermediate energies. Our data are consistent with an approach to $A_1 = 1$ as $x \rightarrow 1$ as required by pQCD. The first extraction of generalized forward spin polarizability γ_0^p was shown to be poorly described by the chiral perturbation theory even at our lowest $Q^2 = 0.05 \text{ GeV}^2/c^2$. The semi-inclusive studies have shown interesting p_T dependence suggesting that the transverse momentum distributions may have non-trivial dependency on the flavor and helicity of quarks. These studies have led to several new proposals with 6 and 11 GeV beams. A recently completed run with 6 GeV beam will provide us with an order of magnitude more π^+, π^-, π^0 for g_1/F_1 on the proton.

The upcoming energy upgrade of Jefferson Lab will allow us to begin a next generation of spin structure studies and will provide a significant increase in kinematic coverage and statistical accuracy in inclusive and semi-inclusive measurements.

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MEASUREMENTS OF G_{Ep}/G_{Mp} TO HIGH Q^2 AND SEARCH FOR 2γ CONTRIBUTION IN ELASTIC ep AT JEFFERSON LAB

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Abstract

The ratio, $\mu_p G_{Ep}/G_{Mp}$, where μ_p is the proton magnetic moment, has been measured extensively over the last decade at the Jefferson Laboratory, using the polarization transfer method. This ratio is extracted directly from the measured ratio of transverse to longitudinal polarizations components of the recoiling proton in elastic electron-proton scattering. The polarization transfer results are of unprecedented high precision and accuracy, due in large part to the small systematic uncertainties associated with the experimental technique. Prior to these measurements, the form factors were empirically observed to exhibit dipole forms, such that $\mu_p G_{Ep}/G_{Mp} \approx 1$ over all regions of momentum transfer studied. With the Hall A results confirming that the ratio $\mu_p G_{Ep}/G_{Mp}$ shows a steady decrease below unity as a function of Q^2 , beginning around $Q^2 \approx 1 \text{ GeV}^2$, discussions revolving around the implication of this deviation from dipole behavior for the structure of the proton have been accompanied by renewed experimental interest in these elastic form factors.

Starting in the fall of 2007, two new experiments, GEp-III and GEp- 2γ in Hall C at JLab, measured the form factor ratio, G_{Ep}/G_{Mp} ; the GEp-III experiment pushed the highest Q^2 limit from 5.6 to 8.49 GeV^2 , with intermediate points at 5.2 and 6.8 GeV^2 , and the GEp- 2γ experiment measured the ratio in three different kinematics at the constant value $Q^2=2.5 \text{ GeV}^2$, by changing beam energy and detector angles. Preliminary results from both experiments are reported.

1 Introduction

The structure of the proton has been investigated experimentally and theoretically over the past 50 years using elastic electron scattering. The two Sachs form factors, G_{Ep} and G_{Mp} , describe the nucleon charge- and magnetization distribution; these form factors have been traditionally obtained by Rosenbluth separation method [1]. The G_{Mp} -data obtained by this method have shown good internal consistency up to 30 GeV^2 . However, the characterization of G_{Ep} has suffered from large inconsistencies in the data base, which are now understood to be the result of the fast decrease of the contribution of G_{Ep} to the cross section.

Indeed, the study of the electromagnetic form factors has become quite central to the rapid progress that is being made in hadronic physics. The results of the GEp-I and GEp-II experiments [2–4] have stimulated a huge amount of theoretical activity, as evidenced

by the nearly 1000 combined citations of these papers and several review papers of nucleon form factors [5–8]. An entirely new picture of the structure of the proton has emerged after the GEp-I and GEp-II experiments showed that the ratio G_{Ep}/G_{Mp} was in fact not constant, and decreased by a factor of 3.7 over the Q^2 range of 1 to 5.6 GeV^2 . These results are illustrated in Fig. 1, where they are also compared with Rosenbluth separation data [9–12].

The meaning of the results seen in Fig. 1 is that the spatial distribution of the electric charge of the proton is “softer”, i.e. larger in extent (in the Breit frame), than its magnetization currents distribution, which is definitively not intuitive. However, the relativistic boost required to transform these spatial distributions back to the laboratory frame are not trivial and only the form factors themselves are relativistic invariants. Recently G.A. Miller [14] has shown that an invariant charge distribution can only be defined on the wave front; the two-dimensional charge density on the wave front is the Fourier transform of the Dirac form factor, F_1 .

It is well known by now that G_{Ep} is difficult to obtain from Rosenbluth separation, a technique which is also especially sensitive to systematics errors and subject to large, ϵ -dependent radiative corrections. The two-photon exchange contribution, neglected in the past, has been shown to be an important term to add to the standard radiative corrections for cross section data; it has a strong ϵ -dependence and brings the Rosenbluth form factor ratio closer to the recoil polarization results [15, 16]. Two-photon contributions are expected to affect the recoil polarization results only very weakly [16].

Following the unexpected results from the polarization experiments, a new experiment, GEp-III [17], was approved to extend the Q^2 -range to 9 GeV^2 in Hall C; and to check the hypotheses of two-photon exchange contribution in ep elastic scattering, the GEp-2 γ experiment [18], using recoil polarization, was approved to measure the ratio at Q^2 of 2.5 GeV^2 for three different kinematics. Two new detectors were built by the collaboration to carry out these experiments; a large solid-angle electromagnetic calorimeter and a double focal plane polarimeter. In both experiments the recoil protons were detected in the high momentum spectrometer (HMS) equipped with the new focal plane polarimeter. The scattered electrons were detected in a new lead glass calorimeter (BigCal) built for this purpose out of 1744 glass bars, 4x4 cm^2 each, giving a total frontal area of 2.6 m^2 , which provides complete kinematical matching. This experiment finished taking data in the spring of 2008. The data analysis is in progress. In this paper, we will describe the recoil polarization method, the experimental results, and discuss the status of the proton elastic electromagnetic form factor data, including the latest results from the GEp-III and GEp-2 γ experiments, and compare them to a number of theoretical predictions.

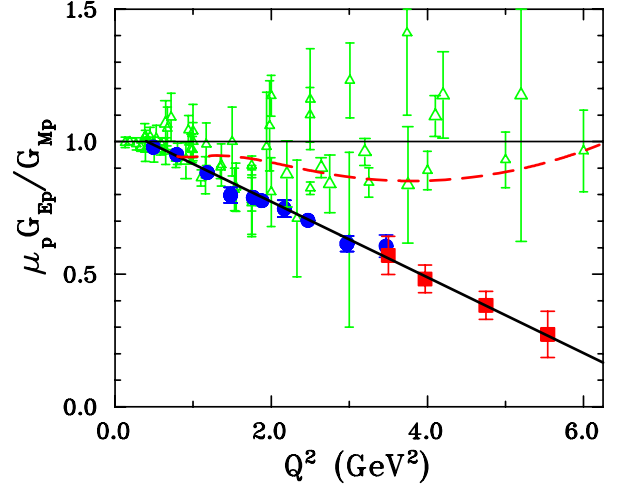


Figure 1: Comparison of $\mu_p G_{Ep}/G_{Mp}$ from the two JLab polarization data (filled circle and square) [3,4], and Rosenbluth separation data (empty triangle) [9–12]; dashed curve is a re-fit of Rosenbluth data [13]; solid curve is an updated form of the fit in ref. [4].

2 Recoil Polarization Method

With a longitudinally polarized electron beam and an un-polarized hydrogen target, the polarization of the incoming electron is transferred to the proton. The non-zero components of the recoil proton polarization are in the reaction plane, parallel, P_ℓ , and perpendicular, P_t , to the proton momentum [19, 20].

In the one-photon exchange process, the form factors depend only on Q^2 and a deviation from constant would indicate a mechanism beyond the Born approximation.

In the general case, elastic ep scattering can be described by three complex amplitudes [16, 21]: $\tilde{G}_M$, $\tilde{G}_E$, and $\tilde{F}_3$, the first two chosen as generalizations of the Sachs electric and magnetic form factors, G_E and G_M , and the last one, $\tilde{F}_3$, vanishing in case of Born approximation. The reduced cross section, σ_{red} , and the proton polarization transfer components P_t and P_l , including two-photon exchange formalism, can be written as [16]:

$$\sigma_{red}/G_M^2 = 1 + \frac{\varepsilon R^2}{\tau} + 2 \frac{\text{Re} \delta \tilde{G}_M}{G_M} + 2R\varepsilon \frac{\text{Re} \delta \tilde{G}_E}{\tau G_M} + 2 \left(1 + \frac{R}{\tau}\right) \varepsilon Y_{2\gamma} \quad (1)$$

$$P_t = -\sqrt{\frac{2\varepsilon(1-\varepsilon)}{\tau}} \frac{G_M^2}{\sigma_{red}} \left(R + R \frac{\text{Re} \delta \tilde{G}_M}{G_M} + \frac{\text{Re} \delta \tilde{G}_E}{G_M} + Y_{2\gamma} \right) \quad (2)$$

$$P_l = \sqrt{(1-\varepsilon^2)} \frac{G_M^2}{\sigma_{red}} \left(1 + 2 \frac{\text{Re} \delta \tilde{G}_M}{G_M} + \frac{2}{1+\varepsilon} \varepsilon Y_{2\gamma} \right), \quad (3)$$

where:

$$\text{Re} \tilde{G}_M(Q^2, \varepsilon) = G_M(Q^2) + \text{Re} \delta \tilde{G}_M(Q^2, \varepsilon) \quad (4)$$

$$\text{Re} \tilde{G}_E(Q^2, \varepsilon) = G_E(Q^2) + \text{Re} \delta \tilde{G}_E(Q^2, \varepsilon) \quad (5)$$

$$R(Q^2) = G_E(Q^2)/G_M(Q^2)$$

$$Y_{2\gamma}(Q^2, \varepsilon) = \sqrt{\frac{\tau(1+\tau)(1+\varepsilon)}{1-\varepsilon}} \frac{\text{Re} \tilde{F}_3(Q^2, \varepsilon)}{G_M(Q^2)} \quad (6)$$

Here $\tau = Q^2/4M_p^2$, and $\varepsilon = [1 + 2(1+\tau) \tan^2 \frac{\theta_e}{2}]^{-1}$, where θ_e is the lab electron scattering angle. While the Sachs form factors depend only on Q^2 , in the general case the amplitudes depend also on ε . The reduced cross section and the transferred proton polarization components are sensitive only to the real part of the amplitudes.

In Born approximation only the first term remains in the reduced cross section, σ_{red} , and the proton polarization transfer components P_t and P_l are :

$$\sigma_{red}/G_M^2 = 1 + \frac{\varepsilon R^2}{\tau} \quad (7)$$

$$P_t = -\sqrt{\frac{2\varepsilon(1-\varepsilon)}{\tau}} \frac{G_M^2 R}{\sigma_{red}}, \quad P_l = \sqrt{(1-\varepsilon^2)} \frac{G_M^2}{\sigma_{red}} \quad (8)$$

Both polarization components can be obtained simultaneously by measuring the azimuthal asymmetry of the recoil protons after re-scattering in an analyzer. The major advantage of this method, compared to cross section measurements, is that in Born approximation the form factor ratio $R = G_{Ep}/G_{Mp}$ is directly proportional to P_t/P_ℓ ,

$R = -P_t/P_l\sqrt{\tau(1+\varepsilon)/2\varepsilon}$; hence only a single measurement is necessary, strongly decreasing the systematic uncertainties.

In a polarimeter one can only measure polarization components normal to the analyzer, however use is made of the fact that the longitudinal component, P_ℓ , of the proton spin precesses in the dipole magnet of the HMS, and transforms into a longitudinal, P_ℓ^{fpp} as well as a normal component, P_n^{fpp} . In first approximation P_ℓ precesses by $\chi_\theta = \gamma\kappa_p(\Theta_{dipole} + \theta_{tar} - \theta_{fp})$, resulting in $P_n^{fpp} \approx P_\ell \times \sin \chi_\theta$ at the polarimeter. In addition, due to the presence of quadrupoles in the HMS, the longitudinal component also produces an additional transverse polarization component, due to the non-dispersive precession angle χ_ϕ . The effect of this precession, by $\chi_\phi = \gamma\kappa_p(\phi_{fp} - \phi_{tar})$, cancels only if the event distribution is symmetric in the angle difference $(\phi_{fp} - \phi_{tar})$.

The azimuthal asymmetry measured in the polarimeters is a function of the scattering angle in the analyzer, ϑ ; for a given bin of the incident proton momentum this distribution has the form:

$$f^\pm(\vartheta, \varphi) = \frac{\eta(\vartheta, \varphi)}{2\pi} (1 \pm P_{beam}(a_1 + \overline{A}_y P_y^{fpp}) \cos \varphi - P_{beam}(b_1 + \overline{A}_y P_x^{fpp} \sin \varphi + 0(n\varphi)) , \quad (9)$$

Here a_1 and b_1 are helicity independent asymmetries, $\eta(\vartheta, \varphi)$ is the polarimeter response efficiency, and $\pm$ stands for the two beam helicities; P_{beam} is the longitudinal beam polarization and $\overline{A}_y$ is the analyzing power averaged over the incident proton momentum and ϑ bin. To obtain the physical asymmetry we form the difference of helicity asymmetry, divided by the sum, $\frac{f^+ - f^-}{f^+ + f^-}$; the systematic asymmetries are canceled in first order with this procedure.

3 Analysis of GEp-III and GEp-2 γ Experiments

For a given beam energy E_e , the scattering angle θ_p and momentum p_p of the recoiling proton in elastic scattering are related by equation:

$$p_p = \frac{2M_p E_e (E_e + M_p) \cos \theta_p}{M_p^2 + 2M_p E_e + E_e^2 \sin^2 \theta_p} \quad (10)$$

The difference between the measured momentum and the momentum predicted by 10 at the measured angle, $\Delta_p = p - p_{el}(\theta_p)$, therefore defines the degree of “inelasticity” of a given event. Fig 2 shows distribution of events $p - p_{el}(\theta_p)$; elastic events are located at $\Delta_p = 0$, with a width determined by the HMS momentum and angular resolution, the beam energy, and the HMS central angle. As seen from Fig. 2, although it is possible to identify elastic and inelastic events based on the reconstructed proton momentum p_p and scattering angle θ_p , the resolution of the HMS is insufficient to achieve a clean separation.

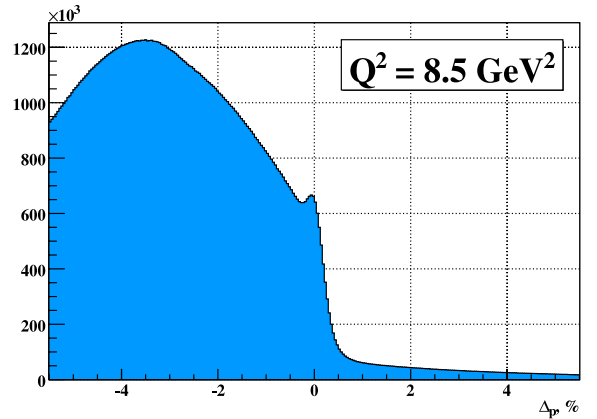


Figure 2: Δp spectrum for $Q^2 = 8.5 \text{ GeV}^2$ with no cuts.

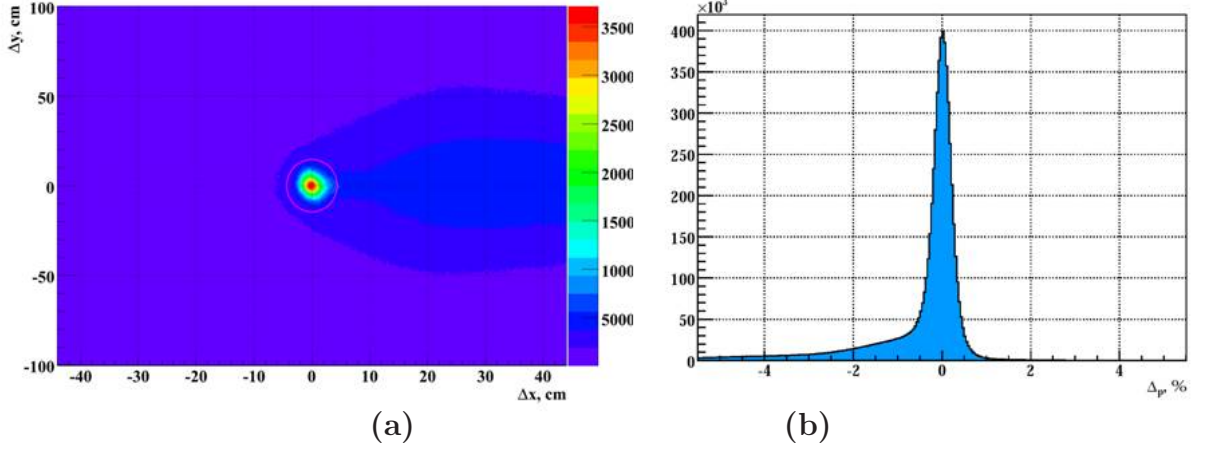


Figure 3: (a) Selection of elastic events using elliptical $(\Delta x, \Delta y)$ cut for $Q^2 = 8.5 \text{ GeV}^2$. (b) Δp spectrum for $Q^2 = 8.5 \text{ GeV}^2$ after $\Delta p - \Delta\theta$ and $(\Delta x, \Delta y)$ cut.

With significant background, kinematic correlation with the scattered electron becomes crucial for elastic ep scattering. A drastic background reduction is achieved by calculating the electron impact point (x, y) in BigCal from the proton momentum, angle and the beam energy assuming two body ep scattering, and comparing it with the observed impact point. Elastic events are found at $(\Delta x, \Delta y) = (0, 0)$ as shown in figure 3a. Figure 3b shows the momentum spectrum after Δp and $(\Delta x, \Delta y)$ cuts.

Figure 4 shows the variation of the extracted form factor ratio as a function of the estimated background within the cuts for the Q^2 of 8.5 GeV^2 point. For this setting, the background increases much more rapidly as a function of the cut width, which is apparent in the rapid decrease of the uncorrected form factor ratio. The corrected form factor ratio in this case is still independent of the background correction up to statistical fluctuations in the polarization of the extra events accepted by wider cuts and the uncertainty in the background polarization used for the correction, demonstrating that

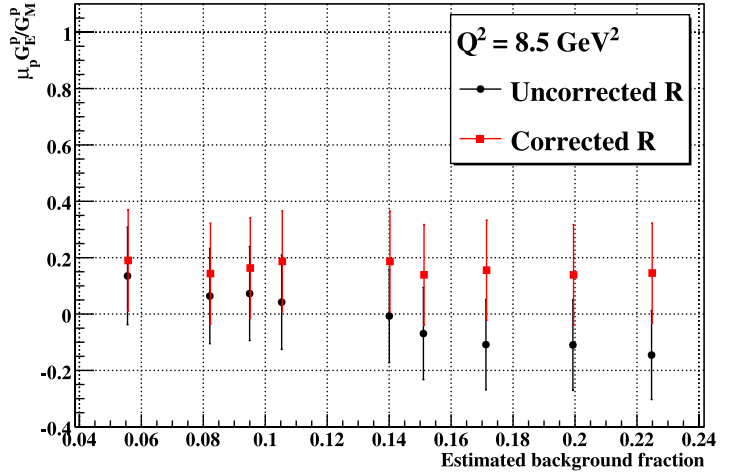


Figure 4: Cut width dependence of the extracted form factor ratio, at $Q^2 = 8.5 \text{ GeV}^2$.

the background subtraction is handled correctly.

4 Results of GEp-III Experiment

In Fig. 5 we show the preliminary results for the GEp-III experiment. Because GEp-III is the first recoil polarization experiment to use a completely different apparatus in a Q^2 range where direct comparison with the Hall A recoil polarization results is possible, it provides an important test of the reproducibility of the recoil polarization technique. Figure 5 shows the results of GEp-III experiment (filled triangles) together with the results of the Hall A experiments [2,3] (filled circles) and [4] (filled squares). The error bars shown are statistical only. The overlap point at 5.2 GeV^2 is in reasonably good agreement with the two surrounding points from [4], and confirms the continuing decrease of G_E^p/G_M^p with Q^2 . The two new data points extend the knowledge of G_E^p/G_M^p to yet higher Q^2 .

Additionally, the preliminary results of the high-statistics survey of the ϵ -dependence of G_E^p/G_M^p at $Q^2 = 2.5 \text{ GeV}^2$, GEp-2 γ experiment, also shown as black triangle in Fig. 5 is in excellent agreement with the Hall A results [3] at $Q^2 = 2.47 \text{ GeV}^2$. Since the GEp-2 γ experiment was performed with identical apparatus and the data were analyzed in exactly the same way as in the GEp-III experiment, a systematic error in the analysis is all but ruled out.

The interesting feature of the results is that the linearly decreasing trend observed in the previous data appears to be slowing in the Hall C data. Although the ratios still decreases with Q^2 , all of the new data points come in higher than the linear fit to the Hall A data.

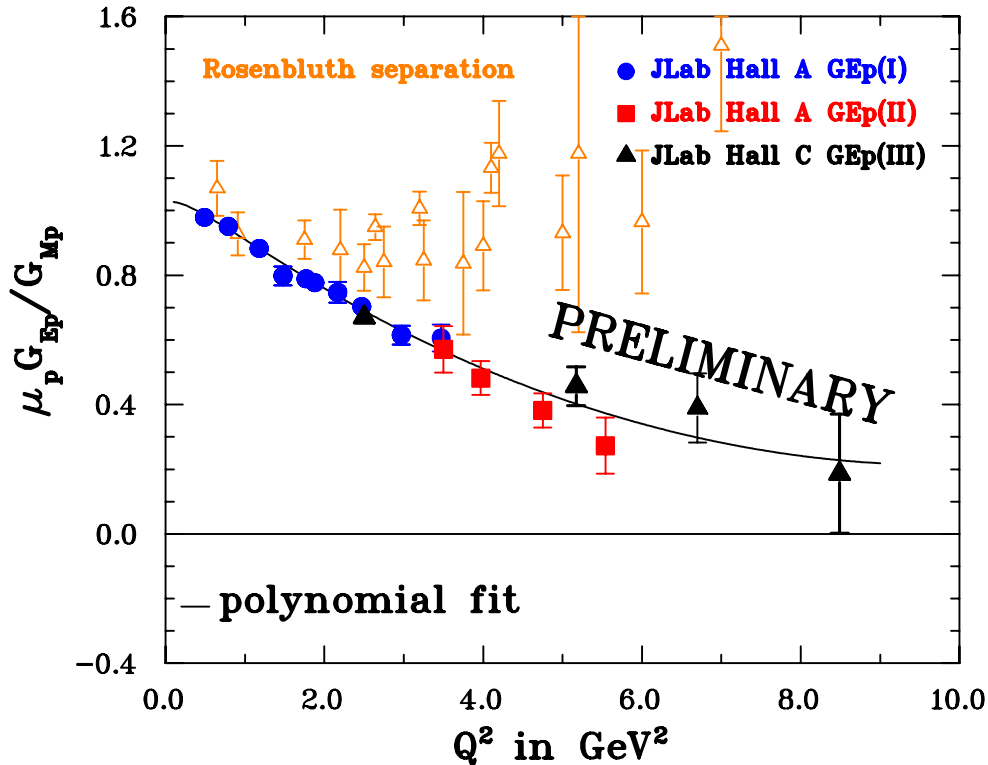


Figure 5: Results of GEp-III experiment (filled triangles) shown as $\mu_p G_E^p / G_M^p$. The results of experiments GEp-I (filled circles) and GEp-II (filled squares) and the Rosenbluth separation data (empty triangles) are shown for comparison. The error bars on the data points are statistical only.

5 Results of GEp-2 γ Experiment

So what are the causes for the different results for $\mu G_{Ep}/G_{Mp}$, from cross section and from polarization measurements? No experimental explanation has been found [13]. Part of the answer may be that radiative corrections at large Q^2 are both large and strongly ϵ dependent, and these corrections are important for cross section data, but not for polarization data.

The JLab polarization results have led to a reexamination of the two-photon contribution, which is included in standard radiative corrections only in the limit of one of the photon energies being small. Recent calculations of the contribution of two-photon exchange with both photons approximately sharing the momentum transfer, first by Guichon *et al* [21], then Afanasev *et al* [16], and Blunden, Melnitchouk and Tjon [15], have shown that it might contribute a correction to the standard radiative correction of several % at $\epsilon = 1$.

The preliminary results from the GEp-2 γ experiment for the proton form factor ratio as extracted from the polarization ratio are shown in Figure 6. The error bars shown are statistical only. The data points lay almost on the same constant line thus showing no deviations from the Born approximation at a percent level. On the same plot several theoretical predictions are shown. These are corrections to the Born approximation. In [16] the two-photon exchange contributions are calculated at the partonic level assuming factorization of the soft nucleon-quark part, and the hard electron-quark interaction where two-photon exchange takes place via box diagram. In another approach [15, 22, 23], the two-photon exchange effects are calculated at the hadronic level.

Similar two-photon exchange box diagrams are used where the quarks are replaced by nucleons. Kivel and Vanderhaeghen [24] estimate the two photon exchange contribution to elastic electron-proton scattering, using several models for the nucleon distribution amplitudes [25, 26].

The key idea of the GEp-2 γ experiment, is to study the ϵ dependence of the recoil proton polarization at a fixed value of Q^2 , hence, (1) the proton momentum is same and as a consequence the analyzing power, A_y , of the reaction used to measure the proton polarization, is the same for all the kinematics; and (2) since the setting of the HMS is fixed, the transport of the spin from the focal plane where it was measured, back to the target is the same for all the data points. This is a significant advantage in the measurement of the polarization ratio P_t/P_l . For these reasons the systematic uncertainty will be quite small for the final GEp-2 γ results.

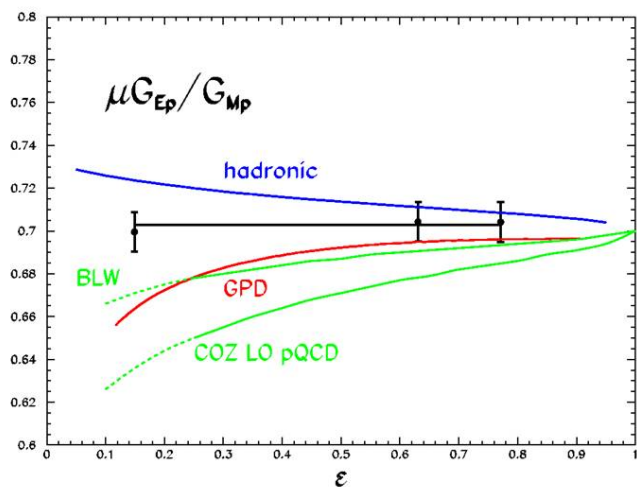


Figure 6: Form factor ratio as measured in GEp-2 γ experiment. Theoretical predictions labeled as GPD are from Afanasev *et al.* [16]; labeled as hadronic are from Blunden *et al.* [15], and labeled as BLW and COZ are calculations done by Kivel and Vanderhaeghen using proton distribution amplitudes from references [25] and [26].

6 Theoretical Developments

The earliest models explaining the global features of the nucleon form factors, such as its approximate dipole behavior, were vector meson dominance (VMD) models. In this picture the photon couples to the nucleon through the exchange of vector mesons. Such VMD models are a special case of more general dispersion relation approach, which allows to relate time-like and space-like form factors. An early VMD fit was performed by Iachello *et al.* [27] and predicted a linear decrease of the proton G_{Ep}/G_{Mp} ratio, which is in basic agreement with the result from the polarization transfer experiments. Such VMD models have been extended by Gari and Krümpelmann [28] to include the perturbative QCD (pQCD) scaling relations [29], which state that $F_1 \sim 1/Q^4$, and $F_2/F_1 \sim 1/Q^2$.

In more recent years, extended VMD fits which provide a relatively good parametrization of all nucleon electromagnetic form factors have been obtained. An example is the fit of Lomon [30], containing 11 parameters. Another such recent parametrization by Bijker and Iachello [31] achieves a good fit by adding a phenomenological contribution attributed to a quark-like intrinsic qqq structure (of *rms* radius ~ 0.34 fm) besides the vector-meson exchange terms. The pQCD scaling relations are built into this fit which has 6 free parameters. In contrast to the early fit of Ref. [27], the new fit of Ref. [31] gives a very good description of the neutron data, albeit at the expense of a slightly worse fit for the proton data. Figure 7 shows predictions from References [27,30,31]; all three fits describe the data, however the best fit is from [30], which goes through all the data points.

Among the theoretical efforts to understand the structure of the nucleon in terms of quark and gluon degrees of freedom, constituent quark models have a long history too. In these models, the nucleon consists of three constituent quarks, which are thought to be valence quarks dressed with gluons and quark-antiquark pairs, and are much heavier than the QCD Lagrangian quarks. All other degrees of freedom are absorbed into the masses of these quarks. To describe the data presented here in terms of constituent quarks, it is necessary to include relativistic effects because the momentum transfers involved are up to 10 times larger than the constituent quark mass.

In the earliest study of the relativistic constituent quark models (RCQM), Chung and Coester [32] calculated electromagnetic nucleon form factors with Poincaré-covariant constituent-quark models and investigated the effect of the constituent quark masses, the anomalous magnetic moment of the quarks, and the confinement scale parameter.

Frank *et al.* [33] have calculated G_{Ep} and G_{Mp} in the light-front constituent quark model and predicted that G_{Ep} might change sign near 5.6 GeV^2 ; this calculation used

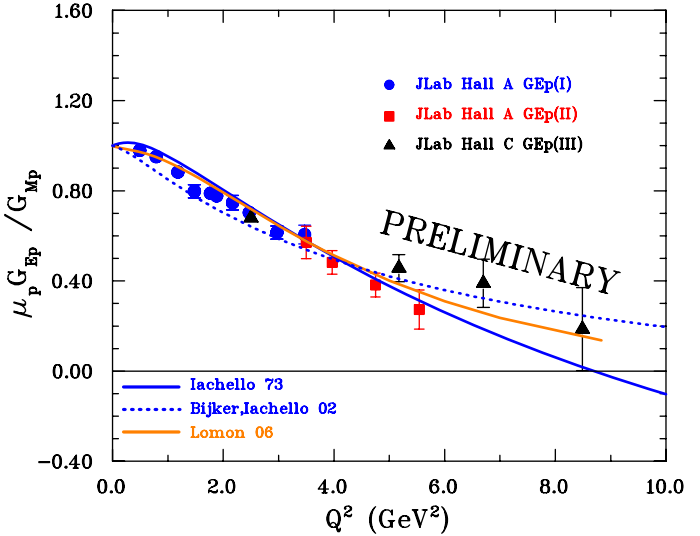


Figure 7: Theoretical calculations with VMD models, compared to the data from GEp-I (filled circles), GEp-II (filled squares), and GEp-III (filled triangles) experiments.

the light-front nucleonic wave function of Schlumpf [34]. Under such a transformation, the spins of the constituent quarks undergo Melosh rotations. These rotations, by mixing spin states, play an important role in the calculation of the form factors.

Cardarelli *et al.* [35] calculated the ratio with light-front dynamics and investigated the effects of SU(6) symmetry breaking. They showed that the decrease in the ratio with increasing Q^2 is due to the relativistic effects generated by Melosh rotations of the constituent quark's spin. In Ref. [36], they pointed out that within the framework of the RCQM with the light-front formalism, an effective one-body electromagnetic current, with a proper choice of constituent quark form factors, can give a reasonable description of pion and nucleon form factors. The chiral constituent quark model based on Goldstone-boson-exchange dynamics was used by Boffi *et al.* [37] to describe the elastic electromagnetic and weak form factors.

More recently Gross and Agbakpe [38] revisited the RCQM imposing that the constituent quarks become point particles as $Q^2 \rightarrow \infty$ as required by QCD; using a covariant spectator model which allows exact handling of all Poincaré transformations, and monopole form factors for the constituent quarks, they obtain excellent ten parameter fits to all four nucleon form factors; they conclude that the recoil polarization data can be fitted with a spherically symmetric state of 3 constituent quarks. Figure 8 shows predictions from References [32, 33, 35–38]; in all predictions the ratio decrease with Q^2 , showing the same trend as the data, with the exception of the predictions of Boffi *et al.* [37] and Frank *et al.* [33].

The nucleon electromagnetic form factors provide a famous test for perturbative QCD. Brodsky and Farrar [39] derived scaling rules for dominant helicity amplitudes which are expected to be valid at sufficiently high momentum transfers Q^2 . A photon of sufficient high virtuality will interact with a nucleon consisting of three mass less quarks moving collinearly with the nucleon. When measuring an elastic nucleon form factor, the final state consists again of three mass less collinear quarks. In order for this process to happen, the large momentum of the virtual photon has to be transferred among the three quarks through two hard gluon exchanges. Because each gluon in such a hard scattering process carries a virtuality proportional to Q^2 , this leads to the pQCD prediction that the helicity conserving nucleon Dirac form factor F_1 should fall as $1/Q^4$ at sufficiently high Q^2 . In contrast to the helicity conserving form factor F_1 , the Pauli form factor F_2 involves a helicity flip between the initial and final nucleons. Hence it requires one helicity flip at the quark level, which is suppressed at large Q^2 . Therefore, for collinear quarks, i.e. moving in a light-cone wave function state with orbital angular momentum projection

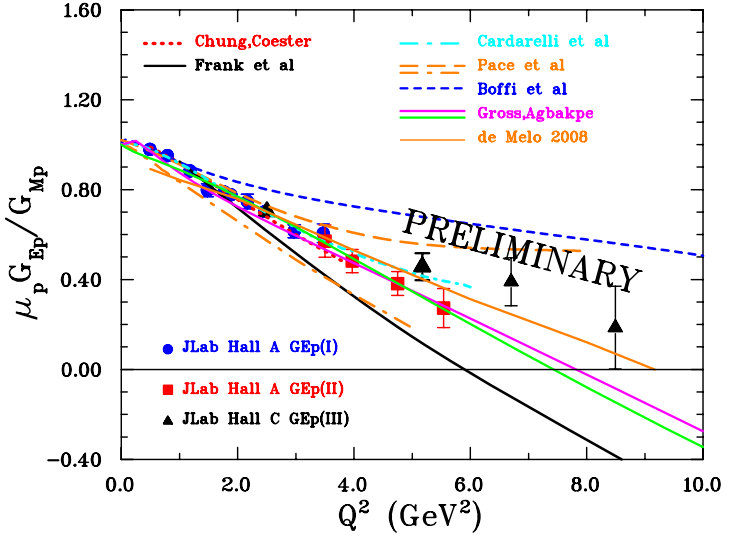


Figure 8: Theoretical calculations with RCQM models, compared to the data from GEp-I (filled circles), GEp-II (filled squares), and GEp-III (filled triangles) experiments.

$l_z = 0$ (along the direction of the fast moving hadron), the asymptotic prediction for F_2 leads to a $1/Q^6$ fall-off at high Q^2 .

In Fig. 9 the JLab recoil polarization data together with cross section data from references [9–11] are shown as $Q^2 F_{2p}/F_{1p}$. The cross section data from Ref. [9] show flattening above Q^2 of 3 GeV²; the recoil polarization data from JLab do not show yet the specific pQCD Q^2 dependence, but the new data from GEp-III suggest a decrease in the slope, hinting that the data might be approaching the pQCD limit.

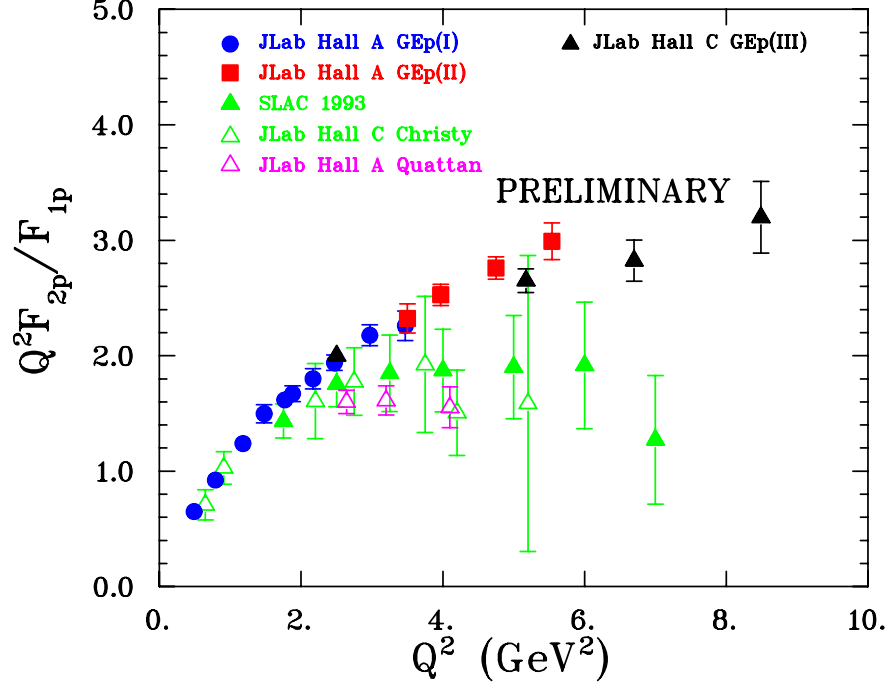


Figure 9: $Q^2 F_{2p}/F_{1p}$ versus Q^2 for data from GEp-I (filled circles), GEp-II (filled squares), and GEp-III (filled triangles) experiments.

In another approach, Belitsky, Ji, and Yuan [40] investigated the assumption of quarks moving collinearly with the proton, underlying the pQCD prediction. It has been shown in [40] that by including components in the nucleon light-cone wave functions with quark orbital angular momentum projection $l_z = 1$, one obtains the behavior $Q^2 F_2/F_1 \rightarrow \ln^2(Q^2/\Lambda^2)$ at large Q^2 , with Λ a non-perturbative mass scale. Choosing Λ around 0.3 GeV, Ref. [40] noticed that the data for $Q^2 F_{2p}/F_{1p}$ support such double-logarithmic enhancement, as can be seen from Figure 10.

Lattice QCD simulations have the potential to calculate nucleon form factors from first principles. This is a rapidly developing field and important progress has been made in the recent past. Nevertheless, lattice calculations are at present still severely limited by available computing power and in practice are performed for quark masses sizably larger than their values in nature.

The generalized parton distributions (GPDs) provide a framework to describe the process of emission and re-absorption of a quark in the non-perturbative region by a hadron in exclusive reactions. The GPDs can be interpreted as quark correlation functions and have the property that their first moments exactly coincide with the nucleon form factors. Precise measurements of elastic nucleon form factors provide stringent constraints on the parametrization of the GPDs. Early theoretical developments in GPDs indicated

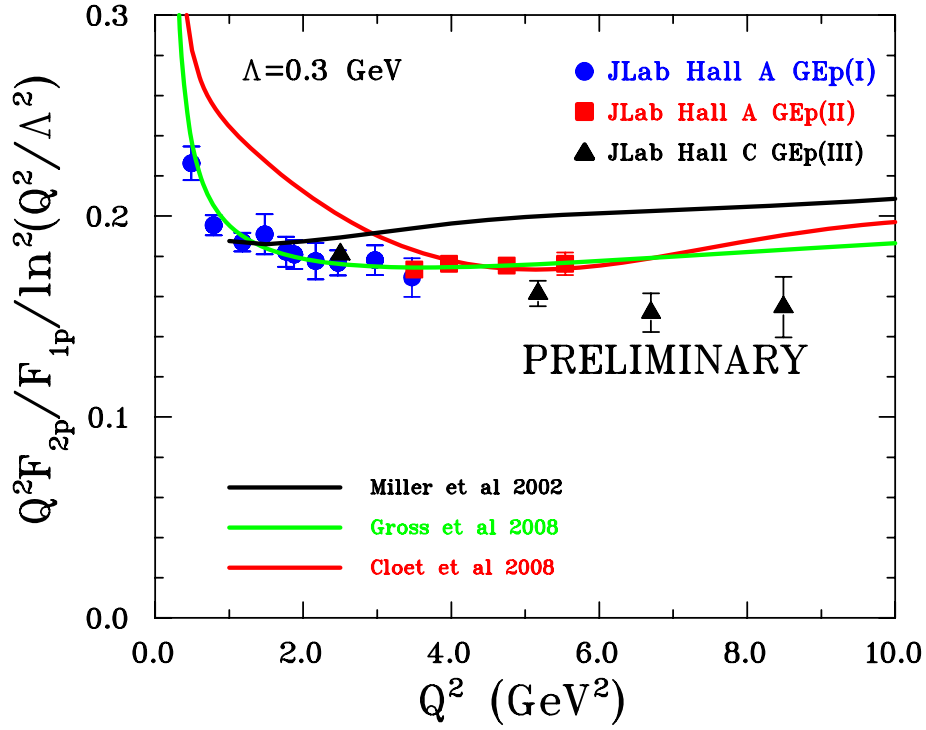


Figure 10: The test of the modified scaling prediction for $Q^2 F_{2p}/F_{1p}$ versus Q^2 [40] for data from GEp-I (filled circles), GEp-II (filled squares), and GEp-III (filled triangles) experiments.

that measurements of the separated elastic form factors of the nucleon to high Q^2 may shed light on the problem of nucleon spin. The first moment of the GPDs taken in the forward limit yields, according to the Angular Momentum Sum Rule [41], a contribution to the nucleon spin from the quarks and gluons, including both the quark spin and orbital angular momentum. The t -dependence of the GPDs has been modeled using a factor corresponding to the relativistic Gaussian dependence of both Dirac [42] and Pauli [43] form factors of the proton. Extrapolation of these GPDs to $t = 0$ leads to the functions entering into the Angular Momentum Sum Rule, and an estimate of the contribution of the valence quarks to the proton spin can then be obtained.

In a recent development of these ideas Guidal *et al* [44] have shown that the difficulties encountered in the Gaussian parametrization used in earlier work could be surmounted with a Regge parametrization. In one version of their parametrization, these authors obtain excellent fits to all 4 nucleon form factors. In a different approach Diehl *et al* [45] use theoretically motivated parametrization of the relevant GPDs in the very small and very large x -domains, and interpolate by fitting the nucleon Dirac form factors F_1^p and F_1^n . They derive the valence contribution to Ji's sum rule [41]. In a related approach, Ji [46] has shown that the GPDs provide a classical visualization of the quark orbital motion.

7 Conclusions

Form factor data are of great interest as a testing ground for QCD, as results from lattice QCD calculations become increasingly accurate and realistic. Phenomenological

models have recently been challenged by the elastic form factors obtained at Jefferson Lab, resulting in an intense discussion of questions related to the shape of the proton [47, 48], and the contribution of the quark orbital angular momentum to its spin, for example see Ref. [41].

The new data from the GEp-III experiment at higher Q^2 , although still preliminary, show a slowing decrease of G_E^p/G_M^p with Q^2 relative to the linear decrease observed in the Hall A data for $Q^2 \leq 5.6 \text{ GeV}^2$. Although the statistical significance of this change in behavior is somewhat marginal, its physical implications are interesting to consider. A constant ratio G_E^p/G_M^p at asymptotically large Q^2 is a signature of the onset of the dimensional scaling expected from perturbative QCD for a nucleon consisting of three weakly interacting quarks.

The final answer to the question of the origin of the discrepancy between the results from Rosenbluth separation and polarization transfer will have to come from experiments. There are several directly- and indirectly related experiments designed to provide answer to the question. For example, (1) the GEp-2 γ experiment measured ϵ -dependence of the G_{Ep}/G_{Mp} ratio at constant Q^2 , obtained from recoil polarization in Hall C [18]. The preliminary results for the $\mu_p G_{Ep}/G_{Mp}$ ratio at $Q^2=2.49 \text{ GeV}^2$, at the 3 values of ϵ 0.14, 0.63 and 0.79, were shown in Fig. 6; no deviation from constancy is seen at the level of 0.01 absolute, showing that the recoil polarization results follow the Born approximation at this Q^2 . (2) Cross section difference in e^+ and e^- proton scattering in Hall B [49] (also approved experiment OLYMPUS at DESY [50]) ; this difference is twice the contribution of a two-photon exchange; and (3) Measuring non-linearity of the Rosenbluth plot; experiment E05-017 measured the cross sections with high precision at several Q^2 values as a function of ϵ to extract the ratio; data taking for this experiment was completed in July 2007 [51].

The measurement of nucleon form factors to the highest possible Q^2 is one of the prime tenets of the JLab 12 GeV upgrade. Two new proposals, GEp-IV [52] and GEp-V [53], were submitted to JLab PAC 34 and 32, respectively to extend the proton form factor ratio measurements, when a 12 GeV beam becomes available. In Experiment GEp-V the ratio will be measured in Hall A to 14.8 GeV^2 , and in the GEp-IV experiment in Hall C, it will be measured to 13 GeV^2 . These planned experiments covering the Q^2 range from 10-15 GeV^2 following the 12 GeV upgrade of the CEBAF accelerator, will answer the question of whether the flattening hinted at by the GEp-III experimental results is real or whether G_E^p/G_M^p continues to decrease and eventually cross zero.

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LONGITUDINAL SPIN TRANSFER TO Λ AND $\bar{\Lambda}$ IN POLARIZED PROTON-PROTON COLLISIONS AT $\sqrt{s} = 200$ GeV

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Abstract

We report our measurement on longitudinal spin transfer, D_{LL} , from high energy polarized protons to Λ and $\bar{\Lambda}$ hyperons in proton-proton collisions at $\sqrt{s} = 200$ GeV with the STAR detector at RHIC. The measurements cover Λ , $\bar{\Lambda}$ pseudorapidity $|\eta| < 1.2$ and transverse momenta, p_T , up to 4 GeV/ c . The longitudinal spin transfer is found to be $D_{LL} = -0.03 \pm 0.13(\text{stat}) \pm 0.04(\text{syst})$ for inclusive Λ and $D_{LL} = -0.12 \pm 0.08(\text{stat}) \pm 0.03(\text{syst})$ for inclusive $\bar{\Lambda}$ hyperons with $\langle \eta \rangle = 0.5$ and $\langle p_T \rangle = 3.7$ GeV/ c . The p_T dependence of D_{LL} for positive and negative η is given.

The longitudinal polarization of Λ hyperons has been studied in e^+e^- annihilation and lepton-nucleon deep inelastic scattering (DIS) with polarized beams and/or targets. Such polarization studies provide access to polarized fragmentation function and the spin content of Λ [1]. Here we report the first measurement on longitudinal spin transfer (D_{LL}) from the proton beam to $\Lambda(\bar{\Lambda})$ produced in proton-proton collisions at $\sqrt{s}=200$ GeV [2],

$$D_{LL} \equiv \frac{\sigma_{p^+p \rightarrow \Lambda^+X} - \sigma_{p^+p \rightarrow \Lambda^-X}}{\sigma_{p^+p \rightarrow \Lambda^+X} + \sigma_{p^+p \rightarrow \Lambda^-X}}, \quad (1)$$

where the superscript + or - denotes the helicity. Within the factorization framework, the production cross sections are described in terms of calculable partonic cross sections and non-perturbative parton distribution and fragmentation functions. The production cross section has been measured for transverse momenta, p_T , up to about 5 GeV/ c and is well described by perturbative QCD evaluations [3]. The spin transfer D_{LL} is thus expected to be sensitive to polarized fragmentation function and helicity distribution function of nucleon, as reflected in different model predictions of D_{LL} at RHIC [4–7].

The spin transfer D_{LL} in Eq. (1) is equal to the polarization of Λ ($\bar{\Lambda}$) hyperons, $P_{\Lambda(\bar{\Lambda})}$, if the proton beam is fully polarized. $P_{\Lambda(\bar{\Lambda})}$ can be measured via the weak decay channel $\Lambda \rightarrow p\pi^-$ ($\bar{\Lambda} \rightarrow \bar{p}\pi^+$) from the angular distribution of the final state,

$$\frac{dN}{d\cos\theta^*} = \frac{\sigma \mathcal{L} A}{2} (1 + \alpha_{\Lambda(\bar{\Lambda})} P_{\Lambda(\bar{\Lambda})} \cos\theta^*), \quad (2)$$

where σ is the (differential) production cross section, $\mathcal{L}$ is the integrated luminosity, A is the detector acceptance, which is in general a function of $\cos\theta^*$ as well as other observables, $\alpha_{\Lambda} = -\alpha_{\bar{\Lambda}} = 0.642 \pm 0.013$ [8] is the weak decay parameter, and θ^* is the angle between the $\Lambda(\bar{\Lambda})$ polarization direction and the (anti-)proton momentum in the Λ ($\bar{\Lambda}$) rest frame.

The data presented here were collected at the Relativistic Heavy Ion Collider (RHIC) with the Solenoidal Tracker at RHIC (STAR) [9] in the year 2005. An integrated luminosity of 2 pb⁻¹ was sampled with average longitudinal beam polarization of $52 \pm 3\%$ and

$48 \pm 3\%$ for two beams. The proton polarization was measured for each beam and each beam fill using Coulomb-Nuclear Interference (CNI) proton-carbon polarimeters [10]. Different beam spin configurations were used for successive beam bunches and the pattern was changed between beam fills. The data were sorted by beam spin configuration.

The analyzed data sample includes three different triggers. One is the minimum bias (MB) trigger sample, defined with a coincidence signal from Beam-Beam Counters (BBC) at both sides of STAR interaction region. The other data samples were recorded with the MB trigger condition and with two additional trigger conditions: a high-tower (HT) and a jet-patch (JP). The HT trigger condition required the BBC proton collision signal in coincidence with a transverse energy deposit $E_T > 2.6$ GeV in at least one Barrel Electromagnetic Calorimeter (BEMC) tower, covering $\Delta\eta \times \Delta\phi = 0.05 \times 0.05$ in pseudorapidity η and azimuthal angle ϕ . The JP trigger condition imposed the MB condition in coincidence with an energy deposit $E_T > 6.5$ GeV in at least one of six BEMC patches each covering $\Delta\eta \times \Delta\phi = 1 \times 1$. The total BEMC coverage was $0 < \eta < 1$ and $0 < \phi < 2\pi$ in 2005.

The Λ ($\bar{\Lambda}$) candidates were identified from the topology of the weak decay channel to $p\pi^-$ ($\bar{p}\pi^+$), which has a branching ratio of 63.9% [8]. Charged particle tracks in the 0.5 T magnetic field were measured with the Time Projection Chamber (TPC), covering $0 < \phi < 2\pi$ and $|\eta| \lesssim 1.3$. The charge tracks after particle identification from specific energy loss dE/dx in TPC were paired to form a Λ ($\bar{\Lambda}$) candidate and topological selections were applied to reduce background [2, 3]. The selections included criteria for the distance of closest approach between the paired tracks and the distance between the point of closest approach and the beam collision vertex, and demanded that the momentum sum of the track pair pointed at the collision vertex. The criteria were tuned to preserve the signal while reducing the background fraction to 10% or less.

Figure 1(a) shows the invariant mass distribution for the Λ (filled circles) and $\bar{\Lambda}$ (open circles) candidates reconstructed from MB data with $|\eta| < 1.2$ and $0.3 < p_T < 3$ GeV/c. The mean values of the Λ and $\bar{\Lambda}$ mass distributions are in agreement with the PDG mass value $m_{\Lambda(\bar{\Lambda})} = 1.11568$ GeV/c² [8]. Figure 1(b) shows the same invariant mass distribution versus $\cos\theta^*$ for the Λ candidates. The number of Λ candidates varies with $\cos\theta^*$ because of detector acceptance. The small variation of the reconstructed invariant mass with $\cos\theta^*$ is understood to originate from detector resolution. In addition to signal, combinatorial background is seen as well as backgrounds of misidentified e^+e^- pairs at low invariant mass values near $\cos\theta^* = -1.0$ and of misidentified K_S^0 in a diagonal band at high invariant mass values and $\cos\theta^* > -0.2$. About 1.2×10^4 Λ and 1.0×10^4 $\bar{\Lambda}$ candidates from the

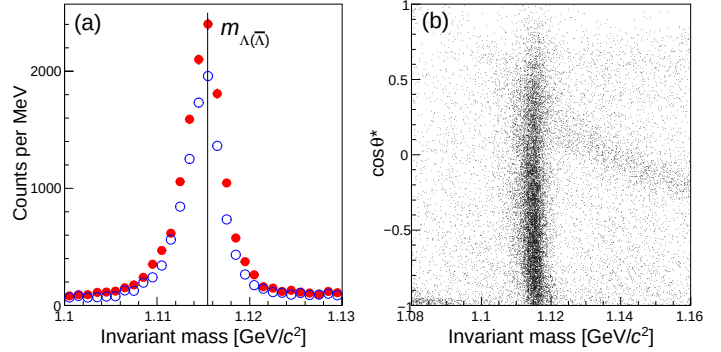


Figure 1: (a) The invariant mass distribution of Λ (filled circles) and $\bar{\Lambda}$ (open circles) candidates from reconstructed $p + \pi^-$ and $\bar{p} + \pi^+$ track pairs in 2005 MB data after topological selections. (b) The invariant mass distribution versus $\cos\theta^*$ for Λ .

MB sample were reconstructed in the invariant mass range $1.109 < m < 1.121 \text{ GeV}/c^2$ and were kept for further analysis.

To minimize the uncertainty associated with acceptance, the longitudinal spin transfer D_{LL} was extracted in small $\cos\theta^*$ intervals as follows:

$$D_{LL} = \frac{1}{\alpha P_{\text{beam}} \langle \cos\theta^* \rangle} \frac{N^+ - RN^-}{N^+ + RN^-} \quad (3)$$

where P_{beam} is the beam polarization, N^+ (N^-) are the Λ ($\bar{\Lambda}$) counts in a small $\cos\theta^*$ interval when the beam is positively (negatively) polarized, $\langle \cos\theta^* \rangle$ is the average value in this $\cos\theta^*$ bin, and $R = L^+/L^-$ is the corresponding luminosity ratio for these two polarization states. Eq.(3) uses the parity conservation in the hyperon production in pp collisions, which leads to a sign flip of hyperon polarization when the beam helicity is flipped. The detector acceptance in a $\cos\theta^*$ interval is taken as a constant and thus canceled. For this measurement, only one polarized beam is needed. At RHIC both beams are polarized and the single spin yields N^+ (N^-) are thus formed from the yields n^{++} , n^{+-} , n^{-+} , and n^{--} by beam helicity configuration weighted with their corresponding relative luminosities. The luminosity ratios were measured with BBC at STAR [11].

The yields N^+ and N^- were determined for each $\cos\theta^*$ interval from the observed Λ and $\bar{\Lambda}$ candidate yields in the mass interval from 1.109 to 1.121 GeV/c^2 . The corresponding raw values D_{LL}^{raw} were averaged over 20 intervals covering the entire $\cos\theta^*$ range. The obtained D_{LL}^{raw} values and their statistical uncertainties were then corrected for (unpolarized) background dilution according to $D_{LL} = D_{LL}^{\text{raw}} / (1 - r)$, where r is the average background fraction. No significant spin transfer asymmetry was observed for the yields in the sideband mass intervals $1.094 < m < 1.103 \text{ GeV}/c^2$ and $1.127 < m < 1.136 \text{ GeV}/c^2$, and thus no further correction was applied to D_{LL} . However, a contribution was included in the estimated systematic uncertainty of the D_{LL} measurement to account for the possibility that the background could nevertheless be polarized.

The combined D_{LL} results from the MB data sample versus $\cos\theta^*$ are shown in Fig. 2(a) for Λ and Fig.2(b) for $\bar{\Lambda}$ hyperons with $0.3 < p_T < 3 \text{ GeV}/c$ and $0 < \eta < 1.2$ and $-1.2 < \eta < 0$. Positive η is defined along the direction of the incident polarized beam. Fewer than 50 counts were observed for $\cos\theta^* > 0.9$ and this interval was discarded for this reason. The extracted D_{LL}

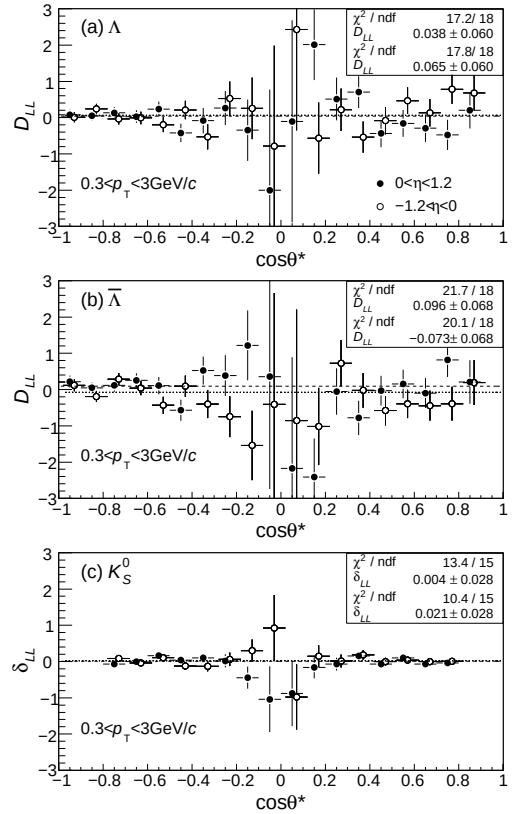


Figure 2: The spin transfer D_{LL} versus $\cos\theta^*$ for (a) Λ , (b) $\bar{\Lambda}$ and (c) the spin asymmetry δ_{LL}^0 for the control sample of K_S^0 mesons versus $\cos\theta^*$. Only statistical uncertainties are shown. The data points with negative η have been shifted slightly in $\cos\theta^*$ for clarity.

is constant with $\cos\theta^*$, as expected and confirmed by the quality of fit. In addition, a null-measurement was performed of the spin transfer for the spinless K_S^0 meson, which has a similar event topology. The K_S^0 candidate yields for $|\cos\theta^*| > 0.8$ were discarded since they have sizable $\Lambda(\bar{\Lambda})$ backgrounds. The result, δ_{LL} , obtained with an artificial weak decay parameter $\alpha_{K_S^0} = 1$, was found consistent with no spin transfer, as shown in Fig. 2(c).

The HT and JP data samples were recorded with trigger conditions that required large energy deposits in the BEMC, in addition to the MB condition. These triggers, however, did not require a highly energetic Λ or $\bar{\Lambda}$. To minimize the effects of this bias, the HT event sample was restricted to Λ or $\bar{\Lambda}$ candidates whose decay (anti-)proton track intersected a BEMC tower that fulfilled the trigger condition. About 50% of the $\bar{\Lambda}$ and only 3% of the Λ candidate events in the analysis pass this selection. This is qualitatively consistent with the annihilation of anti-protons in the BEMC. The $\bar{\Lambda}$ sample that was selected in this way thus directly triggered the experiment read-out. It contains about 1.0×10^4 $\bar{\Lambda}$ candidates with $1 < p_T < 5$ GeV/c.

For the JP triggered sample, events were selected with at least one reconstructed jet that pointed to a triggered jet patch. The same jet reconstruction was used as in Ref. [12]. Jets outside the BEMC acceptance were rejected. The Λ and $\bar{\Lambda}$ candidates whose reconstructed η and ϕ fell within the jet cone of radius $r_{\text{cone}} = \sqrt{(\Delta\eta)^2 + (\Delta\phi)^2} = 0.4$ were retained for further analysis. In most cases, the decay (anti-)proton track is part of the reconstructed jet, whereas the decay pion track is not. No correction was made to the jet finding and reconstruction for this effect. About 1.3×10^4 Λ and 2.1×10^4 $\bar{\Lambda}$ candidates with $1 < p_T < 5$ GeV/c remain after selections.

Figure 3 shows the D_{LL} results of Λ and $\bar{\Lambda}$ versus p_T for positive and negative η respectively. The $\bar{\Lambda}$ results from HT and JP data have been combined. No corrections have been applied for possible decay contributions from heavier baryonic states. The Λ and $\bar{\Lambda}$ results for D_{LL} are consistent with each other and consistent with no spin transfer from the polarized proton beam to the produced Λ and $\bar{\Lambda}$ to within the present uncertainties. The data have p_T up to 4 GeV/c, where $D_{LL} = -0.03 \pm 0.13(\text{stat}) \pm 0.04(\text{syst})$ for the Λ and $D_{LL} = -0.12 \pm 0.08(\text{stat}) \pm 0.03(\text{syst})$ for the $\bar{\Lambda}$ at $\langle\eta\rangle = 0.5$. For reference, the model predictions of Refs. [4, 7], evaluated at $\eta = \pm 0.5$ and $p_T = 4$ GeV/c, are shown as horizontal lines. The expectations of Ref. [4] hold for Λ and $\bar{\Lambda}$ combined and examine different polarized fragmentation scenarios, in which the strange (anti-)quark carries all or only part of the Λ ($\bar{\Lambda}$) spin. The model in Ref. [7] separates Λ from $\bar{\Lambda}$ and otherwise distinguishes the direct production of the Λ and $\bar{\Lambda}$ from the (anti-)quark in the hard scattering and the indirect production via decay of heavier (anti-)hyperons. The evaluations are consistent with the present data and span a range of values that, for positive η , is similar to the experimental uncertainties. The values for negative η are expected to be negligible and thus less sensitive [4, 7]. The current experimental uncertainties are statistics limited, and we expect future data be able to distinguish between different models of polarized fragmentation and parton distribution functions.

The total systematic uncertainty in D_{LL} is varying from 0.02 to 0.04 with increasing p_T and is smaller than the statistical uncertainty, ranging from 0.06 to 0.14, for each of the data points. In estimating the size of the systematic uncertainties, we have combined contributions from the uncertainties in decay parameter α and in the measurements of the proton beam polarization and relative luminosity ratios, as well as uncertainty caused

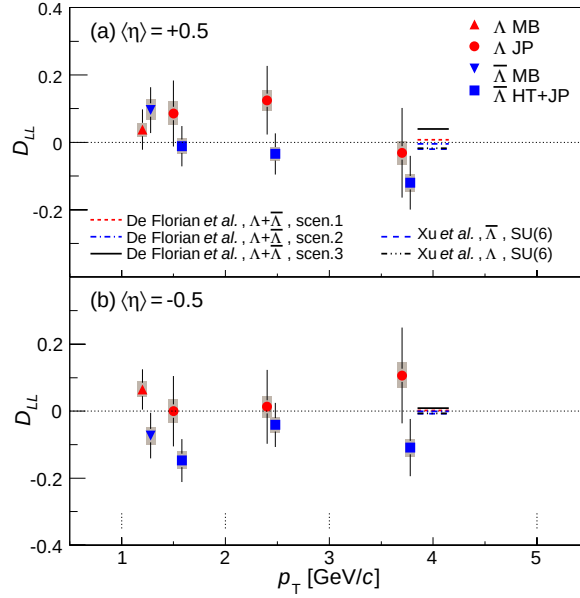


Figure 3: (Color online) Comparison of Λ and $\bar{\Lambda}$ spin transfer D_{LL} in polarized proton-proton collisions at $\sqrt{s} = 200$ GeV for (a) positive and (b) negative η versus p_T . The vertical bars and bands indicate the sizes of the statistical and systematic uncertainties, respectively. The $\bar{\Lambda}$ data points have been shifted slightly in p_T for clarity. The dotted vertical lines indicate the p_T intervals in the analysis of HT and JP data. The horizontal lines show model predictions.

by the aforementioned backgrounds, overlapping events (pile-up), and, in the case of the JP sample, trigger bias studied with Monte Carlo simulation [2]. The effect of Λ ($\bar{\Lambda}$) spin precession in the STAR magnetic field is negligible.

In addition to longitudinal spin transfer, the transverse spin transfer of hyperons from proton is also of particular interest in pp collisions, since it can provide access to the transverse spin content of nucleon, i.e., the transversity distribution, which is still poorly known in experiment. Unlike the polarization is along hyperon's momentum in the case of longitudinal polarization, the azimuthal direction in the transverse plane needs to be determined first to measure transverse hyperon polarization. E704 experiment measured the transverse spin transfer D_{NN} with respect to the production plane [13]. Another choice is to first determine the transverse polarization direction of the fragmenting parton in a hard scattering, which is different from the polarization direction before the scattering, but can be determined by a rotation around the normal of the scattering plane [14]. The fragmenting parton's axis can be obtained via jet reconstruction with charged particle and energy deposits in calorimeters. Then along this direction, the transverse hyperon polarization, and thus the transverse spin transfer D_{TT} can be measured. The transverse polarization of Λ ($\bar{\Lambda}$) is being investigated with hyperons reconstructed at mid-pseudorapidities with TPC at STAR. Sizable transverse spin transfer effect is expected to exist in the large $x_F (\equiv 2p_z/\sqrt{s})$ region. At STAR, a Forward Hadron Calorimeter (FHC) is being proposed to be installed behind the Forward Meson Spectrometer (FMS) in the near future, which may enable the reconstruction of Λ hyperons via the decay channel to $n\pi^0$ with π^0 detected by the FMS and n by the FHC. Simulation studies on Λ reconstruction in the forward region using the FMS and the FHC are underway.

In summary, we made measurements on the longitudinal spin transfer to Λ and $\bar{\Lambda}$ hyperons in $\sqrt{s} = 200$ GeV polarized proton-proton collisions for hyperon p_T up to 4 GeV/ c . The spin transfer is found to be $D_{LL} = -0.03 \pm 0.13(\text{stat}) \pm 0.04(\text{syst})$ for Λ and $D_{LL} = -0.12 \pm 0.08(\text{stat}) \pm 0.03(\text{syst})$ for $\bar{\Lambda}$ hyperons with $\langle \eta \rangle = 0.5$ and $\langle p_T \rangle = 3.7$ GeV/ c . The measurements of the transfer spin transfer may provide access to transversity distribution of the nucleon, and feasibility studies have started.

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LONGITUDINAL SPIN TRANSFER OF THE Λ AND $\bar{\Lambda}$ HYPERONS IN DIS AT COMPASS

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Abstract

The longitudinal spin transfer from muons to Λ and $\bar{\Lambda}$ hyperons has been studied in deep inelastic scattering off an unpolarized isoscalar target at the COMPASS experiment at CERN. The spin transfers to Λ and $\bar{\Lambda}$ produced in the current fragmentation region exhibit different behaviors as a function of x and x_F . The resulting average values are $D_{LL}^{\Lambda} = -0.01 \pm 0.05 \pm 0.02$ and $D_{LL}^{\bar{\Lambda}} = 0.25 \pm 0.06 \pm 0.05$. These results are discussed in the frame of recent model calculations.

Interest to study the polarization of Λ and $\bar{\Lambda}$ in DIS resides in the possibility of accessing the distributions of strange quarks and antiquarks in nucleon ($s(x)$ and $\bar{s}(x)$).

The present results are based on the analysis of data collected by COMPASS during the years 2003-2004. The Λ and $\bar{\Lambda}$ hyperons were produced by scattering 160 GeV polarized μ^+ off a polarized ${}^6\text{LiD}$ target. The COMPASS spectrometer is described in [1]. The average beam polarization was $P_b = -0.76 \pm 0.04$ in the 2003 run and $P_b = -0.80 \pm 0.04$ in the 2004 run. The cryogenic target consists of two 60 cm long cells filled with ${}^6\text{LiD}$, which is characterized by high polarizability. Both cells are oppositely polarized, the polarization vector is reversed every 8 hours.

DIS events are selected by cuts on the photon virtuality ($Q^2 > 1 \text{ (GeV/c)}^2$) and on the fractional energy of the virtual photon ($0.2 < y < 0.9$). The total data sample contains $31.2 \cdot 10^7$ DIS events. We also used a cut on x_F to select the region with an almost constant spectrometer acceptance: $0.05 < x_F < 0.5$. The total amount of hyperon events passing all selection cuts is around 70000 for Λ and 42000 for $\bar{\Lambda}$. It is the world largest $\Lambda(\bar{\Lambda})$ events sample in comparison with other DIS experiments [2–5]. Especially important is the $\bar{\Lambda}$ -hyperon sample which is 25 times larger than in HERMES [5].

For measurement of $\Lambda(\bar{\Lambda})$ polarization we selected the following reactions:

$$\begin{aligned} \mu^+ + d &\rightarrow \mu^+ + \Lambda(\bar{\Lambda}) + X \\ \Lambda &\rightarrow p + \pi^- \\ \bar{\Lambda} &\rightarrow \bar{p} + \pi^+ \end{aligned}$$

The longitudinal hyperon polarization P_L was extracted from analysis of the angular distribution of the hyperon decay products:

$$\frac{1}{N_{tot}} \cdot \frac{dN}{d \cos \theta} = \frac{1}{2} \cdot (1 + \alpha P_L \cos \theta),$$

here N_{tot} is the acceptance corrected amount of reconstructed $\Lambda(\bar{\Lambda})$, P_L is the projection of the polarization vector on the direction of the virtual photon momentum, θ is the

angle between the direction of the hyperon decay daughter baryon (proton for Λ and antiproton for $\bar{\Lambda}$) and the direction of the virtual photon, and $\alpha = +(-)0.642 \pm 0.013$ is the $\Lambda(\bar{\Lambda})$ decay parameter. In the present analysis we used both target cells despite of their polarization vector, which results in averaging of target cells polarization.

It's convenient to use the spin transfer coefficient $D_{LL}^{\Lambda(\bar{\Lambda})}$, which is defined as the part of beam polarization P_b that is transferred to the hyperon:

$$P_L = D_{LL}^{\Lambda(\bar{\Lambda})} P_b D(y),$$

where $D(y) = \frac{1-(1-y)^2}{1+(1-y)^2}$ is the virtual photon depolarization factor.

The dependence of D_{LL}^{Λ} and $D_{LL}^{\bar{\Lambda}}$ on Bjorken scaling variable x is shown in Fig. 1a. The spin transfer values for Λ and $\bar{\Lambda}$ are not equal to each other. D_{LL}^{Λ} is close to zero over the whole x range, whereas $D_{LL}^{\bar{\Lambda}}$ reaches 40-50%.

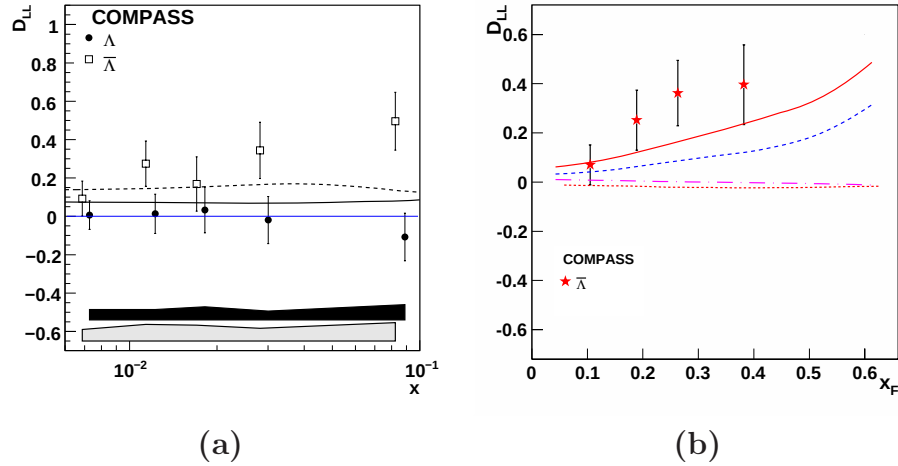


Figure 1: **(a)**: D_{LL} dependence for Λ and $\bar{\Lambda}$ on Bjorken x . Systematic errors are shown at the bottom, black color for Λ , light gray for $\bar{\Lambda}$. Solid line is theoretical calculation of [7] for Λ , dashed line – $\bar{\Lambda}$. The hyperon spin model was SU(6). **(b)**: D_{LL} dependence of $\bar{\Lambda}$ on x_F , calculated in [7] for parton distributions GRV98LO (dashed line), CTEQ5L (solid line) and CTEQ5L with turned off spin transfer from $\bar{s}$ -quark for SU(6) (dash-dotted line) and BJ [11] (dotted line) hyperon spin models.

The average kinematic values for Λ are: $\bar{x}_F = 0.22$, Bjorken $\bar{x} = 0.03$, fractional energy $\bar{z} = 0.27$, $\bar{y} = 0.46$, $\overline{Q^2} = 3.7$ (GeV/c)². In the selected x_F region $\bar{\Lambda}$ kinematic distributions are almost identical to Λ distributions. Spin transfer values averaged over all kinematic range are:

$$\begin{aligned} D_{LL}^{\Lambda} &= -0.012 \pm 0.047_{stat} \pm 0.024_{syst}, \\ D_{LL}^{\bar{\Lambda}} &= 0.249 \pm 0.056_{stat} \pm 0.049_{syst}. \end{aligned}$$

The resulting D_{LL} values are compared with the data from previous experiments in Fig. 2. The present results are in good agreement with others and for $\bar{\Lambda}$ they are the most precise D_{LL} measurements available today. The x_F dependence of D_{LL} for $\bar{\Lambda}$ is in good agreement with the NOMAD [3] result of $D_{LL} = 0.23 \pm 0.15 \pm 0.08$ at $\bar{x}_F = 0.18$.

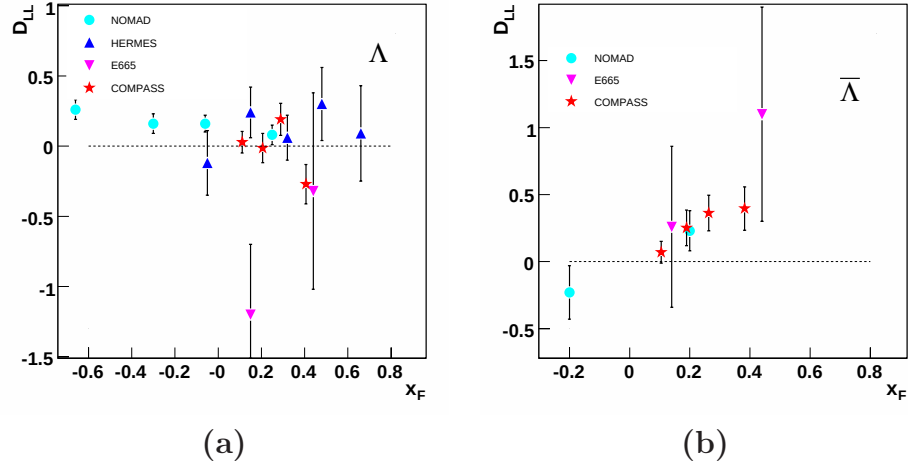


Figure 2: The D_{LL} dependence on x_F for the Λ (a) and $\bar{\Lambda}$ (b) in comparison with the world data [2–5].

It should be noted that the $\bar{\Lambda}$ spin transfer dependence on x_F was measured for the first time.

Different values of Λ and $\bar{\Lambda}$ spin transfer could be explained qualitatively by the following arguments. In the parton model the spin transfer from a polarized lepton scattered off an unpolarized target to the hyperon is given by the equation [6]:

$$D_{LL}^{\Lambda(\bar{\Lambda})}(x, z) = \frac{\sum_q e_q^2 q(x) \Delta D_q^{\Lambda(\bar{\Lambda})}(z)}{\sum_q e_q^2 q(x) D_q^{\Lambda(\bar{\Lambda})}(z)}, \quad (1)$$

where e_q is the quark charge, $q(x)$ is the parton distribution functions, $D_q^{\Lambda(\bar{\Lambda})}(z)$ and $\Delta D_q^{\Lambda(\bar{\Lambda})}(z)$ are the fragmentation functions of unpolarized and polarized quarks respectively. Since the spin of $\Lambda(\bar{\Lambda})$ is carried mainly by the strange quarks, then $\Delta D_q^{\Lambda(\bar{\Lambda})}(z) \sim 0$ for u , d , $\bar{u}$ and $\bar{d}$ quarks. Eq. 1 then results in:

$$D_{LL}^{\Lambda}(x, z) \approx \frac{1}{9} \frac{s(x) \Delta D_s^{\Lambda}(z)}{\sum_q e_q^2 q(x) D_q^{\Lambda}(z)}, \quad (2)$$

$$D_{LL}^{\bar{\Lambda}}(x, z) \approx \frac{1}{9} \frac{\bar{s}(x) \Delta D_s^{\bar{\Lambda}}(z)}{\sum_q e_q^2 q(x) D_q^{\bar{\Lambda}}(z)}. \quad (3)$$

At first sight the difference between (2) and (3) could arise only if $s(x) \neq \bar{s}(x)$. But actually, even if $s(x) = \bar{s}(x)$, equations (2) and (3) could have different denominators, which are accordingly proportional to the production cross-section of Λ and $\bar{\Lambda}$ -hyperons. The production cross-section for the Λ is almost two times bigger than for the $\bar{\Lambda}$ at COMPASS energy range. This leads to a decrease of the magnitude of the Λ -hyperon polarization in comparison with that of $\bar{\Lambda}$.

A comparison of the COMPASS data with theoretical predictions [7] is shown in Fig. 1b. The calculations confirm the high sensitivity of $D_{LL}^{\bar{\Lambda}}$ to the strange quark distributions in the nucleon. Independently of the specific hyperon model, the spin transfer to $\bar{\Lambda}$ vanishes if we turn off the contribution from the strange quarks (dotted line for BJ spin model [11] and dash-dotted line for SU(6) spin model).

The sensitivity of the data to the nucleon parton distribution functions is illustrated by comparison of the results obtained using the parametrizations CTEQ5L [8] (solid line) and GRV98LO [9] (dashed line). The GRV98 parton distributions have been chosen, because the strange sea there is purely perturbative, originating only from the production of $q\bar{q}$ -pairs by gluons. On the contrary CTEQ allows explicit non-perturbative strangeness in the nucleon. The amount of strangeness was extracted from dimuon experiments CCFR and NuTeV [10]. This leads to $\bar{s}(x)$ being nearly two times larger in CTEQ than in GRV98 for COMPASS energies in the $x = 0.001 - 0.01$ range. Therefore the calculation of D_{LL} for $\bar{\Lambda}$ using CTEQ gives twice bigger spin transfer and is closer to the experimental points in comparison with GRV98.

It should be noted that the polarization of Λ and $\bar{\Lambda}$ could be used as a tool to estimate $s(x)$ and $\bar{s}(x)$ independently of each other. More detailed description of the analysis could be found in [12].

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THE GPD PROGRAM AT COMPASS

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Abstract

The COMPASS proposed program of the Generalised Parton Distribution (GPD) studies is reviewed. Various observables for this program and the expected accuracies are discussed. The necessary developments of the experimental setup and the first results from the test run are also presented.

1 Introduction

Generalised Parton Distributions (GPDs) [1–3] provide a comprehensive description of the nucleon’s partonic structure and contain a wealth of new information. In particular, they embody both the nucleon electromagnetic form factors and the parton distribution functions, unpolarised as well as helicity-dependent. But more important, they allow a novel description of the nucleon as an extended object, referred sometimes as 3-dimensional ‘nucleon tomography’ [4]. GPDs also allow access to such a fundamental property of the nucleon as the orbital momentum of quarks [2]. For reviews of the GPDs see Refs [5–7].

The mapping of the nucleon GPDs requires extensive experimental studies of hard processes, Deeply Virtual Compton Scattering (DVCS) and Deeply Virtual Meson Production (DVMP), in a broad kinematic range. The high energies available at the CERN SPS and the availability of both muon beam polarisations make the fixed-target COMPASS set-up a unique place for such studies. In the future program [8] we propose to measure both DVCS and DVMP using an unpolarised proton target during a first period, in order to constrain mainly GPD H , and a transversely polarised ammonia target during another period in order to constrain the GPD E .

2 The proposed setup

The COMPASS apparatus is located at the high-energy (100-200 GeV) and highly-polarized $\mu^\pm$ beam line of the CERN SPS. At present it consists of a two stage spectrometer comprising various tracking detectors, electromagnetic and hadron calorimeters, and particle identification detectors grouped around 2 dipole magnets SM1 and SM2 in conjunction with a longitudinally or transversely polarized target. By installing a recoil proton detector around the target to ensure exclusivity of the DVCS and DVMP events,

COMPASS could be converted into a facility measuring exclusive reactions within a kinematic domain from $x \sim 0.01$ to ~ 0.1 , which cannot be explored at any other existing or planned facility in the near future. Thus COMPASS could explore the uncharted x domain between the HERA collider experiments and the fixed-target experiments as HERMES and the planned 12 GeV extension of the JLAB accelerator. For values of x below 10^{-1} , the outgoing photon (or meson) is emitted at an angle below 10° which corresponds for the photon to the acceptance of the two existing COMPASS electromagnetic calorimeters ECAL1 and ECAL2 and which for charged particles is within the acceptance of the tracking devices and the RICH detector. To access higher x values a large angular acceptance calorimeter ECAL0 is needed, which is presently under a study. Schematic layout of COMPASS is shown in Fig. 1, with only new or upgraded detectors and the spectrometer magnets indicated.

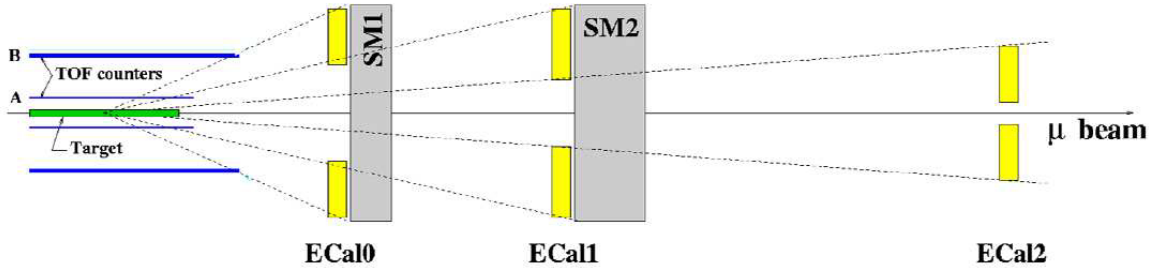


Figure 1: Schematic layout of the proposed setup. Only new and upgraded detectors are shown.

The data will be collected with polarized μ^+ and μ^- beams. Assuming 140 days¹ of data taking and a muon flux of $4.6 \cdot 10^8 \mu$ per SPS spill, reasonable statistics for the DVCS process can be accumulated for Q^2 values up to 8 GeV^2 . It is worth noting that an increase of the number of muons per spill by a factor 4 would result in an increase in the range in Q^2 up to about 12 GeV^2 .

3 Planned measurements

The complete GPD program at COMPASS will comprise the measurements of the DVCS cross section with polarized positive and negative muon beams and at the same time the measurements of a large set of mesons (ρ , ω , ϕ , π , η , ...).

3.1 Deeply Virtual Compton Scattering

DVCS is considered to be the theoretically cleanest of the experimentally accessible processes because effects of next-to-leading order and higher twist contributions are under theoretical control [9]. The competing Bethe-Heitler (BH) process, which is elastic lepton-nucleon scattering with a hard photon emitted by either the incoming or outgoing lepton, has a final state identical to that of DVCS so that both processes interfere at the level of amplitudes.

¹The quoted number corresponds to the assumption that both μ^+ and μ^- beams have the same intensities. In practice, because the intensity of μ^- beam is few times smaller, about 280 days will be required to reach the same precision.

COMPASS offers the advantage to provide various kinematic domains where either BH or DVCS dominates. The collection of an almost pure BH events at small x allows one to get an excellent reference yield and to control accurately the global efficiency of the apparatus. In contrast, the collection of an almost pure DVCS sample at larger x will allow the measurement of the x dependence of the t -slope of the cross section, which is related to the tomographic partonic image of the nucleon. In the intermediate domain, the DVCS contribution will be boosted by the BH process through the interference term. The dependence on ϕ , the azimuthal angle between lepton scattering plane and photon production plane, is a characteristic feature of the cross section [9].

COMPASS is presently the only facility to provide polarized leptons with either charge: polarized μ^+ and μ^- beams. Note that with muon beams one naturally reverses both charge and helicity at once. Practically μ^+ are selected with a polarisation of -0.8 and μ^- with a polarization of +0.8. The difference and sum of cross sections for μ^+ and μ^- combined with the analysis of ϕ dependence allow to isolate the real and imaginary parts of the leading twist-2 DVCS amplitude, and of higher twist contributions.

In the following sections we show projections for DVCS measurements with an unpolarised proton target (3.1.1 and 3.1.2) and with a transversely polarised ammonia target (3.1.3). For each target the integrated muon flux was taken the same as described in Sect. 2 and the value of the global efficiency was assumed to be equal to 0.1.

3.1.1 x -dependence of the t -slope of DVCS

The t -slope parameter $B(x)$ of the DVCS cross section $\frac{d\sigma}{dt}(x) \propto \exp(-B(x)|t|)$ can be obtained from the beam charge and spin sum of the cross sections after integration over ϕ and BH subtraction. The expected statistical accuracy of the measurements of $B(x)$ at COMPASS is shown in Fig. 2. The systematic errors are mainly due to uncertainties involved in the subtraction of the BH contribution. At $x > 0.02$ they are small compared to the statistical errors. For the simulations the simple ansatz $B(x) = B_0 + 2\alpha' \log(\frac{x_0}{x})$ was used. As neither B_0 nor α' are known in the COMPASS kinematics, for the simulations shown in Fig. 2 we chose the values $B_0 = 5.83 \text{ GeV}^2$, $\alpha' = 0.125$ and $x_0 = 0.0012$. The precise value of the t -slope parameter $B(x)$ in the COMPASS x -range will yield new and significant information in the context of ‘nucleon tomography’ as it is expected in Ref. [10].

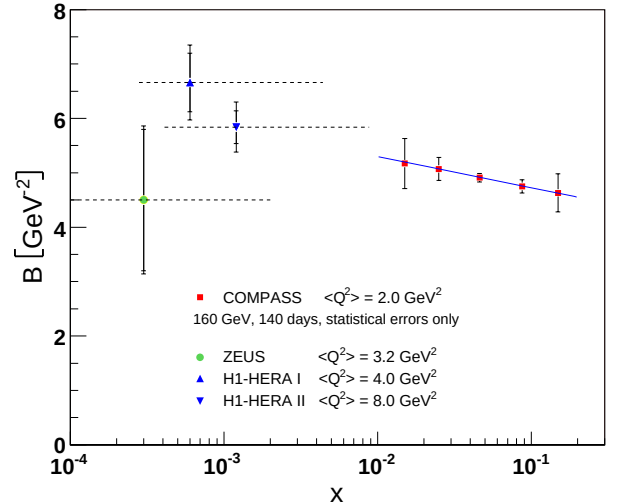


Figure 2: The x dependence of the fitted t -slope parameter B of the DVCS cross section, expressed as $d\sigma/dt \propto e^{-B|t|}$. COMPASS projections are calculated for $1 < Q^2 < 8 \text{ GeV}^2$ and are compared to some HERA results for which the mean value $\langle Q^2 \rangle$ is in this range.

3.1.2 Beam charge and spin asymmetry

Fig. 3 shows the projected statistical accuracy for the beam charge and spin asymmetry $A_{CS,U}$ measured as a function of ϕ in a selected (x, Q^2) bin. The asymmetry is defined as

$$A_{CS,U} = \frac{d\sigma^{\leftarrow} - d\sigma^{\rightarrow}}{d\sigma^{\leftarrow} + d\sigma^{\rightarrow}}, \quad (1)$$

with arrows indicating the orientations of the longitudinal polarisation of the beams. This asymmetry is sensitive to the real part of the DVCS amplitude which is a convolution of GPDs with the hard scattering kernel over the whole range of longitudinal momenta of exchanged quarks. Therefore measurements of this asymmetry provide strong constraints on the models of GPD. Two of the curves shown in the figure are calculated using the 'VGG' GPD model [11]. As this model is meant to be applied mostly in the valence region, typically the value $\alpha' = 0.8$ is used in the 'reggeized' parameterization of the correlated x, t dependence of GPDs. For comparison also the model result for the 'factorized' x, t dependence is shown, which corresponds to $\alpha' \approx 0.1$ in the reggeized ansatz. A recent theoretical development [12] exploiting dispersion relations for Compton form factors was successfully applied to describe DVCS observables at very small values of x typical for the HERA and extended to include DVCS data from HERMES and JLAB. The prediction for COMPASS from this analysis is shown as an additional curve.

As the overall expected data set from the GPD program for COMPASS will allow 9 bins in x vs. Q^2 , each of them expected to contain statistics sufficient for stable fits of the ϕ dependence, a determination of the 2-dimensional x, Q^2 (or x, t) dependence will be possible for the various Fourier expansion coefficients c_n and s_n [9], thereby yielding information on the nucleon structure in terms of GPDs over a range in x . These data are expected to be very useful for future developments of reliable GPD models able to simultaneously describe the *full* x -range.

3.1.3 Predictions for the transverse target spin asymmetry

Transverse target spin asymmetries for exclusive photon production are important observables for studies of the GPD E , and for the determination of the role of the orbital momentum of quarks in the spin budget of the nucleon. The sensitivity of these asymmetries to the total angular momentum of u quarks, J_u , was estimated for the transversely polarised protons in a model dependent way in Ref. [13].

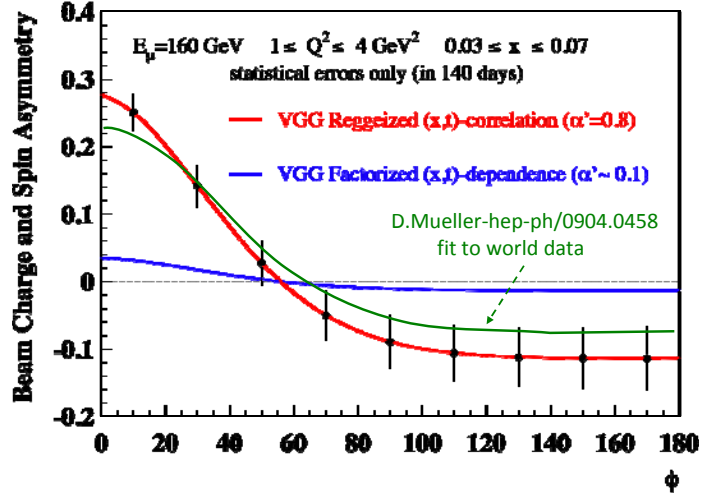


Figure 3: Projections for the Beam Charge and Spin Asymmetry measured at COMPASS in one year for $0.03 \leq x \leq 0.07$ and $1 \leq Q^2 \leq 4 \text{ GeV}^2$. The red and blue curves correspond to different variants of the VGG model [11] while the green curve shows predictions based on the first fit to the world data [12].

The transverse target pin asymmetries for the proton will be measured with the transversely polarised ammonia target, similar to the one used at present by COMPASS. Two options are considered for the configuration of the target magnet and the RPD, each with a different impact on the range of measurable energy of the recoil proton.

The transverse spin dependent part of the cross sections will be obtained by subtracting the data with opposite values of the azimuthal angle ϕ_s , which is the angle between the lepton scattering plane and the target spin vector. In order to disentangle the $|DVCS|^2$ and the interference terms with the same azimuthal dependence, it is necessary to take data with both μ^+ and μ^- beams, because only in the difference and the sum of μ^+ and μ^- cross sections these terms become separated. Both asymmetries for the difference and the sum of μ^+ and μ^- transverse spin dependent cross sections will be analysed. The difference (sum) asymmetry $A_{CS,T}^D$, ($A_{CS,T}^S$) is defined as the ratio of the μ^+ and μ^- cross section difference (sum) divided by the lepton-charge-averaged, unpolarised cross section. Here CS indicates that both lepton charge and lepton spin are reversed between μ^+ and μ^- , and T is for the transverse target polarisation.

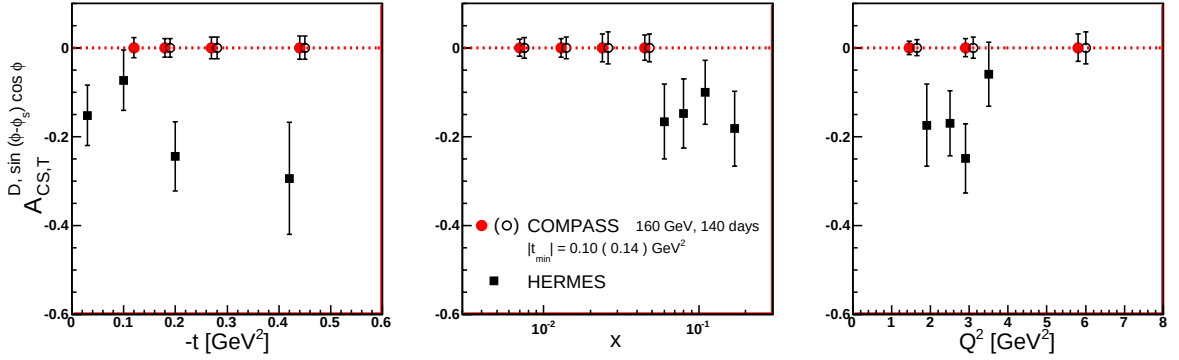


Figure 4: The expected statistical accuracy of $A_{CS,T}^{D, \sin(\phi-\phi_s)\cos\phi}$ as a function of $-t$, x and Q^2 . Solid and open circles correspond to the simulations for the two considered configurations of the target region. Also shown is the asymmetry $A_{UT}^{\sin(\phi-\phi_s)\cos\phi}$ measured by HERMES [13] with its statistical errors.

As an example, the results from the simulations of the expected statistical accuracy of the asymmetry $A_{CS,T}^{D, \sin(\phi-\phi_s)\cos\phi}$ are shown in Fig. 4 as a function of $-t$, x and Q^2 for the two considered configurations of the target region. Here $\sin(\phi-\phi_s)\cos\phi$ indicates the type of azimuthal modulations. This asymmetry is an analogue of the asymmetry $A_{UT}^{\sin(\phi-\phi_s)\cos\phi}$ measured by HERMES with unpolarised electrons, also shown in the figure.

Typical values of the statistical errors of $A_{CS,T}^{D, \sin(\phi-\phi_s)\cos\phi}$, as well as of the seven remaining asymmetries related to the twist-2 terms in the cross section, are expected to be ≈ 0.03 .

3.2 Deeply Virtual Meson Production

Hard exclusive vector meson production is complementary to DVCS as it provides access to various other combinations of GPDs. For vector meson production only GPDs H and E contribute, while for pseudoscalar mesons GPDs $\tilde{H}$ and $\tilde{E}$ play a role. We recall that DVCS depends on the four GPDs. Also in contrast to DVCS, where gluon

contributions enter only beyond leading order in α_s , in DVMP both quark and gluon GPDs contribute at the same order. For example

$$H_{\rho^0} = \frac{1}{\sqrt{2}}\left(\frac{2}{3}H^u + \frac{1}{3}H^d + \frac{3}{8}H^g\right); \quad H_\omega = \frac{1}{\sqrt{2}}\left(\frac{2}{3}H^u - \frac{1}{3}H^d + \frac{1}{8}H^g\right); \quad H_\phi = -\frac{1}{3}H^s - \frac{1}{8}H^g.$$

Therefore by combining the results for various mesons the GPDs for various quark flavours and for gluons could be disentangled.

It was pointed out that vector meson production on a transversely polarised target is sensitive to the nucleon helicity-flip GPD E [5, 14]. This GPD offers unique views on the orbital angular momentum carried by partons in the proton [2] and on the correlation between polarisation and spatial distribution of partons [4]. The azimuthal asymmetry $A_{UT}^{\sin(\phi-\phi_s)}$ for exclusive production of a vector meson off the transversely polarised nucleon depends linearly on the GPD E and at COMPASS kinematic domain it can be expressed as

$$A_{UT}^{\sin(\phi-\phi_s)} \sim \sqrt{t_0 - t} \frac{\text{Im}(\mathcal{E}^M \mathcal{H}_M^*)}{|\mathcal{H}_M|^2}, \quad (2)$$

where t_0 is the minimal momentum transfer. The quantities $\mathcal{H}_M$ and $\mathcal{E}_M$ are weighted sums of integrals over the GPD $H^{q,g}$ and $E^{q,g}$ respectively. The weights depend on the contributions of quarks of various flavours and of gluons to the production of meson M .

We note that the production of ρ , ω , ϕ vector mesons is already being investigated at COMPASS [15]. The recent results from COMPASS on the transverse spin asymmetries for ρ^0 production off transversely polarised protons and deuterons were presented in Ref. [16]. The spin asymmetry $A_{UT}^{\sin(\phi-\phi_s)}$ for the proton is shown in Fig. 5 as a function of Q^2 , x and p_t^2 , where p_t is the transverse momentum of ρ^0 with respect to the virtual photon. The errors shown are statistical ones. Preliminary estimates of systematic errors indicate that they are smaller than the statistical ones. The asymmetry is consistent with zero within the statistical errors.

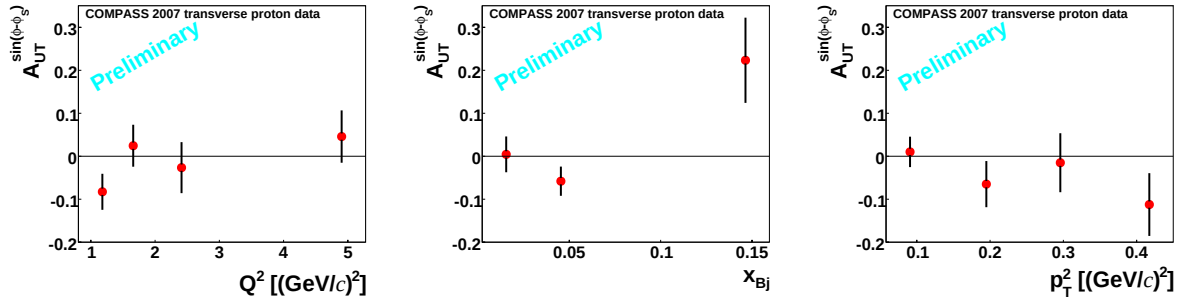


Figure 5: $A_{UT}^{\sin(\phi-\phi_s)}$ for exclusive ρ^0 production off the transversely polarised protons as a function of Q^2 , x and p_t^2 .

A similar analysis of $A_{UT}^{\sin(\phi-\phi_s)}$ for the deuteron was done using the data taken in 2002-2004 with transversely polarised ^6LiD target. Coherent and incoherent contributions to the exclusive production on the deuteron have to be disentangled. The asymmetry for the deuteron is consistent with zero as well. The separation of contributions from longitudinal and transverse virtual photons for both data sets is in progress.

Theoretical calculations of $A_{UT}^{\sin(\phi-\phi_s)}$ for ρ^0 have been performed by Goloskokov and Kroll [17]. At COMPASS kinematics the predicted values for ρ^0 production on the proton

are about -0.02, in agreement with the data. The small values of asymmetry are to a large extent due to a cancellation of E^u and E^d for ρ^0 production. Significantly higher asymmetry, about -0.10, is expected for ω production.

4 Validation tests and outlook

The setup used in 2008 and 2009 for the meson spectroscopy measurements with hadron beams happens to be an excellent *prototype* to perform validation measurements for DVCS. A first measurements of exclusive γ production on a 40 cm long LH target, with detection of the slow recoiling proton in the RPD has been performed during a short (< 2 days) test run in 2008 using 160 GeV μ^+ and μ^- beams. The measurements were performed with the present hadron setup, all the standard COMPASS tracking detectors, the ECAL1 and ECAL2 electromagnetic calorimeters for photon detection and appropriate triggers. An efficient selection of single photon events, and suppression of the background is possible by using the combined information from the forward COMPASS detectors and the RPD.

A way to tag the observed process, $\mu + p \rightarrow \mu' + \gamma + p'$, to which both the DVCS and Bethe-Heitler process contribute, is to look at the angle ϕ between the leptonic and hadronic planes. The Bethe-Heitler contribution, which dominates the sample, show a peak at $\phi \simeq 0$. The observed distribution, after applying all cuts and selections and for $Q^2 > 1$ (GeV/c) ^{2} , is displayed in Fig. 6 with the prediction from the Monte Carlo simulation. The shape of the observed distribution is compatible with the dominant Bethe-Heitler process. The overall detection efficiency can be deduced from the relative normalization of the two distributions. The global efficiency is equal to 0.13 ± 0.05 .

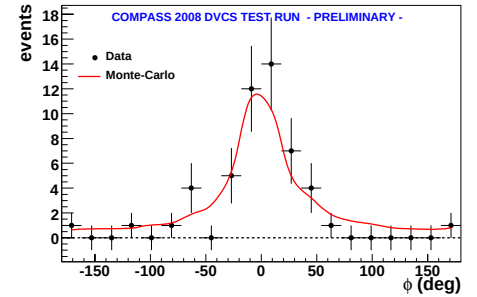


Figure 6: The distribution of the azimuthal angle ϕ for observed exclusive single photon production (points). The line is the Monte Carlo prediction normalized to the data.

In 2009 a two-week DVCS test run was performed with 160 GeV μ^+ and μ^- beams, and with a similar setup as in 2008. In addition the beam momentum station was operational during the test, which allowed momentum measurements of individual beam particles. The main goal for the DVCS test in 2009 is a better understanding of all backgrounds and a first evaluation of the relevant contributions of DVCS and DVCS-BH interference in a kinematic domain at larger x where BH is not dominant.

A proposal to extend the physics program of COMPASS, including the GPD studies, will be submitted in 2010. A possible start of the GPD program, first with an unpolarised proton target, is planned for 2012.

Acknowledgments

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AZIMUTHAL ASYMMETRIES IN PRODUCTION OF CHARGED HADRONS BY HIGH ENERGY MUONS ON POLARIZED DEUTERIUM TARGETS¹

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Abstract

Search for azimuthal asymmetries in semi-inclusive production of charged hadrons by 160 GeV muons on the longitudinally polarized deuterium target, has been performed using the 2002- 2004 COMPASS data. The observed asymmetries integrated over the kinematical variables do not depend on the azimuthal angle of produced hadrons and are consistent with the ratio $g_1^d(x)/f_1^d(x)$. The asymmetries are parameterized taking into account possible contributions from different parton distribution functions and parton fragmentation functions depending on the transverse spin of quarks. They can be modulated (either/or/and) with $\sin(\phi)$, $\sin(2\phi)$, $\sin(3\phi)$ and $\cos(\phi)$. The x -, z - and p_h^T -dependencies of these amplitudes are studied.

1. Introduction. Although the longitudinal spin structure of nucleons has been investigated for more than 20 years and results are very well known, the studies of the transverse spin structure of nucleons have been started recently. Since the pioneering HERMES [1] and CLAS [2] experiments it is known that the signature of the transverse spin effects is an appearance of azimuthal asymmetries (AA) of the hadrons produced in Semi-Inclusive Deep Inelastic Scattering (SIDIS) of leptons on polarized targets.

These asymmetries are related with new Parton Distribution Functions (PDF) and new polarized Parton Fragmentation Functions (PFF), depending on the transverse spin of quarks [3]. AAs on the transversally polarized targets have been already reported by HERMES [4] and COMPASS [5, 6], and on the longitudinally polarized targets - by HERMES [7, 8]. The search for the AA using the COMPASS spectrometer [9] with the longitudinally polarized deuterium target is described below.

In the framework of the parton model of nucleons, the squared modulus of the matrix element of the SIDIS is represented by the type of the diagram in Fig.1a, where an example of one of the new PDF, transversity, $h_1(x)$ and new Collins PFF, $H_1^\perp(z)$, is shown. New PDFs and PFFs, due to their chiral odd structure, always appear in pairs.

The kinematics of the SIDIS is shown in Fig.1b, where ℓ (ℓ') is the 4-momentum of incident (scattered) lepton, $q = \ell - \ell'$, $Q^2 = -q^2$, θ_γ is the angle of the virtual photon momentum $\vec{q}$ with respect to the beam, $P_L(P_T)$ is a longitudinal (transversal) component of the target polarization, P_{II} , with respect to the virtual photon momentum in the laboratory frame, p_h is the hadron momentum with the transverse component p_h^T , ϕ is the azimuthal angle between the scattering plane and hadron production plane, ϕ_S is the angle of the target polarization vector with respect to the lepton scattering

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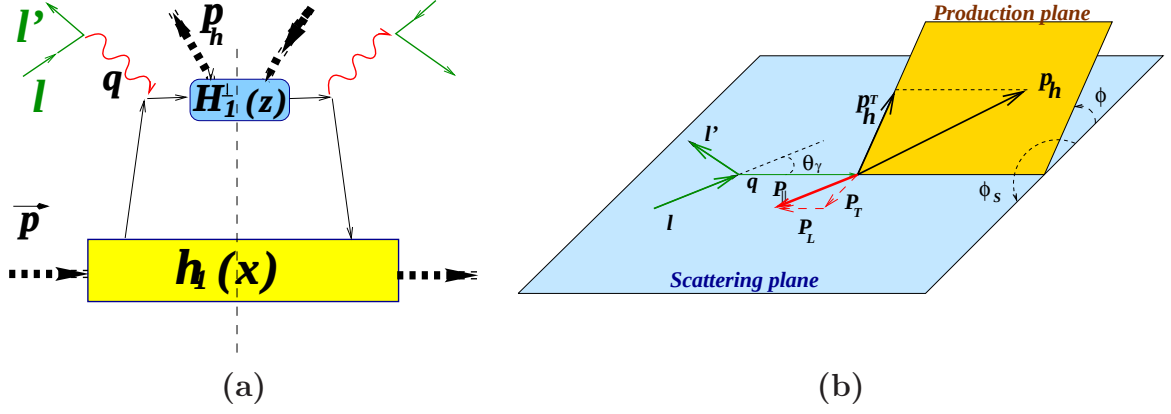


Figure 1: The squared modulus of the matrix element of the SIDIS reaction $\ell + \vec{N} \rightarrow \ell' + h + X$ summed over X states (a) and kinematics of the process (b).

plane (for the longitudinal target polarization $\phi_S = 0$ or π . For the target polarization P_{II} , which is longitudinal with respect to the lepton beam, the transverse component is equal to $|P_T| = P_{II} \sin(\theta_\gamma)$, where $\sin(\theta_\gamma) \approx 2 \frac{M}{Q} x \sqrt{1-y}$, $y = \frac{q \cdot p}{p \cdot \ell}$ and M is the nucleon mass. The Bjorken variable x and the hadron fractional momentum z are defined as $x = Q^2/2p \cdot q$, $z = p \cdot p_h/p \cdot q$, where p is the 4-momentum of the incident nucleon.

In general, the total cross section of the SIDIS reaction is a linear function of the lepton beam polarization, P_μ , and of the target polarization P_{II} or its components:

$$d\sigma = d\sigma_{00} + P_\mu d\sigma_{L0} + P_L (d\sigma_{0L} + P_\mu d\sigma_{LL}) + |P_T| (d\sigma_{0T} + P_\mu d\sigma_{LT}) , \quad (1)$$

where the first (second) subscript of the partial cross sections means the beam (target) polarization.

The asymmetry, $a(\phi)$, in the hadron production from the longitudinally polarized target (LPT), is defined by the expression:

$$a(\phi) = \frac{d\sigma^{\leftarrow\rightarrow} - d\sigma^{\leftarrow\leftarrow}}{d\sigma^{\leftarrow\rightarrow} + d\sigma^{\leftarrow\leftarrow}} \propto P_L (d\sigma_{0L} + P_\mu d\sigma_{LL}) + |P_L| \sin(\theta_\gamma) (d\sigma_{0T} + P_\mu d\sigma_{LT}) . \quad (2)$$

Each of the partial cross sections is characterized by the specific dependence of the definite convolution of PDF and PFF times a function the azimuthal angle of the outgoing hadron. Namely, contributions to Eq. (2) from each quark and antiquark flavor, up to the order (M/Q) , have the forms:

$$\begin{aligned} d\sigma_{0L} &\propto \epsilon x h_{1L}^\perp(x) \otimes H_1^\perp(z) \sin(2\phi) + \sqrt{2\epsilon(1-\epsilon)} \frac{M}{Q} x^2 \left(h_L(x) \otimes H_1^\perp(z) + f_L^\perp(x) \otimes D_1(z) \right) \sin(\phi), \\ d\sigma_{LL} &\propto \sqrt{1-\epsilon^2} x g_{1L}(x) \otimes D_1(z) + \sqrt{2\epsilon(1-\epsilon)} \frac{M}{Q} x^2 \left(g_L^\perp(x) \otimes D_1(z) + e_L(x) \otimes H_1^\perp(z) \right) \cos(\phi), \\ d\sigma_{0T} &\propto \epsilon \{ x h_1(x) \otimes H_1^\perp(z) \sin(\phi + \phi_S) + x h_{1T}^\perp(x) \otimes H_1^\perp(z) \sin(3\phi - \phi_S) - x f_{1T}^\perp(x) \otimes D_1(z) \sin(\phi - \phi_S) \}, \\ d\sigma_{LT} &\propto \sqrt{1-\epsilon^2} x g_{1T}(x) \otimes D_1(z) \cos(\phi - \phi_S), \end{aligned} \quad (3)$$

where $\otimes$ is a convolution in parton's internal transversal momentum, k_T , on which PDF and PFF depend, $\phi_s=0$ for the LPT and $\epsilon \approx \frac{2(1-y)}{2-2y+y^2}$. The structure of the partial cross sections and physics interpretations of the new PDFs and PFFs, entering in $a(\phi)$, are given in [10–12].

So, the aim of this study is to see the AA in the hadron production from LTP, as a manifestation of new PDFs and PFFs and the x , z and p_h^T - dependence of the corresponding amplitudes.

2. Method of analysis. The COMPASS polarized target [9] in 2002-2004 years had two cells, Up- and Down-stream of the beam, placed in the 2.5 T solenoid magnetic field. The target material of the cells (^6LiD or NH_3) can be polarized in opposite directions with respect to the beam, for example in the U-cell along to the beam (positive polarization) and in D-cell – opposite to the beam (negative polarization) and vice versa. Such a configuration can be achieved by means of the microwave field at low temperatures at any direction of the solenoid magnetic field holding the polarization. Suppose that the above configuration of the cell polarizations is realized with the positive (along to the beam) solenoid field, then, to avoid possible systematic effects in acceptance connected with this field, after some time the same configuration of polarizations is realized by means of the microwave field with the negative (opposite to the beam) solenoid field. Microwave polarization reversals are repeated several times while data taking. In order to minimize systematics caused by the time dependent variation of the acceptance between the microwave reversals, the polarizations are frequently reversed by inverting of the solenoid field.

For the AA studies the double ratios of event numbers, R_f , is used in the following form:

$$R_f(\phi) = [N_{+,f}^U(\phi)/N_{-,f}^D(\phi)] \cdot [N_{+,f}^D(\phi)/N_{-,f}^U(\phi)], \quad (4)$$

where $N_{p,f}^t(\phi)$ is a number of events in each ϕ -bin from the target cell t , $t = U, D$, $p = +$ or $-$ is the sign of the target polarization, $f = +$ or $-$ is the direction of the target solenoid field. Using Eqs. (1,2,3) with $P_{\pm}$ as an absolute value of averaged products of the positive or negative target polarization and dilution factor, the number of events can be expressed as

$$N_{p,f}^t = C_f^t(\phi) L_{p,f}^t [(B_0 + B_1 \cos(\phi) + B_2 \sin(\phi) + \dots) \pm P_p (A_0 + A_1 \sin(\phi) + A_2 \sin(2\phi) + \dots)], \quad (5)$$

where $C_f^t(\phi)$ is the acceptance factor (source of false asymmetries), $L_{p,f}^t$ is a luminosity depending on the beam flux and target densities. The coefficients $B_0, B_1, \dots$ and $A_0, A_1, \dots$ characterize contributions of partial cross sections. Substituting Eq.(5) in Eq. (4) one can see that the acceptance factors are canceled, as well as the luminosity factors if the beam muons cross the both cells. So, the ratio $R_f(\phi)$ depends only on physics characteristics of the SIDIS process and it is expressed via asymmetry $a(\phi)$, Eq. (2), in the quadratic equation, approximate solution of which is:

$$a_f = [R_f(\phi) - 1] / (P_{+,f}^U + P_{+,f}^D + P_{-,f}^U + P_{-,f}^D). \quad (6)$$

Since asymmetry should not depend on the direction of the solenoid field, one can expect to have $a_+ = a_-$. Small difference between a_+ and a_- could appear due to the solenoid field dependent contributions non-factorizable in Eq. (5). But these contributions have different signs and canceled in the sum $a(\phi) = a_+(\phi) + a_-(\phi)$. So, the weighted sum $a(\phi) = a_+(\phi) + a_-(\phi)$, calculated separately for each year of data taking and averaged at the end, is obtained for the final results.

3. Data selection. The data selection, aimed at having a clean sample of hadrons, has been performed in three steps, using a preselected sample. This sample contained about 167.5M of SIDIS events with $Q^2 > 1 \text{ GeV}^2$ and $y > 0.1$ in a form of reconstructed vertices with incoming and outgoing muons and one or more additional outgoing tracks.

1. "GOOD SIDIS EVENTS" have been selected out of the preselected ones applying more stringent cuts on the quality of reconstructed tracks and vertices, vertex positions inside the target cells, momentum of the incoming muon (140-180 GeV/c), energy transfer ($y < 0.9$) and invariant mass of the final states ($W > 5 \text{ GeV}$). About 58% of events of the initial sample have survived after these cuts.

2. "GOOD TRACKS" (about 157 M) have been selected out of the total tracks (about 290 M) from GOOD SIDIS EVENTS excluding the tracks identified as muons and tracks with $z > 1$ and $p_h^T < 0.1 \text{ GeV}/c$.

3. "GOOD HADRONS" from GOOD TRACKS have been identified using the information from the hadron calorimeters HCAL1 and HCAL2. Each of the GOOD TRACKS is considered as the GOOD HADRON if: this track hits one of the calorimeter, the calorimeter has the energy cluster associated with this hit with $E_{HCAL1} > 5 \text{ GeV}$, or $E_{HCAL2} > 7 \text{ GeV}$, coordinates of the cluster are compatible with coordinates of the track and the energy of the cluster is compatible with the momentum of the track.

The total number of GOOD HADRONS was about 53 M. Each of the GOOD HADRONS is included to the asymmetry evaluations.

4. Results. The weighted sum of azimuthal asymmetries $a(\phi) = a_+(\phi) + a_-(\phi)$, averaged over all kinematical variables, are shown in Fig. 2 for negative and positive hadrons. They have been fitted by functions

$$a(\phi) = a^{\text{const}} + a^{\sin \phi} \sin(\phi) + a^{\sin 2\phi} \sin(2\phi) + a^{\sin 3\phi} \sin(3\phi) + a^{\cos \phi} \cos(\phi). \quad (7)$$

The fit parameters, characterizing ϕ -modulation amplitudes, are compatible with zero. The ϕ -independent parts of $a(\phi)$ differ from zero and are almost equal for h^- and h^+ . The fits of $a(\phi)$ by constants are also shown in Fig. 2.

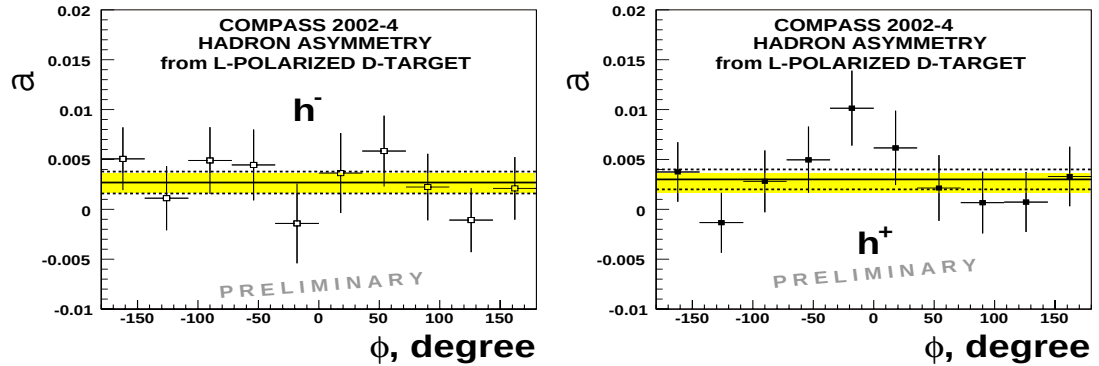


Figure 2: Azimuthal asymmetries $a(\phi)$ for negative (left) and positive (right) hadrons and results of fits by the constants with $\chi^2/d.f.$ equal to 3.4/5 (5.2/5), respectively.

As already specified, the ϕ -independent parts of asymmetries come from the $d\sigma_{LL}$ contributions to the cross sections, which are proportional to helicity PDF times PFF (see Eq. (3)) of non-polarized quarks in the non-polarized hadron. For the deuteron target this contribution is expected to be charge independent.

Dependence of the AA fit parameters on the kinematical variables are shown in Figs. 3-7.

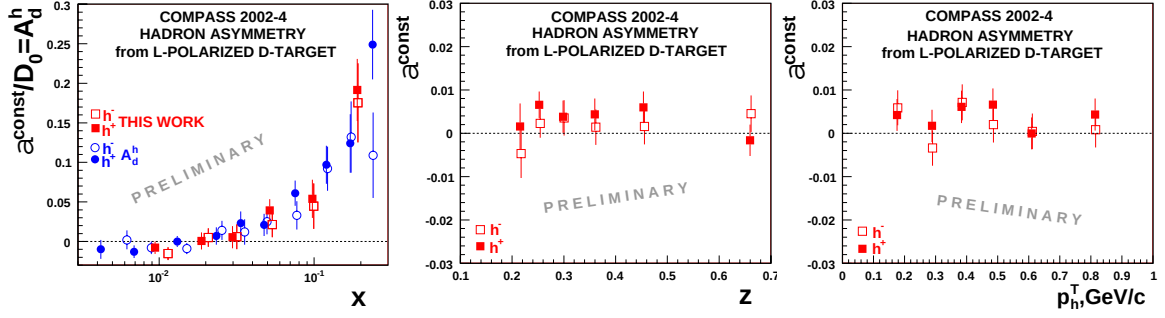


Figure 3: Dependence of the AA fit parameters a^{const} on kinematical variables.

The parameters $a^{\text{const}}(x)$, being divided by the virtual photon depolarization factor D_0 , are equal (by definition) to the asymmetry $A_d^h(x)$, already published by COMPASS [13]. Agreement of these data and data of the present analysis has demonstrated internal consistency of the results.

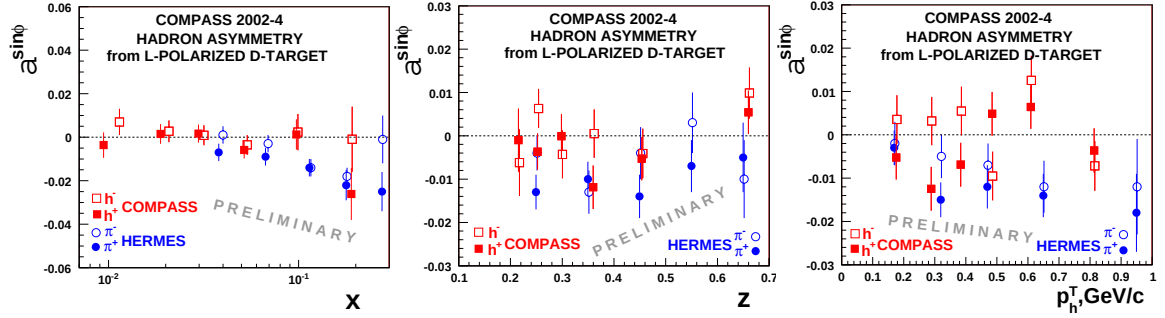


Figure 4: Dependence of the AA fit parameters $a^{\sin \phi}$ on kinematical variables and similar data of HERMES [8] for identified leading pions.

The x -dependence of the $\sin(\phi)$ modulations of the AA, observed by HERMES, is less pronounced at COMPASS. This modulation is due to pure twist-3 PDF's entering from the $d\sigma_{0L}$ contribution to the AA with a factor Mx/Q .

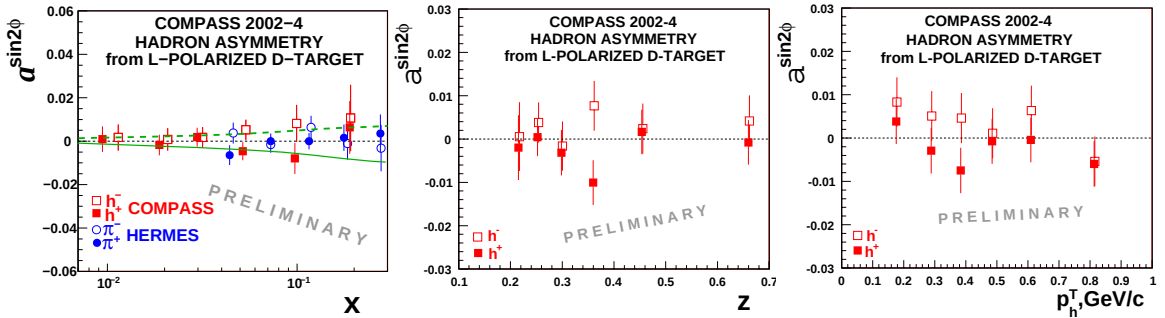


Figure 5: Dependence of the AA fit parameters $a^{\sin 2 \phi}$ on kinematical variables compared to the data of HERMES and calculations by H.Avakian et al. [14]: dashed line – h^- , solid line – h^+ .

The amplitudes of the $\sin(2\phi)$ modulations are small, consistent with zero within the errors. They could be caused by PDF $h_{1L}^\perp$ in $d\sigma_{0L}$.

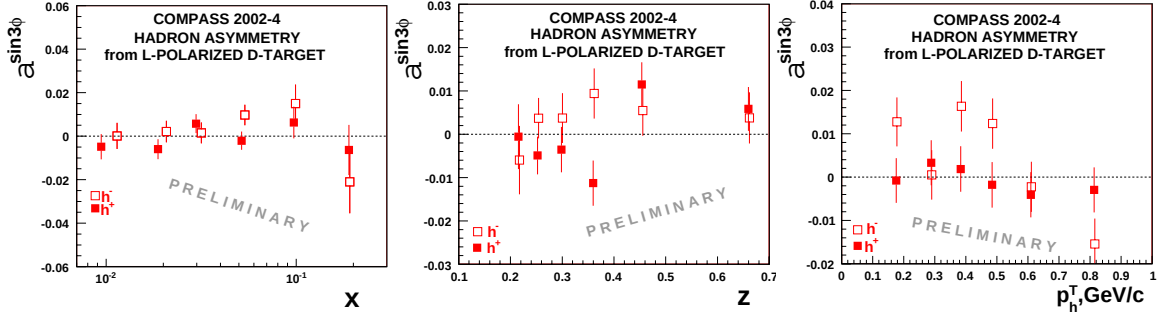


Figure 6: Dependence of the AA fit parameters $a^{\sin 3\phi}$ on kinematical variables.

Some peculiarities of the data on the $a^{\sin 3\phi}$ are seen from Fig. 4, for instance, the points for h^- are mostly positive while for h^+ they are mostly negative like for the COMPASS results from the transversally polarized target [15]. Remind that this modulation could come from the pretzelosity PDF $h_{1T}^\perp$ in $d\sigma_{0T}$, additionally suppressed by $\sin(\theta_\gamma) \sim xM/Q$.

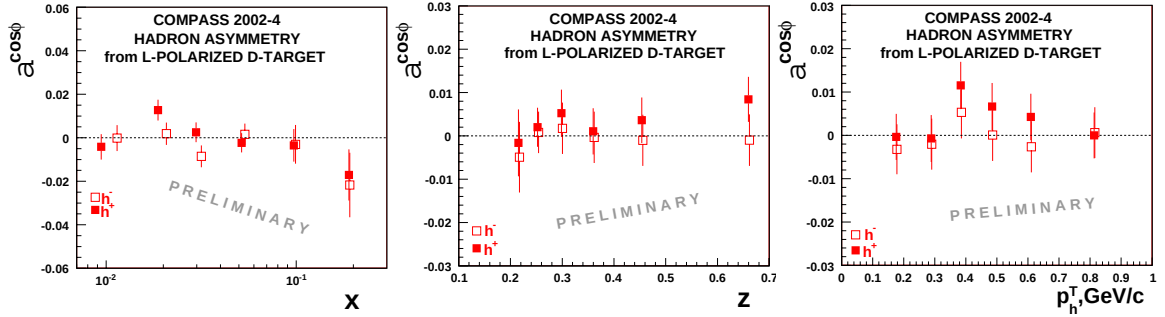


Figure 7: Dependence of the AA fit parameters $a^{\cos\phi}$ on kinematical variables.

The $\cos(\phi)$ modulation of the AA is studied for the first time. It is mainly due to a pure twist-3 PDF $g_L^\perp$ in $d\sigma_{LL}$, an analog to the Cahn effect [16] in unpolarized SIDIS.

5. Conclusions and prospects.

1. The azimuthal asymmetries ($a(\phi)$) in the SIDIS ($Q^2 > 1 \text{ GeV}^2$, $y > 0.1$) production of negative (h^-) and positive (h^+) hadrons by 160 GeV muons on the longitudinally polarized deuterium target, have been studied with the COMPASS data collected in 2002 – 2004.
2. For the integrated over x , z and p_h^T variables all ϕ -modulation amplitudes of $a(\phi)$ are consistent with zero within errors, while the ϕ -independent parts of the $a(\phi)$ differ from zero and are almost equal for h^- and h^+ .
3. The amplitudes as functions of kinematical variables are studied in the region of $x = 0.004 - 0.7$, $z = 0.2 - 0.9$, $p_h^T = 0.1 - 1 \text{ GeV}/c$. It was found that:
 - ϕ -independent parts of the $a(\phi)$, $a^{\text{const}}(x)/D_0 = A_d^h$, where D_0 is a virtual photon depolarization factor, are in agreement with the COMPASS published data [13] on A_d^h , calculated by another method and using different cuts;

- the amplitudes $a^{\sin\phi}(x, z, p_h^T)$ are small and in general do not contradict to the HERMES data [8], if one takes into account the difference in x , Q^2 and W between the two experiments. One can also note, that in the HERMES experiment the asymmetries are calculated for identified leading pions, while in this analysis every hadron is included in the asymmetry evaluations;
 - the amplitudes $a^{\sin 2\phi}$, $a^{\sin 3\phi}$ and $a^{\cos\phi}$ are consistent with zero within statistical errors of about 0.5% (only statistical errors are shown in the plots while systematic errors are estimated to be much smaller).
4. The results of this analysis are obtained with restriction $z > 0.2$ of the energy fraction of the hadron in order to assure that it comes from the current fragmentation region. This request removes almost one half of statistics. The tests have shown that with a lower cut, $z > 0.05$, the results are identical.
 5. The reported data are preliminary. New data of 2006 from the deuterium target will be added. These data will increase the statistics by about a factor of 2. New data of 2007 from the hydrogen target will be very interesting in comparison with the effects already observed by the COMPASS and HERMES on the transversally polarized targets.

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TRANSVERSE SPIN AND MOMENTUM EFFECTS IN THE COMPASS EXPERIMENT

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Abstract

The study of the spin structure of the nucleon is part of the scientific program of COMPASS, a fixed target experiment at the CERN SPS. COMPASS investigates transverse spin and transverse momentum effects by studying the azimuthal distributions of the hadrons produced in deep inelastic scattering of naturally polarized 160 GeV/c muons on transversely polarized target. The results obtained using a ${}^6\text{LiD}$ from the data collected in 2002-2004, and new results obtained using a NH_3 target, from the data collected in 2007, are presented.

1 Introduction

A powerful method to access the nucleon structure is the measurement of the azimuthal modulations of the distribution of the hadrons inclusively produced in the deep inelastic scattering. On the basis of general principles of quantum field theory in the one photon exchange approximation, the semi inclusive deep inelastic scattering (SIDIS) cross section can be written in a model independent way [1].

In its expression there are 18 structure functions, 8 of them are leading order. They can be accessed measuring the corresponding modulation in different combinations of the azimuthal angle of the hadron and of the nucleon spin vector, all of them independent. Most of them have a clear interpretation in terms of the Parton Model and can be written as the convolution of a distribution function of the quarks inside the nucleon (PDF) and a function which describes the fragmentation of the struck quark into a specific hadron (FF).

Recent data on single spin asymmetries in SIDIS off transversely polarized nucleon targets [2], [3] triggered a lot of interest towards the transverse momentum dependent and spin dependent distributions and fragmentation functions. Correlations between spin and transverse momentum give rise to effects, which would be zero in the absence of intrinsic motion of the quarks inside the nucleon.

One of the most famous PDF is the transversity distribution which gives the difference of the number of quarks with momentum fraction x with their transverse spin parallel and anti-parallel to the nucleon spin [4]. The transversity distribution, together with the better known spin-averaged distribution and helicity distribution functions, is necessary to fully specify the nucleon structure at leading order. Since it is chiral odd it can be measured only coupled to another chiral odd function. In SIDIS one candidate is the Collins fragmentation function which is the spin dependent part of the FF of a polarized

quark in an unpolarized hadron. The convolution of the two functions generates the so called Collins asymmetry in the hadrons distribution. A second SIDIS channel to access the transversity distribution is in conjunction with the interference fragmentation function of the polarized quark in two hadrons giving rise to an azimuthal modulation of the plane of inclusively produced hadrons pairs with respect to the scattering plane.

Of special interest among the transverse momentum dependent distribution functions are: the Sivers function [5] which describes a possible deformation of the quark intrinsic transverse momentum distribution in a transversely polarized nucleon, and the Boer-Mulders function [6] which gives the correlation of the quark transverse momentum and its transverse spin in an unpolarized nucleon. The measurements of these functions can give further insights on the connection between the transverse spin and transverse momentum and are needed to progress towards a more structured picture of the nucleon structure, beyond the collinear partonic representation.

2 The COMPASS experiment

COMPASS is an high energy experiment at the CERN SPS with a wide physics program focused on the study of the nucleon spin structure and hadron spectroscopy using muon and hadron beams. The data used to investigate the transverse spin and transverse momentum structure of the nucleon have been taken with a positive muon beam of 160 GeV/c on a transversely polarized target. The scattered muon and the produced hadrons are detected in a two stage spectrometer with a wide angular acceptance and an excellent particle identification [2]. For the data taking with a transversely polarized target, in the years 2002 2003 and 2004, deep inelastic scattering data have been collected using a ${}^6\text{LiD}$ target material, while in 2007 a NH_3 target material was used.

The target material was placed in 2 (during 2004) or 3 (during 2007) cells oppositely transversely polarized and the direction of the target polarization has been reversed every 5 days in each target cell.

Many results have been determined from these data, in this report i will only discuss the results for unpolarized asymmetries from 2004 data and the results for the Sivers and Collins asymmetries from 2002-2004 and 2007 data.

3 Unpolarized target azimuthal asymmetries

The cross section for the hadron production of the lepton nucleon deep inelastic scattering off an unpolarized is [1]:

$$\begin{aligned} \frac{d^5\sigma}{dx dy dz d\phi_h dp_T^2} = & \frac{2\pi\alpha^2}{xyQ^2} \cdot \left\{ \frac{1 + (1-y)^2}{2} F_{UU} + (2-y) \sqrt{1-y} \cos\phi_h F_{UU}^{\cos\phi_h} \right. \\ & \left. + (1-y) \cos 2\phi_h F_{UU}^{\cos 2\phi_h} + \lambda_l y \sqrt{1-y} \sin\phi_h F_{LU}^{\sin\phi_h} \right\} \quad (1) \end{aligned}$$

There are three independent modulations on the azimuthal angle of the hadron with respect to the scattering plane. The $\sin\phi_h$ modulation is proportional to the beam polarization and the structure function $F_{LU}^{\sin\phi_h}$ has no clear explanation in term of the parton

model. Both the $\cos\phi_h$ and the $\cos 2\phi_h$ modulations can be explained in terms of the Cahn effect and the already introduced Boer Mulders distribution function and are related to the transverse momentum of the quark inside the nucleon. One has to notice that the magnitude of the Boer Mulders function can depend upon the quark flavor while the Cahn effect is a pure kinematical effect which depends only from the struck quark $\vec{k}_T$ and is expected to be the same for positive and negative hadrons. Moreover the Cahn effect is the main contribution to the $\cos\phi_h$ asymmetry [7].

Also pQCD gives rise to azimuthal modulations in the hadron production but it has been shown [7] that these effects are negligible for the range of transverse momentum p_T of the hadrons produced in COMPASS.

3.1 Analysis of unpolarized asymmetries

The analyzed data have been taken during 2004 with a polarized deuteron target. The kinematical cuts which have been applied to select the DIS region are:

- the squared four momentum transfer $Q^2 > 1 \text{ (GeV/c)}^2$,
- the hadronic invariant mass $W > 5 \text{ GeV/c}^2$ and
- the fractional energy transfer of the muon $0.1 < y < 0.9$.

The analyzed hadrons sample include all charged hadrons coming from the interaction with:

- transverse momentum $p_T > 0.1 \text{ GeV/c}$ to have a good p_T resolution and
- fraction of the available energy $z > 0.2$ to exclude hadrons coming from the target fragmentation region.
- $p_T < 1.5 \text{ GeV/c}$ and $z < 0.85$, specific for the unpolarized asymmetries analysis.

Since the data were taken with the polarized target (longitudinally or transversely), both the two spin configurations (up/down) collected from the transverse data taking and the two spin configurations (forward/backward) collected from the longitudinal data taking were individually combined in order to cancel spin related effects.

In the extraction of the unpolarized asymmetries there is no cancellation of the apparatus acceptance and a Monte Carlo simulation is needed to evaluate it. The Monte Carlo chain used in COMPASS starts from Lepto for the generation of SIDIS events, while the particle interactions in the materials and the detectors response simulation is performed by COMGEANT, based on GEANT3. The Monte Carlo data are finally reconstructed by the same software used for the real data.

The experimental apparatus acceptance as a function of the azimuthal angle ϕ_h is then calculated as the ratio between the generated and the reconstructed hadrons azimuthal distribution performed in a given kinematical range. An example of the apparatus acceptance correction is shown in figure 1.

The measured unpolarized hadron azimuthal distributions are divided by the acceptance and the resulting distributions are fit with the function:

$$N_{corr}(\phi_h) = N_0 \cdot (1 + A_{\cos(\phi_h)}^D \cdot \cos(\phi_h) + A_{\cos(2\phi_h)}^D \cdot \cos(2\phi_h) + A_{\sin(\phi_h)}^D \cdot \sin(\phi_h)).$$

The $A_{\cos(\phi_h)}^D$, $A_{\cos(2\phi_h)}^D$, $A_{\sin(\phi_h)}^D$ amplitudes are extracted for each considered bin of x , z and p_T ; when extracting the asymmetries as function of one of these three variables, the other two are integrated over all their SIDIS kinematical range. The Monte Carlo simulation plays an important role in the unpolarized asymmetries measurement. Two different sets of Lepto parameters have been used to evaluate a possible source of systematic errors.

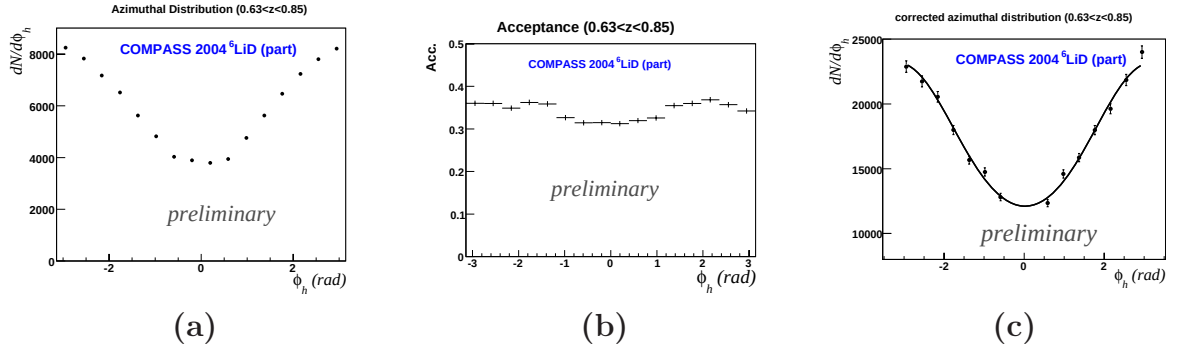


Figure 1: (a) Measured spin averaged azimuthal distribution of the hadrons for $0.63 < z < 0.85$. (b) Azimuthal acceptance calculated from the ratio between Monte Carlo generated and Monte Carlo reconstructed hadrons distribution in the same z bin. (c) Corrected azimuthal distribution corrected by the acceptance with the fit superimposed.

Both of them give a satisfactory description of the COMPASS data in the main kinematical variables, both for data collected with transversely and longitudinally polarized deuteron target. All the asymmetries extracted with different Monte Carlo have been compared in order to estimate the systematic errors. Further systematic tests, like splitting the data sample according to the event topology and the time of the measurement gave no significant contributions.

3.2 Results

The $\sin(\phi_h)$ asymmetries measured by COMPASS are compatible with zero over the full range of x , z and p_T both for positive and for negative hadrons and suggest a beam polarization effect compatible with zero on the deuteron target.

The COMPASS results for all $\cos(\phi_h)$ asymmetries are shown in figure 2a, for pos-

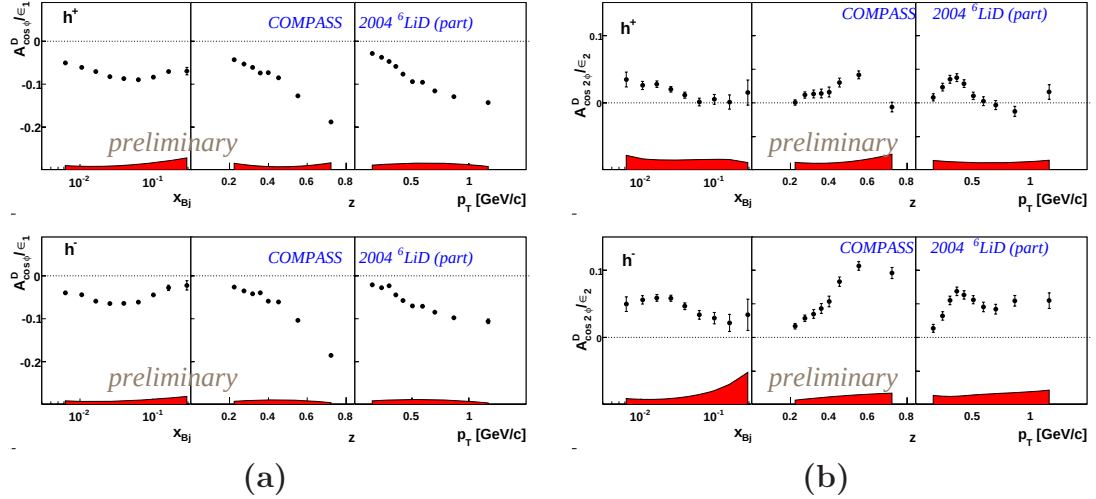


Figure 2: (a) Top picture: $\cos(\phi_h)$ asymmetries extracted from positive hadrons. Bottom picture: $\cos(\phi_h)$ asymmetries extracted from negative hadrons. (b) Top picture: $\cos(2\phi_h)$ asymmetries extracted from positive hadrons. Bottom picture: $\cos(2\phi_h)$ asymmetries extracted from negative hadrons. All asymmetries shown in each picture are extracted as a function x (left plot), z (middle plot) and p_T (right plot), the red bands are the systematic errors. The shown quantities are the amplitudes of each modulation divided by the y dependent corresponding kinematical factor.

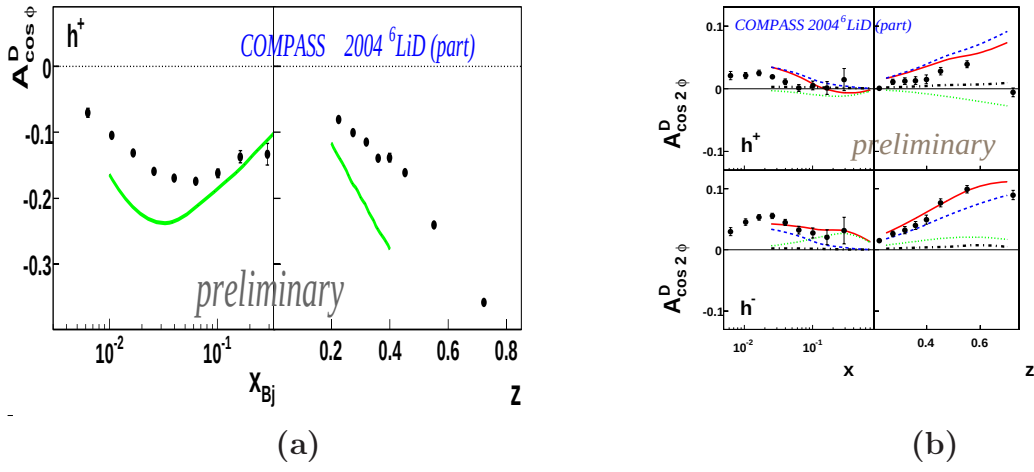


Figure 3: **(a)** Comparison with theoretical predictions. Black points are the $\cos(\phi_h)$ amplitude extracted from COMPASS data, positive hadrons, as a function of x and z and in green are shown the values predicted in [7] for COMPASS kinematics without considering the Boer Mulders effect. **(b)** Comparison with theoretical predictions. Black points are the $\cos(2\phi_h)$ amplitude extracted from COMPASS data, for positive hadrons (upper two plots) and negative hadrons (lower row), as a function of x and z . The red line shows the predictions from [8], and is the sum of Cahn contribution (blue line), Boer Mulders contribution (green dashed line) and perturbative QCD (black dotted line). These values have been calculated for the COMPASS kinematical region.

itive (upper plot) and negative (bottom plot) hadrons. The measured asymmetries are large and slightly different for positive and negative hadrons. The COMPASS results for all $\cos(2\phi_h)$ asymmetries are shown in figure 2b, for positive (upper plot) and negative (bottom plot) hadrons.

A comparison of the $\cos(\phi_h)$ asymmetries with the theoretical predictions for the COMPASS kinematics [7] is shown in figure 3a. In [7] only Cahn effect was considered. The differences between positive and negative hadrons measured from COMPASS data could hint to a possible Boer-Mulders PDF contribution.

The results for the $\cos(2\phi_h)$ asymmetry are compared to theoretical predictions [8], made for the COMPASS kinematical region, in figure 3b. In [8] different contributions were taken into account: the Cahn effect (blue dashed line in the picture), the Boer-Mulders PDF (green dashed line) and the first order pQCD contribution which is negligible as shown by the black dashed-dotted line. In this work the Boer-Mulders PDF was assumed to be proportional to the better known Siverts function. Comparing the results obtained for positive and negative hadrons one can see that only the contribution coming from the Boer-Mulders PDF changes significantly with the hadron charge. COMPASS measurements confirm this trend and agree quite well with these predictions.

4 Transverse spin dependent azimuthal asymmetries

4.1 The Collins asymmetry

In semi inclusive deep inelastic scattering the transversity distribution function $\Delta_T q(x)$ can be measured combined with the Collins fragmentation function $\Delta_T^0 D_q^h(z, p_T)$. According to Collins, the fragmentation of a transversely polarized quark into an unpolarized hadron generates an azimuthal modulation in the hadron distribution with respect to the

scattering plane [12]. The hadron yield $N(\Phi_{Coll})$ can be written as:

$$N(\Phi_{Coll}) = N_0 \cdot (1 + \dots + f \cdot P \cdot D_{NN} \cdot A_{Coll} \cdot \sin(\Phi_{Coll})), \quad (2)$$

where N_0 is the average hadron yield, f is the fraction of polarized material in the target, P is the target polarization, $D_{NN} = (1-y)/(1-y+y^2/2)$ is the depolarization factor and y is the fractional energy transfer of the muon. The dots stand for the other terms of the SIDIS cross section, all of them are proportional to independent azimuthal modulation. The Collins angle is $\Phi_{Coll} = \phi_h + \phi_s - \pi$ where ϕ_h is the hadron azimuthal angle and ϕ_s is the azimuthal angle of the nucleon spin, with respect to the scattering plane [4]. The Collins asymmetry can be expressed as:

$$A_{Coll} = \frac{\sum_q e_q^2 \cdot \Delta_T q(x) \cdot \Delta_T^0 D_q^h(z, p_T)}{\sum_q e_q^2 \cdot q(x) \cdot D_q^h(z, p_T)}, \quad (3)$$

where the sum runs over the quark flavors and e_q is the quark charge, $D_q^h(z, p_T)$ is the unpolarized fragmentation function, $q(x)$ is the unpolarized parton distribution function.

4.2 Transversely polarized two hadrons asymmetries

The transversity distribution can also be measured in SIDIS combined with the interference fragmentation function $H_1^\lessgtr(z, M_{inv}^2)$ where M_{inv} is the invariant mass of the hadrons pair [13]. The fragmentation of transversely polarized quarks into two hadrons can give an azimuthal modulation in the produced hadrons distribution which, for two hadrons of opposite charge, can be written as:

$$N_{h^+h^-} = N_0 \cdot (1 + f \cdot P \cdot A_{RS} \cdot \sin(\Phi_{RS}) \cdot \sin(\theta)), \quad (4)$$

where the angle $\Phi_{RS} = \phi_R + \phi_s - \pi$ is given by the sum of the azimuthal angle of the spin vector ϕ_s and the azimuthal angle of $\vec{R}_T$ with respect to the lepton scattering plane. $\vec{R}_T$ is the transverse component, with respect to the virtual photon direction, of the vector $\vec{R}$ defined as:

$$\vec{R} = (z_2 \cdot \vec{p}_1 - z_1 \cdot \vec{p}_2) / (z_1 + z_2), \quad (5)$$

where $\vec{p}_1$ and $\vec{p}_2$ are the momenta in the laboratory frame of h^+ and h^- respectively. θ is the angle between the momentum vector of h^+ in the center of mass frame of the hadrons pair and the momentum vector of the two hadron system. The measured amplitude is proportional to the product of the transversity distribution and the polarized two-hadrons interference fragmentation function:

$$A_{RS} \propto \frac{\sum_q e_q^2 \cdot \Delta_T q(x) \cdot H_1^\lessgtr(z, M_{inv}^2)}{\sum_q e_q^2 \cdot q(x) \cdot D_q^{2h}(z, M_{inv}^2)}, \quad (6)$$

where $D_q^{2h}(z, M_{inv}^2)$ is the unpolarized two-hadron interference fragmentation function.

The hadrons used in the analysis have $z > 0.1$ and $x_F > 0.1$ to exclude hadrons coming from the target fragmentation region; $z_1 + z_2 < 0.9$ to exclude pairs from exclusively produced ρ^0 mesons and $R_T > 0.07$ GeV/c to have a good resolution on the azimuthal angle ϕ_R .

4.3 The Sivers asymmetry

According to Sivers, the coupling of the intrinsic transverse momentum $\vec{k}_T$ of unpolarized quarks with the spin of a transversely polarized nucleon can give an asymmetry in the hadrons azimuthal distribution [5]. The correlation between the transverse nucleon spin and the quark transverse momentum is described by the transverse momentum dependent PDF $\Delta_0^T q(x, \vec{k}_T)$ called the Sivers function. The hadron yield can be written as:

$$N(\Phi_{Siv}) = N_0 \cdot (1 + \dots + f \cdot P \cdot D_{NN} \cdot A_{Siv} \cdot \sin(\Phi_{Siv})), \quad (7)$$

where the Sivers angle $\Phi_{Siv} = \phi_h - \phi_s$. The Sivers asymmetry can be expressed as:

$$A_{Siv} = \frac{\sum_q e_q^2 \cdot \Delta_0^T q(x, \vec{k}_T) \cdot D_q^h(z, p_T)}{\sum_q e_q^2 \cdot q(x) \cdot D_q^h(z, p_T)}, \quad (8)$$

Since Collins and Sivers asymmetries are given by independent azimuthal modulations in the SIDIS hadrons distribution, they can be measured from the same data set. The kinematical selection is similar to the one applied for unpolarized asymmetries analysis, previously described in section 3.1.

4.4 Results

The Collins and Sivers asymmetries measured from the 2004 transversely polarized deuteron data are compatible with zero [9], within the statistical and systematical errors, over the full range of x , z and p_T . A global fit to the COMPASS results on deuteron and of the HERMES measurements on proton of the Collins asymmetries and to the BELLE results on the Collins fragmentation function, led to the first extraction of the transversity distribution function [10]. The emerging picture is that the transversity distributions are different from zero, of the same magnitude and opposite sign for the u and d quark leading to a cancellation of the Collins asymmetry for an isoscalar target like deuteron. A global fit to HERMES proton data and COMPASS deuteron data suggested that a similar cancellation for the deuteron target occurs for the Sivers function [11].

Obviously there was a great interest for the results on polarized proton at the COMPASS energy. In 2007 a polarized NH_3 (proton) target was used and the first preliminary results were produced in 2008.

In figure 4 the results for the Collins asymmetries on proton are shown. The measured asymmetry is small and compatible with zero up to $x = 0.05$ and then increases up to about 10% in the x valence region, with opposite sign for positive and negative hadrons. The asymmetry has the same sign and size of the one measured by HERMES. The COMPASS measurement confirms the HERMES result at an higher energy, i.e. in a region where higher twist effects are less important.

In figure 5 the two-hadrons asymmetries are shown. At variance with the deuteron results the asymmetry is evident in the x -valence region and there is a strong signal in M_{inv} around the ρ^0 mass.

In figure 6 the Sivers asymmetries are shown. They are small and compatible with zero over all the x range both for positive and negative hadrons. This is a surprising result since no strong Q^2 dependence is expected for the Sivers distribution function.

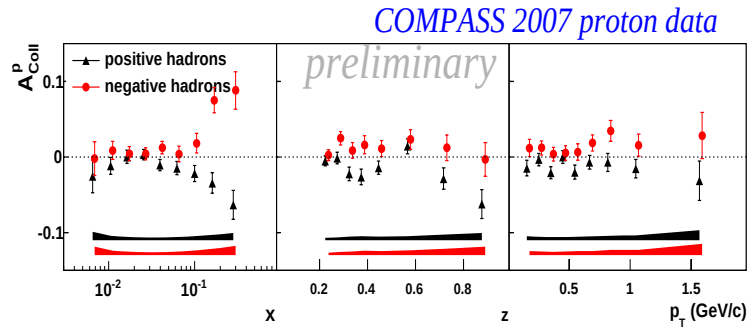


Figure 4: COMPASS Collins asymmetries on polarized proton for positive (black triangles) and negative (red circles) hadrons as a function of x , z , and p_T . The red bands stand for the systematic errors.

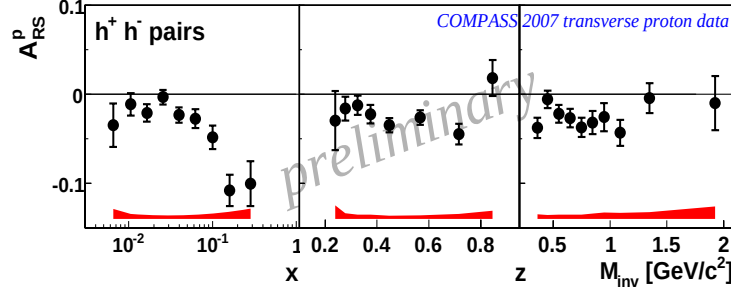


Figure 5: COMPASS two-hadrons asymmetries on polarized proton for positive (black triangles) and negative (red circles) hadrons as a function of x , z , and M_{inv} . The red bands stand for the systematic errors.

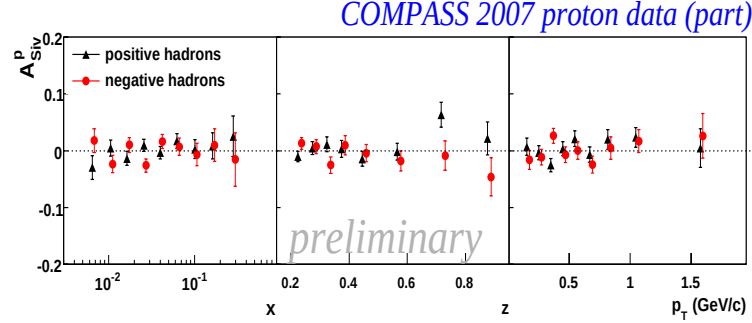


Figure 6: COMPASS Sivers asymmetries on polarized proton for positive (black triangles) and negative (red circles) hadrons as a function of x , z , and p_T . The systematic errors are about 0.5 of the statistical ones.

5 Outlook

An important contribution to the understanding of the nucleon structure comes from these SIDIS results. The work is still ongoing and more precise high energy data, complementary to the measurement performed at JLAB, are needed and COMPASS will have a full year dedicated data taking with polarized proton in 2010.

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DELTA-SIGMA EXPERIMENT - THE RESULTS OBTAINED.
 $\Delta\sigma_T(np)$ MEASUREMENTS PLANNED AT THE NUCLOTRON-M

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Abstract

Under the research program of "Delta-Sigma experiment" project the energy dependences of the $-\Delta\sigma_L(np)$ observable and ratio R_{dp} have been obtained over the neutron beam energy region of several GeV . Further measurements of the $-\Delta\sigma_{L,T}(np)$ and $A_{ookk}(np)$, $A_{oonn}(np)$ energy dependences using L and T orientations of beam and target polarizations will be possible in the near future when both the new high intensity source of polarized deuterons will be put in operation at the Nuclotron and the T mode of target polarization will be ready. An intensity of the polarized deuteron beam from new creating polarized deuteron ions source at the Nuclotron-M will be more then ten times higher then from the existing source "Polaris".

1. Introduction. The nucleon beam energy region of $1 - 5 GeV$ is considered as transitional one between the usual boson exchange mechanism of NN interaction and nonperturbative QCD one. So the new experimental NN results to be obtained will promote to create an adequate phenomenological and theoretical description of NN interaction over this region. The measurements of the energy behavior of elastic np spin-dependent observables over this energy region have been proposed [1], successfully started and carried out at the JINR LHE under the research program of project "Delta-Sigma experiment" [2].

The aim of the Delta-Sigma experimental program is to extend investigations of properties of NN interaction over a new high energy region of free polarized neutron beams. This neutron beam energy region from 1.2 up to $3.7 GeV$ is provided in the last years only by the JINR LHE accelerators. The main task of these studies is to determine for the first time the imaginary and real parts of all the $np \rightarrow np$ forward ($\theta_{n,CM} = 0^\circ$) and backward ($\theta_{n,CM} = 180^\circ$) spin dependent np scattering amplitudes over the specified energy region. To reach this aim a sufficient data set on energy dependences of np spin dependent observables has to be obtained for direct and unambiguous reconstruction of these amplitudes.

Under the Delta-Sigma research program [2] the energy dependences of the $-\Delta\sigma_L(np)$ observable and ratio $R_{dp} = (d\sigma/d\Omega)(nd)/(d\sigma/d\Omega)(np)$ have been obtained over the neutron beam energy region of several GeV . Further measurements of $-\Delta\sigma_{L,T}(np)$ and $A_{ookk}(np)$, $A_{oonn}(np)$ energy dependences using L and T orientations of beam and target polarizations will be possible in the near future when both the new high intensity source of polarized deuterons will be put in operation at the Nuclotron and the T mode of target polarization will be ready.

2. Delta-Sigma Experimental Set-up. Transmission measurements over 1.2 - 3.7 GeV of the energy dependence of $\Delta\sigma_L(np)$ and $\Delta\sigma_T(np)$ have been proposed [1] and started in Dubna in collaboration with the groups from FSU, France and U.S. To speed up the proposed measurements the main part of the set-up elements was provided by the project collaborators. A large (140cm^3) Argonne-Saclay polarized proton target (PPT) [3, 4] was reconstructed, shipped to Dubna, installed at the neutron beam line and tested. The new polarized neutron beam lines [5, 6] with suitable parameters were prepared, tuned, and tested by the JINR LHE specialists. The set of neutron detectors with corresponding electronics (Saclay, Kurchatov Institute and JINR LHE), data acquisition system (JINR LHE), deuteron beam line polarimeters (PINP, Gatchina and JINR LHE), and other equipment needed were also prepared, tuned, and tested. The set-up for the transmission measurements is fully functional but some parts of it needed to be upgraded. More comprehensive description of the Delta-Sigma set-up for transmission measurements elements is given in [7 - 8, 11].

The magnetic spectrometer with two sets of multiwire proportional chambers [9 - 10, 11] for detecting protons from $np \rightarrow pn$ elastic and $nd \rightarrow p(nn)$ quasi-elastic charge-exchange processes at $\theta_{p,\text{Lab}} = 0^\circ$, was installed and tested at the free neutron-beam line of the JINR VBLHE Nuclotron facility. The spectrometer is a part of the DeltaSigma setup.

3. Main Results Obtained under the Project.

Measurements of the $-\Delta\sigma_L(np)$ energy dependence were made using the L - polarized deuteron beam from the JINR LHE Synchrophasotron and the Dubna L - polarized proton target. Three data taking runs (1995, 1997 and 2001) were carried out for the $\Delta\sigma_L(np)$ measurements with an adequate intensity up to 5×10^9 particles/cycle of polarized deuterons. An enough detailed data set at 10 energy points for the $-\Delta\sigma_L(np)$ energy behavior over 1.2 II 3.7 GeV was obtained [7 - 8]. See also fig.1.

The measurements of the $-\Delta\sigma_L(np)$ energy dependence were in the main completed. The Dubna group of the participants of this investigation has gained wide experience in carrying out transmission measurements and is ready to measure the $\Delta\sigma_T(np)$ spin dependent observable. The detectors for transmission measurement, electronics and data acquisition system are also ready.

In the last years, the significant efforts were made to test and tune the magnetic spectrometer equipment and there were provided the possibilities for

$A_{ookk}(np)$, $A_{oonn}(np)$ and R_{dp} measurements using prepared magnetic spectrometer. A number of data taking runs were successfully carried out using an intense quasimonochromatic unpolarized neutron beam and liquid H_2 , liquid D_2 , C , CH_2 and CD_2 targets, placed in the neutron beam line instead of the PPT. Measurements of the ratio R_{dp} of the

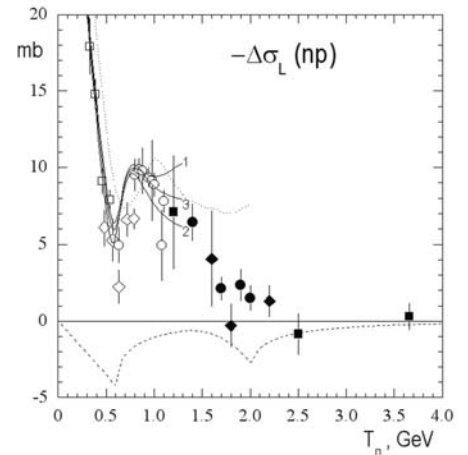


Figure 1: Energy dependence of the $-\Delta\sigma_L(np)$: $\bullet$, $\blacksquare$, $\blacklozenge$ JINR results [7 - 8]; $\square$ PSI; $\diamond$ LAMPF; $\circ$ Saturne II; (solid curves 1, 2, and 3) FA85, SP99, and SP03 ED GW/VPI PSA solutions, respectively; (dotted curve) meson-exchange model; (dashed curve) contribution from nonperturbative QCD interaction induced by instantons [7, 8].

charge-exchange quasi-elastic differential cross-section for the reaction $nd \rightarrow p(nn)$ at the outgoing proton angle $\theta_{p,\text{Lab}} = 0^\circ$ to the elastic charge-exchange differential cross-section $np \rightarrow pn$, were carried out at the neutron beam kinetic energies of 0.55, 0.8, 1.0, 1.2, 1.4, 1.8 and 2.0 GeV. The obtained $R_{dp} = d\sigma/d\Omega(nd \rightarrow p(nn))/d\sigma/d\Omega(np \rightarrow pn)$ results were published in [9 - 10]. See also fig.2.

Analysis of the experimental R_{dp} results obtained proves that this data can be used for the Delta-Sigma experimental program to reduce the total ambiguity in the extraction of the amplitude real parts.

The testing of the spectrometer elements and carrying out of the R_{dp} measurements were a powerful test for the future measurements of the spin correlation parameters $A_{ookk}(np)$ and $A_{oonn}(np)$. These measurements will be made using the same method and the existing described spectrometer. We have gained wide experience in carrying out such charge-exchange measurements.

4. Readiness of the Dubna Polarized Proton Target. The equipment (superconducting coils) needed for the providing of T mode of target polarization has been prepared and complex tests of coils together with systems MPT were begun. But since a middle of 2007 the financing of the MPT modernization works was completely stopped. Works on modernization of the Saclay-Argonne-JINR polarized proton target (setup PPT) will be foreseen by the new JINR theme 02-1-1097-2010/12 "Study of Polarization Phenomena and Spin Effects at the JINR Nuclotron-M Facility".

5. High Intensity Polarized Deuteron Beam at the Nuclotron-M. A creation of new source of polarized deuteron (SPD) ions [12] is foreseen now for the Nuclotron-M accelerator under the JINR theme 02-0-1065-2007/2009. The proposed project of source assumes the development of a universal high-intensity source of polarized deuterons (protons) using a charge-exchange plasma ionizer. The output current of the source under design will be up to 10 mA for $\uparrow D^+(\uparrow H^+)$ and polarization will be up to 90% of the maximal vector (-1) and tensor (+1, -2) polarization. The project is based on the equipment which was supplied within the framework of an agreement between JINR and IUCF (Bloomington, USA).

Recent status of the source for polarized deuteron creation is shown in [13]. The SPD putting into operation is in the end of 2010.

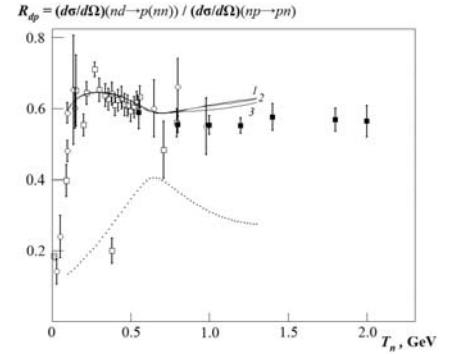


Figure 2: The R_{dp} energy dependence. Black squares our data [9 - 10], open squares and circles for the existing data at energies below 1 GeV. The open squares represent the R_{dp} data obtained in the $nd \rightarrow p(nn)$ reaction measurements and the open circles represent the R_{dp} data obtained in either the $pd \rightarrow n(pp)$ or the $dp \rightarrow (pp)n$ reactions measurements. The curves represent the energy behaviour of the R_{dp} calculated using NN invariant amplitude data sets from the energy-dependent phase-shift analyses. The solid curves are the calculations with the amplitude data sets for the elastic $np \rightarrow pn$ charge-exchange at $\theta_{p,\text{CM}} = 0^\circ$: curve 1 corresponds to the SP07 solution, curve 2 to SP00, and curve 3 to SM97. The dotted curve represents the R_{dp} values with the amplitudes data set for the elastic $np \rightarrow np$ scattering at $\theta_{n,\text{CM}} = \pi$ (for the SP07 solution).

6. Estimate of the count rate in transmission measurements using new SPD.

An expression for the measured $-\Delta\sigma_{L,T}(np)$ value is following

$$-\sigma_{L,T} = \ln [(M^-T^+)/ (M^+T^-)] / (|P_B P_T| n_H), \quad (6.1)$$

where the M^+ , M^- and T^+ , T^- are the statistics for monitor and transmission neutron detectors with the P_B^+ and P_B^- neutron beam polarizations, respectively. The P_T is the polarized proton target polarization value and n_H is the number of polarizable hydrogen atoms at polarized proton target.

A systematic uncertainty of the $-\Delta\sigma_{L,T}$ value is mainly caused by the errors in $|P_B|$, $|P_T|$, and n_H . The statistical error of $-\Delta\sigma_{L,T}$ is given by the formula

$$\delta_{stat} = \sqrt{(1/M^+ + 1/M^- + 1/T^+ + 1/T^-) / (|P_B P_T| n_H)}. \quad (6.2)$$

where the M^+ , M^- and T^+ , T^- are the statistics for monitor and transmission neutron detectors with the P_B^+ and P_B^- neutron beam polarizations, respectively.

To estimate the required statistics for the monitor neutron detectors M in order to reach desired value of the $\Delta\sigma_{L,T}$ statistical error δ_{stat} , we substitute the experimentally obtained n_H , $|P_B|$, $|P_T|$, and T/M values [7 - 8] in to this formula:

$$n_H = 9.14 \times 10^{23}; \quad |P_B| = 0.5; \quad |P_T| = 0.7; \quad T/M = 0.82.$$

To reach the δ_{stat} value of $\sim 1 mb$ it is required to accumulate the M statistics equal to $\sim 4.5 \times 10^7$ counts. From our experience obtained during the np charge-exchange measurements [9 - 10] using high intensity ($\sim 2 \times 10^{10}$ particles per cycle) unpolarized deuteron beam, such M statistics will be accumulated within 8 hours approximately.

Therefore to measure the $-\Delta\sigma_T$ values at 10 points of the neutron beam energy T_n it is required about ~ 80 hours of the "pure" run time for statistics accumulation. The count rate which will be accessible with new source of polarized deuteron ions at the Nuclotron-M estimated above, allows obtaining the $-\Delta\sigma_T$ values with very high precision [14].

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COMPASS RESULTS ON GLUON POLARISATION FROM HIGH p_T HADRON PAIRS

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Abstract

One of the goals of the COMPASS experiment is the determination of the gluon polarisation $\Delta G/G$, for a deep understanding of the spin structure of the nucleon. In DIS the gluon polarisation can be measured via the Photon-Gluon-Fusion (PGF) process, identified by open charm production or by selecting high p_T hadron pairs in the final state. The data used for this work were collected by the COMPASS experiment during the years 2002-2004, using a 160 GeV naturally polarised positive muon beam scattering on a polarised nucleon target. A new preliminary result of the gluon polarisation $\Delta G/G$ from high p_T hadron pairs in events with $Q^2 > 1 \text{ (GeV/c)}^2$ is presented. In order to extract $\Delta G/G$, this analysis takes into account the leading process γq contribution together with the PGF and QCD Compton processes. A new weighted method based on a neural network approach is used. A preliminary $\Delta G/G$ result for events from quasi-real photoproduction ($Q^2 < 1 \text{ (GeV/c)}^2$) is also presented.

1 Introduction

The COMPASS experiment is located in the Super Proton Synchrotron (SPS) accelerator at CERN. For a more complete description of the experimental apparatus the reader is addressed to [1]. In 2007, the COMPASS collaboration estimated the quark contribution to the nucleon spin with high precision [2], using a NLO QCD fit with all world data available. This contribution confirms that approximately 1/3 of the nucleon spin is carried by the quarks, as demonstrated by earlier experiments [3].

The nucleon spin can be written as:

$$\frac{1}{2} = \frac{1}{2}\Delta\Sigma + \Delta G + L \quad (1)$$

$\Delta\Sigma$ and ΔG are, respectively, the quark and gluon contributions to the nucleon spin and L is the orbital angular momentum contribution coming from the quarks and gluons.

The aim of this analysis is to estimate the gluon polarisation, $\Delta G/G$, using the high transverse momentum (high p_T) hadron pairs sample. The analysis is performed in two complementary kinematic regions: $Q^2 < 1 \text{ (GeV/c)}^2$ (low Q^2) and $Q^2 > 1 \text{ (GeV/c)}^2$ (high Q^2) regions. The present work is mainly focused on the analysis for high Q^2 . However, the analysis for the low Q^2 region is summarised in sec. 6.

For completeness, the slides of the presentation can be found in [4].

2 Analysis Formalism

Spin-dependent effects can be measured experimentally using the helicity asymmetry

$$A_{LL} = \frac{\Delta\sigma}{2\sigma} = \frac{\sigma^{\uparrow\downarrow} - \sigma^{\uparrow\uparrow}}{\sigma^{\uparrow\downarrow} + \sigma^{\uparrow\uparrow}} \quad (2)$$

defined as the ratio of polarised ($\Delta\sigma$) and unpolarised (σ) cross sections. $\uparrow\uparrow$ and $\uparrow\downarrow$ refer to the parallel and anti-parallel spin helicity configuration of the beam lepton ($\uparrow$) with respect to the target nucleon ($\uparrow$ or $\downarrow$).

According to the factorisation theorem, the (polarised) cross sections can be written as the convolution of the (polarised) parton distribution functions, $(\Delta)q_i$, the hard scattering partonic cross section, $(\Delta)\hat{\sigma}$, and the fragmentation function D_f .

The gluon polarisation is measured directly via the Photon-Gluon Fusion process (PGF); which allows to probe the spin of the gluon inside the nucleon. To tag this process directly in DIS a high p_T hadron pairs data sample is used to calculate the helicity asymmetry. Two other processes compete with the PGF process in leading order QCD approximation, namely the virtual photo-absorption leading order (LO) process and the gluon radiation (QCD Compton) process. In Fig. 1 all contributing processes are depicted.

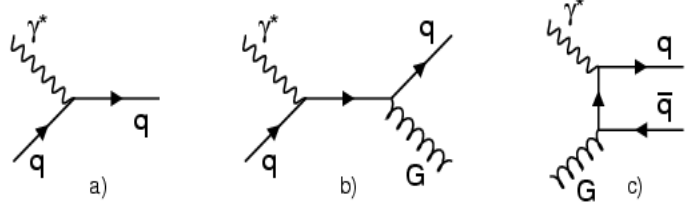


Figure 1: The contributing processes: a) DIS LO, b) QCD Compton and c) Photon-Gluon Fusion.

The helicity asymmetry for the high p_T hadron pairs data sample can thus be schematically written as:

$$A_{LL}^{2h}(x_{Bj}) = R_{PGF} a_{LL}^{PGF} \frac{\Delta G}{G}(x_G) + R_{LO} D A_1^{LO}(x_{Bj}) + R_{QCDC} a_{LL}^{QCDC} A_1^{LO}(x_C) \quad (3)$$

The R_i (the index i refers to the different processes) are the fractions of each process. a_{LL}^i represents the partonic cross section asymmetries, $\Delta\hat{\sigma}^i/\hat{\sigma}^i$, (also known as analysing power). D is the depolarisation factor¹. The virtual photon asymmetry A_1^{LO} is defined as $A_1^{LO} \equiv \frac{\sum_i e_i^2 \Delta q_i}{\sum_i e_i^2 q_i}$.

To extract $\Delta G/G$ from eq. (3) the contribution from the physical background processes LO and QCD Compton needs to be estimated. This is done using Monte Carlo (MC) simulation to calculate R_i fractions and a_{LL}^i . The virtual photon asymmetry A_1^{LO} is estimated using a parametrisation based on the A_1 asymmetry of the inclusive data [5]. Therefore a similar equation to (3) can be written to express the inclusive asymmetry of a data sample, A_{LL}^{incl} .

Using eq. (3) for the high p_T hadron pairs sample and a similar eq. for the inclusive sample the following expression is obtained:

¹The depolarisation factor is the fraction of the muon beam polarisation transferred to the virtual photon.

$$\begin{aligned}
\frac{\Delta G}{G}(x_G^{av}) &= \frac{A_{LL}^{2h}(x_{Bj}) + A^{corr}}{\beta} \\
A^{corr} &= -A_1(x_{Bj})D \frac{R_{LO}}{R_{LO}^{incl}} - A_1(x_C)\beta_1 + A_1(x'_C)\beta_2
\end{aligned} \tag{4}$$

and

$$\begin{aligned}
\beta_1 &= \frac{1}{R_{LO}^{incl}} \left[a_{LL}^{QCDC} R_{QCDC} - a_{LL}^{incl, QCDC} R_{QCDC}^{incl} \frac{R_{LO}}{R_{LO}^{incl}} \right] \\
\beta_2 &= a_{LL}^{incl, QCDC} \frac{R_{QCDC}^{incl}}{R_{LO}^{incl}} \frac{R_{QCDC}}{R_{LO}^{incl}} \frac{a_{LL}^{QCDC}}{D} \\
\alpha_1 &= a_{LL}^{PGF} R_{PGF} - a_{LL}^{incl, PGF} R_{LO} \frac{R_{PGF}^{incl}}{R_{LO}^{incl}} \\
\alpha_2 &= a_{LL}^{incl, PGF} R_{QCDC} \frac{R_{PGF}^{incl}}{R_{LO}^{incl}} \frac{a_{LL}^{QCDC}}{D} \\
\beta &= \alpha_1 - \alpha_2.
\end{aligned} \tag{5}$$

The term A^{corr} comprises the correction due to the other two processes, namely the LO and the QCD Compton processes. α_1 , α_2 , β_1 , β_2 , x_C , x'_C and x_G^{av} are estimated using high p_T and inclusive MC samples.

3 Data Selection

Data from 2002 to 2004 years is used. The selected events have an interaction vertex containing an incoming muon beam and a scattered muon. As mentioned in sec. 2 the data samples are divided into two data sets: the high p_T hadron pairs and the inclusive data samples.

Both data sets have the $Q^2 > 1 \text{ (GeV/c)}^2$ kinematic cut applied. Another cut is applied on the fraction of energy taken by the virtual photon, y : $0.1 < y < 0.9$. These cuts described previously are used to select the *inclusive sample*.

In the high p_T hadrons data sample, besides the inclusive selection, events with (at least) two outgoing high p_T hadrons are considered. These so-called hadron candidates must fulfill the following requirement: the two hadrons with the highest transverse momentum must have $p_T > 0.7 \text{ GeV/c}$. This requirement constitutes the high p_T cut. All these cuts additionally correspond to the *high p_T sample* selection.

4 Monte Carlo simulation

Important information to be used in the $\Delta G/G$ extraction is obtained from MC simulation. In this analysis it is fundamental that the simulation describes well the experimental data. Two MC samples were produced: one for the high p_T sample and another for the inclusive sample, to estimate the terms in the set of eq. (5).

The MC production comprises three steps: first the events are generated, then the particles pass through a simulated spectrometer using a program based on GEANT3 [6] and finally the events are reconstructed using the same procedure applied to real data.

For the first step the LEPTO 6.5 [7] event generator is used together with a leading order parametrisation of the unpolarised parton distributions. The MRST04LO set of parton distributions is used in a fixed-flavour scheme generation. This set of parton distributions has a good description of F_2 in the COMPASS kinematic region.

NLO corrections are simulated partially by including gluon radiation in the initial and final states (parton shower – PS).

The fragmentation is based on the Lund string model [9] implemented in JETSET [10]. In this model the probability that a fraction z of the available energy will be carried by a newly created hadron is expressed by the Lund symmetric function $f(z) = z^{-1}(1-z)^a e^{-bm_\perp^2/z}$, with $m_\perp^2 = m^2 + p_\perp^2$, where m is the hadron mass.

To improve the agreement between MC and data, the parameters (a,b) in the fragmentation function are modified from their default values (0.3,0.58) to (0.6, 0.1).

The transverse momentum of the hadrons, p_T , at the fragmentation level is given by the sum of the p_T of each hadron quarks. Then the p_T of the newly created hadrons is described by three steering parameters JETSET parameters: PARJ(21), PARJ(23) and PARJ(24). The default values of these three parameters are (0.36, 0.01, 2.0), and were modified to (0.30, 0.02, 3.5).

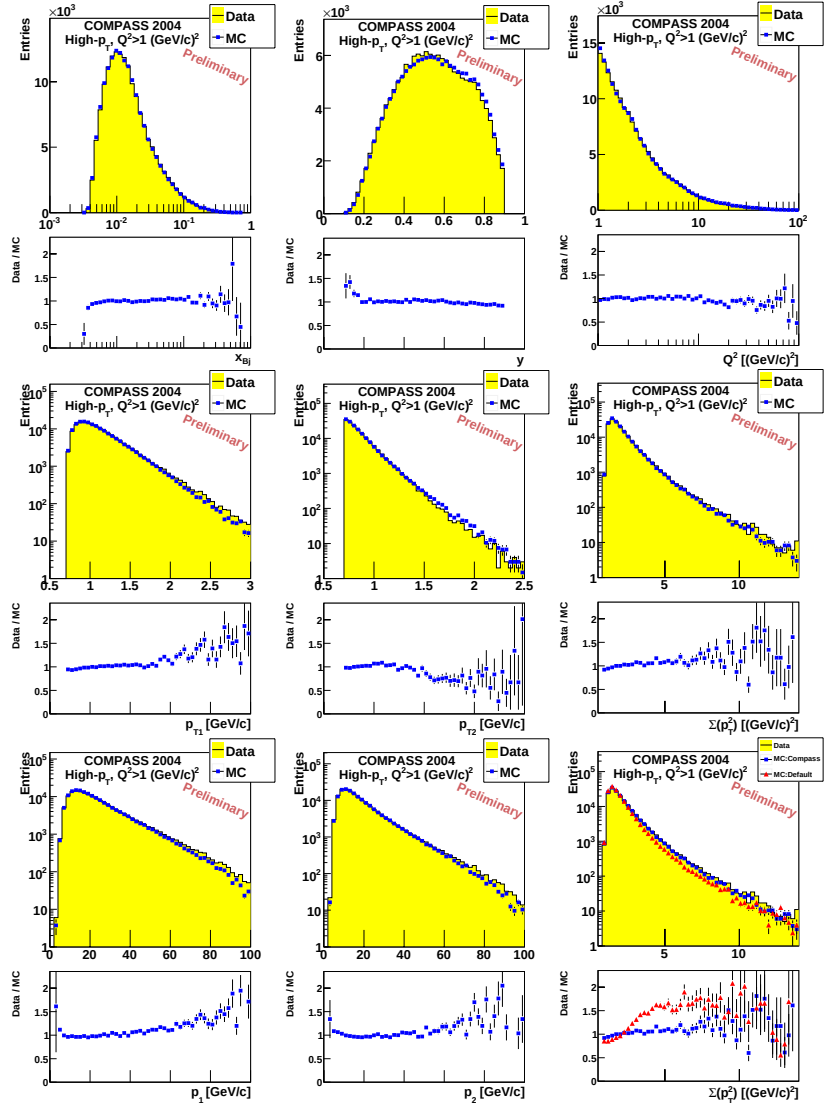


Figure 2: Comparison between data and MC simulations – The distributions and ratios of Data/MC for: inclusive variables: x_{Bj} , Q^2 , y (1st row). For hadrons p_T (2nd row). For hadron momenta (3rd row), and also the comparison of MC with LEPTO default tuning.

The remarkable agreement of the MC simulation with the data is illustrated in Fig. 2; this figure shows the data–MC comparison of the kinematic variables: x_{Bj} , y and Q^2 (1st row), the hadronic variables, p_T for the leading and sub-leading hadrons, together with the sum of p_T^2 , i.e. $\sum p_{T1}^2 + p_{T2}^2$ (2nd row), and the momentum p of those hadrons (3rd row), also two comparisons of the $\sum p_T^2$ variable one using the COMPASS tuning and another using the default LEPTO tuning. In this example, it is evident that the COMPASS tuning describes better our data sample than the LEPTO default one.

5 The $\Delta G/G$ extraction method

In the original idea of the high p_T analysis, the selection was based on a very tight set of cuts to suppress LO and QCD Compton. This situation results in a dramatic loss of statistics. A new approach was found, in which a loose set of cuts applied, combined with the use of a neural network [11] to assign a probability to each event to be originated from each of the three processes. The main goal of this method is to enhance the PGF process in the events sample, which accounts for the gluon contribution to the nucleon spin.

The neural network is trained using MC samples. In this way the neural network is able to learn about the three processes in order to be disentangled. A parametrisation of the variables R_i , x^i and a_{LL}^i for each process type are estimated by the neural network using as input the kinematic variables: x_{Bj} and Q^2 , and the hadronic variables: p_{T1} , p_{T2} , p_{L1} , and p_{L2} .

As the fractions of the three processes sum up to unity, we need two variables to parameterise them: o_1 and o_2 . The relations between the two neural network outputs o_1 and o_2 and the fraction are $R_{PGF} = 1 - o_1 - 1/\sqrt{3} \cdot o_2$, $R_{QCDC} = o_1 - 1/\sqrt{3} \cdot o_2$ and $R_{LO} = 2/\sqrt{3} \cdot o_2$.

A statistical weight is constructed for each event based on these probabilities. In this way we do not need to remove events that most likely do not came from PGF processe, because the weight will naturally reduce their contribution in the gluon polarisation, thus enhancing the sample of events that have a PGF likelihood.

The resulting neural network outputs for the fractions are presented in Fig. 3 in a 2-dimensional plot. The triangle limits the region where all fractions are positive. For the inclusive sample the average value of o_2 is quite large, which means that the LO process is the dominant one. The situation is different for the high p_T sample, in which the average outputs are $\langle o_1 \rangle \approx 0.5$ and $\langle o_2 \rangle \approx 0.35$. Note also that the spread along o_2 is larger than along o_1 .

This means that the neural network is able to select a region where the contribution of PGF and QCDC is significant compared to LO, although it can not easily distinguish between the PGF and QCDC processes themselves.

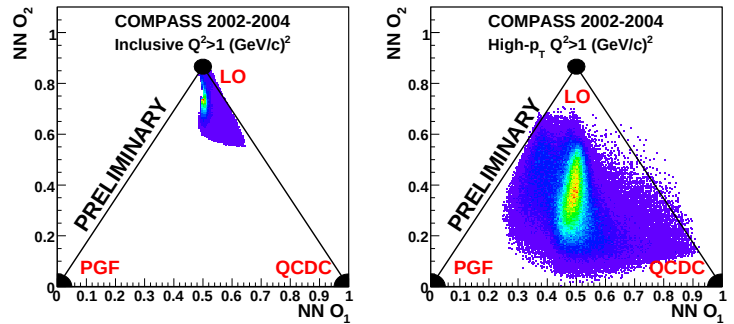


Figure 3: 2-d output of neural network for estimation that the given event is PGF, QCDC or LO; (left) for the inclusive sample and (right) for high p_T sample.

6 High p_T hadron pair analysis for low Q^2 region

The reason for splitting the Q^2 range in two complementary regions is that for the low Q^2 region the resolved photon contributions are considerably higher ($\approx 50\%$) than in the high Q^2 region, which contains practically only the three processes previously mentioned. This means that the QCD hard scale is also different for both Q^2 regions: for low Q^2 the scale is given by the high p_T hadrons, while for the high Q^2 is given by the Q^2 value itself.

A more complicated description of the physics than pure QCD in lowest order needs to be included in the MC simulation for this case. Therefore the event generator used in this analysis is PYTHIA 6.2 [12] which covers the physical processes for quasi-real photoproduction.

In this analysis the selection is essentially the same as in high Q^2 region plus a slightly strict set of cuts: $x_F > 0.1$, $z > 0.1$, and $\sum p_T^2 > 2.5 \text{ (GeV/c)}^2$. The data sample in this region is 90 % of the whole data for all Q^2 range. The weighting method used in the high Q^2 analysis is not applied in this case.

The MC simulated and real data samples of high p_T events are compared in Fig. 4 for Q^2 , y (1st row), and for the total and transverse momenta of the high p_T hadron (2nd row), showing a good agreement.

The gluon polarisation in the low Q^2 region is extracted using averaged values as shown by this expression:

$$\left\langle \frac{A_{LL}}{D} \right\rangle = R_{\text{PGF}} \left\langle \frac{\hat{a}_{LL}^{\text{PGF}}}{D} \right\rangle \left\langle \frac{\Delta G}{G} \right\rangle + R_{\text{QCDC}} \left\langle \frac{\hat{a}_{LL}^{\text{QCDC}}}{D} A_1 \right\rangle + \sum_{f, f^\gamma} R_{ff^\gamma} \left\langle \hat{a}_{LL}^{ff^\gamma} \frac{\Delta f}{f} \frac{\Delta f^\gamma}{f^\gamma} \right\rangle \quad (6)$$

R_{ff^γ} is the fraction of events in the high p_T sample for which a parton f from the nucleon interacts with a parton f^γ from a resolved photon; A_1 is the virtual photon deuteron asymmetry measured in an inclusive sample; $\Delta f/f$ ($\Delta f^\gamma/f^\gamma$) is the polarisation of quarks or gluons in the deuteron (photon).

This analysis was performed using a data sample from the years 2002 to 2004. For more details about this analysis the reader is invited to look into ref. [13].

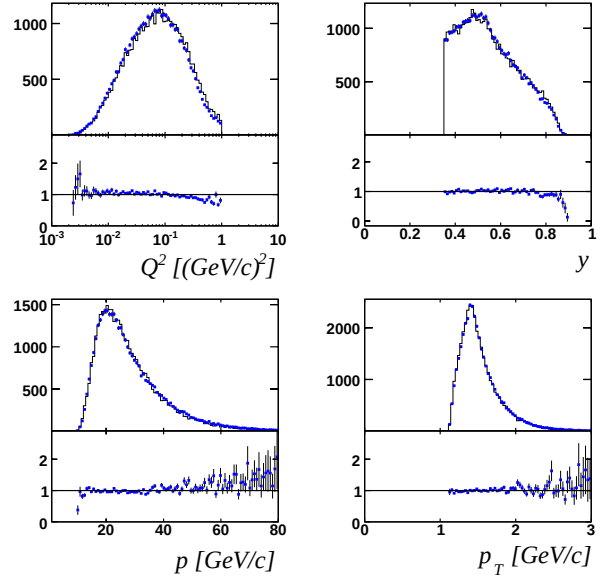


Figure 4: Comparison between data and MC simulations – The distributions and ratios of Data/MC for: kinematic variables: Q^2 , y (1st row). Total and transverse momentum p and p_T of the high p_T hadron (2nd row).

7 Results

The preliminary measurements of the gluon polarisation in low and high Q^2 regions, using data from the years 2002 to 2004, are:

$$\begin{aligned} (\Delta G/G)_{\text{low } Q^2} &= 0.02 \pm 0.06_{(\text{stat.})} \pm 0.06_{(\text{syst.})} & \text{with } x_G &= 0.09^{+0.07}_{-0.04} \\ (\Delta G/G)_{\text{high } Q^2} &= 0.08 \pm 0.10_{(\text{stat.})} \pm 0.05_{(\text{syst.})} & \text{with } x_G &= 0.08^{+0.04}_{-0.03} \end{aligned}$$

The average of the hard scale, μ^2 , for low and high Q^2 is about 3 (GeV/c)². x_G is the momentum fraction carried by the probed gluons obtained from the MC parton kinematics. The result of the measurement for low Q^2 using data from 2002 and 2003 can be found in [13].

Fig. 5 shows these new values of $\Delta G/G$ together with the preliminary value from the open charm analysis. Also the figure shows the measurements from SMC collaboration, from the high p_T analysis for the $Q^2 > 1$ (GeV/c)² region [14] and also the measurements from HERMES collaboration, for single hadrons and high p_T hadron pairs analyses [15]. The curves in the figure are the parametrisation of $\Delta G/G(x)$ using a NLO QCD analysis done by COMPASS [2] in the $\overline{MS}$ scheme with a renormalisation scale $\langle\mu^2\rangle = 3$ (GeV/c)². The dashed line curve is the QCD fit assuming that $\Delta G > 0$, the dotted line is the QCD fit assuming $\Delta G < 0$. It is seen that both results from high p_T analyses, for high and low Q^2 regions, are compatible with each other and also, within their x_G region, in agreement with the NLO QCD fits.

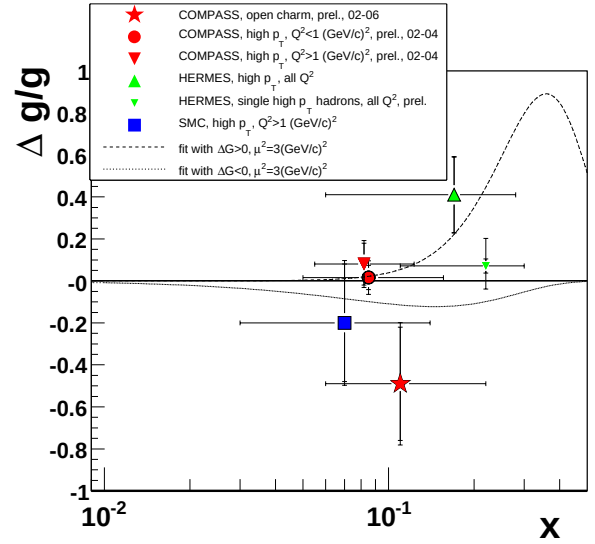


Figure 5: Comparison of $\Delta G/G$ measurements from COMPASS [2], SMC [14], and HERMES [15]. The two curves correspond to the parametrisation from the NLO QCD analysis in the $\overline{MS}$ scheme with scale at 3 (GeV/c)².

8 Conclusions

The preliminary values of the gluon polarisation for low and high Q^2 regions were presented. The gluons were probed at an average scale $\langle\mu^2\rangle \approx 3$ (GeV/c)². Both measurements show that the gluon contribution to the nucleon spin for $x_G \approx 0.1$ is compatible with zero. In that region of x_G the presented measurements are in agreement with other well known results.

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STATUS OF THE PAX EXPERIMENT

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Abstract

The PAX experiment aims at the production of a spin-polarized antiproton beam by means of Spin-Filtering, and the study of the spin-dependent pp cross sections. The PAX collaboration is preparing for detailed investigations of the spin-filtering process of protons at COSY (FZ Jülich) and of antiprotons at the AD (CERN). The physics program with polarized antiprotons includes the study of the Drell-Yan process as a direct measurement of transversity, a leading twist structure function of the nucleon. In the talk, a recent measurement of the depolarization of an initially polarized proton beam by co-moving electrons at low relative velocity is presented. Implications for the role of electrons in the polarization buildup of a stored proton beam are discussed.

1 The PAX experiment

There is consensus that the QCD physics potential of experiments with high energy polarized antiprotons would be enormous, but up to now high-luminosity experiments were not possible. The situation would change dramatically with the production of stored polarized antiproton beams, and the subsequent realization of a double-polarized high-luminosity antiproton-proton collider. The list of fundamental physics issues to be addressed with such a collider includes the measurement of transversity, the quark transverse polarization inside a transversely polarized proton, which constitutes the last missing leading twist piece of the QCD description of the partonic structure of the nucleon. The transversity can be directly accessed only via double-polarized antiproton-proton Drell-Yan production. It should be noted, that without a measurement of the transversity, our knowledge of the the spin structure of the proton will remain incomplete. Other items of interest are the measurement of the phases of the timelike form factors of the proton, and double-polarized hard antiproton-proton scattering.

The PAX collaboration (**P**olarized **A**ntiproton **eX**periments) has formulated its ambitious physics program [1]. A Technical Proposal [2] has been submitted for the new Facility for Antiproton and Ion Research (FAIR) to be built at GSI in Darmstadt, Germany. The uniqueness and the strong scientific merits that would become available with the advent of stored beams of polarized antiprotons have been well received [3], and there is now an urgency to convincingly demonstrate experimentally that a high degree of antiproton beam polarization can be reached. The suggested collider aims at luminosities in excess of $10^{31} \text{ cm}^{-2}\text{s}^{-1}$. An integral part of such a machine is a dedicated large-acceptance Antiproton Polarizer Ring (APR) [4], for which a basic design has been developed [5].

2 Towards polarized Antiprotons

For more than two decades, physicists have tried to produce beams of polarized antiprotons [6], generally without success. Conventional methods like atomic beam sources, appropriate for the production of polarized protons and heavy ions cannot be applied, since antiprotons annihilate with matter. Polarized antiprotons have been produced from the decay in flight of $\bar{\Lambda}$ hyperons at Fermilab. The intensities achieved with antiproton polarizations $P > 0.35$ never exceeded $1.5 \cdot 10^5 \text{ s}^{-1}$ [7]. Scattering of antiprotons off a liquid hydrogen target could yield polarizations of $P \approx 0.2$, with beam intensities of up to $2 \cdot 10^3 \text{ s}^{-1}$ [8]. Small angle $\bar{p}C$ scattering has been studied at LEAR, but the observed antiproton polarizations are negligibly small [9,10]. Unfortunately, all the above mentioned approaches do not allow efficient accumulation in a storage ring, which would greatly enhance the luminosity. Spin splitting using the Stern–Gerlach separation of the given magnetic substates in a stored antiproton beam was proposed in 1985 [11]. Although the theoretical understanding has much improved since then [12], spin splitting using a stored beam has yet to be observed experimentally. In contrast to that, a convincing proof of the spin–filtering principle was provided by the FILTEX experiment at the TSR–ring in Heidelberg [13].

3 Do unpolarized electrons affect the polarization of a stored beam?

The original idea of PAX to use polarized electrons to produce a polarized beam of antiprotons [4] has triggered further theoretical work on the subject, which led to a new suggestion by a group from Mainz [14,15] to use co-moving electrons (or positrons) at slightly different velocities than the orbiting protons (or antiprotons) as a means to polarize the stored beam. The cross section for $e\vec{p}$ spin–flip predicted by the Mainz group in a *numerical* calculation is as large as about $2 \cdot 10^{13}$ barn, if the relative velocities between proton (antiproton) and electron (positron) are adjusted to $v/c \approx 0.002$. At the same time, *analytical* predictions for the same quantity by a group from Novosibirsk [16] range well below a mbarn. Thus prior to the experiment, the two theoretical estimates differed by about 16 orders of magnitude. It should be also emphasized, that the practical use of the $\vec{e}p$ or $e^+\vec{p}$ processes to polarize anything is excluded if the spin–flip cross sections are smaller than about 10^7 barn.

In order to provide an experimental answer for this puzzle, the ANKE and PAX collaborations joined forces at COSY and mounted an experiment, where for the first time, electrons in the electron cooler have been used as a target, and in which the effect of electrons on the polarization of a 49.3 MeV proton beam orbiting in COSY was determined. Instead of studying the buildup of polarization in an initially unpolarized beam, here the inverse situation was investigated by observation of the depolarization of an initially po-

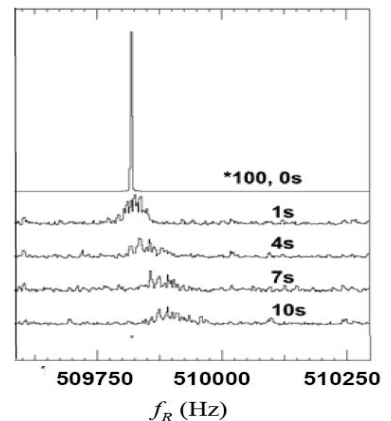


Figure 1: Shown are Schottky spectra of the stored electron-cooled protons at different delay times at which the cooler electrons were detuned. The peaks correspond to the revolution frequency of the protons, from which their velocity change can be derived.

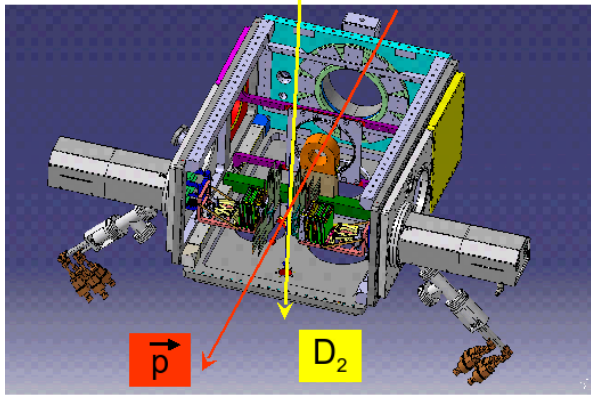


Figure 2: Using events from elastic $\bar{p}d$ scattering, two Silicon detector telescopes mounted near the deuterium beam (D_2) of the ANKE cluster jet target measure the change of the proton ($\bar{p}$) beam polarization induced by the depolarizing $e\bar{p}$ spin-flips in the COSY electron cooler.

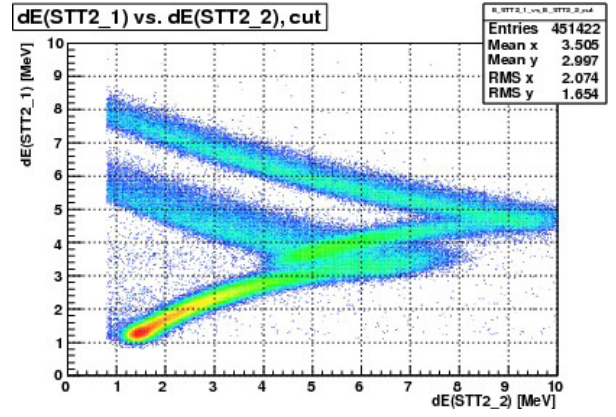


Figure 3: Identification of pd elastic scattering in the detector system.

larized beam. A key question was how fast the stored protons follow the velocity change of the cooler electrons. The result of a test measurement is shown in Fig. 1. It can be concluded that for periods of about 5s the change of the proton velocity is negligible, which guarantees a well defined relative velocity during the measurement cycles [17].

The proton beam polarization has been measured by making use of the analyzing power of $\bar{p}d$ elastic scattering on a deuterium cluster jet target. The experimental setup of the polarimeter consisting of two telescopes, each of them containing two $300\ \mu\text{m}$ thick layers of Silicon, is depicted in Fig. 2.

The energy deposit in the first detector layer is plotted as a function of the energy deposit in the second layer in Fig. 3. The upper band corresponds to deuterons and the lower one to protons from pd elastic scattering. The depolarizing cross section is determined from the ratio of the measured beam polarizations P_{detuned} and P_{tuned} (Fig. 4). These polarizations correspond to well-defined changes of the electron velocity with respect to the protons, which were achieved by detuning the accelerating voltage in the electron cooler by a specific amount.

The depolarizing cross section is plotted in Fig. 5 as a function of the magnitude of the electron velocity in the proton rest frame. The measurement shows that the predictions by the Mainz group were too large by at least six orders of magnitude. It should be noted that very recently, the Mainz group has submitted two errata to their original publications stating that because of numerical problems in the calculation their theoretical estimates were too large by about 15 orders of magnitude.

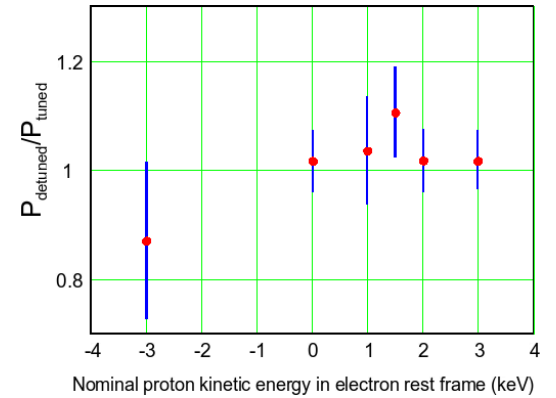


Figure 4: Ratio of the measured polarizations $\frac{P_{\text{detuned}}}{P_{\text{tuned}}}$ as function of the proton kinetic energy in the electron rest frame.

4 Spin–Filtering Experiments at COSY and AD

at COSY, shedding light on the ep spin–flip cross sections when the target electrons are unpolarized. The experimental finding rules out the practical use of polarized leptons to polarize a beam of antiprotons with present–day technologies. This leaves us with the only proven method to polarize a stored beam in situ, namely *spin filtering* by the strong interaction. At present, we are lacking a complete quantitative understanding of all underlying processes, therefore the PAX collaboration is aiming at high–precision polarization buildup studies with transverse and longitudinal polarization using stored protons in COSY [18].

The buildup process itself can be studied in detail, because in this situation the spin–dependence of the pp interaction is completely known. The polarized internal target required for these investigations was previously used at the HERMES experiment at HERA/DESY. Including the target polarimeter, it has meanwhile been relocated to Jülich [19], where it is presently set up to be installed together with a large–acceptance detector system for the determination of target and beam polarizations in a dedicated low– β section at COSY.

In contrast to the pp system, the experimental basis for predicting the polarization buildup in a stored antiproton beam by spin filtering is practically non–existent. Therefore, it is of high priority to perform, subsequently to the COSY experiments, a series of dedicated spin–filtering experiments using stored antiprotons. The AD–ring at CERN is a unique facility at which stored antiprotons in the appropriate energy range are available with characteristics that meet the requirements for the first–ever antiproton polarization buildup studies. At the same time, the equipment required for the dedicated spin–filtering experiments at the AD, i.e. the polarized internal target and the new low– β section, an efficient polarimeter to determine target and beam polarizations, and a Siberian snake to maintain the longitudinal beam polarization, must be extensively commissioned and tested at COSY, prior to the experimental investigations at CERN.

Already in 2005, the PAX collaboration suggested in a Letter–of–Intent [20] to the SPS committee of CERN to study the polarization buildup via spin filtering of stored antiprotons by multiple passage through a polarized internal hydrogen gas target, because only through these investigations, one can obtain direct access to the spin dependence of the total $\bar{p}p$ cross sections. Apart from the obvious interest for the general theory of $\bar{p}p$ interactions, the knowledge of these cross sections is necessary for the interpretation of unexpected features of the $\bar{p}p$, and other antibaryon–baryon pairs, contained in final states in J/Ψ and B –decays. Once these experiments have provided an experimental data base, the design of a dedicated APR can be targeted, as the next major milestone.

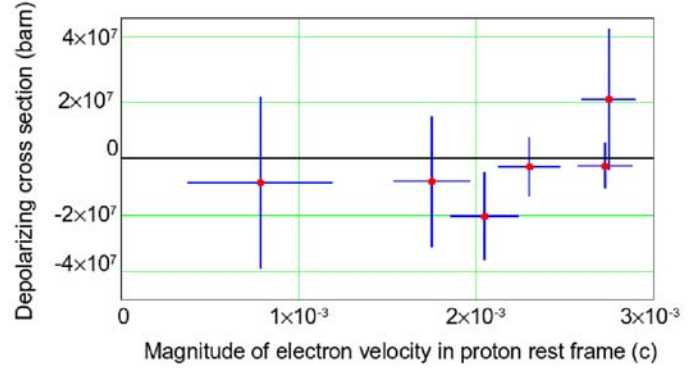


Figure 5: The depolarizing cross section is plotted as a function of the magnitude of the electron velocity in the proton rest frame, indicating an upper limit of a few 10^7 b. Prior to the experiment, the two theoretical estimates of this cross section differed by about 16 orders of magnitude at a relative velocity of $v/c = 2 \cdot 10^{-3}$ [15,16]. The experiment rules out the higher estimate.

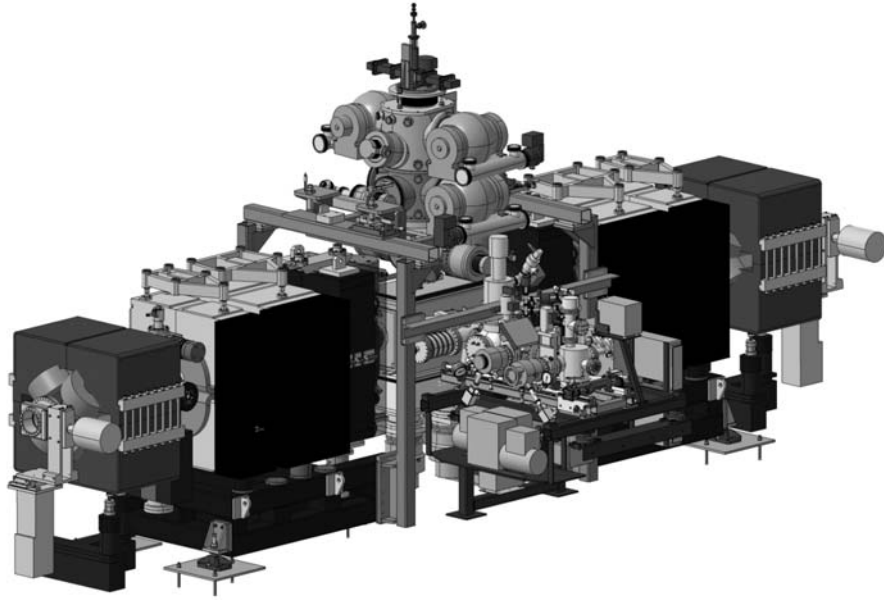


Figure 6: Target section required for the CERN test with low-beta quadrupoles, target chamber and target cell (center, not visible), source of polarized atoms (top) feeding the target cell, and target polarimeter (center).



Figure 7: Preparatory work for the installation of low- β quadrupoles at the PAX target point in COSY, as of July 2009

The development of the optimum spin filtering conditions requires a parallel interlaced program of studies at the cooler synchrotron COSY (FZ Jülich) and at the antiproton decelerator AD (CERN). At COSY, the target insertion (see Fig.6) will be tested and commissioned, and precision studies on pp spin filtering will be performed which serve

to test the theoretical models of the polarizing process. The time schedule includes four years at COSY for the various steps of installation and experiment, and in parallel a period of four years at the AD starting with a delay of one year. The final step at the AD will be $p\bar{p}$ spin filtering in the energy range up to $T = 450\text{MeV}$ with longitudinal spins at the $\vec{H}$ -target section. As a first step (see Fig. 7), a set of four low- β quadrupoles is being installed at COSY in order to test the machine performance with a modified lattice.

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ASYMMETRY MEASUREMENTS IN THE ELASTIC PION-PROTON SCATTERING IN THE RESONANCE ENERGY RANGE

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Abstract

The asymmetry parameter P was measured for the elastic pion-proton scattering in the very backward angular region of $\theta_{cm} \approx 150 - 170^\circ$ at several pion beam energies in the invariant mass range containing most of the pion-proton resonances. The general goal of the experimental program was to provide new data for partial wave analysis in order to allow the unambiguous light baryon spectrum reconstructions. The experiment was performed at the ITEP U-10 proton synchrotron by the ITEP-PNPI collaboration in the latest 5 years.

Current situation in light baryon spectroscopy is far from perfect both from theoretical and experimental points of view. Most data on light baryons is taken from the partial wave analyses of πN scattering KH80 [1] and CMB80 [2], both performed about 30 years ago. Later solutions by GWU group [3, 4] do not reveal several resonant states seen by [1, 2], though sufficient improvements were made to their analysis in the recent years. This work was aimed at providing new experimental data to the partial wave analyses in kinematic areas, where their behavior is most unstable and where they are most suspicious for various kinds of ambiguities.

Polarization parameter P in the elastic scattering is determined by the direct measurement of the azimuthal asymmetry of the reaction produced by pions on a proton target polarized normally to the scattering plane.

The experimental setup SPIN-P02 is schematically shown in fig. 1. The polarized target with vertical orientation of the polarization vector $\vec{P}_T$ was located at the focus of the secondary pion beam line of the ITEP proton synchrotron. The incident and scattered pions and the recoil proton were tracked with blocks of multi-wire chambers CH1-CH14 and the full event reconstruction was performed allowing good selection of the elastic events. In case of positive beam TOF technique was used to separate pions from protons at low beam energies, while at beam momenta above 1.8 GeV/c the aerogel Cherenkov counter AGCC was additionally introduced into the beam for the pion tagging. Further details on the SPIN-P02 setup could be found in [5]. The asymmetry is determined by measuring the normalized event counts for two opposite directions of the target polarization vector.

The procedure of the elastic event selection is illustrated by fig. 2. For each event the deviation from the elastic kinematics was calculated in terms of two variables: $\Delta\theta$

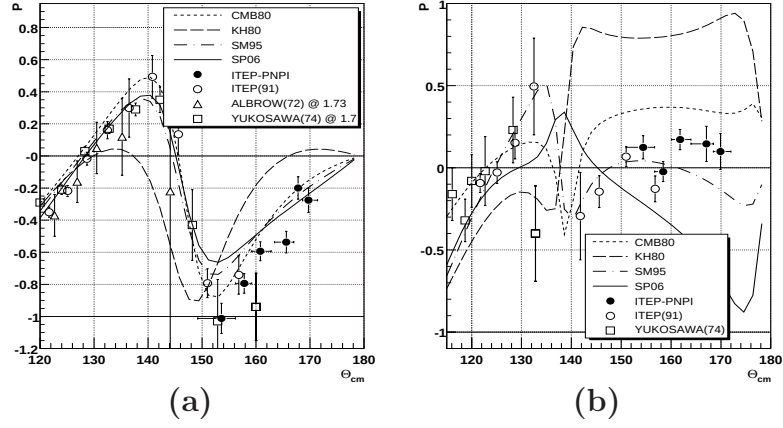


Figure 3: Asymmetry in π^-p elastic scattering at 1.78 GeV/c (a) and 2.07 GeV/c (b).

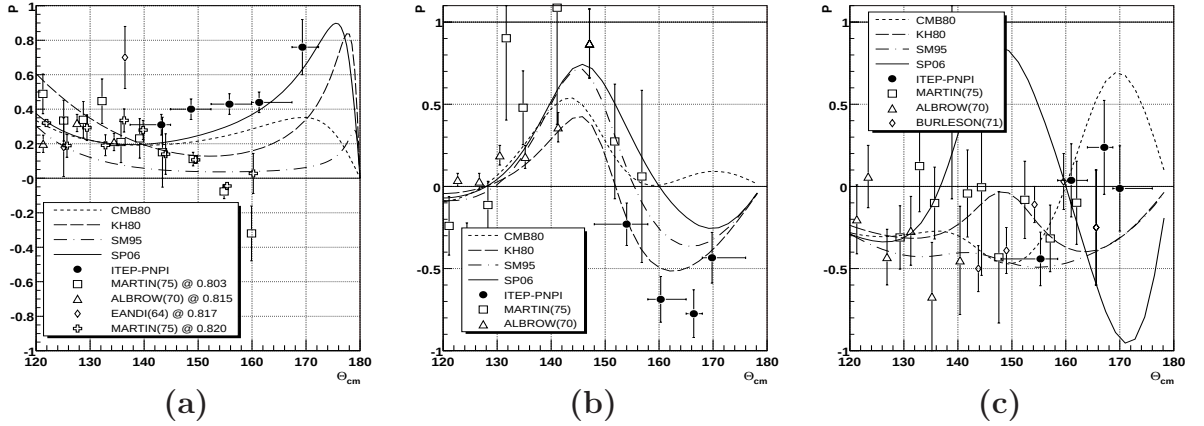


Figure 4: Asymmetry in π^+p elastic scattering at 0.80 GeV/c (a), 1.94 GeV/c (b) and 2.07 GeV/c (c).

on the actual cross section values the background conditions are also very much different (compare figures 2a and 2b) and the ratio of the background to the elastic signal lays in the range $Bg/S = 0.15 - 1.2$ in various kinematic areas of the experiment.

Figures 3,4 present the results of the measurements (closed circles) as a function of the c.m. scattering angle. The errors are only statistical and account for the elastic event numbers, the background error matrix and the intensity normalization uncertainty. The overall scale error due to the target polarization measurement is below 3% (typically 1.5%). Continuous lines show the PWA predictions [1]- [4]. Previous experimental data at smaller angles and in the overlapping regions are also shown with open symbols.

The points at 1.78 GeV/c in π^-p (fig. 3a) were taken to check the setup. Partial wave analysis do not show large variations from each other. The data show best agreement with CMB80, while the deviation of KH80 from the data points is obviously statistically significant. The point at 153.6° coincides with earlier data and supports even deeper minimum than any of the analysis show.

The data at 2.07 GeV/c with negative pions (fig. 3b) are not exactly followed by any of the PWA solutions. Yet the closest two are SM95 and CMB80, while SP06 and KH80 do

not even resemble the data behavior. Good agreement with earlier ITEP results [6] in the overlapping region proofs the data quality. A narrow local minimum at 157° confirmed by two independent measurements indicates the significant presence of a partial wave with high orbital momentum $L \approx 8 - 9$. The sharp step at 167° in 1.78 GeV/c data may have similar interpretation.

In both cases with π^-p scattering the latest solution SP06 of the GWU group is not the closest to the new data. In the lower energy domain (see fig. 4a for π^+p at 0.80 GeV/c) the data are best described by this very solution. CMB80 does not show the sharp peak at 175° implied by the data, while SM95 gives much smaller asymmetry values. Our results obviously contradict to the 3 rightmost points from [7] at two adjacent energies.

π^+p asymmetry at 1.94 GeV/c shows large negative values around 165° (fig. 4b). Neither of the solutions manifest so deep a minimum though all but CMB80 have qualitatively similar behavior. The closest prediction in this case is from KH80.

The cross section for backward angles at 2.07 GeV/c is extremely low for positive pions. Our data are not much better statistically than that from previous works and feature high background levels. The angular dependence of the data most resemble the curve from CMB80 (fig. 4c). All other solutions show qualitatively different behavior though KH80 and SM95 are not beyond 3σ boundary of the data.

The obtained results show that in some kinematic areas one or both of the "classic" partial wave analysis CMB80 and KH80 are in disagreement with the new data. In some cases even the qualitative behavior of mentioned PWA does not correspond to that of the data, which may indicate the wrong choice of the solution branch by these analysis and, consequently, wrong extraction of baryon properties. The latest solution SP06 of GWU group seems to be consistent with the data in the lower energy domain, while in the momentum region above 1.8 GeV/c it's behavior looks unstable.

The ITEP-PNPI team believes that their new data notably improves the experimental database for partial wave analysis and becomes a significant step towards reliable light baryon spectrum.

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CHARGE ASYMMETRY AND SYMMETRY PROPERTIES

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Abstract

Applying general symmetry properties of electromagnetic interaction, information from electron proton elastic scattering data can be related to charge asymmetry in the annihilation channels $e^+ + e^- \leftrightarrow \bar{p} + p$ and to the ratio of the cross section of elastic electron and positron scattering on the proton. A compared analysis of the existing data allows to draw conclusions on the reaction mechanism.

Elastic and inelastic electron scattering has been considered the most direct way to learn about the internal structure of hadrons. If one assumes that the underlying mechanism is the exchange of one virtual photon of mass $q^2 = -Q^2$ (OPE), an elegant and simple formalism allows to express the electromagnetic current of a spin S hadron in terms of $2S + 1$ electromagnetic form factors (FFs).

In this contribution, we discuss the proton structure and the reaction mechanism for annihilation and scattering reactions in the energy range of a few GeV. Unpolarized electron proton elastic scattering has been considered the simplest way to determine FFs, using the Rosenbluth separation: measurements at fixed Q^2 , for different angles (which requires to change the beam energy and the spectrometer setting) allow to extract the electric G_E and the magnetic G_M proton form factors, through the slope and the intercept of the Rosenbluth plot.

In recent years it has been possible to measure the polarization of the outgoing proton, scattered by a longitudinally polarized beam. The ratio of the transverse to longitudinal polarization of the scattered proton P_T/P_L allows to access the ratio of the electric to magnetic form factor, not only their squared values as the cross section does. This method, suggested by A.I. Akhiezer and M.P. Rekalo [1], is more sensitive to a small contribution of the electric FF, especially at large values of Q^2 , where the magnetic contribution is dominant. The surprising result, which was obtained by the GEp collaboration [2], gave rise to a huge number of theoretical and phenomenological papers, and to a large experimental activity. Polarization experiments show that over $Q^2 = 1$ (GeV/c)² the Q^2 dipole approximation of FFs does not hold anymore and the FFs ratio follows a straight line: $R = \mu G_E/G_M = 1.059 - 0.143 Q^2$ [(GeV/c)²] at least up to a $Q^2 \sim 6$ (GeV/c)².

As no bias has been found in the experiments, the reason for the discrepancy is possibly related to radiative corrections (RC). In unpolarized experiments, RC can reach 40% and high order radiative corrections are typically not included in the data analysis. No RC are applied to the polarization ratio, as they are assumed to be negligible, and indeed, first order RC cancel. The magnetic FF can be considered well known from cross section measurements, so the common interpretation of the present results is that the electric distribution in the proton is different from what previously assumed. Being G_E related

to the slope of the Rosenbluth plot, it has been noticed [3] that radiative corrections are largely responsible for this slope, in particular the C-odd corrections which are due to bremsstrahlung, more exactly, to the interference between electron and proton soft photon emission and two photon exchange (TPE).

TPE can not be calculated in model independent way when the target is a proton. Exact calculations for a lepton target, in a pure QED framework, have been done and show that TPE can not exceed 1-2% [4]. Moreover, the lepton case is, by definition, an upper limit of the case where the target is a proton and the intermediate state in the box diagram is also a proton [5]. TPE calculations for a proton target require modelization and different calculations have been done, with no quantitative agreement (see [6] and refs therein). Let us mention that in [7] it has been argued that elastic and inelastic intermediate states in the box diagram essentially cancel, due to analytical properties of the reaction amplitude.

Standard radiative corrections take into account the contribution of TPE where most of the transferred momentum is carried by one photon, while the other photon has very small momentum. The TPE contribution is larger when the two photons share equal momentum, and, due to the steep decrease of FFs, it might compensate the α counting rule. If such mechanism is important, it should be more visible at larger values of momentum transfer and for hadrons heavier than proton.

Crossing symmetry, which holds at the level of Born diagram, allows to relate the matrix elements $\mathcal{M}$ of the crossed processes, through the amplitude $f(s, t)$:

$$\overline{|\mathcal{M}(eh \rightarrow eh)|^2} = f(s, t) = \overline{|\mathcal{M}(e^+e^- \rightarrow \bar{h}h)|^2}. \quad (1)$$

The line over $\mathcal{M}$ denotes the sum over the polarizations of all particles (in initial and final states), s and t are the Mandelstam variables, which span different kinematical regions for annihilation and scattering channels.

The presence of a single virtual photon in the reaction $e^+e^- \rightarrow \gamma^* \rightarrow \bar{h}h$ constrains the total angular momentum $\mathcal{J}$ and the P -parity for the $\bar{h}h$ -system, to take only one possible value, $\mathcal{J}^P = 1^-$, the quantum number of the photon. Therefore, in the framework of one-photon exchange (OPE), the $\cos\theta$ -dependence of $\overline{|\mathcal{M}(eh \rightarrow eh)|^2}$ can be predicted in a general form (θ is the CMS angle between the momenta of the electron and the detected antinucleon):

$$\overline{|\mathcal{M}(eh \rightarrow eh)|^2} = a(t) + b(t) \cos^2 \theta, \quad (2)$$

where $a(t)$ and $b(t)$ are definite quadratic combinations of the electromagnetic form factors for the hadron h . The C-invariance of the electromagnetic hadron interaction allows only even powers of $\cos\theta$ and the degree of the $\cos\theta$ -polynomial is limited to the second order, due to the spin of the virtual photon. One can show [8] that there is a one-to-one correspondence between $\cos^2\theta$ and $\cot^2(\theta_e/2)$ (θ_e is the LAB angle of the emitted electron in the scattering channel) which explains the origin of the linear $\cot^2 \frac{\theta_e}{2}$ -dependence of the differential cross section for any eh -process in OPE approximation.

The presence of TPE in the intermediate state: $e^+ + e^- \rightarrow 2\gamma \rightarrow \bar{h} + h$ can induce any value of the total angular momentum and space parity in the annihilation channel, but the $\bar{h}h$ -system, produced through OPE and TPE mechanisms has different values of C-parity, because $C(\gamma) = -1$ and $C(2\gamma) = +1$. Therefore the interference contribution to the differential cross section in the annihilation channel must be an **odd** function of $\cos\theta$.

These model independent statements allow to sign the presence of TPE: non linearities in the Rosenbluth plot in the scattering channel, and odd $\cos\theta$ contributions in the differential cross section for the annihilation channel. The illustration of different sets of data is given below, for $e^- + {}^4\text{He}$ elastic scattering, $e^+ + e^- \rightarrow \bar{p} + p$ annihilation and $e^\pm + p$ elastic scattering.

1 Scattering channel

The search of model independent evidence of TPE in the experimental data which should appear as a non linearity of the Rosenbluth fit was firstly done in case of deuteron in Ref. [8] and in case of proton in Ref. [3]. No evidence was found, in the limit of the precision of the data.

Let us note, that these experiments are sensitive to the real part of the interference between OPE and TPE. Very precise measurements of the transverse beam spin asymmetry in elastic electron scattering on proton and ${}^4\text{He}$ suggest a non zero imaginary part of the TPE amplitude [9, 10]. Of particular interest is the case of ${}^4\text{He}$ target:

$$e^-(p_1) + {}^4\text{He}(q_1) \rightarrow e^-(p_2) + {}^4\text{He}(q_2), \quad (3)$$

as the spin structure of the matrix element is highly simplified for a spinless target. Using the general properties of the electron-hadron interaction, such as the Lorentz invariance and P-invariance, taking into account the identity of the initial and final states and the T-invariance of the strong interaction, the scattering of a spin 1/2 particle on a spin zero target is described by two independent amplitudes and the general form of the matrix element can be written independently from the reaction mechanism, as :

$$\mathcal{M}(s, q^2) = \frac{e^2}{Q^2} \bar{u}(p_2) \left[m F_1(s, q^2) + F_2(s, q^2) \hat{P} \right] u(p_1) \varphi(q_1) \varphi(q_2)^* = \frac{e^2}{Q^2} \mathcal{N}, \quad (4)$$

where $\varphi(q_1)$ and $\varphi(q_2)$ are the wave functions of the initial and final helium, with $P = q_1 + q_2$ and $u(p_1)$, $u(p_2)$ are the spinors of the initial and final electrons, respectively. Here F_1 and F_2 are two invariant amplitudes, which are, generally, complex functions of two variables $s = (q_1 + p_1)^2$ and $q^2 = (q_2 - q_1)^2 = -Q^2$ and m is the electron mass. The matrix element (4) contains the helicity-flip amplitude F_1 proportional to the electron mass which is explicitly singled out. The single-spin asymmetry, of interest here, is proportional to F_1 . In OPE approximation one has:

$$F_1^{\text{Born}}(s, q^2) = 0, \quad F_2^{\text{Born}}(s, q^2) = F(q^2), \quad (5)$$

where $F(q^2)$ is the helium electromagnetic charge form factor depending only on q^2 , with normalization $F(0) = Z$, where Z is the helium charge.

To separate the effects due to the Born and the two-photon exchange contributions, let us define the following decompositions of the amplitude [11]

$$F_2(s, q^2) = F(q^2) + f(s, q^2) \text{ where } F_1(s, q^2) \sim \alpha, \quad f(s, q^2) \sim \alpha, \text{ and } F(q^2) \sim \alpha^0. \quad (6)$$

Since the terms F_1 and f are small in comparison with the dominant one, one can safely neglect the bilinear combinations of these small terms multiplied by the factor m^2 .

The differential cross section of the reaction (3), for the case of unpolarized particles, has the following form in the Born approximation

$$\frac{d\sigma_{un}^{Born}}{d\Omega} = \frac{\alpha^2 \cos^2 \frac{\theta}{2}}{4E^2 \sin^4 \frac{\theta}{2}} \left[1 + 2 \frac{E}{M} \sin^2 \frac{\theta}{2} \right]^{-1} F^2(q^2), \quad (7)$$

where θ is the electron scattering angle in Lab system and M is the helium mass.

In the Born approximation, the ${}^4He$ FF depends only on the momentum transfer squared, Q^2 . The presence of a sizable TPE contribution should appear as a deviation from a constant behavior of the reduced cross section measured at different angles and at the same Q^2 . In case of ${}^4He$ few data exist at the same $\bar{Q}^2$ value, for $Q^2 < 8 \text{ fm}^{-2}$ [12]. No deviation of these data from a constant value is seen, from a two parameter linear fit. The slope for each individual fit is always compatible with zero (see Fig. 1).

The TPE contribution leads to new terms in the differential cross section :

$$\frac{d\sigma_{un}}{d\Omega} = \frac{\alpha^2 \cos^2 \frac{\theta}{2}}{4E^2 \sin^4 \frac{\theta}{2}} \left[1 + 2 \frac{E}{M} \sin^2 \frac{\theta}{2} \right]^{-1} \left\{ F^2(q^2) + 2F(q^2) \text{Re } f(s, q^2) + |f(s, q^2)|^2 + \right. \quad (8)$$

$$\left. + \frac{m^2}{M^2} \left[\frac{M}{E} + \left(1 + \frac{M}{E} \right) \tan^2 \frac{\theta}{2} \right] F(q^2) \text{Re } F_1(s, q^2) \right\}, \quad (9)$$

and to a non-zero asymmetry, in the case of the elastic scattering of transversally polarized electron beam. Let us define a coordinate frame with $z \parallel p$, $y \parallel \vec{p} \times \vec{p}'$, where $\vec{p}(\vec{p}')$ is the initial (scattered) electron momentum, and the x axis directed to form a left-handed coordinate system. The transverse asymmetry, due to the interference between OPE and TPE, is determined by the polarization component perpendicular to the reaction plane:

$$A_y = \frac{\sigma^\uparrow - \sigma^\downarrow}{\sigma^\uparrow + \sigma^\downarrow} \sim \vec{s}_e \cdot \frac{\vec{p} \times \vec{p}'}{|\vec{p} \times \vec{p}'|} \equiv s_y, \quad (10)$$

where $\sigma^\uparrow(\sigma^\downarrow)$ is the cross section for electron beam polarized parallel (antiparallel) to the normal of the scattering plane and $\vec{s}_e$ is the spin vector of the electron beam. In terms of the amplitudes, it is expressed as:

$$A_y = 2 \frac{m}{M} \tan \frac{\theta}{2} \frac{\text{Im } F_1(s, q^2)}{F(q^2)}. \quad (11)$$

Being a T-odd quantity, it is completely determined by the TPE contribution through the the imaginary part of the spin-flip amplitude $F_1(s, q^2)$ and, therefore, it is proportional to the electron mass.

For elastic $e^- + {}^4He$ scattering, a value of $A_y^{exp}({}^4He) = -13.51 \pm 1.34(stat) \pm 0.37(syst)$ ppm for $E = 2.75 \text{ GeV}$, $\theta = 6^\circ$, and $Q^2 = 0.077 \text{ GeV}^2$ has been measured [10], to be compared to a theoretical prediction $A_y^{th}({}^4He) \approx 10^{-10}$ [13]. The difference (by five orders of magnitude) was possibly explained by a significant contribution of the excited states of the nucleus.

The measured value of the asymmetry allows to determine the size of the imaginary part of the spin-flip amplitude F_1 [10]. From Eq. (11) we obtain $Im F_1 \approx -F(q^2)$ for $\theta = 6^\circ$. Assuming that $Re F_1 \approx Im F_1$, then the contribution of F_1 to the differential cross section is negligible due to the small factor m^2/M^2 . One may expect that the imaginary part of the non-spin-flip amplitude, namely, its TPE part, is of the same order as $Im F_1$ since we singled out the small factor m/M from the amplitude F_1 . In this case we obtain an extremely large value for the TPE mechanism, of the same order as the OPE contribution itself, at such low q^2 value. We can conclude that either our assumption, about the magnitudes of $Im f$ and $Im F_1$, is not correct, or the experimental results on the asymmetry are somewhat large.

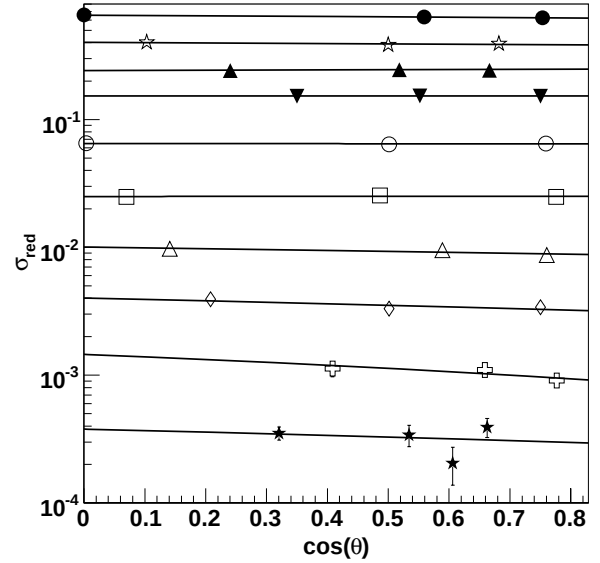


Figure 1: Reduced cross section as a function of $\cos \theta$, at $Q^2=0.5, 1, 1.5, 2, 3, 4, 5, 6, 7, 8 \text{ fm}^{-2}$ (from top to bottom). The data are from Ref. [12]. Lines are two parameter linear fits.

2 Annihilation channel

The general analysis of the polarization phenomena in the reaction

$$e^+(p_+) + e^-(p_-) \rightarrow p(p_1) + \bar{p}(p_2) \quad (12)$$

and in the time reversal channel, taking into account the TPE contribution, was done in Refs. [14].

The TPE contribution, if present, should also manifest itself in the time-like region. From first principles, as the C-invariance of the electromagnetic interaction and the crossing symmetry, the presence of TPE would create a forward backward asymmetry in the differential angular distribution of the emitted particle.

Such angular distributions have been recently measured by Babar [15], for different ranges of the invariant mass of the $p\bar{p}$ pair, after selection of the reaction (12) by tagging a hard photon from initial state radiation (ISR). The distributions have been built with the help of a Monte Carlo (MC) simulation, which takes into account the properties of the detection and allows to subtract the background.

For each $\cos \theta$ bin and each invariant mass interval, the angular asymmetry is defined as:

$$\mathcal{A}(c) = \frac{\frac{d\sigma}{d\Omega}(c) - \frac{d\sigma}{d\Omega}(-c)}{\frac{d\sigma}{d\Omega}(c) + \frac{d\sigma}{d\Omega}(-c)}, \quad c = \cos \theta \text{ and } \mathcal{A}(0) = 0. \quad (13)$$

The dependence of the asymmetry as a function of $\cos \theta$ is rather flat and can be fitted by a constant, in each mass range. The experimental values of the asymmetry are compatible

with zero for all mass ranges, with a typical error is $\sim 5\%$. As no systematic effect over $M_{p\bar{p}}$ appears, one can calculate the global average: $\mathcal{A} = 0.01 \pm 0.02$ (Fig. 2).

Note that radiative corrections of C-odd nature could also contribute to an eventual asymmetry in the data. Other odd contributions to the reaction (12), with respect to $\cos\theta$, may arise due to Z -boson exchange and C-odd interference of radiative amplitudes (including the emission of virtual and real photons). For energies smaller than the Z -boson mass, $\sqrt{t}/M_Z \ll 1$, the Z -boson exchange can be neglected.

The largest contribution to the asymmetry is represented by a factor which depends on the soft photon energy ΔE and is partially compensated by hard photon emission, with energy $\omega > \Delta E$. The hard contribution to the asymmetry A^{hard} was explicitly calculated in Ref. [16]. The total contribution to the asymmetry is expected to be $A^{tot} = A^{soft} + A^{hard} \leq 2\%$ [17].

The total contribution from radiative corrections to the angular asymmetry is not expected to exceed 2%. Moreover, radiative corrections have been applied to the data, therefore part or all of the asymmetry arising by soft and hard photon emission is already taken into account in the differential cross section.

The analysis of the available data shows no asymmetry, within an error of 2%. Such error is of the order of the asymmetry expected from radiative corrections as calculated from QED. As no systematic deviations are seen, we can conclude that these data do not give any hint of the presence of TPE, in all the considered kinematical range.

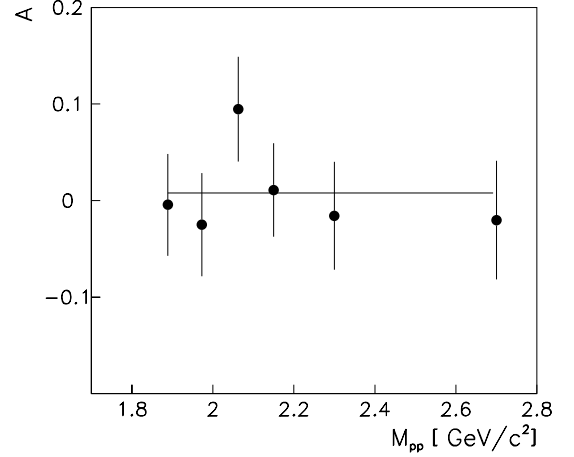


Figure 2: Average forward backward asymmetry as a function of $M_{p\bar{p}}$.

3 Electron and positron scattering

In the Born approximation, the elastic cross section is identical for positrons and electrons. A deviation of the ratio:

$$R = \frac{\sigma(e^+h \rightarrow e^+h)}{\sigma(e^-h \rightarrow e^-h)} = \frac{1 + A^{odd}}{1 - A^{odd}} \quad (14)$$

from unity would be a clear signature of processes beyond the Born approximation. Those processes include the interference of OPE and TPE, and all the photon emissions which bring a C-odd contribution to the cross section.

A model for TPE in $e^\pm p$ scattering was derived in Ref. [7], and the charge asymmetry:

$$A^{odd} = \frac{d\sigma^{e+p} - d\sigma^{e-p}}{d\sigma^{e+p} + d\sigma^{e-p}} = -\frac{2\alpha}{\pi} \left[\ln \frac{1}{\rho} \ln \frac{Q^2 x}{2\rho(\Delta E)^2} + \text{Li}_2 \left(1 - \frac{1}{\rho x} \right) - \right. \quad (15)$$

$$\left. \text{Li}_2 \left(1 - \frac{\rho}{x} \right) \right], \quad x = \frac{\sqrt{1+\tau} + \sqrt{\tau}}{\sqrt{1+\tau} - \sqrt{\tau}}, \quad \tau = \frac{Q^2}{4M^2} \quad (16)$$

was expressed as the sum of the contribution of two virtual photon exchange, (more exactly the interference between the Born amplitude and the box-type amplitude) and

a term which depends on the maximum energy of the soft photon, which escapes the detection, ΔE (E is the initial energy and ρ is the fraction of the initial energy carried by the scattered electron). It turns out that it is namely this term which gives the largest contribution to the asymmetry and contains a large ϵ dependence ($\epsilon^{-1} = 1 + 2(1 + \tau) \tan^2(\theta_e/2)$). Note that Eq. (16) holds at first order in α and does not include multi-photon emission.

Let us note that a C-odd effect is enhanced in the ratio (14) with respect to the asymmetry (16). Experiments on elastic and inelastic scattering of e^+ and e^- beams in identical kinematical conditions have been performed and recently reviewed in [18].

The elastic data are shown in Fig. 3 as a function of ϵ and compared to the model of Ref. [7] plotted at three values of Q^2 . The comparison of the calculation with the data has to be performed point by point for the corresponding (ϵ, Q^2) values. The agreement turns out to be very good, in the limit of the experimental errors, showing that a possible deviation of this ratio from unity is related to soft photon emission. One can conclude that the data on the cross section ratio are compatible with the assumption that the hard two-photon contribution is negligible.

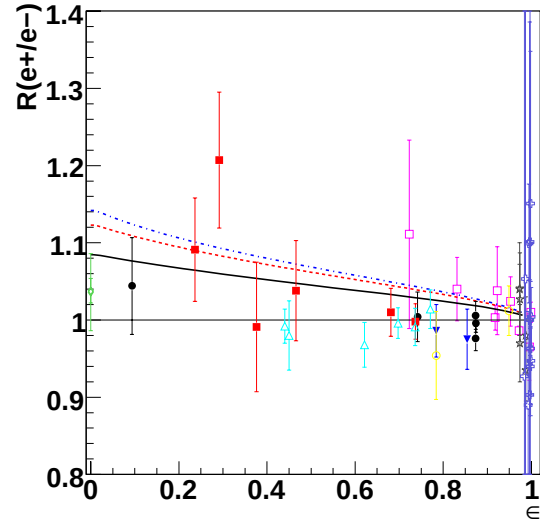


Figure 3: Ratio of cross sections $R = \sigma(e^+p)/\sigma(e^-p)$, as a function of ϵ , for $c = 0.97$ and $Q^2 = 1 \text{ GeV}^2$ (solid line, black) $Q^2 = 3 \text{ GeV}^2$ (dotted line, red) and $Q^2 = 5 \text{ GeV}^2$ (dash-dotted line, blue).

4 Conclusions

In the scattering and annihilation channels involving the electron proton interaction, the presence of TPE can be parameterized in a model independent way. In the scattering channel, the additional terms induced by TPE depend on the angle of the emitted particle, and manifest as an angular dependence of the reduced differential cross section at fixed Q^2 . The TPE contribution could also be detected using a transversally polarized electron beam, through a T-odd asymmetry of the order of the electron mass. An analysis of the existing data does not allow to reach evidence of the presence of the TPE mechanism for ^4He , as well as for other reactions involving protons and deuterons.

In the annihilation channel, the analysis of the BABAR data on $e^+ + e^- \rightarrow \bar{p} + p$, in terms of $\cos \theta$ asymmetry of the angular distribution of the emitted proton, does not show evidence of TPE, in the limit of the uncertainty of the data.

The difference in the cross section for $e^\pm p$ scattering can be explained by odd terms, which are present in standard radiative corrections.

One can conclude that the data do not show evidence for the presence of the TPE at the level of their precision. TPE is expected to become larger when the momentum transfer increases. Its study in the kinematical range covered by the present experiments requires

more precise and dedicated measurements. In this respect, future antiproton beams at the FAIR facility in Darmstadt will provide very good conditions for the measurement of time-like proton form factors and for the study of the reaction mechanism.

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FIRST MEASUREMENT OF THE INTERFERENCE FRAGMENTATION FUNCTION IN e^+e^- AT BELLE

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Abstract

A first measurement of the di-hadron interference fragmentation function of light quarks in pion pairs with the Belle detector is presented. The chiral odd nature of this fragmentation function allows the use as a quark polarimeter sensitive to the transverse polarization of the fragmenting quark. Therefore it can be used together with data taken at fixed target and collider experiments to extract the quark transversity distribution. A sample consisting of 711×10^6 di-hadron pairs was extracted from 661 fb^{-1} of data recorded near the $\Upsilon(4S)$ resonance delivered by the KEKB e^+e^- collider.

1 Introduction

The di-hadron interference fragmentation function (IFF), suggested first by Collins, Hoppmann and Ladinsky [1] describes the production of unpolarized hadron pairs in a jet from a transversely polarized quark. The transverse polarization is translated into an azimuthal modulation of the yields of hadron pairs around the jet axis. In addition the IFF is chiral odd and can therefore act as a partner for the likewise chiral odd quark transversity function. The resulting amplitude is chiral even and therefore leads to observable effects in semi deep inclusive scattering off a transversely polarized target [2, 3] or in proton-proton collisions [4] in which one beam is transversely polarized. Besides the fragmentation into two hadrons, the fragmentation of a transversely polarized quark into one unpolarized hadron can be used to extract transversity via the Collins effect. Measurement of this effect at Belle [5] made the first extraction of transversity possible [6]. However as compared with the Collins fragmentation function, using the IFF to extract transversity exhibits a number of advantages. These are connected to the additional degree of freedom provided by the second hadron. It allows to define the azimuthal angle between the two hadrons as an observable in the transverse plane and at the same time integrate over transverse momenta of the quarks and hadrons involved. Because transverse momenta are integrated over collinear schemes in factorization and evolution can be used which are known and which do not need assumptions of the intrinsic transverse momenta [7]. Since the IFF is not a transverse momentum dependent function (TMD) it is universal and therefore directly applicable to SIDIS and proton proton data. From

both types of experiment results are available [2–4]. The results from SIDIS are indicating a non-zero IFF while the first analysis of the IFF effect at PHENIX [4] opens a way to disentangle transverse spin effects in proton proton collisions. Because the IFF is not a TMD function, extraction of transversity times IFF is not dependent on a model of the transverse momentum dependence, which is the case in the measurement of the Collins effect [6]. Here transverse momentum in the final state originates from a convolution of quark distribution and fragmentation function. This leads also to the Sudakov suppression of the effect [8]. Technically, the extraction of the IFF from electron colliders is also easier, as the signal is not competing with asymmetries from QCD radiation. Also acceptance effects are smaller for the relative azimuthal angles of hadron pairs.

2 Observables in e^+e^- collisions

The transverse polarization of the fragmenting quark leads to a cosine modulation of the azimuthal angle of the plane spanned by the two hadrons h_1, h_2 which is described by the vector $\mathbf{R} = \mathbf{P}_{h1} - \mathbf{P}_{h2}$ lying in the plane. Since the electron beams are unpolarized any effect that one would measure in one hemisphere of an event would average out. Instead one can make use of the fact that the spins of quark and anti quark in electron-positron annihilation are 100% correlated. Thus the correlation of the azimuthal angles of the vectors $\mathbf{R}^\alpha$ in the hemispheres $\alpha \in \{1, 2\}$ around the thrust axis with regard to the event plane is sensitive to the IFF. The measurement uses the center of mass system and defines the event plane as the plane which contains the beam axis $\hat{\mathbf{z}}$ and thrust axis $\hat{\mathbf{n}}$. Figure 1 shows the coordinate system with the relevant quantities. With these the angles ϕ_α between $\mathbf{R}_\alpha$ and the event plane can be expressed as

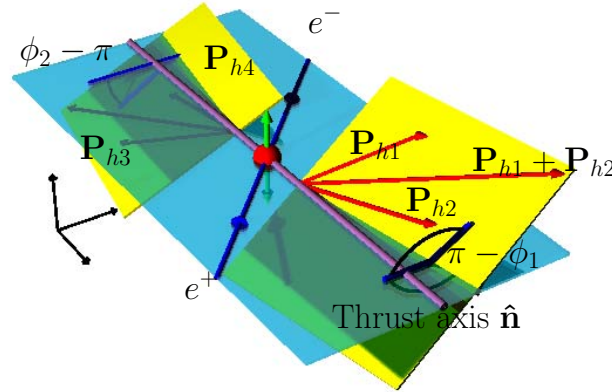


Figure 1: Azimuthal angle definition. Azimuthal angles ϕ_1 and ϕ_2 are defined relative to the thrust axis.

$$\begin{aligned} \phi_{\{1,2\}} &= \text{sgn} [\hat{\mathbf{n}} \cdot (\hat{\mathbf{z}} \times \hat{\mathbf{n}} \times (\hat{\mathbf{n}} \times \mathbf{R}_{1,2}))] \\ &\times \arccos \left(\frac{\hat{\mathbf{z}} \times \hat{\mathbf{n}}}{|\hat{\mathbf{z}} \times \hat{\mathbf{n}}|} \cdot \frac{\hat{\mathbf{n}} \times \mathbf{R}_{1,2}}{|\hat{\mathbf{n}} \times \mathbf{R}_{1,2}|} \right) . \end{aligned} \quad (1)$$

The product of the quark and anti quark interference fragmentation functions $H_1^\triangleleft \cdot \bar{H}_1^\triangleleft$ is then proportional to the amplitude a_{12} of the modulation $\cos(\phi_1 + \phi_2)$ of the di-hadron

pair yields [9]. The di-hadron fragmentation functions are dependent on the kinematic variables m_{Inv} , the invariant mass of the hadron pair, and $z = \frac{2E_h}{Q}$ the normalized energy of the hadron pair. Here E_h is the energy of the hadron and Q the absolute energy transferred by the virtual photon. Therefore the normalized azimuthal yield of di-hadron pairs can be described as

$$N(y, z_1, z_2, m_1, m_2, \phi_1 + \phi_2) \propto a_{12}(y, z_1, z_2, m_1, m_2) \cdot \cos(\phi_1 + \phi_2) \quad (2)$$

with

$$a_{12}(y, z_1, z_2, m_1, m_2) \sim B(y) \cdot \frac{\sum_q e_q^2 H_1^{\triangleleft}(z_1, m_1) \bar{H}_1^{\triangleleft}(z_2, m_2)}{\sum_q e_q^2 D_1(z_1, m_1) \bar{D}_1(z_2, m_2)}. \quad (3)$$

Here the sum goes over all quark flavors q and $B(y) = \frac{y(1-y)}{\frac{1}{2}-y+y^2} \stackrel{\text{CM}}{=} \frac{\sin^2 \theta}{1+\cos^2 \theta}$ is the kinematic factor describing the transverse polarization of the quark-anti quark with regard to its momentum. The polar angle θ is defined between the electron axis and the thrust axis as shown in fig. 1. The yields are dependent on z_1, z_2, m_1, m_2 , the fractional energies and invariant masses of the hadron pairs in the first and second hemisphere, and y . In the following the yields are integrated over y in the limits of the fiducial cuts introduced in section 2.3. The labeling of the hemisphere is at random, and it is experimentally not possible to distinguish between quark and anti quark fragmentation. Further information about the IFF can be learned from the dependence on the decay angle θ_h in the CMS of the two hadrons produced. Using a partial wave decomposition to isolate components that contain the interference between waves with one unit difference in angular momentum, one expects a dependence on $\sin \theta_h$ which gives the interference term between s and p waves. This term is expected to dominate at Belle kinematics and is favored by the acceptance. Since the acceptance is symmetric around $\theta_h = \frac{\pi}{2}$ the p-p contribution proportional to $\cos \theta_h$ should average out. Table 3 shows that this is approximately the case.

2.1 Models

As described in the previous section the IFF is an interference effect between hadrons, here pions, created in partial waves with a relative angular momentum difference of one. The dominant contribution being from the s-p interference term [10,11]. Therefore information about the IFF can be gained from a partial wave analysis for di-pion production [12]. Here the data from [13] suggest a phase shift around the ρ mass, leading to a sign change in the fragmentation function. In some models [14] the location of the phase shift around the mass of the ρ meson is caused by the interference of pion pairs produced in a p wave coming from the decay of the spin one ρ which interferes with the non-resonant background. Based on this model and estimation of particle yields from simulations, Bacchetta, Ceccopieri, Mukherjee and Radici [15] made model predictions for the magnitude of IFF Asymmetries at Belle. Again, a strong dependence on the invariant mass is predicted, with a maximum around the rho mass. The asymmetries are expected to rise with z due to the preservation of spin information early in the fragmentation.

2.2 The Belle experiment

The Belle detector, located at the KEKB asymmetric energy e^+e^- collider is described in detail in [16]. For the purpose of this study it is important that a high number of events

is recorded, a good particle identification allows to identify pions up to high values of z and that the detector is hermetic to minimize acceptance effects. Both is well fulfilled by the Belle detector. It is located at the collision point of the 3.5 GeV e^+ and 8 GeV e^- beam and is almost symmetric in the center of mass system of the beams. It is a large solid angle magnetic spectrometer consisting of barrel and end caps parts. For a more homogeneous acceptance function only the barrel part was used for this analysis. It comprises a silicon vertex detector, a 50-layer drift chamber and an electromagnetic calorimeter (CsTI) in a magnetic field of 1.5T. Particle identification over a wide range of momenta is done using an array of aerogel Cherenkov counters, a barrel-like arrangement of time-of-flight scintillation counters and an instrumented iron flux return yoke outside of the coil to detect K_L^0 mesons and muons.

2.3 Data selection and Asymmetry extraction

From the 661 fb^{-1} data sample roughly 589 fb^{-1} were taken on the $\Upsilon(4S)$ resonance and 73 fb^{-1} taken in the continuum 60 MeV below. Because the thrust cut used to select events with a two jet topology containing light and charm quarks also rejects events in which B mesons were produced, on-resonance and continuum data can be combined. The thrust is defined as $T = \frac{\sum_i |p_i \cdot \hat{n}|}{\sum_i |p_i|}$ and $\hat{n}$ is direction of the thrust chosen such that T is maximal. B meson events produced on the Υ resonance have a spherical shape, since their high mass does not allow for large kinetic energy. On the other side, light quark anti quark pairs have a more two jet like topology.. An applied thrust cut of $T > 0.8$ reduces the contamination with B events to an order of 2% [5]. The inversion of the thrust cut selects events that don't have a clear two jet topology and that are contaminated by Υ decays into B mesons. This leads to a decrease of the asymmetry.

As described earlier, only the barrel region of the detector is used in the analysis. Therefore the thrust axis is required to lie in a region well contained within it. This translates into a cut on the z component of the thrust axis $|\mathbf{n}_z| < 0.75$, which also allows for the particles of the jet around the axis to be reconstructed in the barrel. Further cuts on the event level are a reconstructed energy of at least 7 GeV to reject $e^+e^- \rightarrow \tau^+\tau^-$ and to reliably reconstruct the thrust axis. The later is computed using all charged tracks and photons passing some minimum energy cuts. A mean deviation of 135 mrad with a RMS of 90 mrad of the thrust axis from the real quark-anti quark axis is computed from simulations.

Only events were selected which satisfy a vertex cut of 2 cm in the radial and 4 cm in the beam direction. On the track level the fiducial cuts to reduce acceptance effects are a constraint to the barrel region of the detector using a cut on the polar angle θ in the laboratory system of $-0.6 < \cos(\theta) < 0.9$ which translates to an almost symmetric cut in the CMS. To make sure that the azimuthal range of tracks around the thrust axis is not biased, only tracks are chosen that have at least 80% of their energy along the thrust axis. This restricts tracks to be within a cone that is entirely contained in the acceptance and reduces false asymmetries considerably as shown later in sec. 2.4. No false asymmetries from this cut is expected and since most of the energy of the jet is contained within the jet no significant dilution of the asymmetries either. Only tracks above a minimal fractional momentum $z > 0.1$ are considered. They have to be positively identified as pions. These tracks are then sorted into two hemispheres according to the sign of their momentum

projection on the thrust axis. All possible pairs of $\pi^+\pi^-$ in the same hemisphere are selected and the angle $\phi_\alpha, \alpha \in \{1, 2\}$ computed. To this end the vector $\mathbf{R}_\alpha = \mathbf{P}_1 - \mathbf{P}_2$ for each pion pair in hemisphere α is formed and the azimuthal angle around the thrust axis with respect to the event plane computed according to eq. 1. Charge ordering of h_1, h_2 in the computation of $\mathbf{R}$ is always the same so that the effect is not averaged out. A weighting of the hadron momentum vector with the inverse fractional energy z as suggested by [17] only led to differences in the thrust axis within numerical uncertainties. For the computation of the yields in a specific $\phi_1 + \phi_2$ bin all combinations of pairs in the two hemispheres are considered. Due to the cone cut described earlier, the number of hadrons with a false hemisphere assignment is negligible. From Monte Carlo studies a signal purity for di-pion pairs (4 particles) of better than 90% over the whole kinematic range is obtained.

2.4 Systematics studies

In order to determine the systematic error on the measurement, simulation and real data was used to determine the contribution of detector effects and competing physical processes to false asymmetries or a dilution of the measurement. The studies that lead to the biggest contribution to the systematic error are the study of false asymmetries in simulation and real data in which no asymmetry is expected due to a wrong assignment of the thrust axis or the hemispheres of the particle pairs. Any false asymmetries, together with their statistical errors were added to the systematic error.

Since the IFF effect is not included in the Pythia event generator used, checking for the asymmetries in fully simulated data allows to estimate effects of the detector acceptance and efficiencies on the asymmetry. Table 1 shows the false asymmetries extracted from a full simulation of the detector using GEANT.

Another source for false asymmetries are events reconstructed from real data in which the asymmetries are averaged out. This is the case for mixed events in which the angles ϕ_1 and ϕ_2 are taken from different events or events in which the angles are computed for pion pairs in the same hemisphere. The later case leads to asymmetries due to phase space restrictions, which could be reproduced, the former to asymmetries in the order of one per-mille, shown in table 2, which has been added to the systematic error. Contributions from higher harmonics in the fits to the cosine modulation are under one per-mille.

Due to the smearing of the thrust axis reconstruction with respect to the true quark-anti quark axis, the extracted asymmetries are diluted. Since this dilution could be reproduced in weighted Monte Carlo using the observed smearing of the thrust axis, it was corrected for.

More physics motivated systematic checks concern the dependence on the various kinematic factors. The dependence on the kinematic factor $\frac{\sin^2 \theta}{1 + \cos^2 \theta}$ should be linear, as should be the dependence on $\sin \theta_h$, if the effect is dominated by the s-p interference term. Both could be validated.

Even though the Belle detector is very stable over time, checks have been done to determine the compatibility of data taking periods and data taken on and off the Υ resonance. To this end, the χ^2 of each fit has been computed and all values have been fitted by an appropriate χ^2 distribution. The result of these tests show very good compatibility.

We checked for correlations in the data that might lead to false error estimates by breaking up weighted Monte Carlo data in 500 small chunks and comparing the mean

error of the extracted asymmetries to their variation. The results are compatible, so no systematic error was assigned.

For the interpretation of the results the respective fraction of processes contributing to the asymmetries are very important. These have been estimated from Monte Carlo simulations and are shown in fig. 2. It is evident that the contribution from Υ decays has been almost eliminated by the thrust cut. There are also marginal contributions from τ pairs. For these the false asymmetries are compatible with zero and they have been added with their statistical error to the overall systematic error. Besides the contribution from light quark pairs there is a considerable contribution from charm quarks. This contributions decreases with increasing z which can be understood, since charmed mesons have to undergo an additional decay before they can contribute to the pion asymmetries. This decay can also cause the invariant mass dependence namely a general decrease with higher masses. However, the highest bin shows a very high charm contribution, in some bins more than half of the events, which is not yet understood. The asymmetry results for these bins indicate that the IFF for charm quarks is non-vanishing and of similar magnitude as that of lighter quarks.

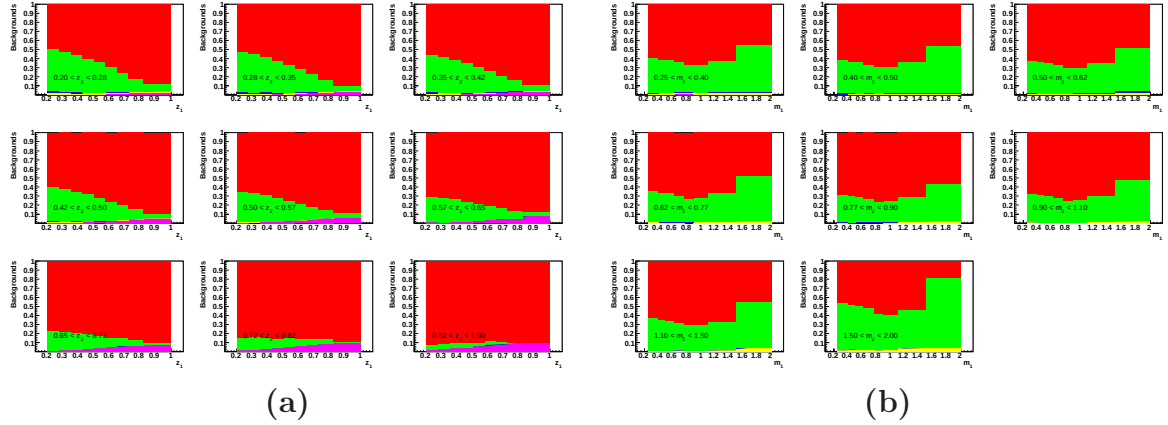


Figure 2: Relative process contributions from light quark-anti quark events (red), charm events (green), charged B meson pairs (blue), neutral B meson pairs (yellow) and τ pairs (purple) as a function of z_2 for all z_1 bins (a) and as a function of m_2 for all m_1 bins (b)

2.5 Results

The results obtained for the a_{12} asymmetry as defined in eq. 2 are shown in figs. 3 binned in the invariant masses m_1 , m_2 of the hadrons pairs in the first and second hemisphere and their fractional energies z_1 , z_2 , respectively. Table 3 shows the integrated asymmetries and the averaged kinematic observables. The extracted asymmetries are large, especially when considering that a product of the IFF for quark and anti quarks is measured. As expected the magnitude of the effect rises with z . However, the invariant mass behavior does not match model predictions from [15]. But these model predictions are only available in leading order and heavily dependent on simulations which were tuned for the SIDIS experiment HERMES at a center of mass energy of roughly 7 GeV. The asymmetries rise up to around the mass of the ρ but then plateau instead of decreasing again. A

sign change of the IFF can therefore not be confirmed. However bins with high invariant masses receive also considerable contributions from charm quarks as shown before.

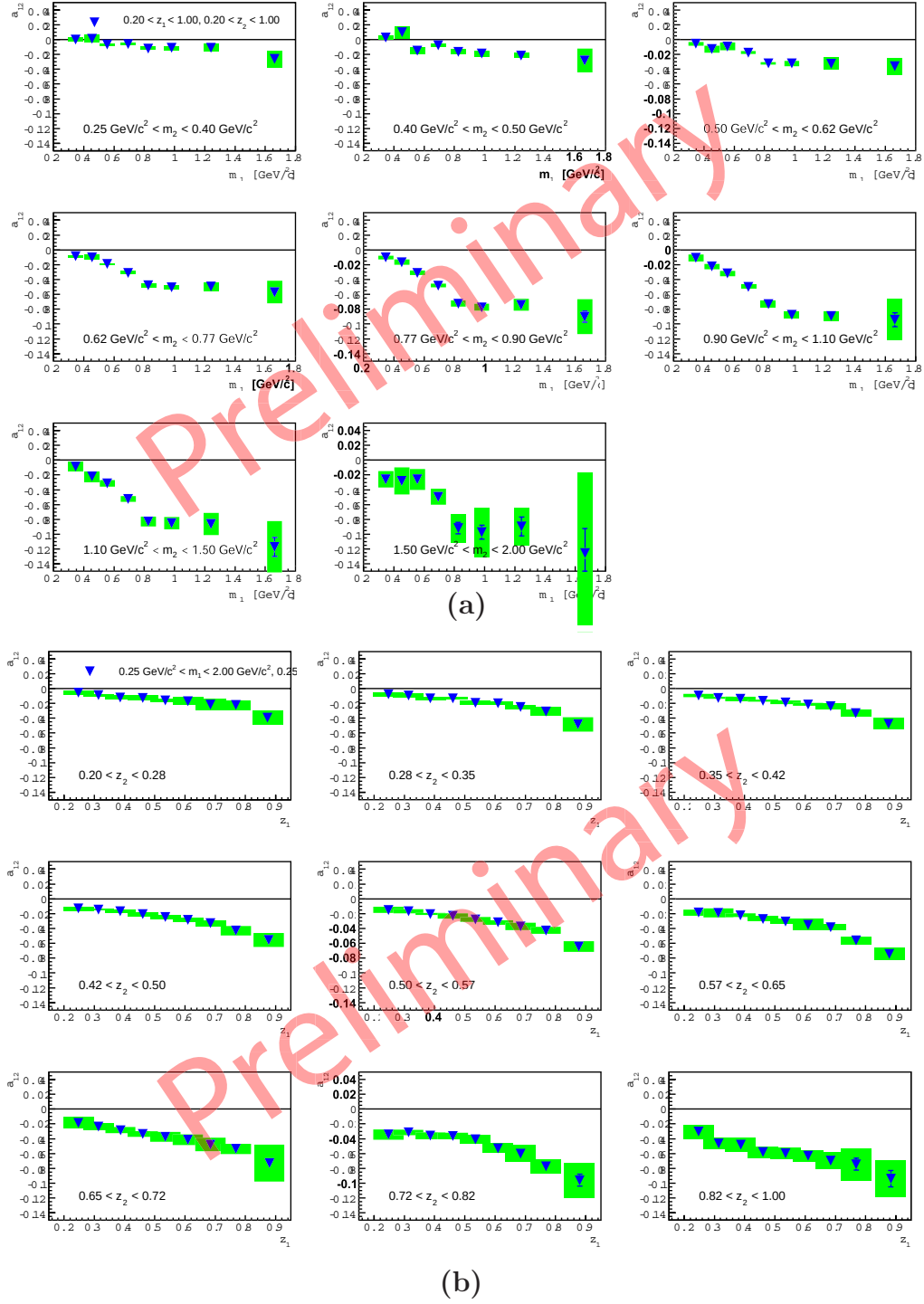


Figure 3: Results for the a_{12} modulations in a symmetric 8×8 binning in m_1 , m_2 for $m_{\{1,2\}}$ between 0.25 GeV and 2 GeV (a) and for a symmetric 9×9 binning in z_1 , z_2 for $z_{\{1,2\}}$ between 0.2 and 1 (b). The statistical error is shown in blue, the systematic error in green

2.6 Summary and Outlook

The first direct measurement of the interference fragmentation function using 661 fb^{-1} of data recorded at the Belle experiment has been presented. The asymmetries are large, up to 10 %, which would correspond to an IFF contribution of over 30%. Models predicting a sign change or a decrease for invariant masses higher than the ρ meson's could not be confirmed. Charm quarks play a significant role at high invariant masses. They seem to introduce an IFF asymmetry of similar magnitude as light quarks and further studies are under way to determine the charm quark contribution to the asymmetries. The analysis presented here should enable a combined analysis to extract the transversity distribution of data taken in SIDIS, proton-proton and e^+e^- . This is very desirable due to the various advantages as compared with the extraction via the Collins effect and its complementarity. In proton-proton collisions the use of the IFF effect to access transversity is especially helpful, since it can help disentangle different contributions to the measured transverse single spin asymmetries A_N . Our plans for the future contain also an extraction of other particle combinations, namely those including neutral pions and charged kaons. Furthermore an extraction of the unpolarized yields is planned to facilitate the extraction of the IFF.

Table 1: MC results averaged over all z bins in %.

sample	species	z_1, z_2 -Asymmetries	
		$\langle a_{12} \rangle$	$\langle a_{12R} \rangle$
No opening cut			
uds 4π	$\pi\pi$	-0.089 ± 0.008	-0.108 ± 0.008
uds acceptance	$\pi\pi$	-0.488 ± 0.011	-0.490 ± 0.011
uds MC rec.	$\pi\pi$	-0.394 ± 0.013	-0.418 ± 0.013
charm rec.	$\pi\pi$	-0.446 ± 0.041	-0.388 ± 0.044
With opening cut of 0.8			
uds 4π	$\pi\pi$	-0.038 ± 0.013	-0.035 ± 0.013
uds acceptance	$\pi\pi$	-0.112 ± 0.016	-0.113 ± 0.016
uds MC rec.	$\pi\pi$	0.012 ± 0.019	0.008 ± 0.019
charm rec.	$\pi\pi$	0.006 ± 0.040	0.027 ± 0.040

Table 2: Mixing results averaged over the z binning in %. The results integrated over other binnings are nearly identical.

sample	z_1, z_2 -Asymmetries	
	$\langle a_{12} \rangle$	$\langle a_{12R} \rangle$
uds 4π	0.070 ± 0.013	0.030 ± 0.013
uds acceptance	0.020 ± 0.016	-0.021 ± 0.016
uds rec.	0.091 ± 0.019	0.087 ± 0.019
charm rec.	-0.024 ± 0.040	-0.017 ± 0.040
Data	-0.019 ± 0.017	-0.012 ± 0.017

Table 3: Integrated asymmetries and average kinematics.

$\langle z_1 \rangle, \langle z_2 \rangle$	0.4313
$\langle m_1 \rangle, \langle m_2 \rangle$	0.6186
$\langle \sin^2 \theta / (1 + \cos^2 \theta) \rangle$	0.7636
$\langle \sin \theta_h \rangle$	0.9246
$\langle \cos \theta_h \rangle$	0.0013
a_{12}	$-0.0199 \pm 0.0002 \pm 0.0009$

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TECHNICS
and
NEW DEVELOPMENTS

PROTON BEAM POLARIZATION MEASUREMENTS AT RHIC

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Abstract

We discuss the beam polarization measurement strategy at RHIC with achieved relative precision of better than 5%.

1 Introduction

Polarimetry is one of the crucial part of the RHIC Spin Program [1]. It provides precise measurements of beam polarizations which is a necessary component of all high precision spin dependent measurements at RHIC.

Two types of polarimeters are used at RHIC, which are based on small angle elastic scattering, with sensitivity to the proton beam polarization coming from interference between electromagnetic and hadronic amplitudes (Coulomb-Nuclear Interference (CNI) region).

One type of polarimeter (pC) uses an ultra-thin ($10\text{--}20\text{ }\mu\text{g}/\text{cm}^2$) carbon ribbon target, and provides fast relative polarization measurements with a few percent statistical uncertainty in 20–30 sec. The other type polarimeter (H-Jet) uses a polarized hydrogen gas target and utilizes the analyzing power in proton-proton elastic scattering in the CNI region. It accumulates data over the entire store and provides absolute polarization measurements, with $\sim 10\%$ statistical uncertainty in a store (6–8 hours). H-Jet data accumulated over several stores are used for the absolute calibration of the pC polarimeters.

In addition, two experiments STAR and PHENIX employ local polarimeters which are used to set up and to monitor the spin direction at collision. They also proved to be a valuable tool for beam polarization monitoring when spin orientation at collision was set up to be transverse.

2 H-Jet Polarimeter

The polarized hydrogen jet target polarimeter locates at one of the collision points in the RHIC accelerator and measures polarization of one of the two RHIC beams at a time. In 2008 the the H-Jet polarimeter was tested running both beams separated vertically, and both incident on the jet. In 2009 RHIC run it allowed the H-Jet to simultaneously monitor the polarization of both beams on store by store basis.

H-Jet polarimeter consists mainly of three parts: atomic beam source, scattering chamber and Breit-Rabi polarimeter (BRP). Atomic beam source produces polarized atomic hydrogen jet with velocity 1560 ± 20 m/c [2] which crosses RHIC proton beam in scattering chamber perpendicularly. The FWHM of the jet in the center of scattering chamber is ~ 6.5 mm. The total atomic beam intensity in the scattering chamber was measured to be $(12.4 \pm 0.2) \cdot 10^{16}$ atoms/cm² [3]. Along the RHIC beam axis the target thickness is $(1.3 \pm 0.2) \cdot 10^{12}$ atoms/cm² [2]. BRP measured the atomic hydrogen polarization. It is monitored during the run and showed a stable polarization of $P_{target} = 0.924$ after correction for $\sim 3\%$ molecular hydrogen contamination, which diluted jet polarization [2].

Recoil protons from elastic pp scattering are detected using an array of silicon detectors located on the left and right of the beam at a distance ~ 80 cm (Fig. 1a)). The elastically scattered protons are identified with small background contamination ($\sim 4\%$) using the kinematical correlation between recoil proton kinetic energy and time of flight, and kinetic energy and scattering angle (or silicon detector strip number in Fig. 1a)). Then asymmetry between scattering on the left and on the right are measured relative to beam polarization direction or target polarization direction. Since beam polarization direction alternates every bunch (separated by 106 ns) or pair of bunches, and target polarization flips every ~ 10 min, we simultaneously collect data for all beam-target spin configurations: $N_{\uparrow}^L$, $N_{\downarrow}^L$, $N_{\uparrow}^R$ and $N_{\downarrow}^L$, where subscript denotes the beam or target polarization direction, and superscript relates to left or right detector, defined relative to beam direction. Using so called square-root formula:

$$\epsilon = \frac{\sqrt{N_{\uparrow}^L \cdot N_{\downarrow}^R} - \sqrt{N_{\uparrow}^R \cdot N_{\downarrow}^L}}{\sqrt{N_{\uparrow}^L \cdot N_{\downarrow}^R} + \sqrt{N_{\uparrow}^R \cdot N_{\downarrow}^L}}, \quad (1)$$

we cancel effects from left-right detector acceptance asymmetry and spin up-down beam luminosity asymmetry, when measuring raw asymmetry relative to beam polarization, ϵ_{beam} , or relative to target polarization ϵ_{target} .

Beam polarization P_{beam} is then defined from the measured asymmetries ϵ_{beam} and ϵ_{target} and known target polarization P_{target} :

$$P_{beam} = -\frac{\epsilon_{beam}}{\epsilon_{target}} \cdot P_{target}. \quad (2)$$

The details of H-Jet setup and Data Acquisition system are described elsewhere [4].

The asymmetry measurements were performed for the recoil proton kinetic energy range $T_R = 1 - 4$ MeV (corresponding to a momentum transfer $0.002 < -t < 0.008$ (GeV/c)²), in which the proton-proton elastic scattering analyzing power is slightly T_R dependent and on the average is $\langle A_N^{pp} \rangle \sim 4\%$. It also doesn't show any sizable beam energy dependence (Fig. 1b), which makes pp CNI polarimetry particularly valuable for measuring and monitoring proton beam polarization in wide energy range.

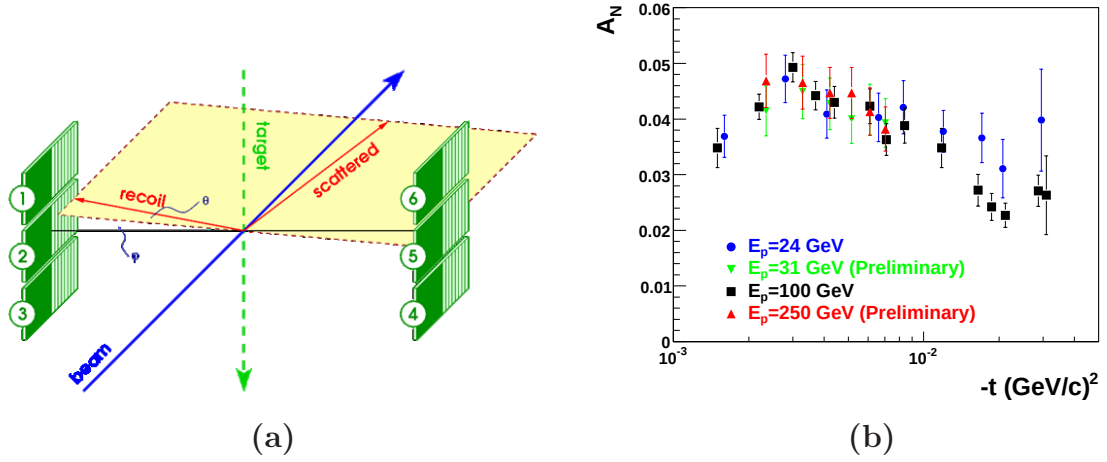


Figure 1: (a) H-Jet layout of silicon detectors with strips oriented perpendicular to the beam line; atomic hydrogen jet goes from the top to the bottom; forward strips are used to measure recoil carbon from the interaction of jet with one beam, and backward strips - with the other beam. (b) Analyzing power (A_N) for elastic pp scattering defined as $\epsilon_{target}/P_{target}$; 24 GeV and 100 GeV data are from RHIC Run 2004 ([5] and [6]), 31 GeV and 250 GeV preliminary data are from Run 2006 and 2009, correspondingly.

The accumulated statistics in H-Jet in the first long polarized proton RHIC run in 2006 provided 2.4% uncertainty for the pC polarimeter absolute normalization for each of 100 GeV beams. Other uncertainties came from the molecular hydrogen background, $\sim 2\%$, and from background in the identification of elastically scattered recoil protons, $\sim 1.5\%$. The latter background didn't show any beam or target polarization dependence, so that it diluted both ϵ_{beam} and ϵ_{target} in the same way, and according to Eq. (2) didn't affect beam polarization measurements.

3 Proton-Carbon Polarimeter

Both RHIC rings are equipped with pC polarimeters, which provide fast relative measurements of beam polarizations, which subsequently is normalized to the absolute polarization measurements performed by H-Jet polarimeter. Repeated measurements during a store also provide a handle on the beam polarization vs time as well as measurements of beam polarization profile in the transverse plane. The latter directly enters the average beam polarization observed in collisions at the interaction regions of the RHIC experiments.

The pC polarimeters consist of a carbon target and six silicon strip detectors mounted in a vacuum chamber at 45, 90 and 135 degrees relative to vertical axis in both left and right sides with respect to the beam direction, Fig. 2a. It allows to monitor not only vertical component of beam polarization through left-right asymmetry, but also horizontal component through up-down asymmetry using 45 and 135 degree detectors. The details of pC polarimeter setup are described elsewhere [7].

pC polarimeters utilize the analyzing power for the elastic polarized proton-carbon scattering in the CNI region (Fig. 2b). For the the recoil carbon kinetic energy range $T_R = 400\text{--}900 \text{ keV}$ used for the asymmetry measurements the average analyzing power

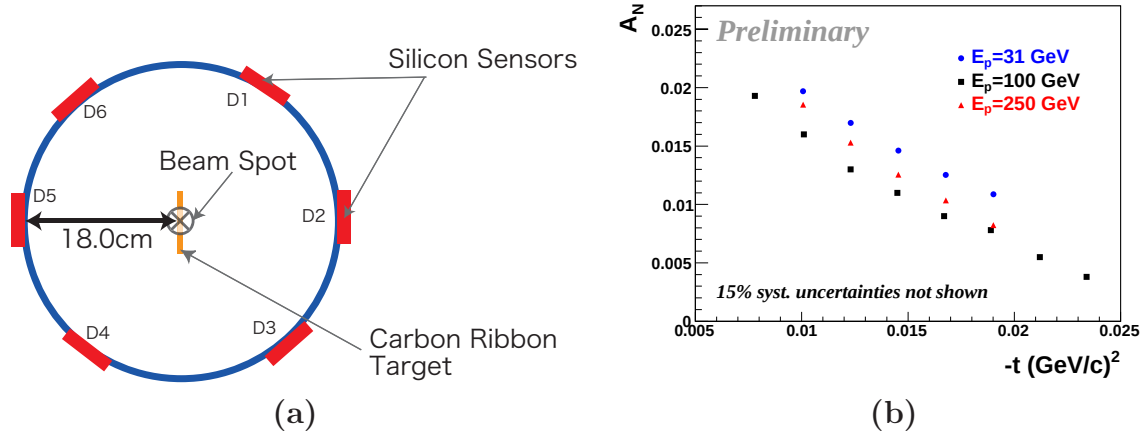


Figure 2: (a) The RHIC pC-polarimeter setup; silicon detectors are aligned 45, 90 and 135 degrees azimuthally in both left and right side with respect to the beam direction; the beam is pointing into the figure perpendicularly. (b) Analyzing power (A_N) for elastic pC scattering measured by RHIC pC polarimeters, for proton beam energy 31 GeV from RHIC Run 2006, 100 GeV from Run 2004 and 250 GeV from Run 2009.

for 100 GeV protons was $\langle A_N^{pC} \rangle \sim 1.2\%$. Within this energy range, corresponding to a momentum transfer $0.009 < -t < 0.020$ (GeV/c) ^{2} , A_N is T_R dependent and drops from $\sim 1.5\%$ at 400 keV to $\sim 0.7\%$ at 900 keV.

The identification of recoil carbon is performed through the correlation between recoil carbon kinetic energy and time of flight measured by silicon detectors. The recoil carbon energy includes corrections due to energy losses [8]. After applying 3σ mass cut corresponding to carbon mass, where σ is reconstructed carbon mass peak width, the background under carbon mass peak was estimated on the level of 1% [7].

A typical carbon ribbon target is 2.5 cm long, 10–20 $\mu\text{g}/\text{cm}^2$ thick and 10–20 μm wide, which is long, thin and narrow compared to typical beam width of 1 mm. Scanning target across the beam allows to measure beam intensity and polarization profiles in both vertical and horizontal directions in transverse plane. For non-flat polarization profile the beam polarization will be seen differently by H-Jet polarimeter, pC polarimeter (in fixed target measurement mode) and when colliding beams in the interaction regions of RHIC experiments. Due to that the jet target in H-Jet polarimeter is much wider than beam width, the polarization profile is weighted with beam intensity profile in the average beam polarization. For colliding beams, the polarization profile is weighted with a product of two beam intensity profiles in transverse plane. The polarization measurement in pC polarimeter will depend on the positioning of the carbon target, and on whether horizontal or vertical target is used for the measurements (in case beam horizontal and vertical profiles are different).

Assuming gaussian profiles ($I(x, y)$ for intensity and $P(x, y)$ for polarization), the average beam polarization seen by H-Jet polarimeter can be expressed:

$$\langle P \rangle_{HJet} = \frac{\iint P(x, y) I(x, y) dx dy}{\iint I(x, y) dx dy} = \frac{P_{max}}{\sqrt{(1 + R_X) \cdot (1 + R_Y)}}, \quad (3)$$

by pC polarimeter for the case with vertical target positioned at beam maximum intensity

(and polarization) along axis X :

$$\langle P_{max-X} \rangle_{pC} = \frac{\int P(y)I(y)dy}{\int I(y)dy} = \frac{P_{max}}{\sqrt{(1+R_Y)}}, \quad (4)$$

and by experiments in two beam collision, for $I_{1,2}$ relating to intensity profiles for two beams respectively:

$$\langle P \rangle_{Exp} = \frac{\iint P(x,y)I_1(x,y)I_2(x,y)dxdy}{\iint I_1(x,y)I_2(x,y)dxdy} = \frac{P_{max}}{\sqrt{(1+\frac{1}{2}R_X) \cdot (1+\frac{1}{2}R_Y)}}, \quad (5)$$

with P_{max} the polarization at beam maximum intensity and polarization in transverse plane, R_X and R_Y the squared ratio of the intensity profile width and polarization profile width $(\sigma_I/\sigma_P)^2$, for transverse X and Y projections, respectively.

These relations between average polarizations are taken into account when normalizing pC measurements to H-Jet absolute polarization measurements and when providing polarization values for RHIC experiments. For example for the typical at RHIC values of $R \sim 0.1 - 0.2$, beam polarization seen by H-Jet is lower than polarization seen by experiments in beam collisions by $\sim 5 - 10\%$.

Profile parameter R can be extracted either directly from the measurements of intensity and polarization profile widths σ_I and σ_P , as shown in Fig. 3a, or from the correlation between polarization and intensity at each target position (for gaussian intensity and polarization profiles), for example for the measurements with vertical target:

$$\frac{P}{P_{max-X}} = \left(\frac{I}{I_{max-X}} \right)^{R_X}, \quad (6)$$

which doesn't require knowledge on target position. I_{max-X} and P_{max-X} are beam intensity and polarization¹ (averaged along vertical direction Y), respectively, when target is positioned in the beam center along axis X . A typical measurement for 100 GeV beam is shown in Fig. 3b. In this approach the average over beam transverse cross section polarization for the measurements with vertical target is

$$\langle P \rangle_{pC} = \frac{P_{max-X}}{\sqrt{1+R_X}}, \quad (7)$$

and similarly if horizontal target is used. The value $\langle P \rangle_{pC}$ is compared to H-Jet polarization measurements, when the normalization for pC polarization measurements is derived.

The main sources of systematic uncertainties in polarization measurements by pC polarimeter comes from the drift of the recoil carbon energy correction (2.5% in Run6), and from corrections due to polarization profile, particularly in runs when one or both transverse projections of the profile are not measured regularly (2% in Run6).

Combining these uncertainties with uncertainties from polarization normalization using H-Jet, the final systematic uncertainties for the proton beam polarization measurements were 4.7% and 4.8% for two RHIC beams and 8.3% for a product of two beam polarizations in 2006 for 100 GeV beam running.

¹Compared to notation in Eq. (4) we here omitted angle brackets for brevity.

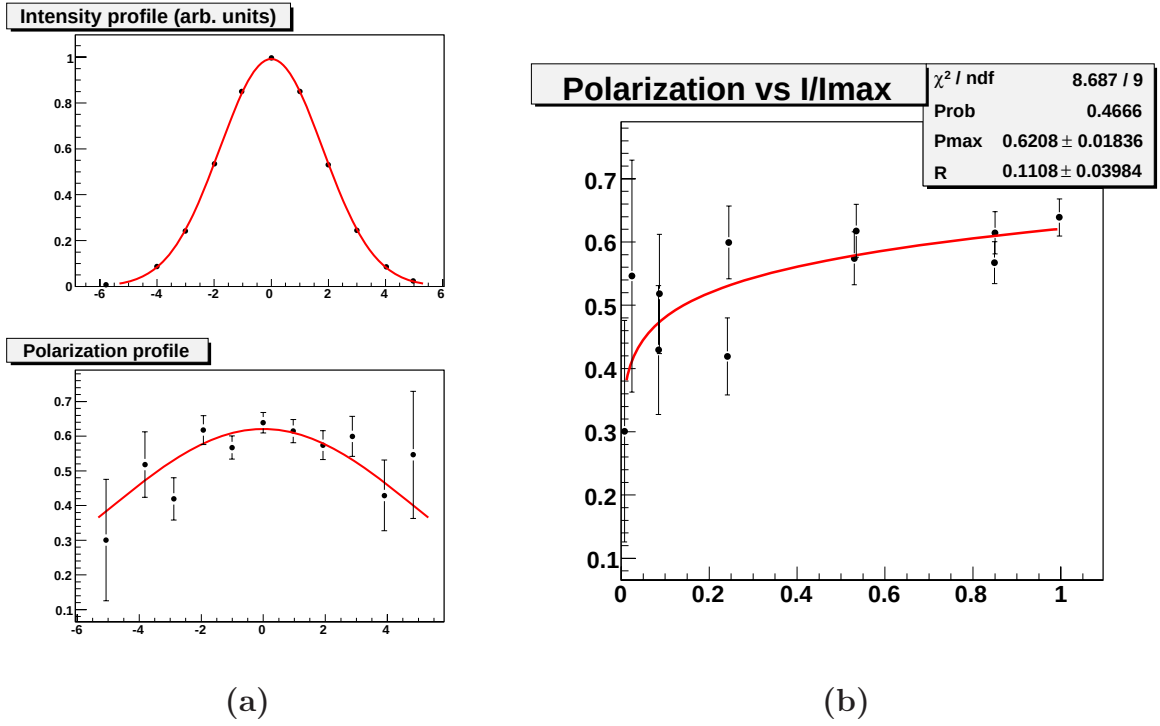


Figure 3: (a) Intensity (upper plot) and polarization (bottom plot) vs target position (arb. units), measured in one of the RHIC stores for 100 GeV beam; gaussian fits are shown. (b) Polarization vs intensity normalized to the intensity at beam center (I/I_{max}); fit of Eq. (6) to the data points is shown, with extracted parameters P_{max} and R .

4 Local Polarimeters

Local polarimeters employed by PHENIX and STAR also make a vital contribution to the measurements and monitoring of the beam polarizations at RHIC. They are sensitive to beam transverse polarizations. Their main role is to set up and to monitor the spin direction at collision. For the physics measurements with longitudinally polarized beams they are to confirm the absence of the polarization in transverse plane. In addition, they also proved to be a precise tool to monitor beam polarization bunch by bunch and versus time.

PHENIX local polarimetry is based on the analyzing power of the very forward neutron production discovered in the first polarized proton RHIC run of 2001-2002 [9]. Forward neutrons in PHENIX are measured by Zero-Degree Calorimeters (ZDC) with position sensitive Shower-Max Detectors which cover ± 2.8 mrad of the forward and backward directions [10]. The analyzing power increases nearly linearly with center of mass energy of colliding beams in the fixed ZDC geometry, which may indicate on p_T scaling of A_N (Fig. 4a, [11]), taking into account the x_F scaling of the forward neutron production cross section ([11]).

STAR local polarimetry is based on the sizable analyzing power of forward hadron production, first measured at $\sqrt{s} < 20$ GeV ([12]) and then shown to persist at RHIC energies, Fig. 4b ([13], [14]). Forward hadrons are detected by STAR Beam-Beam Counters (BBC), which are scintillator annuli mounted around the beam pipe. The inner rings of the BBCs, covering $\Delta\phi = 2\pi$ and $3.4 < |\eta| < 5.0$ are used for local polarimetry. In this geometry, the BBCs average spin asymmetry for hadrons in wide x_F range, so the

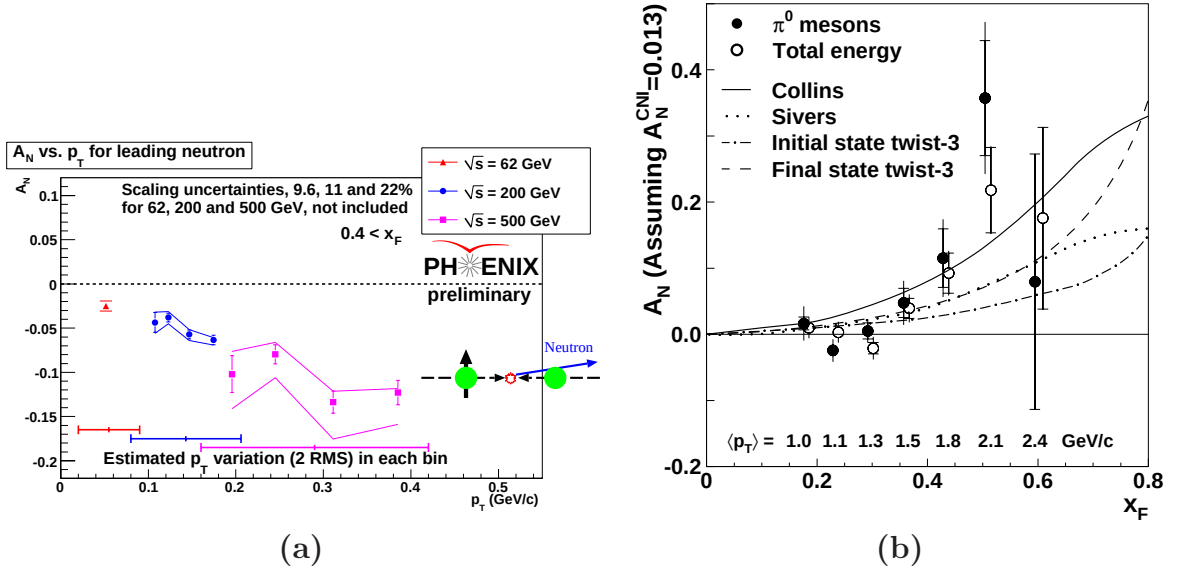


Figure 4: (a) The A_N for forward neutrons produced in polar angle < 2.8 mrad as a function of p_T ; data from PHENIX at $\sqrt{s} = 62, 200$ and 500 GeV are shown ([11]). (b) The A_N for π^0 mesons produced in the interval $3.3 < \eta < 4.1$ with polarized proton collisions at $\sqrt{s} = 200$ GeV from STAR ([13]).

effective analyzing power for the polarized pp scattering at $\sqrt{s} = 200$ GeV is of the order of 0.8% [15].

5 Summary

RHIC polarimetry consists of several complementary subsystems. Atomic hydrogen jet polarimeter provides the absolute polarization measurements averaged over the beam transverse profile, with limited statistical uncertainty on a store level ($\sim 10\%$). Proton-carbon polarimeters provide measurement of polarization profiles, which directly enters the average beam polarization observed in collisions at the interaction regions of RHIC experiments. After normalization to the absolute measurements by hydrogen jet, proton-carbon polarimeters provide store by store polarization measurements with good accuracy (better than 5%). Local polarimeter measurements serve to set up and monitor the spin direction at collision.

The pC polarimeter upgrade is under way, which will further improve polarization measurements at RHIC particularly in high event rate environment anticipated in future high luminosity RHIC runs [16, 17]. PHENIX and STAR collaboration are planning to add other forward detectors to local polarimetry task.

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DEVELOPMENT OF HIGH ENERGY DEUTERON POLARIMETER BASED ON dp -ELASTIC SCATTERING AT THE EXTRACTED BEAM OF NUCLOTRON-M.

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Abstract

The selection of the dp -elastic scattering events at the energy of 2 GeV using scintillation counters has been performed. The procedure of the CH_2-C subtraction has been established. This method can be used to develop the efficient high energy deuteron beam polarimetry.

The investigations with the use of high energy polarized deuteron beams proposed at different facilities require the efficient polarimetry at these energies to deduce values of polarization observables reliably. However, only some reactions can be used to provide efficient polarimetry of high energy deuteron. Since deuteron is spin-1 particle deuteron beam has vector and tensor polarizations. The polarimetry should have a capability to determine both components of polarization, if possible, simultaneously.

The polarimeters based on deuteron inclusive breakup with the proton emission at zero degree [1] and pp - quasi-elastic scattering [2] are currently used at LHEP-JINR Accelerator Complex to provide the polarimetry of the deuterons at high energies. But these analyzing reactions cannot be used for the simultaneous measurements of the tensor and vector polarizations of the beam, because deuteron inclusive breakup at zero degree and pp - quasi-elastic scattering have no vector and tensor analyzing powers, respectively. The dp -elastic scattering at backward angles has been successfully used for the deuteron polarimetry at RIKEN at a few hundreds of MeV [3]. This reaction has several advantages as a beam-line polarimetry over the others. Firstly, both vector and tensor analyzing powers have large values. Secondly, a kinematical coincidence measurement of deuteron and proton with simple plastic scintillation counters suffices for event identification. The same polarimeter for higher energies has been proposed [4], constructed and calibrated at 880 MeV [5] at Internal Target Station at Nuclotron.

The goal of present investigation is to study the possibility to use high energy dp -elastic scattering at forward angles with the selection by means of the kinematical coincidences of deuteron and proton with simple plastic scintillation counters for extracted beam polarimetry at Nuclotron.

The schematic view of the experiment at extracted beam of Nuclotron is shown in Fig. 1. The setup consists of two scintillation detectors based on FEU-85 photomultiplier tubes. Measurements were performed using deuteron beam of 1.6 and 2.0 GeV and polyethylene and carbon targets. The deuteron detector was placed at 8° in lab. The proton detector was placed in the kinematical coincidence.

The amplitudes of the signals and timing information from the both P and D detectors were recorded for each event. The distributions of the amplitudes for scattered deuterons and recoil protons for polyethylene target are presented in Fig. 2a and Fig. 2b, respectively. The correlation of the amplitudes and time difference are shown in Fig. 2c and Fig. 2d, respectively. One can see the clean correlation between the amplitudes from P and D detectors and well pronounced peak in time difference spectrum corresponding to the dp -elastic scattering events.

The selection of the dp -elastic scattering events was done by the applying of the graphical cut on the amplitudes correlation (see Fig. 3a). This cut allows to reduce significantly background in the time difference spectrum show in Fig. 3b.

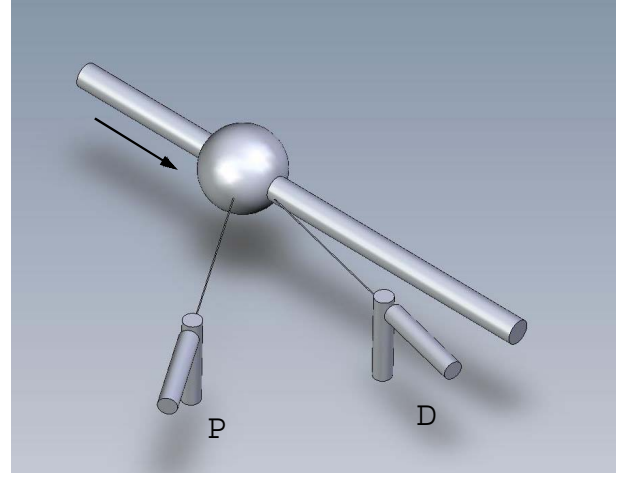


Figure 1: Schematic view of the experiment. P-proton detector, D-deuteron detector.

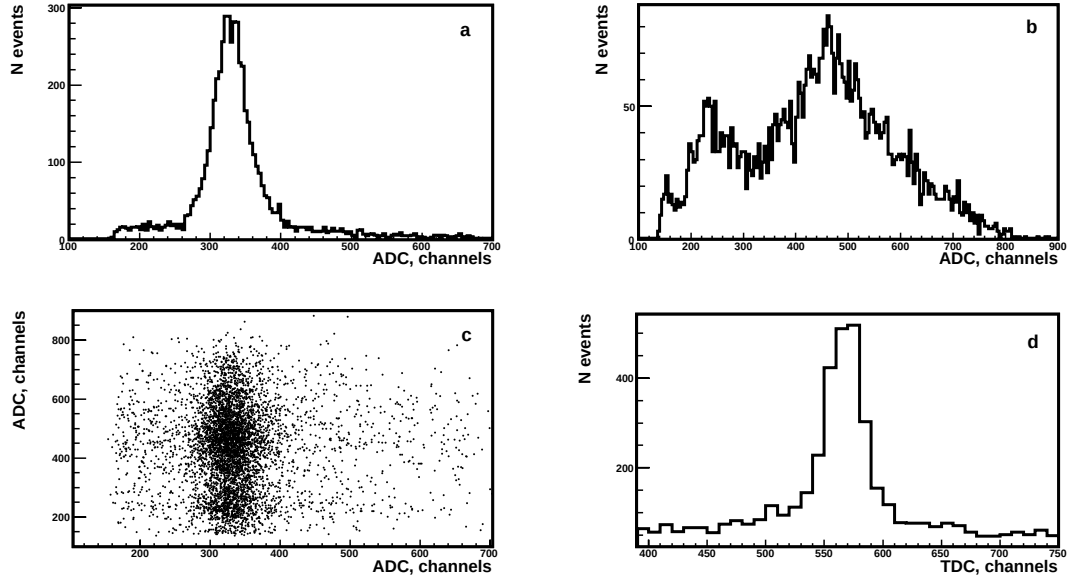


Figure 2: The results obtained at 2 GeV deuteron beam on polyethylene target for 8° deuteron scattering angle in lab: a) deuteron energy losses, b) proton energy losses, c) correlation of proton and deuteron energy losses, d) time difference between the signals for deuteron and proton detectors.

The measurements on carbon target were also performed to estimate the carbon contribution from polyethylene target. The corresponding distributions of the amplitudes for scattered deuterons and recoil protons, the correlation of the amplitudes and time difference are shown in Fig. 4. One can see that the distribution for recoil proton energy losses for $d + C$ interaction is much wider than for $d + CH_2$ scattering.

The time difference for P and D detectors for polyethylene and carbon targets after applying of the graphical cut on the amplitudes correlation are shown in Fig. 5a. The relative normalization of the spectra was obtained from the ratio of the background events

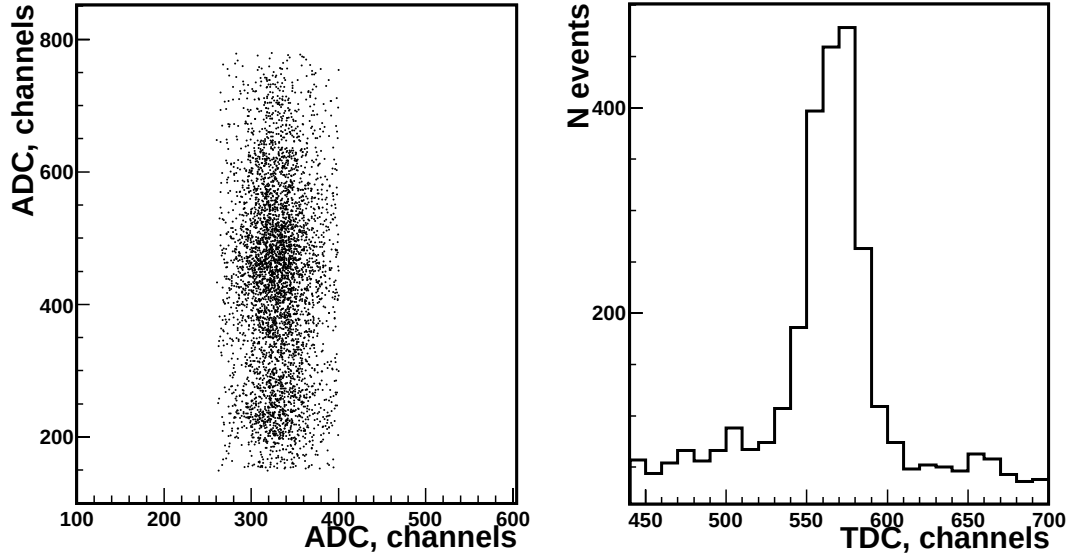


Figure 3: The selection of the dp -elastic scattering events by the energy losses correlation a), time difference after applying of the energy losses correlation cut b).

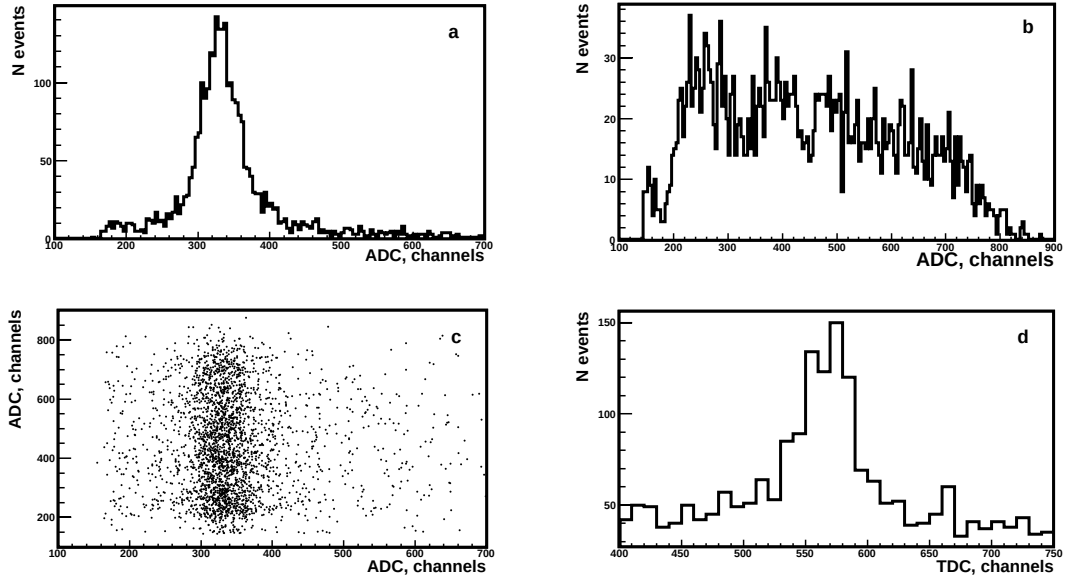


Figure 4: The results obtained at 2 GeV deuteron beam on carbon target for 8° deuteron scattering angle in lab: a) deuteron energy losses, b) proton energy losses, c) correlation of proton and deuteron energy losses, d) time difference between the signals for deuteron and proton detectors.

placed on the left and right from the peak. For this purpose the time difference spectra in Fig. 5a were fitted by the sum of gaussian and constant. The ratio of the obtained constants was considered as a normalization factor.

The results of the $CH_2 - C$ subtraction of the time difference spectra are presented in Fig. 5b. The lines correspond to the prompt time windows for dp -elastic events. The background placed outside of the window is about 3% from the peak height.

The following results has been obtained: The possibility of the dp-elastic scattering events selection at high energies and small scattering angles using scintillation counters techniques has been demonstrated. The procedure of the $CH_2 - C$ subtraction has been established. This method can be used to develop the efficient high energy deuteron beam polarimetry.

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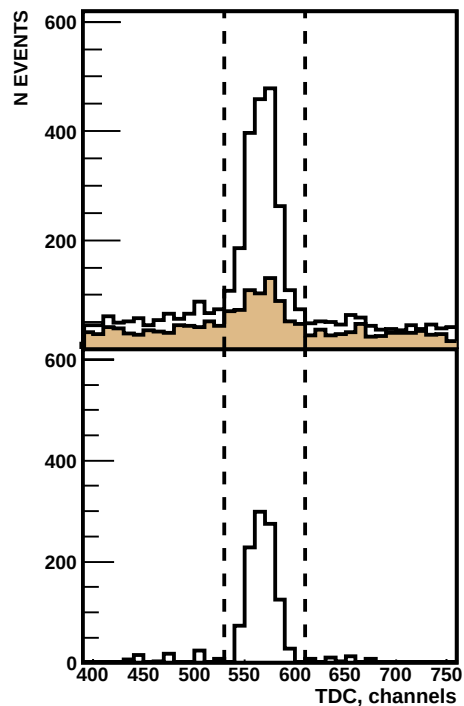


Figure 5: $CH_2 - C$ subtraction: a) time difference spectra obtained on polyethylene (white) and carbon (gray) targets after applying graphical cut, b) dp -elastic events.

SPIN-ISOTOPIC ANALYSIS AT SUPERLOW TEMPERATURES

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Abstract

Free electrons produced by beam irradiation of any dielectrics, are able to provide studying its nuclear surroundings. At superlow temperatures, the nuclear signals can be enhanced due to electron-nuclear dipolar interactions by the Dynamic Nuclear Polarization method. Electron spins induce the Fermi indirect interactions between nuclear spins in surrounding molecules. This allows one to obtain the "squeezed" image of quadrupole nuclei in the proton spectra. Herein we discuss the application of this method for early cancer detection using the irradiated samples of the human blood. The method is demonstrated on the simplest irradiated LiD and NH₃ materials [1] investigated by SMC and COMPASS collaborations at CERN.

1. Introduction.

The spin-isotopic analysis of biological structures is a new endeavour for early cancer detection applying the technique of strong magnetic fields, superlow temperatures, and the dynamic nuclear polarization (DNP). Unlike the "warm" nuclear magnetic resonance (NMR) at room temperatures, the proposed method operates with hundreds-fold DNP-enlarged signals. DNP requires the presence of paramagnetic dopant (F-centres) that can be induced by a human blood sample exposure to the x-ray, electron, or other ionizing radiation [2]. As compared with the chemical analysis, the method distinguishes isotopes, allows one to detect the routine NMR spectra and the quadrupole spectra called by us "squeezed". In the case considered, we use the data taken after completing DNP-process and cooling the samples below 0.1 K. It is difficult to judge and compare the results because there are no experiments having applied this advanced tool in biological studies.

2. Isotopic analysis at superlow temperatures.

The isotope density (N_i) is proportional to the integral of the NMR-signal absorption part $v_i(\omega)$. The relative densities of two spin species (N_i/N_j) are as follows [3]:

$$\frac{N_i}{N_j} = \frac{I_j \gamma_j^2 B(\gamma_j h H_0 / k_B T_j)}{I_i \gamma_i^2 B(\gamma_i h H_0 / k_B T_i)} \frac{\int_0^\infty v_i(\omega) d\omega}{\int_0^\infty v_j(\omega) d\omega}, \quad (1)$$

where h is the Plank constant, indices i and j stand for the two species, $B(i, j)$, $I_{i,j}$, $\gamma_{i,j}$ are the Brillouin function, the spins, the gyromagnetic ratios, respectively.

At superlow temperatures, the nuclear spin lattice relaxation times reach days and for precise measurements it is convenient to keep the fixed tuning of the spectrometer varying the magnetic field to any isotope resonances. Typical measurements are illustrated with LiD material having the cubic crystalline lattice. In a perfect cubic crystal the quadrupole

interaction vanishes and a single Zeeman line is observed. If the NMR spectrometer is tuned to the Larmor frequency $\nu_c = \gamma_{Li}H_{Li} = \gamma_D H_D = 16.38$ MHz, then the D, ^6Li and ^7Li isotopes are detected at the resonant fields listed in the Table. Fig. 1a [4] shows the D and ^7Li spectra taken by frequency scanning across the ν_c . The spin-isotopic composition obtained by NMR method was found in good agreement with the data of the mass-spectrometric analysis. It was shown that for the most polarized dielectrics, the

Material	Isotopes	Spin	γ (Hz/G)	Field (G)	ν_c MHz	Quadr. mom $10^{-26}(\text{cm}^2)$
LiD	D	1	653.6	25060	16.38	+0.287
LiD	^6Li	1	626.6	26141	16.38	+0.05
LiD	^7Li	3/2	1654.8	9898	16.38	-1.2 ?
NH_3	^1H	1/2	4257.7	25060	106.7	0
NH_3	^{14}N	1	307.8	25060	7.713	+1.56
Potassium	^{39}K	3/2	198.9	25060	4.984	+14
Sodium	^{23}Na	3/2	1126.2	25060	28.22	+11

DNP mechanism ensured the equal spin temperature (EST) for all nuclear species (i.e. $T_i = T_j$ in Eq.(1)). At these conditions Brillouin functions in Eq.(1) are reduced and the error of $\simeq 2\%$ comes mainly from the accuracy of integration.

The spectrum of nitrogen (Fig. 1b) [5] in irradiated ammonia (NH_3) illustrates the problems of detection of ^{14}N , ^{39}K , ^{23}N and similar quadrupole nuclei in the amorphous biological structures. ^{14}N -spin ($I_n = 1$) has a small gyromagnetic ratio, a large quadrupole moment (see Table) and the broadened spectrum due to strong quadrupole interactions with electric field gradients in the lattice. The ability of the spectrometer to detect NMR spectrum requires the signal width to be a small fraction of the Larmor frequency. At the fixed central frequency of the spectrometer this condition is fulfilled only for the two small pieces of the ^{14}N spectrum (hashed in Fig. 1b). As a result, the total line shape (Fig. 1b) needs to be reconstructed from these spectral pieces that actual spectra are shown in Fig. 1c. It is clear that the direct integration by Eq.(1) becomes doubtful.

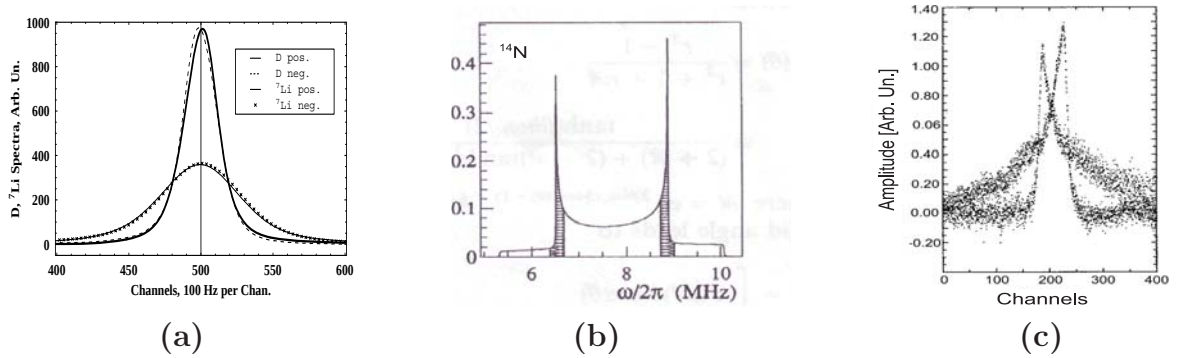


Figure 1a. D and ^7Li line shapes at positive and negative spin temperatures. Spectra of the negative polarization were inverted and aligned with the positive spectra; $P_D = \pm 39, 2\%$.

Figure 1b. The reconstructed signal of ^{14}N at about +10% polarization. The hashed areas represent the detected regions and the actual spectra are shown in Fig. 1(c).

Fig. 1c. Two signals at +10% polarization detected in the hashed regions in Fig. 1b.

3. Squeezing of quadrupole spectra.

The electric field gradients and isotopic composition have to identify any blood sample for a noninvasive early cancer detection. In fact, at low temperatures any molecular movements are stopped and the electrical interactions reveal themselves only over the electric field gradients determined by the specific broadening of the spectral line shape. Magnetic interactions and, as we have seen, the isotopic composition are also withdrawn from the spectral analysis. In general they would have been understood as the complete set of parameters for the blood control for the early cancer detection, if the signal amplitudes did not reduce so much due to the quadrupole broadening of their spectra (see Fig. 1c). This broadening comes from random angle (θ) orientations of the electric field gradients relatively to the magnetic field direction. The eigen-states of ^{14}N spins in NH_3 are [6]

$$E_n = -h\nu_c m_n + h\nu_Q \{3 \cos^2(\theta) - 1\} \{3(m_n)^2 - I(I+1)\} , \quad (2)$$

where ν_c is the nitrogen Larmor frequency, eq is the quadrupole moment, $\nu_Q = (e^2 q Q / h) / 8$ and eQ is an electric field gradient value, $m_n = (+1, 0, -1)$ and $m_p = (+1/2, -1/2)$ are the nitrogen and proton (see below) magnetic quantum numbers.

The problem can be solved if the wave function of free electrons couples the nuclear spins situated nearby to F-centers with scalar Fermi interaction [7]. In this case the removed and nearest spins have energies E_d and E_F , correspondingly:

$$E_d = -h\gamma_p(H_0 + H_{loc})m_p, \quad E_F = -h\gamma_p H_0 m_p + hJ \cdot m_n m_p . \quad (3)$$

The first of Eq.(3) describes the energy of remote spins situated far away from F-centres. Their Larmor frequency equals to $\nu_0 = \gamma_p H_0$ and it is shifted by local field (H_{loc}) of polarized protons to ν_+ or ν_- depending on the sign of H_{loc} (see the left and middle-hand diagrams in Fig. 2a). The second Eq.(3) approximately describes the energy of the nearest spins (the right-hand diagram in Fig. 2a). These protons are coupled with ^{14}N spins by the Fermi indirect interaction. J -constant of the Fermi interaction is determined by the density of the electron wave function inside nuclei [8]. The total effect from heterogeneous magnetization is demonstrated in Fig. 2b [5] which shows the proton NMR signals in

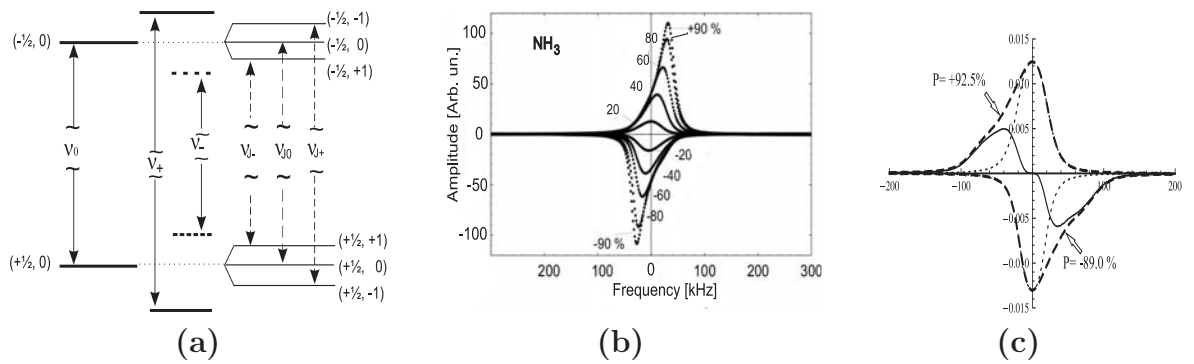


Figure 2a. (m_p, m_n) -sublevels in NH_3 : ν_0 -Larmor proton frequency, ν_+ , ν_- are the same frequency at opposite signs of polarizations, ν_{J+} , ν_{J0} and ν_{J-} are J -coupling transitions.

Figure 2b. Proton NMR line shapes for different values of polarizations.

Figure 2c. Proton lines in NH_3 centred at the tops. They are subdivided into the symmetrical dipolar part and asymmetrical fractions owing to J -coupling between H_3 and N spins.

an amorphous NH_3 for low and high, positive and negative nuclear polarizations (spin temperatures). Since magnetic transitions obey the $\Delta(m_p + m_n) = \pm 1$ selection rule, for protons $\Delta(m_p) = \pm 1$, then $\Delta(m_n) = 0$ in Eq.(3). At parallel orientation ($m_p = 1/2$, $m_n = 1$ or $m_p = -1/2$, $m_n = -1$), the proton energy will be increased by amount $hJ_{nm}/2$; for antiparallel ($m_p = -1/2$, $m_n = 1$ or $m_p = 1/2$, $m_n = -1$) it will be decreased by the same amount. In the order of increasing frequency of the spectrometer, the positive spectra will include ν_{J-} , ν_{J0} , ν_{J+} Fermi transitions and the ν_+ transition. For negative spectra we have ν_- and ν_{J-} , ν_{J0} , ν_{J+} transitions, respectively. Since the remote spins interact only by the dipolar interaction which is independent of the spin permutations, they generate the symmetrical part of a spectra i.e. ν_+ and ν_- transitions in Fig. 2a. The proton spins surrounding F-centres are undergone by shorter-acting and, hence, the stronger J -interaction [8] which produces the residual i.e. asymmetrical part of the line. This logic allows one to explain the left-hand asymmetry of the spectral line at the positive and the right-hand asymmetry at the negative polarizations. It is clear that the asymmetrical parts of signals, shown with the solid lines in Fig. 2c, are the squeezed images of nitrogen spins into the proton line shape. Using Eq. (1), the Table data and the spectral bandwidths in Fig. 1b and Fig. 2c it isn't difficult to estimate the amplitude enhancement given by the method:

$$\frac{A_{NH}}{A_N} = \xi \frac{\Delta_N}{\Delta_H} \left(\int_0^\infty v_H d\omega / \int_0^\infty v_N d\omega \right) \approx 0.3 \frac{2.3 \text{ MHz}}{0.1 \text{ MHz}} \cdot (2150) \approx 1.5 \cdot 10^4, \quad (4)$$

where $\xi \approx 0.3$ is the asymmetrical contribution in the proton spectrum estimated from Fig. 2c, Δ_N/Δ_H is the ratio of nitrogen and proton bandwidths estimated from Fig. 1b and Fig. 2c and the ratio in the brackets was estimated by Eq. (1) for the proton of 90% ($B_H \approx 0.9$) and the nitrogen of 11% ($B_N \approx 0.11$) polarizations both having about the equal temperatures. The effect can be visualized by a comparison between the noisy routine spectra in Fig. 1c and the noiseless solid curves in Fig. 2c obtained by our method.

The above consideration has shown that the relative isotope densities of different quadrupole nuclei in frozen amorphous biological samples can be compared by their squeezed images in the proton spectrum. The method results in the better accuracy for isotopic comparison between the normal and cancerous blood samples due to large amplitude enhancement and the detection at fixed tuning of the NMR-spectrometer.

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Abstract

In papers [1–4] the new technique of spin resonance crossing has been offered, which essentially expands opportunities of well-known methods for fast or adiabatic crossing. Technique of transparent spin resonance crossing takes into account interference of spin precession phases inside resonance region. In this paper the received results are summarized with the help of spin precession axis and generalized spin tune concept.

1 The spin motion description under stationary conditions

The description of spin motion in cyclic accelerators and storage rings essentially does not differ from the orbital beam motion one. Under stationary conditions the main characteristics of orbital motion are canonically conjugated variables of actions and phases. Particles acceleration relates to nonstationary conditions and requires the special description of beam polarization behavior which is similar to the description of orbital motion. Therefore it is natural to study first spin motion under stationary conditions and then to investigate nonstationary conditions.

At nonequilibrium orbits the spin field is periodic function of all particle oscillations phases near a closed orbit under stationary conditions

$$\vec{W}(\theta, \Psi_i, I_i) = \vec{W}(\theta + 2\pi, \Psi_i + 2\pi, I_i),$$

where θ is a generalized azimuth, i.e. length along the closed orbit in terms of its radius, I_i, Ψ_i are action-phase variables of particle orbital motion near closed orbit in accelerator. It is shown, that for multifrequency system there is a precession axis with similar spin field properties of periodicity [5, 6]: $\vec{n}(\theta, \Psi_i, I_i) = \vec{n}(\theta + 2\pi, \Psi_i + 2\pi, I_i)$,

The precession axis $\vec{n}$ is the solution of the Thomas-BMT equation

$$d\vec{n}/d\theta \equiv \vec{n}' = [\vec{W} \times \vec{n}]$$

The generalized spin tune describing spin motion in a perpendicular to the $\vec{n}$ axis plane is determined by the general expression

$$\nu = \vec{W} \cdot \vec{n} - \vec{\ell}_2 \cdot \vec{\ell}_1' \quad (1)$$

and depends on a choice of $\vec{\ell}_1$ and $\vec{\ell}_2$ orts in a perpendicular to the $\vec{n}$ axis plane. It is shown, that $\vec{\ell}_1(\theta, \Psi_i, I_i)$ and $\vec{\ell}_2(\theta, \Psi_i, I_i)$ orts can be chosen with similar to field properties of periodic dependence phases Ψ_i and azimuth θ . Moreover, there is a “natural” system of orts $\{\vec{\ell}_1, \vec{\ell}_2, \vec{n}\}$ when the generalized spin tune doesn’t depend on phases Ψ_i and azimuth θ and remains constant at each particles trajectory: $\nu = \nu(I_i) = \text{const}$.

The spin, directed along $\vec{n}$ axis, is stable to small field deviations. On the contrary, motion of perpendicular to an $\vec{n}$ axis spin components is unstable to small field deviations which cause a small deviation of spin tune ν .

In natural system of orts $\{\vec{\ell}_1, \vec{\ell}_2, \vec{n}\}$, spin rotates in a constant field:

$$\vec{h}(I_i) = \nu(I_i) \vec{n}.$$

The generalized spin tune which depends on periodic orts $\vec{\ell}_1$ and $\vec{\ell}_2$ can be changed for an integer combination from orbital motion tunes $\nu_k = \Psi'_k = \sum k^i \nu_i$. The generalized spin tune is changed for value ν_k ($\nu \rightarrow \nu - \nu_k$) if periodic orts $\vec{\ell}_1 + i \vec{\ell}_2$ are replaced by periodic orts $(\vec{\ell}_1 + i \vec{\ell}_2) \exp(i \Psi_k)$.

Under stationary conditions the investigation of beam polarization behavior in the accelerator is reduced to search of “natural” system orts which can be found, for example, by numerical methods.

The vector of polarization $\vec{P}$ is determined by average meaning of spin $\vec{S}$ over particles distribution and is a function of azimuth θ :

$$\vec{P} = \langle \vec{S} \rangle = \left\langle J \vec{n} + \sqrt{1 - J^2} \operatorname{Re} \left[(\vec{\ell}_1 + i \vec{\ell}_2) \exp(-i(\nu\theta + \alpha)) \right] \right\rangle,$$

where $J = \vec{S} \cdot \vec{n} = \text{const}$ is a spin variable of action or spin adiabatic invariant which is kept under stationary conditions. In this formula the spin is normalized to unit. Averaging over the distribution of particles under stationary conditions is reduced to independent averaging over all orbital phases Ψ_i and orbital actions I_i . Because of spin tune $\nu(I_i)$ and orbital tunes Ψ'_i spreading the beam polarization under stationary conditions relaxes to following value: ¹

$$\vec{P} = \langle J \rangle \langle \vec{n} \rangle. \quad (2)$$

The greatest possible degree of polarization is achieved at the maximum value of spin adiabatic invariant, i.e. when $J = \pm 1$ (for example, for a particle with 1/2 spin the adiabatic invariant in dimensional units is equal $\pm \hbar/2$). In this case $\vec{P}_{\max} = \pm \langle \vec{n} \rangle$. The deviation of $|\vec{P}_{\max}|$ value from unit is connected with spread of precession axes directions at particles trajectories. This “dynamical” beam depolarization is reversible. The beam polarization can be restored by means of the organization of an appropriate spin field in the accelerator. The $\vec{P}_{\max}$ value depends on azimuth θ and in a place of carrying out of experiment can be equaled to unit. For example, in an opposite section of the accelerator with one Siberian Snake precession axes of all particles have a longitudinal direction ($|\vec{P}_{\max}(0)| = 1$). In other places of an orbit there is a spread of precession axes and consequently in these places there is the reduction of a degree of polarization ($|\vec{P}_{\max}| < 1$).

¹Here the effects connected with loss of beam polarization at injection in accelerator are not considered.

2 The spin motion description under nonstationary conditions

The special analysis is required to determine beam polarization behavior under nonstationary conditions. Thus, in the traditional accelerator the spin field $\vec{W}$ depends on particles energy and is changed during acceleration. Orbital motion of particles is changed during acceleration too. However, a wide class of problems can be solved in adiabatic approximation when the “nonstationary” parameter is changed relatively little. For example, the energy deviation of a particle per a turn in the accelerator is rather small. In this approximation action variables of orbital motion remain constant during all cycle of acceleration. It is necessary to pay attention only to small region of “critical” energy during acceleration in which orbital action variables I_i can be changed and synchrotron tune is tending to zero.

Dependence of a spin field $\vec{W}$ on time under nonstationary conditions can be described as a function from additional parameter $\lambda(\theta)$ which is slowly changing during acceleration: $\vec{W}(\theta, \Psi_i, I_i, \lambda)$.

The beam polarization behavior is determined by means of solution for an precession axis and the generalized tune which are found under stationary conditions. Procedure is the following: if $\lambda = \text{const}$ a spin field $\vec{W}$, as well as under stationary conditions, is a periodic function of all phases Ψ_i and azimuth θ . We find spin bases $\{\vec{\ell}_1(\theta, \Psi_i, I_i, \lambda), \vec{\ell}_2(\theta, \Psi_i, I_i, \lambda), \vec{n}(\theta, \Psi_i, I_i, \lambda)\}$ with periodic properties of each phase and azimuth for each value of λ . Under this conditions the spin action variable J and the generalized spin tune ν is a function of λ too: $J = J(\lambda)$, $\nu = \nu(I_i, \lambda)$.

From definition of the generalized spin tune the basic statement follows: tune ν , as well as natural system of orts $\{\vec{\ell}_1, \vec{\ell}_2, \vec{n}\}$, is a continuous function of λ parameter. That follows from continuous dependence of a field $\vec{W}$ on λ parameter. “Points” of spin resonances mean only, that near these points spin motion is very sensitive to spin field value on particles orbits.

The spin action variable J remains constant during the changing of $\lambda(\theta)$ in adiabatic approximation, if the speed of λ changing is small enough i.e. spin field changing is relatively small during characteristic time of the precession axis changing. Characteristic time of $\vec{n}$ changing is determined by the nearest to ν combination from orbital motion tunes ν_k , i.e. by the nearest spin resonance. Therefore, the condition of a spin action variable J conservation while $\lambda(\theta)$ is changing, will be the following:

$$\left| \frac{d\lambda}{d\theta} \frac{\partial}{\partial \lambda} [(\nu - \nu_k) \vec{n}] \right| \ll (\nu - \nu_k)^2, \quad (3)$$

where $\nu_k = \sum k^i \nu_i$ is the nearest combination of orbital motion tunes. In this case spin, rotating around $\vec{n}$ axis, “has time” to follow up $\vec{n}(\lambda)$ direction changing and the vector of polarization is still determined by an average precession axis over distribution of particles in a beam according to the formula (2).

The condition (3) indicates that there are only small regions of parameter in which condition of a spin action variable conservation can be violated. These are the regions of spin resonances of rather small strength.

3 The conditions for transparent spin resonance crossing

Taking into account the above conception one can get a complete compensation of depolarization degree at resonance crossing. The conditions for transparent crossing are given by

$$J(\lambda_-) = \pm J(\lambda_+). \quad (4)$$

where λ_- and λ_+ are λ values before ($\theta = \theta_-$) and after ($\theta = \theta_+$) resonance crossing.

Let us consider the simple example of circular accelerator with only one Fourier harmonic of the field $\vec{W}$ due to vertical betatron oscillations for the purposes of illustration:

$$\vec{W} = \nu_0 \vec{e}_z + w \vec{e}_1(\Psi_z) \quad (5)$$

where $\vec{e}_1(\Psi_z) = \vec{e}_x \cos \Psi_z + \vec{e}_y \sin \Psi_z$, Ψ_z is the phase, $\nu_z = \Psi'_z$ is the vertical betatron tune, $\vec{e}_x$, $\vec{e}_y$, $\vec{e}_z$ denote the corresponding radial, longitudinal and vertical orts of “accelerator” basis, $\nu_0 = \gamma G = \gamma(g - 2)/2$ is spin tune in ideal accelerator at $w = 0$.

Thus, the “natural” spin basis is:

$$\vec{\ell}_1 = -\frac{w}{h} \vec{e}_z + \frac{\varepsilon}{h} \vec{e}_1(\Psi_z), \quad \vec{\ell}_2 = [\vec{n} \times \vec{\ell}_1] = \vec{e}_y \cos \Psi_z - \vec{e}_x \sin \Psi_z, \quad \vec{n}(\Psi_z) = \frac{\varepsilon}{h} \vec{e}_z + \frac{w}{h} \vec{e}_1(\Psi_z),$$

here are $h = \sqrt{\varepsilon^2 + w^2}$, $\varepsilon = \nu_0 - \nu_z$ — spin detuning.

The general spin tune ν is equal to h in this natural basis $\{\vec{\ell}_1, \vec{\ell}_2, \vec{n}\}$, near the isolated spin resonance.

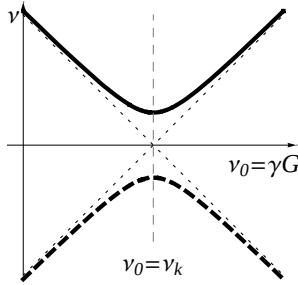


Figure 1: The general spin tune dependence vs tune ν_0 (vs the energy)

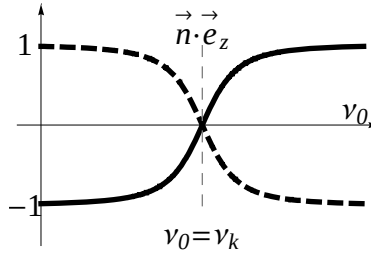


Figure 2: The vertical component of precession axis plotted vs the tune ν_0

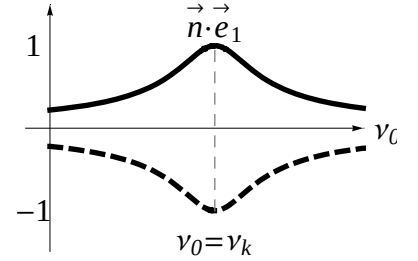


Figure 3: The transverse component of precession axis plotted vs the tune ν_0

The dependencies of general spin tune as well as vertical and transverse $\vec{n}$ axis projections on “energy” are shown in Fig. 1-3. The mirror-reflected decision $\{\vec{n}, \nu\} \rightarrow \{-\vec{n}, -\nu\}$ is shown by a dashed line.

The variable λ coincide with the $\nu_0 = \gamma G$ value at the particles accelerations. The condition (3) can be fulfilled at any moment of adiabatic crossing if $(\nu - \nu_k = h, w = \text{const})$

$$|d\nu_0/d\theta| \ll w^2. \quad (6)$$

The condition (6) can be violated for very small resonance strength (at $w \rightarrow 0$). Spin action variable $J(\lambda)$ can be changed only in the resonance region. The beam’s polarization is decrease in this case. The condition for transparent crossing is the restoration of the spin action variable value after spin resonance crossing.

4 Examples of transparent crossing

Examples of transparent resonance crossing is given in papers [1–3]. The condition (3) can be fulfilled only if detuning changed in spin resonance region. The SAI dependence is shown schematically in Fig. 4 and Fig. 5. SAI changing is plotted symbolic by the dashed line inside spin resonance region. The detuning shape defines SAI changing in resonance region.

When condition (3) is violated the width of resonance region is defined by crossing speed and is equal approximately to $\theta_{\text{res}} \sim 1/\sqrt{d\varepsilon/d\theta}$.

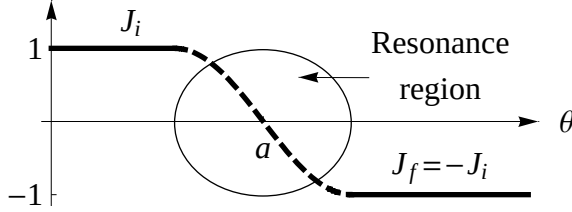


Figure 4: Transparent resonance crossing in the case of SAI-flip

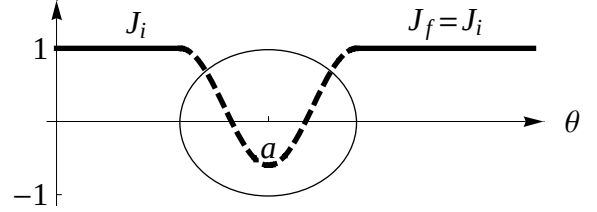


Figure 5: Transparent resonance crossing in the case of SAI keeps sign

Another possibility to restore SAI appears, if one uses multiple resonance crossing [4]. At single crossing with constant speed $\varepsilon' = d\varepsilon/d\theta$ the SAI values before and after resonance crossing are connected by general FS formulae [7]:

$$J_f = \left(1 - 2 \exp\left(-\frac{w^2}{2|\varepsilon'_a|}\right)\right) J_i = \cos 2\alpha J_i, \quad \sin \alpha = \exp\left(-\frac{w^2}{4|\varepsilon'_a|}\right)$$

where angle α is defined by speed of detuning changing at resonance crossing moment $\theta = a$. It should be noted that it is impossible to restore SAI after single crossing with constant speed. The restoration of SAI after resonance crossing can be obtained due to detuning of changing in outside resonance region.

The SAI dependence is shown schematically in Fig. 6 and Fig. 7 for the cases of two times transparent crossing.

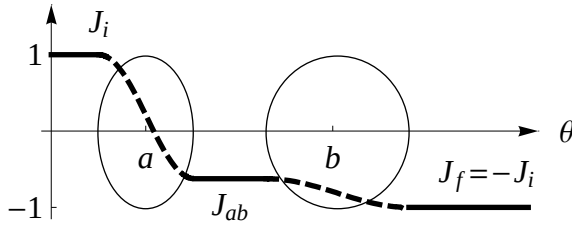


Figure 6: Two times transparent resonance crossing in the case of SAI-flip

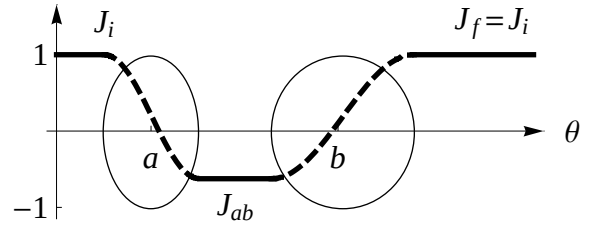


Figure 7: Two times transparent resonance crossing in the case of SAI keeps sign

Two times transparent crossing can be used for spin-flip equipment in the accelerator ring directly before beam is extracted to external target.

The conditions of two times transparent resonance crossing are the following:

- a) in the case of SAI-flip (Fig. 6): $\alpha = \beta, \quad \Psi_{ab} = 2\pi m;$
- b) in the case of SAI keeps sign (Fig. 7): $\alpha = \pi/2 - \beta, \quad \Psi_{ab} = 2\pi m + \pi,$

where angles α and β are defined by the speed of detuning changing at resonance crossing moments $\theta = a$ and $\theta = b$, the spin precession phase is $\Psi_{ab} = \int_a^b h d\theta + \varphi_a + \varphi_b$ on area $a < \theta < b$ is defined by general spin tune and phase shift φ_a and φ_b at moments of resonance crossing.

The SAI dependence is shown schematically in Fig. 8 and Fig. 9 for three times transparent crossing.

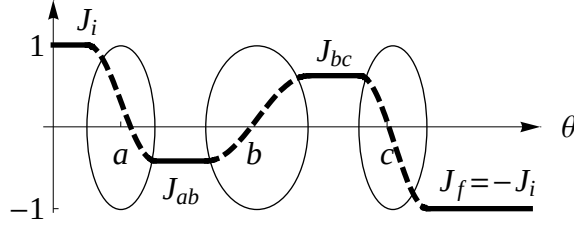


Figure 8: Three times transparent resonance crossing in the case of SAI-flip

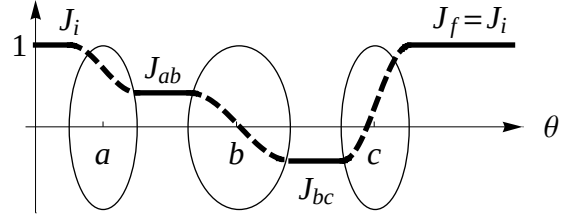


Figure 9: Three times transparent resonance crossing in the case of SAI keeps sign

Let us write the conditions of SAI restoration for symmetrical three times crossing, when phase Ψ_{ab} is equal to phase Ψ_{bc} and crossing speeds at moments a and c are the same. For symmetrical three times crossing the transparent condition is the following:

- a) in the case of SAI-flip (Fig. 8): $\Psi_{ab} = 2\pi m, \quad \beta = 2\alpha - \pi/2, (\alpha > \pi/4),$
 $\Psi_{ab} = 2\pi m + \pi, \quad \beta = \pi/2 - 2\alpha, (\alpha < \pi/4);$
- b) in the case of SAI keeps sign (Fig. 9): $\cos \Psi_{ab} = \frac{\text{tg}\beta}{\text{tg}2\alpha}.$

The first experimental test of this technique to overcome spin resonance at COSY accelerator has shown the potential benefit of transparent crossing [8].

Note that in a real accelerator the orbital parameters spread and slow synchrotron modulation of the particle phase and energy, the problems of crossing nonisolated resonances should be taken into account.

We thank participants of SPIN@COSY teleconference for useful discussions of some questions, considered in this paper.

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DEUTERON BEAM POLARIMETRY AT THE INTERNAL TARGET STATION AT NUCLOTRON-M AT GeV ENERGIES

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Abstract

A polarimeter based on the asymmetry measurement of the deuteron-proton elastic scattering has been constructed at Internal Target Station of Nuclotron-M(JINR). It allows to measure vector and tensor components of the deuteron polarization simultaneously. We have measured also the analyzing powers A_y , A_{yy} and A_{xx} in the $d-p$ elastic scattering at 880 and 2000 MeV. The analyzing powers values at angles in the c.m. are large enough to provide the efficient polarimetry at these energies.

Measurement of the beam polarization is an important element in the experiments on the spin physics studies. The deuteron beam has vector and tensor polarization in general. As far as possible it is necessary to conduct the simultaneous measurement both components of the deuteron beam polarization.

The polarimetry based on zero-degree inclusive deuteron breakup from a proton target and pp quasi-elastic scattering are currently used at LHEP Accelerator Complex to provide the polarimetry of the deuterons at GeV-energies. These polarimeters do not fit the requirement of the vector-tensor mixed polarimetry, because the first and the second one have no vector and tensor analyzing powers, respectively.

On the other hand, $d-p$ elastic scattering at forward scattering angles is traditionally used for the tensor and vector polarimetry at intermediate energies. The polarization of the deuteron beam was measured using the $d-p$ elastic scattering at 1600 MeV at

forward angles, where vector A_y and tensor A_{yy} analyzing powers have large values with two-arm magnetic spectrometer ALPHA[1]. It was demonstrated recently at COSY at ANKE spectrometer at the initial deuteron energy of 1170 MeV that the vector and tensor analyzing powers for $d-p$ elastic scattering are also large. $d-p$ elastic scattering at large angles has been successfully used for deuteron beam polarimetry at RIKEN[2].

The aim of the experiment is to obtain the A_y , A_{yy} , and A_{xx} analyzing powers in $d-p$ elastic scattering at large angles at the 270–2000 MeV energy range to study the possibility to use this reaction for the efficient polarimetry of deuterons in the wide energy range.

The analyzing powers data at the energies in a GeV region are necessary for the DSS project at Nuclotron–M. New facility RIBF at RIKEN will have polarized deuterons at 880 MeV and, therefore, also requires the development of the high energy polarimetry.

The polarimeter was constructed at the Internal Target Station(ITS) of Nuclotron-M.

A schematic view of the polarimeter is shown in Fig.1. A detector support to install deuteron and proton detectors is placed downstream of the ITS. Each plastic scintillation counters coupled to a photo-multiplier tube Hamamatsu H7416MOD is arranged on the left, right, up and down of the beam axis to allow determination of the spin-dependent asymmetry of the reaction. A detector for quasielastic- pp scattering is used as the monitor of beam intensity.

Analyzing powers for the $d-p$ elastic scattering at the energies 880 and 2000 MeV were measured with this newly-constructed polarimeter in June 2005.

Polarized deuterons were provided by PIS("polarized ion source") POLARIS. In the present experiment the data were taken for unpolarized, "2-6" and "3-5" spin modes of POLARIS which had the following theoretical maximum polarization $(p_z, p_{zz})=(0,0), (1/3,1)$ and $(1/3,-1)$, respectively. The quantization axis is perpendicular to the beam circulation plane of Nuclotron. The spin states of POLARIS were changed spill by spill.

The typical beam intensity in Nuclotron ring was $2-3 \times 10^7$ deuterons per spill. The $10\mu\text{m}$ CH_2 foil was used as a proton target, while the measurements with carbon target were made to estimate the background originating from quasifree scattering on carbon. Prior to the analyzing powers measurement at 880 and 2000 MeV, the beam polarization measurement was carried out at 270 MeV, where well established data on the analyzing powers exist [2]-[5].

The beam polarization was determined by using asymmetry of the $d-p$ elastic scattering yields and known analyzing powers of the reaction at 270 MeV[2,4,5]. Data for several scattering angles were used in the polarization determination: 75° , 86.5° , 95° , 105° , 115° , 126.3° , and 135° . The analyzing powers values A_y , A_{yy} , A_{xx} and A_{xz} and their errors at these angles were obtained by the cubic spline interpolation of the data from Refs.[2,4,5].

The values of the tensor p_{yy} and vector p_y polarizations of the beam for "2-6" and "3-5" spin modes of POLARIS obtained at different angles are presented in Fig.2. The

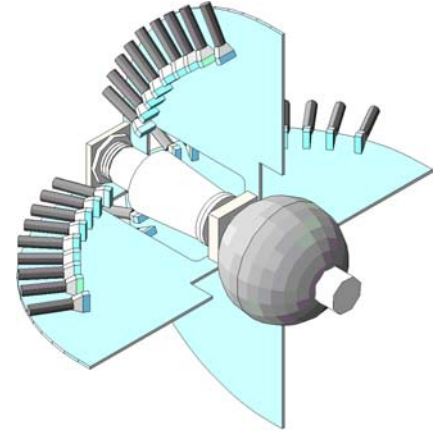


Figure 1: Schematic view of the detector setup.

error bars include both statistical and systematical uncertainties. The weighted average values of the tensor polarizations obtained at 270 MeV are 0.605 ± 0.025 and -0.575 ± 0.02 for the spin modes "2-6" and "3-5", respectively. They are in a good agreement with the values obtained by the low energy polarimeter based on ${}^3\text{He}(d, p(0^\circ)){}^4\text{He}$ reaction at 10 MeV which were 0.69 ± 0.13 and -0.67 ± 0.16 for the spin modes "2-6" and "3-5", respectively[6]. The value of the angle between the beam direction and polarization axis was found $\beta = -91.0 \pm 1.3^\circ$, therefore, the polarization axis is perpendicular to the plane containing the mean beam orbit in the accelerator.

The $d - p$ elastic events have been selected using the information on the energy losses in the plastic scintillators by protons and deuterons, their time-of-flight difference and interaction point position. The values of the analyzing powers A_y and A_{yy} at 880 and 2000 MeV were obtained from the following expressions

$$A_y = \frac{3(L - R)}{2p_y} \quad (1)$$

$$A_{yy} = \frac{L + R - 2}{p_{yy}} \quad (2)$$

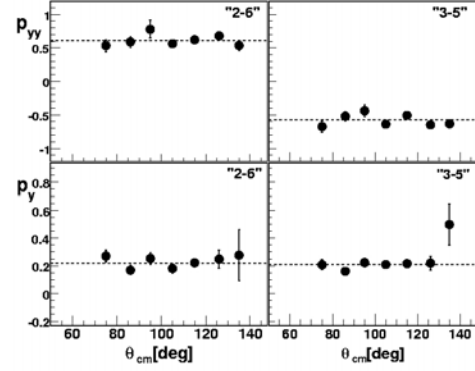


Figure 2: The vector p_y and tensor p_{yy} polarizations of the deuteron beam for "2-6" and "3-5" spin modes of POLARIS obtained at different angles. The error bars include both statistical and systematical uncertainties.

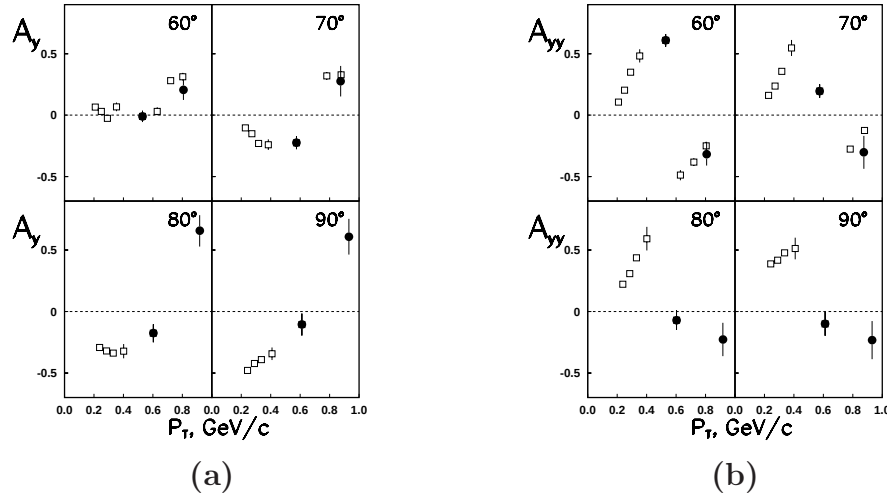


Figure 3: (a) The dependence of the vector A_y analyzing power in dp - elastic scattering at the fixed angles in the c.m.s. as a function of transverse momentum p_T .

(b) The dependence of the tensor A_{yy} analyzing power in dp -elastic scattering at the fixed angles in the c.m.s. as a function of transverse momentum p_T . The open and solid symbols represent data obtained at RIKEN, Saclay, ANL and at Nuclotron(Dubna), respectively.

Here L and R are the normalized yields of the selected events to the yield for unpolarized beam for scattering left and right, respectively.

The results on the vector A_y and tensor A_{yy} analyzing powers at the fixed angles in c.m.s are plotted as the functions of transverse momentum p_T in Fig.3a and in Fig.3b,

respectively. The open symbols represent the data obtained at RIKEN, Saclay and ANL. Nuclotron data obtained at the energies 880 and 2000 MeV are shown by the solid symbols. Both vector A_y and tensor A_{yy} analyzing powers change the sign at $p_T \sim 600\text{-}700$ MeV/c. The additional measurement are necessary to understand such behavior.

The A_y is consistent with zero at the angle of 60° in c.m.s, while it is about -0.2 at large angles at the energy of 880 MeV. The A_{yy} is sizeable only at the angles 60° and 70° at this energy. The absolute values of the vector A_y and tensor A_{yy} analyzing powers at 2000 MeV are sufficiently large to provide the efficient polarimetry of the deuteron beam.

The obtained results indicate the possibility of the efficient deuteron beam polarimetry using the $d - p$ elastic scattering at large angles in c.m.s. in the energy region from 880 to 2000 MeV. However, the measurements with larger statistics are needed. On the other hand, theoretical calculations[7]-[9] predict the large values of A_y , A_{yy} and A_{xx} at the angles $\theta_{c.m.} \leq 60^\circ$ at the energy of 880 MeV. It is very desirable to conduct the analyzing powers measurements at these angles.

Acknowledgments

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SPIN-MANIPULATING POLARIZED DEUTERONS AND PROTONS

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Abstract

We made the first systematic study of the spin resonance strengths induced by rf dipoles and solenoids using 2.1 GeV/c polarized protons and 1.85 GeV/c polarized deuterons stored in COSY. We found huge disagreements between the strengths measured in Froissart-Stora sweeps and the theoretical values calculated using the well-known formulae. These data resulted in correction of these formulae. We also tested Chao's matrix formalism for describing the spin dynamics near and inside a spin resonance, which allows analytic calculations of the beam polarization's behavior. Our measurements agreed precisely with the Chao formalism's predicted oscillations. We also tested Kondratenko's proposal to overcome depolarizing resonances by ramping through them with a crossing pattern that forces the depolarizing contributions to cancel themselves. Our tests with an rf bunched deuteron beam gave an ~ 20 -fold reduction in the depolarization. We recently used an rf solenoid to study rf spin resonances with both unbunched and bunched polarized beams of both protons and deuterons. We found narrowing of the bunched deuteron resonance, indicating that the beam's polarization behaved as if there was no momentum spread. However, we found widening of the proton resonance due to bunching.

Polarized scattering experiments require frequent spin-direction reversals to reduce systematic errors; one must overcome many spin resonances to maintain the polarization. In flat circular rings, each beam particle's spin precesses around the vertical fields of the ring's bending magnets. The spin tune ν_s (the number of spin precessions during one turn around the ring) is given by $\nu_s = G\gamma$, where G is the particle's gyromagnetic anomaly and γ is its Lorentz energy factor. The vertical polarization can be perturbed by any horizontal magnetic field. RF fields can induce an rf spin resonance when $f_r = f_c(k \pm \nu_s)$, where f_r is the resonance frequency, f_c is the circulation frequency and k is an integer.

Ramping an rf magnet's frequency through f_r can flip each particle's spin. The Froissart-Stora (FS) equation [1] relates the beam's final polarization after crossing the resonance P_f , to its initial polarization P_i

$$P_f = P_i \left\{ 2 \exp \left[\frac{-(\pi \mathcal{E}_{FS} f_c)^2}{\Delta f / \Delta t} \right] - 1 \right\}; \quad (1)$$

$\mathcal{E}_{FS}$ is the spin resonance strength, Δf and Δt are the rf frequency's ramp range and ramp time. For a flat ring, with a short rf magnet causing the only spin perturbation, $\mathcal{E}_{Bdl}$ due to rf solenoid's or an rf dipole's B -field was thought to be given by [2]

$$\text{RF solenoid (longitudinal } B): \quad \mathcal{E}_{Bdl} = \frac{1}{\pi 2\sqrt{2}} \frac{e(1+G)}{p} \int B_{rms} dl, \quad (2)$$

$$\text{RF dipole (transverse } B): \quad \mathcal{E}_{Bdl} = \frac{1}{\pi 2\sqrt{2}} \frac{e(1+G\gamma)}{p} \int B_{rms} dl, \quad (3)$$

where e is the particle's charge, p is its momentum, and $\int B_{rms} dl$ is the rf magnet's rms magnetic field integral in its rest frame.

We analyzed all available data on spin-flipping stored beams of protons, deuterons and electrons [3]. We calculated the rf-induced spin resonance strength ratios $\mathcal{E}_{FS}/\mathcal{E}_{Bdl}$, where $\mathcal{E}_{FS}$ was obtained by fitting the measured polarization to Eq. (1); $\mathcal{E}_{Bdl}$ was calculated using Eqs. (2) and (3). We found that $\mathcal{E}_{FS}/\mathcal{E}_{Bdl}$ was 7 times lower than predicted for deuterons, and 12 to 170 times higher for protons. We systematically studied $\mathcal{E}_{FS}/\mathcal{E}_{Bdl}$ ratios with vertically polarized protons and deuterons in COSY.

Identical apparatus was used for both protons and deuterons, including the COSY ring, the EDDA detector, the low energy polarimeter, the electron cooler, the injector cyclotron, the polarized ion source, and the rf dipole or solenoid [3–5, 8, 10–12]. All available $\mathcal{E}_{FS}/\mathcal{E}_{Bdl}$ data are shown in Fig. 1. With an rf dipole, we studied the dependence of $\mathcal{E}_{FS}/\mathcal{E}_{Bdl}$ on the beam's size, momentum spread $\Delta p/p$, and distance from the nearest 1st-order intrinsic spin resonance for both protons [3] and deuterons [4], and on Δf for deuterons. We observed no dependence of $\mathcal{E}_{FS}/\mathcal{E}_{Bdl}$ on the beam's size or $\Delta p/p$ for either protons or deuterons, and no dependence of $\mathcal{E}_{FS}/\mathcal{E}_{Bdl}$ on Δf for deuterons. We varied the vertical betatron tune ν_y near a 1st-order intrinsic spin resonance and observed an enhancement of $\mathcal{E}_{FS}/\mathcal{E}_{Bdl}$ with a hyperbolic dependence on distance from the 1st-order intrinsic spin resonance for both protons and deuterons. This can explain the enhancement for protons; however it can not explain the deuteron's very small $\mathcal{E}_{FS}/\mathcal{E}_{Bdl}$.

With an rf-solenoid and stored polarized deuterons, we measured an $\mathcal{E}_{FS}/\mathcal{E}_{Bdl}$ of 1.02 ± 0.05 over a wide range of parameters [5]. These new data agree precisely with Eq. (2). These rf-solenoid data together with earlier rf-dipole results [3, 4] indicate that Eq. (2) is correct for longitudinal rf fields in an ideal accelerator, while Eq. (3) for radial rf fields is incorrect. Moreover, Eq. (3) must be replaced by more complex calculations, which depend on each ring's properties. One must properly include the modification of $\mathcal{E}_{Bdl}$ due to coherent beam oscillations caused by the rf dipole, which then interact with all the ring's radial and longitudinal magnetic fields [2, 6].

Equation (1) is only valid if Δf is significantly larger than the spin resonance's width. Chao's new matrix formalism [7] deals with conditions where the FS formula is not valid. The Chao formalism can be used to calculate the spin dynamics anywhere inside a piecewise linear resonance crossing. Our measurements [8] of the deuteron's polarization near and inside the resonance agree with the Chao formalism's predicted oscillations.

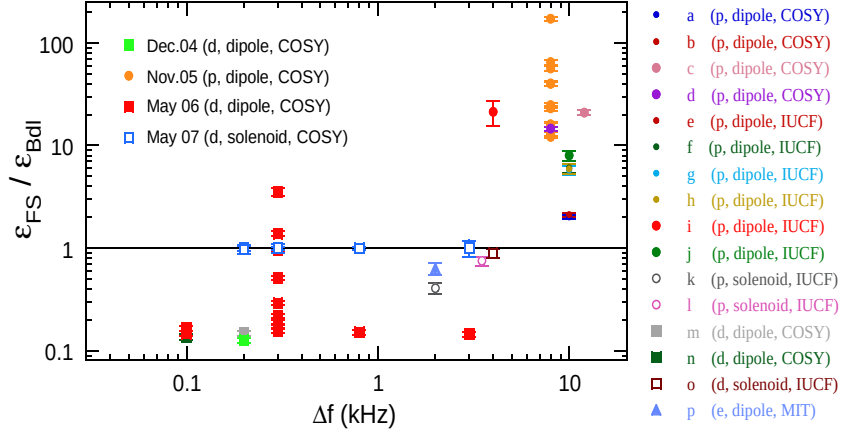


Figure 1: Ratio of $\mathcal{E}_{FS}$ to $\mathcal{E}_{Bdl}$ is plotted vs. rf magnet's frequency sweep range Δf .

Several techniques [9] are used to overcome spin resonances. Siberian snakes [10] are most useful at very high energies. In the 1-25 GeV region one can use harmonic correction of the imperfection resonances and fast crossing (FC) of the intrinsic resonances; these are less effective than Siberian snakes and often leave some depolarization at each crossed resonance. Kondratenko proposed a new technique for overcoming depolarizing resonances [11], by crossing each resonance using the Kondratenko Crossing (KC) pattern. We tested this proposal using stored 1.85 GeV/c vertically polarized deuterons [12]. We crossed an rf depolarizing resonance by ramping an rf solenoid's frequency through the KC pattern shown in Fig. 2, while varying its parameters. As shown in Fig. 3, with its optimal parameters, KC gave measured polarization losses of $3.3 \pm 0.2\%$ and $0.8 \pm 0.3\%$ with unbunched and bunched beams, respectively; while fast crossing (FC) at the same crossing rate, gave measured losses of $15.6 \pm 0.2\%$ and $15.0 \pm 0.3\%$, respectively. This clearly shows KC's potential value.

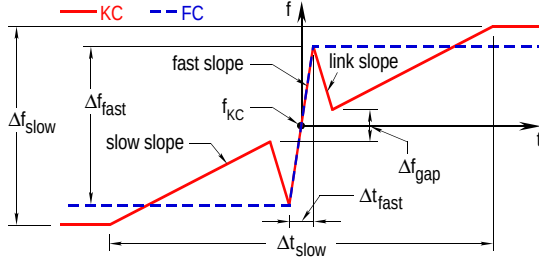


Figure 2: Kondratenko Crossing (KC) [solid line] and Fast Crossing (FC) [dashed line] patterns. This figure defines the KC pattern parameters.

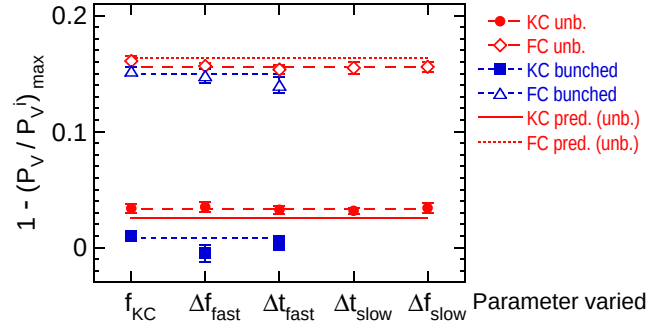


Figure 3: Summary of depolarizations at KC peak for both KC and FC, with unbunched and bunched beams.

We used an rf solenoid to study rf spin resonances with both bunched and unbunched stored beams of 1.85 GeV/c polarized deuterons [13]. With the unbunched beam, we set many different fixed rf solenoid frequencies near the resonance; we found only partial depolarization. However, we found that the bunched beam's polarization was almost fully flipped, and its resonance was about 5 times narrower (see Fig. 4). We then used Chao's equations [7] to explain this behavior and to calculate the polarization's dependence on various rf-solenoid and beam parameters. Our data and calculations show that a bunched deuterons' polarization can behave as if the beam has zero momentum spread.

We recently used an rf solenoid to study rf spin resonances with both unbunched and bunched stored beams of 2.1 GeV/c polarized protons [14]. For the unbunched beam at different fixed rf-solenoid frequencies, we found only a possible very shallow depolarization dip; it was greatly enhanced by making 400 Hz frequency sweeps centered at similar frequencies (see Fig. 5). With the bunched proton beam, both the fixed-frequency and frequency-sweep maps were similar; however, they were more than twice as wide as the unbunched maps. Moreover, the bunched maps showed full depolarization in the wide resonance region. This widening of the proton resonance due to bunching is exactly opposite to the narrowing of deuteron resonances due to bunching.

I thank the COSY staff for the successful operation of COSY, and my SPIN@COSY colleagues for their help and advice. This research was supported by grants from the German BMBF Science Ministry, its JCHP-FFE program at COSY.

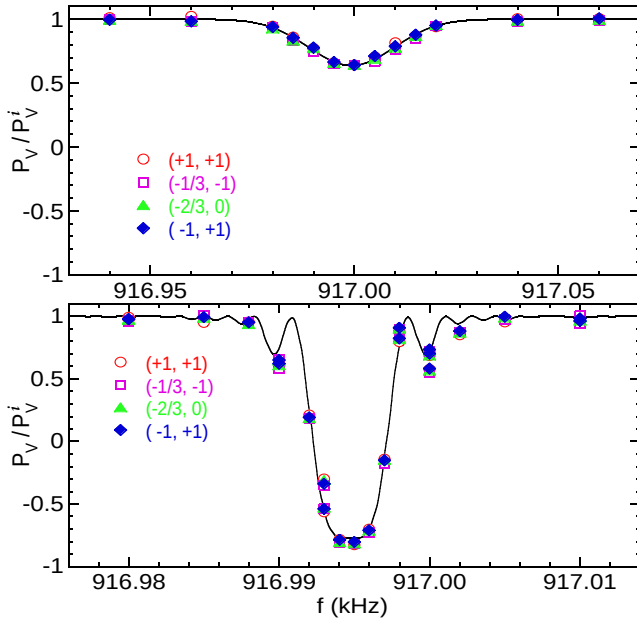


Figure 4: Unbunched (Top) and bunched (Bottom) deuteron resonance maps.

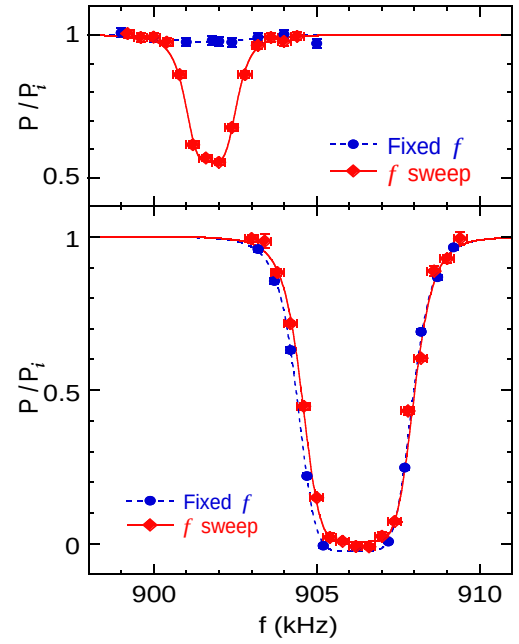


Figure 5: Unbunched (Top) and bunched (Bottom) proton resonance maps.

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A STUDY OF POLARIZED METASTABLE HELIUM-3 ATOMIC BEAM PRODUCTION

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Abstract

The problems of metastable ^3He atomic beam formation are considered. The difference between metastable ^3He and hydrogen or deuterium atoms concerning the radio-frequency transitions of atomic states is essential. The Schroedinger equation in the uncoupled basis $|\psi_e, \psi_h\rangle$ and also in the basis of stationary states are received. The results of computer simulations agree with published data. The possibility to get the positive and negative values of helion polarization by using two types of the weak field transitions in the metastable helium-3 atom is shown.

1 Introduction

The spin-dependent part of the Hamiltonian for helium-3 atoms in the metastable state 2^3S_1 with electron spin moment one is

$$\hat{H} = -\mu_J \vec{J} \vec{B}(t) - \mu_h \vec{\sigma}_h \vec{B}(t) - \frac{1}{3} \Delta W \vec{\sigma}_h \vec{J}, \quad (1)$$

where $\vec{\sigma}_h$ are the Pauli spin matrices of the helion, $\vec{J}$ are the electron spin matrices ($J = 1$), $B(t)$ is magnetic field strength, ΔW is hyperfine splitting; $\Delta W = 4.4645 \times 10^{-24}$ J = $\hbar \times 4.2335 \times 10^{10}$ rad/s, $\mu_J = 2\mu_e = -1.85695275 \times 10^{-23}$ J/T = $-\hbar \times 1.76085977 \times 10^{11}$ rad/s T, $\mu_h = -1.07455 \times 10^{-26}$ J/T = $-\hbar \times 1.0189 \times 10^8$ rad/s T.

We consider that a static magnetic field B directed along a z-axis. The wave functions of the hyperfine states $\Psi(F, M_F)$ for this magnetic field and in a high field limit are $\Psi_1(1/2, +1/2) = -\sin \beta \psi_h^+ \psi_J^0 + \cos \beta \psi_h^- \psi_J^+ \Rightarrow \psi_h^- \psi_J^+$, $\Psi_2(3/2, +3/2) = \psi_h^+ \psi_J^+$, $\Psi_3(1/2, -1/2) = -\sin \alpha \psi_h^+ \psi_J^- + \cos \alpha \psi_h^- \psi_J^0 \Rightarrow \psi_h^- \psi_J^0$, $\Psi_4(3/2, +1/2) = \cos \beta \psi_h^+ \psi_J^0 + \sin \beta \psi_h^- \psi_J^+ \Rightarrow \psi_h^+ \psi_J^0$, $\Psi_5(3/2, -1/2) = \cos \alpha \psi_h^+ \psi_J^- + \sin \alpha \psi_h^- \psi_J^0 \Rightarrow \psi_h^+ \psi_J^-$, $\Psi_6(3/2, -3/2) = \psi_h^- \psi_J^-$, where $\sin \beta = \sqrt{A_+}$, $\cos \beta = \sqrt{1 - A_+}$; $\sin \alpha = \sqrt{A_-}$, $\cos \alpha = \sqrt{1 - A_-}$;

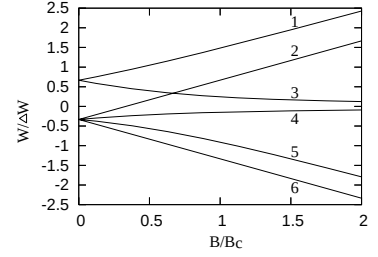
$$A_+ = \frac{1}{2} \left(1 - \frac{x + 1/3}{\sqrt{1 + \frac{2}{3}x + x^2}} \right), \quad A_- = \frac{1}{2} \left(1 - \frac{x - 1/3}{\sqrt{1 - \frac{2}{3}x + x^2}} \right);$$

$$x = \frac{B}{B_c}, \quad B_c = \frac{\Delta W}{-\mu_J/J + \mu_I/I} = \frac{\Delta W}{-\mu_J + 2\mu_h} = 0.2407 \text{ T}.$$

The Breit-Rabi diagram of six Zeeman hyperfine components of this metastable state is shown at Fig. 1, where the numbers correspond to those of the wave functions.

The energies of the states $\Psi_1 - \Psi_6$ are

$$\begin{aligned} W_1 &= \frac{\Delta W}{6} - \mu_J \frac{B}{2} + \frac{\Delta W}{2} \sqrt{1 + \frac{2}{3}x + x^2}, \quad W_2 = -\frac{\Delta W}{3} - \\ &\mu_J B - \mu_h B, \quad W_3 = \frac{\Delta W}{6} + \mu_J \frac{B}{2} + \frac{\Delta W}{2} \sqrt{1 - \frac{2}{3}x + x^2}, \quad W_4 = \\ &\frac{\Delta W}{6} - \mu_J \frac{B}{2} - \frac{\Delta W}{2} \sqrt{1 + \frac{2}{3}x + x^2}, \quad W_5 = \frac{\Delta W}{6} + \mu_J \frac{B}{2} - \\ &\frac{\Delta W}{2} \sqrt{1 - \frac{2}{3}x + x^2}, \quad W_6 = -\frac{\Delta W}{3} + \mu_J B + \mu_h B. \end{aligned}$$



2 The Schroedinger equation in the uncoupled state basis

Figure 1: The scheme of $^3\text{He } 2^3\text{S}_1$ hyperfine structure and Zeeman splitting, $B_c = 0.2407$ T.

The adiabatic transitions for the hyperfine states of hydrogen were considered by Antishev and Belov [1]. Oh [2] published detailed results for the weak field transitions (WFT) in deuterium. Here we solve this problem for ^3He in the metastable state 2^3S_1 .

In the uncoupled $|m_h, m_J\rangle$ state basis

$$\Psi(t) = C_1(t)\psi_h^+\psi_J^+ + C_2(t)\psi_h^+\psi_J^0 + C_3(t)\psi_h^+\psi_J^- + C_4(t)\psi_h^-\psi_J^+ + C_5(t)\psi_h^-\psi_J^0 + C_6(t)\psi_h^-\psi_J^-, \quad (2)$$

and we obtain the following equations for the amplitudes:

$$\begin{aligned} dC_1/dt &= -i/\hbar\{C_1[-(\mu_h + \mu_J)B_z + \Delta W/3] - C_2\mu_J B_x/\sqrt{2} - C_4\mu_h B_x\} \\ dC_2/dt &= -i/\hbar\{-C_1\mu_J B_x/\sqrt{2} - C_2\mu_h B_z - C_3\mu_J B_x/\sqrt{2} + C_4\sqrt{2}\Delta W/3 - C_5\mu_h B_x\}. \\ dC_3/dt &= -i/\hbar\{-C_2\mu_J B_x/\sqrt{2} + C_3[-(\mu_h + \mu_J)B_z - \Delta W/3] + C_5\sqrt{2}\Delta W/3 - C_6\mu_h B_x\}. \quad (3) \\ dC_4/dt &= -i/\hbar\{-C_1\mu_h B_x + C_2\sqrt{2}\Delta W/3 + C_4[(\mu_h - \mu_J)B_z - \Delta W/3] - C_5\mu_J B_x/\sqrt{2}\}. \\ dC_5/dt &= -i/\hbar\{-C_2\mu_h B_x + C_3\sqrt{2}\Delta W/3 - C_4\mu_J B_x/\sqrt{2} + C_5\mu_h B_z - C_6\mu_J B_x/\sqrt{2}\}. \\ dC_6/dt &= -i/\hbar\{-C_3\mu_h B_x - C_5\mu_J B_x/\sqrt{2} + C_6[(\mu_h + \mu_J)B_z + \Delta W/3]\}. \end{aligned}$$

3 The Schroedinger equation in the basis of the stationary states

Another way is to solve the Schroedinger equation in the basis of the stationary states, as it was made by Beijers [3] for hydrogen. But he used the "static" Hamiltonian slowly changing with time because atoms move through a changing magnetic field. Hasuyama and Wakuta [4] used the "static" Hamiltonian for strong field transitions in deuterium. This approach is not always correct.

We use, as the basis, the stationary states existing at the value of the magnetic field, $B_z = B_0$, at the enter into the region of a RF field.

For WFT 2-6 we consider only the four-level system of 2,4,5 and 6 substates of the $F = 3/2$ state, since the levels of the $F = 1/2$ state are sufficiently distant as to have no significant effect in our problem. For the amplitudes of these states, we use the notations $c_2 - c_6$.

The Schroedinger equation is written as

$$i\hbar \frac{d\Psi}{dt} = [\hat{H} + \hat{H}'(t)]\Psi, \quad (4)$$

where $\hat{H}$ is the time-independent Hamiltonian whose eigenfunctions satisfy the equations $\hat{H}\psi_n = W_n\psi_n$. The exact wave function is written in the form

$$\Psi = \sum c_n(t)\psi_n e^{-iW_n t/\hbar}. \quad (5)$$

The coefficients $c_n(t)$ must satisfy the differential equations

$$i\hbar \frac{dc_k(t)}{dt} = \sum H'_{kn}(t)c_n(t)e^{i\omega_{kn}t}, \quad (6)$$

where $\omega_{kn} = (W_k - W_n)/\hbar$, and $H'_{kn}(t) = \langle \psi_k | \hat{H}'(t) | \psi_n \rangle$.

For the weak field transitions, one may use

$$\hat{H}'(t) = -\mu_h \sigma_{hx} B_x(t) \sin \omega t - \mu_J S_{Jx} B_x(t) \sin \omega t - \mu_h \sigma_{hz} b_z(t) - \mu_J S_{Jz} b_z(t), \quad (7)$$

where $b_z(t) = (dB_z/dx)vt$.

The matrix elements are

$$\begin{aligned} H'_{22} &= (-\mu_h - \mu_J)b_z(t), \\ H'_{24} &= [-\mu_h \sin \beta - (\mu_J/\sqrt{2}) \cos \beta]B_x(t) \sin \omega t, \\ H'_{44} &= [\mu_h(\sin^2 \beta - \cos^2 \beta) - \mu_J \sin^2 \beta]b_z(t), \\ H'_{45} &= [-\mu_h \sin \alpha \cos \beta - \mu_J/\sqrt{2}(\sin \alpha \sin \beta + \cos \alpha \cos \beta)]B_x(t) \sin \omega t \\ H'_{55} &= [\mu_h(\sin^2 \alpha - \cos^2 \alpha) + \mu_J \cos^2 \alpha]b_z(t), \\ H'_{56} &= [-\mu_h \cos \alpha - (\mu_J/\sqrt{2}) \sin \alpha]B_x(t) \sin \omega t, \\ H'_{66} &= (\mu_h + \mu_J)b_z(t). \end{aligned} \quad (8)$$

The values of $\sin \alpha$, $\cos \alpha$, $\sin \beta$, $\cos \beta$, ω_{ik} are taken at the initial value of $B_z(x_{init})$. To find the final amplitude we must multiply the resulting wave function by Ψ_n^* at the final value of the field, $B_z(x_{final})$.

Also, we need the matrix elements for the transition 1 – 3. We use the notations c_1 and c_3 .

$$\begin{aligned} H'_{11} &= [\mu_h(\cos^2 \beta - \sin^2 \beta) - \mu_J \cos^2 \beta]b_z(t), \\ H'_{13} &= [\mu_h \sin \beta \cos \alpha - \mu_J/\sqrt{2}(\sin \alpha \sin \beta + \cos \alpha \cos \beta)]B_x(t) \sin \omega t \\ H'_{33} &= [\mu_h(\cos^2 \alpha - \sin^2 \alpha) + \mu_J \sin^2 \alpha]b_z(t), \end{aligned} \quad (9)$$

We note that at a weak magnetic field ($x \ll 1$) the level distance $W_1 - W_3 \approx -\frac{4}{3}\mu_J B$, and $W_2 - W_4 = W_4 - W_5 = W_5 - W_6 \approx -\frac{2}{3}\mu_J B$. This is different from the case of deuterium where all distances between the levels at $F = 3/2$ and $F = 1/2$ at the weak magnetic field are equal to $\approx -\frac{2}{3}\mu_e B$.

Let we have a system of two sextupoles with the space between them. Then, realizing the transition $1 \rightarrow 3$ in the space between the sextupoles, after the second sextupole we get the pure state 2 with $F = 3/2$, $m_F = 3/2$ and after ionization in a strong magnetic

field $P \approx +1$. If we add the transition $2 \rightarrow 6$ after the second sextupole, we produce pure state 6 with $F = 3/2$, $m_F = -3/2$ and $P \approx -1$ after ionization.

So, for the polarized metastable ^3He atomic beam, we need two weak field transition units that should be placed between and after the sextupole magnets.

It is interesting to note that the pure states of $^3\text{He}(2\text{S})$ may be transferred into the pure states of $^3\text{He}^+$ after stripping one electron in 2S -state. The results of Slobodrian [5] tentatively confirm this point. In their scheme the adiabatic transition in a weak magnetic field ($m_F \rightarrow -m_F$) transforms components 1-2 by turns into 3-6. In a magnetic field the wave function of the hyperfine substate 3 of the $^3\text{He}(2\text{S})$ atom is

$$\psi(F = 1/2, m_F = -1/2) = -\sin \alpha \psi_h^+ \psi_J^- + \cos \alpha \psi_h^- \psi_J^0 \Rightarrow \psi_h^- \psi_J^0.$$

With $B = 0.2$ T, $(x = 0.8309)$ $\cos \alpha = 0.8564$ and $\sin \alpha = 0.5164$.

It can be easily shown that if the second ionization is effected in zero magnetic field the expected value of P for the pure state of $^3\text{He}^+$ would be $P = -0.68$, and for the mixed state $^3\text{He}^+$ $P = -0.44$. The measured value is $P = -(0.6 \div 0.8)$.

A tapered electromagnet produces a static magnetic field $B_z(x)$ perpendicular to the beam path with a field gradient dB_z/dx along $x = vt$.

$$B_z(x) = B_0 + \frac{dB_z}{dx}x, \quad B_x(x) = B_1(x) \sin \omega x. \quad (10)$$

We accepted that some parameters have the same values as in the paper by Oh [2]:

$B_0 = 1.17 \times 10^{-3}$ T, $dB_z/dx = -1.4 \times 10^{-2}$ T/m for a negative static field gradient (or $B_0 = 4.7 \times 10^{-4}$ T, $dB_z/dx = 1.4 \times 10^{-2}$ T/m for a positive gradient), $l = 5 \times 10^{-2}$ m, $\omega = 9.63 \times 10^7$ rad/s for the $2 \rightarrow 6$ transition and $\omega = 1.93 \times 10^8$ rad/s for the $1 \rightarrow 3$ transition. The atomic beam velocity $v = 1.2 \times 10^3$ m/sec. The RF amplitude $B_1(x)$ is a quadratic function of x with zero values at $x = 0$ and $x = l$; $B_1^{max} = B_1(l/2) = (1 - 2) \times 10^{-4}$ T.

The results of the computer calculations for an atom velocity of 1200 m/s give practically 100% probability of the transition.

4 Conclusion

Some aspects of developing a polarized helion source for JINR Accelerator Complex was discussed. The possibility to get positive and negative values of the helion polarization by using two types of the weak field transitions in the metastable helium-3 atom is shown.

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PROTON POLARIMETER AT 200 MEV ENERGY

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Abstract

The precise measurements of proton beam polarization at energy 200 MeV (at the exit of the linac before injection into Booster) are required for the polarization adjustment (monitoring and optimization) out of the source, for the spin-rotator tune of the vertical polarization and for the study of the polarization losses in the Booster and AGS. For this purpose, we have investigated two of possible solutions: polarimeters based on elastic pp and pD scattering. We compare these two polarimeters by estimating their factor of merit, statistical and systematical precisions and quality of the experimental apparatus. The absolute polarization measurements of about $\pm 1\%$ can be achieved at 200 MeV by using very accurate known value of pD analyzing power.

RHIC is the first collider where the "Siberian snake" technique was successfully implemented to avoid the resonance depolarization during beam acceleration in RHIC. Polarimetry is another essential component of the polarized collider facility. A complete set of polarimeters includes: Lamb-shift polarimeter at the source energy, a 200 MeV polarimeter after the linac, and polarimeters in AGS and RHIC based on proton-Carbon scattering in Coulomb-Nuclear Interference region. A polarized hydrogen jet polarimeter was used for the absolute polarization measurements in RHIC. Also local polarimeters for tuning of longitudinal polarization are installed at STAR and PHENIX detectors.

The polarized proton source operates at about 1 Hz repetition rate and additional source pulses were directed to the 200 MeV proton-Carbon polarimeter for the source polarization measurements and optimization, spin-rotator tuning, and continuous polarization monitoring. The p-Carbon polarimeter was calibrated by the elastic p-Deuteron scattering where analyzing power was measured to less than 0.5% absolute accuracy at IUCF. This polarimeter is described in the papers [1], [2].

The polarimeter schematic layout is shown in Fig.1. The targets are exchangeable carbon filaments and CD₂ polyethylene foils. The measurements of elastic proton-Deuteron collisions and complete collision kinematics

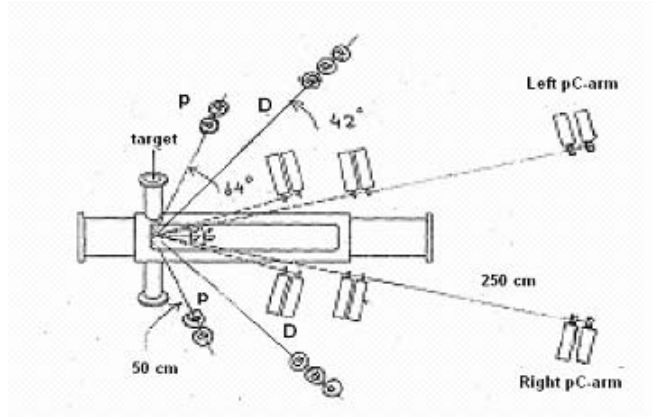


Figure 1: Layout of 200 GeV proton polarimeter.

reconstruction provide more accurate absolute polarization determination. However low counting rate requires a long measurement time (typically about 8 hrs for $\pm 2\%$ measurement accuracy). This limits the accuracy of calibration since the source polarization can vary during the measurements.

The proton-Carbon inclusive polarimeter was used because of higher cross-section of this process comparatively to pp or pD processes. The use of carbon nuclei as a scattering target material has the following shortcomings, stated in [3]. The cross-section and analyzing power of pC polarimeter sharply depends on proton scattering angle that leads to significant false asymmetries if the beam traverses the polarimeter not precisely in the center or with some angular displacement. Therefore to achieve the accuracy of absolute polarization measurements $\pm 1\%$ it is required systematic continuous polarimeter recalibration.

To choose the optimal process for pC polarimeter calibration we shall compare pp and pD polarimeters by such characteristics as Factor-of-Merit (FoM), and absolute accuracy of polarization measurements.

Factor-of-Merit M is defined as [3]:

$$M = \frac{d\sigma}{d\Omega} A_y^2, \quad (1)$$

where $\frac{d\sigma}{d\Omega}$ is differential elastic scattering cross-section and A_y is the polarimeter analyzing power.

To define the angle dependence of cross-section the equation from [4] was used:

$$\begin{aligned} \sigma(\theta) = & \left[\frac{\eta}{2k \sin^2(\theta/2)} \right]^2 + \frac{\eta A_I(0)}{k \sin^2(\theta/2)} \sin(\eta \ln \sin^2(\theta/2)) - \\ & - \frac{\eta A_R(0)}{k \sin^2(\theta/2)} \cos(\eta \ln \sin^2(\theta/2)) + \sum_{l=0}^{l_{\max}} a_{2l} \cos^{2l} \theta. \end{aligned} \quad (2)$$

Here $\eta = e^2/\hbar v$, v is the velocity in the laboratory system, $\hbar k$ and θ are momentum and scattering angle in the center-of-mass system. The imaginary part of the amplitude $A_I(0)$ is defined by optical theorem from the total scattering cross-section σ_t :

$$A_I(0) = k\sigma_t/4\pi. \quad (3)$$

The real part of the amplitude $A_R(0)$ and the values a_{2l} also are taken from [4]. The derived dependence is shown in Fig. 2. The points show experimental data from [5].

Experimental angular dependence of analyzing power An for the beam proton energy 210 MeV is taken from [6], where it is fitted by function:

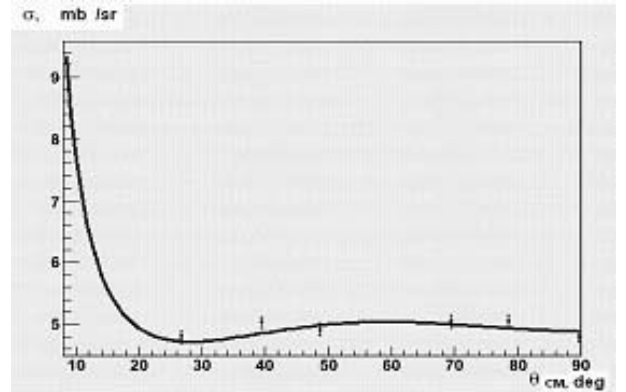


Figure 2: Elastic pp-scattering cross-section vs scattering angle.

$$P(\theta) \sigma(\theta) = \sin \theta \cos \theta \sum_{n=0}^2 b_{2n} \cos^{2n} \theta, \quad (4)$$

where $b_0=2.78\pm0.16$, $b_2=-2.24\pm0.17$, $b_4=2.96\pm 0.77$.

Using equation (1) factor of merit (FoM) for the polarimeter on pp scattering was found. It's angular dependence in center-of-mass (CM) frame is shown in Fig. 3.

After recalculation we see that maximal value of Factor-of-Merit, achieved at 40° in CM frame corresponds to 19.1° of the scattered proton in the laboratory frame.

After usual transformations we defined, that maximal value of Factor-of-Merit, achieved in CM frame at 40° corresponds to 69.1° of the recoil proton angle in the laboratory frame and this recoil angle corresponds to the recoil proton energy 20 MeV.

Now we shall consider the polarimeter based on pD scattering. The pD scattering cross-section at 200 GeV is two times higher than pp scattering cross-section [7].

We shall use the analyzing power dependence on the scattering angle in the laboratory frame, measured with high precision in [3] and shown in Fig. 3 for the proton beam energy 200 MeV.

In [8] is given the value of an angle in the laboratory frame with maximal analyzing power equal to 42.6° and is given the maximal value of the analyzing power, equal to $A_{pD} = 0.507\pm0.002$, that gives relative statistic uncertainty better, than 0.5%. This high precision allows use the pD polarimeter to improve the absolute accuracy of polarization measurements.

The dependence of recoil deuteron kinetic energy on recoil angle for pD reaction was calculated.

Maximum of Factor-of-Merit being reached at 42.6° corresponds to 110 MeV of recoil deuteron kinetic energy. Factor-of-Merit for pD polarimeter in maximum is equal 2.5 and approximately 5 times larger, than for pp polarimeter.

So we choose pD polarimeter for further development. In this case we shall choose the laboratory angles for scattered protons and recoil deuterons as in [3] being 64.5° and 42.6° respectively.

As was mentioned above, pp and pD targets are composed of polyethylene (CH_2) or deuterated (CD_2) polyethylene, containing carbon.

We estimate proton rate from the scattering on carbon in the target 2-3 times higher than on deuterons. This background can be discriminated by time-of-flight measurements for scattered proton and recoil deuteron and deuteron energy measurements.

We estimate the counting rate of pD polarimeter as several MHz during the bunch (mostly from p-Carbon, not pD). So there is a significant problem to detect particles with this counting rate. Even with the good selection factors the pD events rate is limited

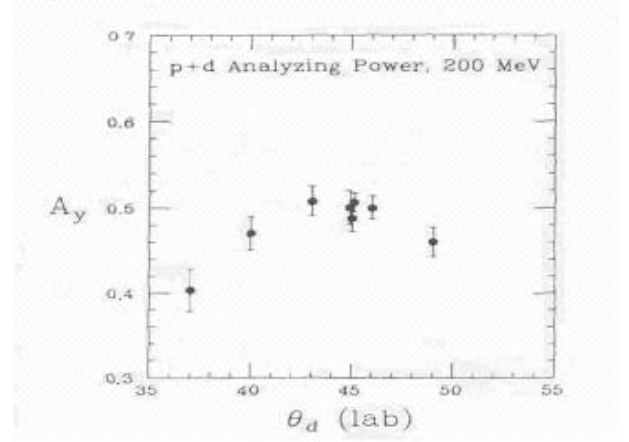


Figure 3: Deduced values for the p+D analyzing power at 200 MeV as a function of the deuteron scattering angle observed in the laboratory frame [3]. Error bars represent only the statistical uncertainties.

by detector saturation from pC collisions. Nevertheless, with bunch duration 300 us and bunch repetition rate 0.5 Hz we shall detect only some events per second in coincidence. This is insufficient for good statistical precision for a short measurement time.

So we are planning in the future to analyze the possibility to use proton-carbon polarimeter with separation of elastically scattered protons.

The use of contemporary fast detectors (and DAQ) would allow higher rate and better statistical error. For application in pD polarimeter we shall investigate as scintillators new materials like LYSO ($\text{Lu}_{1.8}\text{Y}_{0.2}\text{SiO}_5$) with high light output, better energy resolution (8%) and short rise (10 ns) and decay (40 ns) time. High density of this material ($\sim 7 \text{ g/cm}^3$) will provide full absorption of deuteron in the scintillator with the thickness less than 1 cm. Also we shall examine new plastic scintillators with exclusively short timing properties. As light sensors will be investigated Silicon photomultipliers (SiPM).

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SPIN-ORBIT POTENTIALS IN NEUTRON-RICH HELIUM ISOTOPES

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Abstract

We measured the vector analyzing power of the proton elastic scattering from neutron-rich helium isotopes ${}^6\text{He}$ and ${}^8\text{He}$ at 71 MeV/u. A solid polarized proton target specially constructed for radioactive-ion beam experiments was successfully applied. Deduced spin-orbit potentials between protons and ${}^{6,8}\text{He}$ are found to be remarkably shallow. Dominance of the interaction between a proton and an α core in ${}^{6,8}\text{He}$ is also indicated.

1 Introduction

Recently, renewed interest has been focused on the spin-dependent interaction in unstable nuclei. One of the most direct methods to extract the information on the spin-dependent interaction is the scattering experiment induced by spin-polarized light ions. In order to investigate the unstable nuclei with spin polarization, a polarized proton solid target, which was specially designed for radioactive-ion beam experiments, has been developed at CNS [1, 2]. The most prominent advantage of this target is in its modest operation condition, *i.e.*, a low magnetic field of 0.1 T and a high temperature of 100 K. This condition allows us to detect low energy recoiled protons, which are essential for event selection. Making use of this target, we measured the vector analyzing power of the proton elastic scattering from ${}^6\text{He}$ and ${}^8\text{He}$ at 71 MeV/A in 2005 and 2007. The aim of the measurements is to extract the feature of the shape of spin-orbit potentials between protons and ${}^{6,8}\text{He}$, and discuss the effect of valence neutrons on the spin-orbit potential from the systematics between isotopes. Details of the measurement and the optical model analysis are reported.

2 $\vec{p}+^8\text{He}$ analyzing power measurement

The experiment was carried out at the RIKEN Nishina Center Accelerator Research Facility (RARF) using the RIKEN Projectile-fragment Separator (RIPS). The ^8He beam was produced by a fragmentation reaction of an ^{18}O beam with an energy of 100 MeV/A bombarded onto a 13 mm² Be target. The energy and the intensity of a ^8He beam were 71 MeV/A and 1.5×10^5 pps, respectively. The purity of the beam was 77 %.

As a secondary target, we used the polarized proton solid target (Fig. 1). The material of the secondary target was a single crystal of naphthalene with a thickness of 4.3×10^{21} protons/cm². Protons in the crystal were polarized under a low magnetic field of 0.09 T at 100 K by the optical excitation of electrons and the cross-polarization method.

The target polarization was $11.0 \pm 2.5\%$ on average. Recoil protons were detected using multiwire drift chambers and CsI (TI) scintillators placed on the left and right sides of the beam line. Each MWDC was located 180 mm away from the target and covered a scattering angle of $50^\circ - 70^\circ$ in the laboratory system. A CsI scintillator with a sensitive area of $135 \text{ mm}^H \times 60 \text{ mm}^V$ was placed just behind the MWDC. Leading particles $^6,^8\text{He}$ were detected using another MWDC and $\Delta E - E$ plastic scintillator hodoscopes with thicknesses of 5 and 100 mm.

The measured differential cross sections and analyzing powers are shown in Fig. 2 by closed circles. They are consistent with the previous data of the differential cross section [3] plotted by open circles.

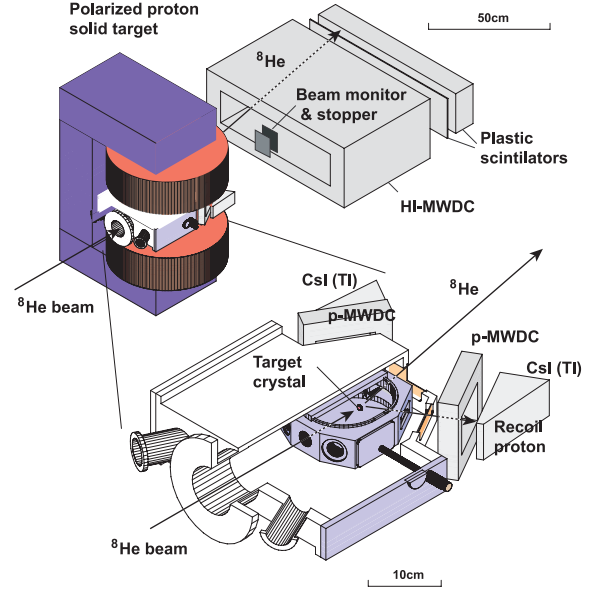


Figure 1: Schematic view of the experimental setup.

3 Optical model analysis

3.1 Phenomenological optical potential

In order to extract the global nature of $\vec{p}-^6,^8\text{He}$ interactions, the data were phenomenologically analyzed with an optical model potential. For the central and spin-orbit terms, we assumed Woods-Saxon and Thomas type functions, respectively. We searched a parameter set that reproduces the data using a fitting code ECIS79. As an initial potential, a parameter set for $\vec{p}+^6\text{Li}$ elastic scattering at 72 MeV/A [4] was used. The dashed lines in Fig. 2 show the calculation with initial parameters. Results of the best-fit parameters are presented by solid curves. Both differential cross section and analyzing power data are well reproduced except for the scattering at backward angles.

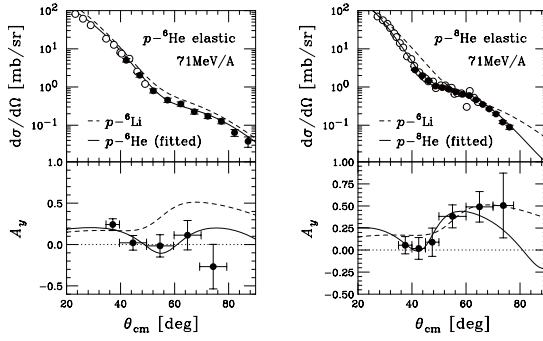


Figure 2: The differential cross section and the analyzing power of the $p+{}^6,{}^8\text{He}$ elastic scattering at 71 MeV/A.

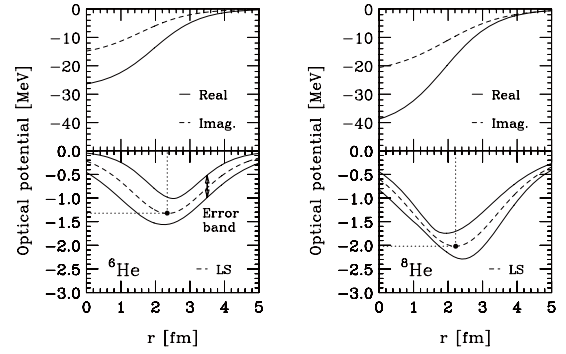


Figure 3: Phenomenological optical potentials between a proton and ${}^6,{}^8\text{He}$ as functions of the radius.

The potentials obtained are shown in Fig. 3 as a function of radius. The upper panel displays real and imaginary parts of the central term, while the lower one presents a spin-orbit term with error bands resulted from fitting uncertainty. Due to the target polarization uncertainty, there is an additional scale error of 28% for ${}^6\text{He}$ and 23% for ${}^8\text{He}$ in the depth of the spin-orbit potentials.

3.2 Feature of the spin-orbit potentials

The shape of the spin-orbit potentials in neutron-rich helium isotopes is discussed here. In order to extract the gross feature of the potentials, we focus on the radius and the amplitude of the peak of spin-orbit potential as shown by the dotted lines in Fig. 3. We call the former “LS radius” and the latter “LS amplitude”. Since the spin-orbit potential is usually approximated by the radial derivative of density distribution, LS radius and LS amplitude should be closely related to the radius and the gradient of the density distribution, respectively.

The LS radii and LS amplitudes of ${}^6\text{He}$ and ${}^8\text{He}$ are presented in Fig. 4 by closed squares. Those of neighboring even-even stable nuclei [5, 6] and a global optical potential [7] are also plotted by closed and open circles. It is clearly demonstrated that the LS amplitudes of ${}^6\text{He}$ and ${}^8\text{He}$ are remarkably smaller than those of stable nuclei. The LS amplitudes of stable nuclei are almost constant and distributed between 4–5.5 MeV, whereas those of ${}^6\text{He}$ and ${}^8\text{He}$ are as small as 1.3 and 2.0 MeV. It can be concluded that the spin-orbit potentials in neutron-rich helium isotopes are characterized by the significantly shallow shape. This can be intuitively explained from the largely diffused density distribution of ${}^6\text{He}$ and ${}^8\text{He}$, whose gradient is more than twice as small as that of ${}^4\text{He}$. Results of more detailed analysis with microscopic optical model calculation will be reported elsewhere.

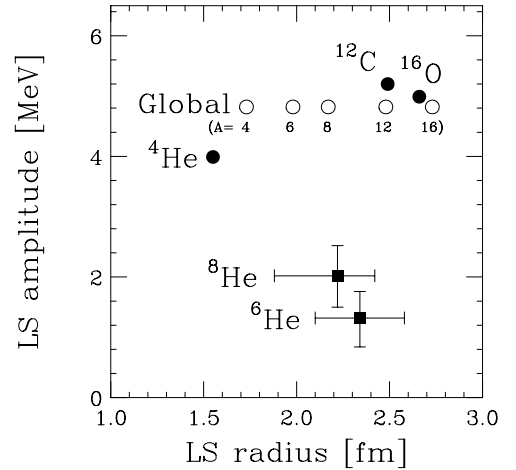


Figure 4: Two-dimensional distribution of LS radius and LS amplitude.

3.3 Effect of valence neutrons on spin-orbit potential

Here, we discuss the effect of valence neutrons on the spin-orbit potential from the systematics in helium isotopes. For this purpose, an $\alpha + xn$ cluster folding calculation [5] could be a useful method, since the contribution of an α core and valence neutrons can be separately evaluated. It has been expected from the calculation that the contribution of valence neutrons to the spin-orbit potential is one-order of magnitude smaller than that of an α core. Consequently, the spin-orbit potentials of ${}^6,8\text{He}$ are expected to be related to the spatial distribution of the α core in a nucleus, *i.e.*, the LS radius should increase with the diffuseness of the α core distribution, whereas the LS amplitude should decrease.

The LS radius and LS amplitude of helium isotopes are plotted in Fig. 5 as a function of the proton orbital radius, which represents the diffuseness of the α core distribution. Although there are large experimental uncertainties, almost linear relations can be seen. From this we can derive a picture that the valence neutrons affect the spin-orbit potential rather strongly through the $p - \alpha$ interaction than via the direct $p - n$ interaction. The recoil motion of the α core induced by valence neutrons can be important in the explanation of the large difference between spin-orbit potentials in ${}^4\text{He}$ and neutron-rich helium isotopes.

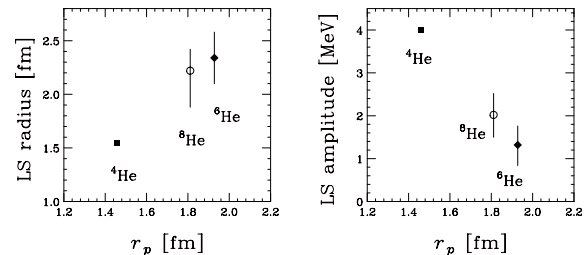


Figure 5: LS radius and LS amplitude of helium isotopes as functions of proton radius.

4 Summary

The vector analyzing power of the proton elastic scattering from ${}^6\text{He}$ and ${}^8\text{He}$ have been measured at 71 MeV/A. We found that the spin-orbit potentials in ${}^6\text{He}$ and ${}^8\text{He}$ are significantly shallower than those in stable nuclei. This shallowness can be naturally explained by the large diffuseness of the neutron-rich helium isotopes. The dominance of $p - \alpha$ core interaction in the $p - {}^6,8\text{He}$ spin-orbit potentials is also indicated.

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SPIN CONTROL BY RF FIELDS AT ACCELERATORS AND STORAGE RINGS

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Abstract

The first experiments to apply RF fields for resonant beam depolarization and spin flip at the VEPP-2M storage ring were carried out more than 30 years ago [1]. Later this technique was used at VEPP-2M in the experiment for comparison of electron and positron anomalous magnetic moments. Recently, interest in RF spin control has appeared at proton machines. This paper describes a general approach for consideration of RF influence on spin dynamics at electron (positron) and hadron accelerators. Some practical applications of RF fields are discussed.

The spin motion of a particle in electromagnetic fields is described by the BMT equation: [2]

$$\begin{aligned} \frac{d\mathbf{S}}{dt} &= \dot{\mathbf{S}} = [\mathbf{\Omega} \times \mathbf{S}] \\ -\mathbf{\Omega} &= \left(\frac{q_0}{\gamma} + q' \right) \mathbf{B}_\perp + \frac{q_0 + q'}{\gamma} \mathbf{B}_\parallel + \left(\frac{q_0}{\gamma + 1} + q' \right) [\mathbf{E} \times \mathbf{V}], \end{aligned} \quad (1)$$

where q_0 and q' are normal and anomalous parts of the particle gyro-magnetic ratio; $\mathbf{B}_\perp$ and $\mathbf{B}_\parallel$ are magnetic field components along and transverse to the particle velocity $\mathbf{V}$. Since in circular accelerators the one-turn energy change is relatively small, we can neglect, in the first approximation, the electric field: $\mathbf{E} = 0$. Following the usual approach for orbital motion, we use as the independent variable the generalized azimuth θ and subdivide the spin precession vector in two parts: $\mathbf{\Omega} = \mathbf{W}_0(\theta) + \mathbf{w}(\theta)$, where $\mathbf{W}_0(\theta)$ contains only fields on the Closed Orbit $\mathbf{R}_0$, while $\mathbf{w}(\theta)$ denotes *all* of the other terms (contributions from closed orbit imperfection and orbital oscillations). One can treat $\mathbf{w}$ as a small perturbation for the spin motion. ([3]- [5]) In the accelerator vector triad $\mathbf{e}_x$, $\mathbf{e}_y = \mathbf{V}/V$, $\mathbf{e}_z = [\mathbf{e}_x \times \mathbf{e}_y]$ components of the precession vector $\mathbf{W}_0$ can be presented in the linear approximation in the next form: ¹

$$\begin{aligned} W_{0x} &= \nu_0 K_x; & K_x &= \frac{B_x}{B_0}; \\ W_{0y} &= (1 + a) K_y; & K_y &= \frac{B_y}{B_0}; \\ W_{0z} &= \nu_0 K_z; & K_z &= \frac{B_z}{B_0}, \end{aligned} \quad (2)$$

where we introduce the particle magnetic anomaly $a = q'/q_0$ and denote $\nu_0 = \gamma \cdot a$. On the reference orbit the equation has one solution $\mathbf{n}_0$, which is periodic around the ring,

¹We use dimensionless units: fields are normalized to mean guiding field $B_0 = 1/2\pi \oint B_z d\theta$; length and time are measured correspondent in units of mean radius R and revolution time.

i.e. $\mathbf{n}_0(\theta + 2\pi) = \mathbf{n}_0(\theta)$. There are also two other linearly independent solutions: the orthogonal complex vectors $\boldsymbol{\eta}$ and $\boldsymbol{\eta}^*$, which rotate around $\mathbf{n}_0$ with spin tune ν :

$$\boldsymbol{\eta}(\theta + 2\pi) = e^{i2\pi\nu}\boldsymbol{\eta}(\theta); \quad \boldsymbol{\eta}^*(\theta + 2\pi) = e^{-i2\pi\nu}\boldsymbol{\eta}^*(\theta)$$

We restrict consideration to planar machines and start with a proton ring, where the spin tune is $\nu = \nu_0$ and $\mathbf{n}_0 = \mathbf{e}_z$; $\boldsymbol{\eta} = (\mathbf{e}_x - i\mathbf{e}_y)e^{i\nu_0\theta}$. Let's apply on a short piece of the orbit Δl a longitudinal RF-field $\tilde{K}_y \cos(\nu_d\theta)$ (RF-solenoid) with a frequency ν_d , which is an external spin perturbation. It can be presented as a number of circular harmonics $w = \sum_k w_k e^{i(k \pm \nu_d)\Theta}$ with equal amplitudes $w_k = (1 + a)\tilde{K}_y \frac{\Delta l}{4\pi R}$.

A more complicated situation occurs for the application of a radial RF-field $\tilde{K}_x \cos(\nu_d\theta)$ (RF-dipole). This field disturbs not only spin motion, but it excites also forced vertical oscillations. As a result, spin gets additional kicks from off-orbit fields in dipole and quadrupole magnets around the machine. The tight frame of this paper doesn't allow a mathematical description of this process and we refer readers to papers devoted to this topic. [7]- [8] The linear spin response formalism has been developed for *simultaneous* treating of the orbit and spin dynamics. Based on this formalism the code "ASPIRRIN" ([10]) calculates 5 complex response functions $F_1(\theta) - F_5(\theta)$ which satisfy the periodicity condition: $|F_i(\theta + 2\pi)| = |F_i(\theta)|$. These response functions characterize the sensitivity of the spin precession axis $\mathbf{n}$ to kicks of orbital variables ($X^T = (p_x, x, p_z, z, \delta\gamma/\gamma, 0)$) for specific machines optics and working points. In the case of an ideal flat machine, $\frac{\partial \mathbf{n}}{\partial p_z} = \nu_0 \text{Re}(i F_3(\theta) \boldsymbol{\eta}^*)$. Thereby the RF-dipole, applied at azimuth Θ , creates a set of perturbing harmonics with amplitudes $w_k = \nu_0 \tilde{K}_y F_3(\Theta) \frac{\Delta l}{4\pi R}$.

Next, it's necessary to take into account synchrotron oscillation of the particle energy, inherent to accelerators: $\gamma = \gamma_0 + \Delta\gamma \cos(\nu_s\theta)$. Usually, the synchrotron tune $\nu_s \sim 10^{-2} - 10^{-3}$ is much less than other orbital frequencies. Hence, the spin tune is modulated by the synchrotron tune. It results in a modification of the RF-field spectrum. Each line of the spin perturbation is transformed to a set of side bands resonances $k \pm \nu_d \pm m\nu_s$. The side band harmonics are given by the expression: [3]

$$|w_k^m| = |w_k| J_m(\lambda), \quad (3)$$

where $J_m(\lambda)$ are the Bessel functions. As a rule, the modulation index $\lambda = \frac{\nu_0}{\nu_s} \left(\frac{\Delta\gamma}{\gamma} \right) \ll 1$, at proton and electron accelerators. For RF spin resonances not all sidebands overlap ($w_k^m \ll \nu_s$). It's important to emphasize that the central line of the spectrum corresponds to the mean energy γ_0 of the beam. Moreover, the width of this spectrum line, averaged over the synchrotron oscillations, shrinks to a size of $\sigma_\nu \sim (\delta\gamma/\gamma)^2$. [14] Assuming a Gaussian particle distribution in the longitudinal phase space, mean value of the central resonance strength ($w_k^0 \equiv w$) and it's rms deviation σ_w are given by the expressions:

$$w^2 = w_k^2 I_0(\Lambda) e^{-\Lambda}; \quad \sigma_w^2 = w_k^2 \frac{\Lambda}{2} I_1(\Lambda) e^{-\Lambda}, \quad (4)$$

where I_0, I_1 are the modified Bessel functions from the argument $\Lambda = \frac{\nu_0}{\nu_s} \sigma_\gamma$; ($\sigma_\gamma^2 = \overline{\left(\frac{\Delta\gamma}{\gamma} \right)^2}$).

Evidently, spin control operations have to be done at the central resonance. The choice of frequency ν_d is determined by concrete conditions of an experiment, but always one must satisfy the resonance condition: $|\nu_d \pm k - \nu_0| \ll 1$. It's more visual to consider this situation in a frame rotating at the resonant frequency, where the spin motion is simply

precession in a "field" $\mathbf{h} = \epsilon \mathbf{e}_z + w \mathbf{e}_y$, where $\epsilon = \nu_0 - \nu_d \pm k$ is the resonance tune, see Fig. 1. The precession frequency $h = \sqrt{\epsilon^2 + w^2}$ in the resonance case ($\epsilon = 0$) decays to w_k , which can be taken as a strength of the resonance.

The spin resonance crossing was described by Froissart-Stora [12], who found that in result of passing from $\epsilon = \infty$ up to $\epsilon = -\infty$ with a tune rate $\dot{\epsilon} = \text{const}$ the residual projection averaged over spin phases $S_z = (\mathbf{S} \cdot \mathbf{n}_0)$ is described by a formula:

$$S_z = S_z(0)(2e^{-\frac{\pi w^2}{2\dot{\epsilon}}} - 1). \quad (5)$$

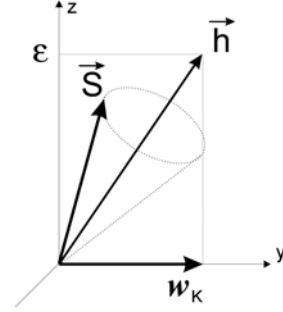


Figure 1: Resonant frame.

The final result depends on the spin phase advance near the resonance center: $\Psi = \int h dt \sim w^2/\dot{\epsilon}$. The polarization is preserved by an adiabatic change of parameters ($\Psi \gg 1$) and flips down together with the $\mathbf{n} = \mathbf{h}/h$ -axis. In the opposite case of a fast crossing ($\Psi \ll 1$), the spin only tilts slightly from its initial direction with a depolarization $\Delta S_n \simeq w^2/\dot{\epsilon}$. Polarization losses may attain $\sim 100\%$ in intermediate situations $\Delta\phi \leq 1$.

At electron accelerators there are other limitations for RF-device usage. Quantum fluctuations of the synchrotron radiation lead to so called, "spin diffusion", which is enhanced at spin resonances. [4] In a general case, a quantum emission in the resonant frame brings at the same instant jumps of the resonance tune ($\delta\epsilon$), the resonant harmonic (δw_k) and it's phase ($\delta\phi$), see Fig. 2. At that, the precession axis $\mathbf{n} = \mathbf{h}/h$ gets a kick $\delta\mathbf{n}$:

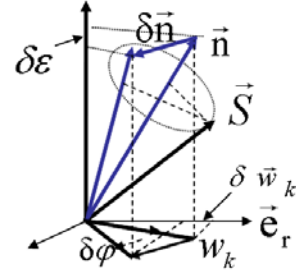


Figure 2: Spin diffusion.

$$(\delta\mathbf{n})^2 = (\delta \arctan(\frac{|w_k|}{\epsilon}))^2 + \frac{w_k^2}{h^2}(\delta\phi_k)^2, \quad (6)$$

but spin vector does not change (we neglect here "spin-flip" quantum emissions). However, the projection S_n undergoes a change $\delta S_n = -1/2 (\delta\mathbf{n})^2 S_n$. A consequence of random kicks results in a polarization loss with a decay time $\tau_d = 1/2 \langle \frac{d}{dt}(\delta\mathbf{n})^2 \rangle$, where the angle brackets $\langle \rangle$ denote an average over time and an ensemble.

In the case of the RF-resonance ($\delta\phi = 0$), we assume the time of the resonance crossing $\Delta t \sim w/\dot{\epsilon}$ is much longer than the radiative damping time τ_0 and the characteristic times of the orbital motion. Therefore we can consider, in average, numerous jumps of $\delta\epsilon$ and δw only as a diffusion around mean values $\bar{\epsilon}$ and $\bar{w}$ with corresponding rms deviations σ_ν , and σ_w . So, from (6) we find the depolarization time, caused by the quantum fluctuations:

$$\tau_d^{-1} = \frac{1}{2} \left\langle \frac{d}{dt} \left(\frac{\epsilon \delta w - w \delta \epsilon}{w^2 + \epsilon^2} \right)^2 \right\rangle \simeq \frac{\bar{\epsilon}^2 \sigma_w^2 + \bar{w}^2 \sigma_\nu^2}{(\bar{w}^2 + \bar{\epsilon}^2)^2} \cdot \tau_0^{-1}. \quad (7)$$

To obtain above expression, we ignored the interference term by the averaging over time ($\overline{\delta\epsilon \cdot \delta w} = 0$) and used usual formulas for an equilibrium state: $\sigma_\epsilon^2 = 1/2 \langle \frac{d}{dt}(\delta\epsilon)^2 \rangle \cdot \tau_0$; $\sigma_w^2 = 1/2 \langle \frac{d}{dt}(\delta w)^2 \rangle \cdot \tau_0$ with σ_w from (4), where $\sigma_\gamma^2 = 1/2 \langle \frac{d}{dt}(\delta\gamma)^2 \rangle \cdot \tau_0$. At the next step we employ the FS-formula for electron machines, taking into account the spin diffusion under the condition of adiabatic crossing ($\psi \gg \frac{w^2}{\dot{\epsilon}}$). We calculate the beam polarization $\zeta(t) = \langle S_n \rangle = \zeta(0) \cdot \int_0^t e^{-t/\tau_d} dt$, while resonance crossing with different

w and $\dot{\epsilon}$. For example, we use parameters of VEPP-2M storage ring ($E=700$ MeV). The initial polarization is always $\zeta(0) = 1$. Fig.3 shows three curves $\zeta(t)$ versus $\epsilon(t)$ for different RF resonance amplitudes and spin tune spreads but for the same crossing rate $\dot{\epsilon} \cdot f_0 = 400$ Hz/sec: (solid line) $w = 1 \cdot 10^{-5}$; $\sigma_\nu = 3 \cdot 10^{-7}$; (dotted line) $w = 4 \cdot 10^{-5}$; $\sigma_\nu = 1 \cdot 10^{-6}$; and (dashed line) $w = 1 \cdot 10^{-4}$; $\sigma_\nu = 3 \cdot 10^{-6}$.

These calculations demonstrate clearly the influence of the spin diffusion and spin tune spread. Even increasing the RF amplitude by ten times does not help to avoid some depolarization, when the spin tune spread grows up to three times (compare curves in Fig.3). So, using such "simulations", it's possible to choose RF device parameters for successful spin flip. Moreover, the measurement of a residual polarization in the case of a reasonable polarization loss appears as a way to measure the spin tune spread and minimize it, if that is necessary. [14]

It's interesting also to study the opposite case of a small RF amplitude. Fig.4 presents three other curves, where we fixed the spin tune spread $\sigma_\nu = 1 \cdot 10^{-6}$ and the crossing rate $\dot{\epsilon} \cdot f_0 = 2$ Hz/sec, but changed the amplitude w : (solid line) $w = 2 \cdot 10^{-7}$; (dotted line) $w = 5 \cdot 10^{-7}$ and (dashed line) $w = 1 \cdot 10^{-6}$. One can see from Fig.4 the resonant depolarization by the RF-field. Decreasing the RF power provides a measurement of the spin tune with accuracy up to its spread σ_ν . In turn, the spin tune determination is simultaneously the absolute mean energy measurement, since the magnetic anomalies are

well known. [13] For instance: $a_e = 1.159652193 \times 10^{-3}$. Figure 4: Resonant depolarization

Beam energy calibration has been routinely used at electron-positron colliders in precise experiments for secondary particle mass measurements. [15] The coherent spin rotation by 90 degrees and full spin flip were crucially important at VEPP-2M in the experiment for electron and positron anomalous magnetic moments comparison. [16]

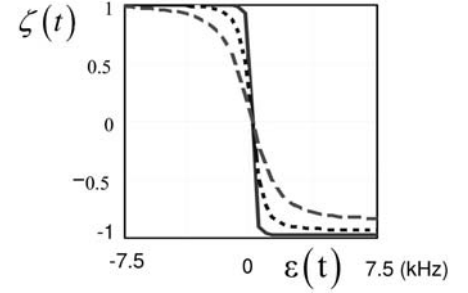


Figure 3: Spin flip by RF.

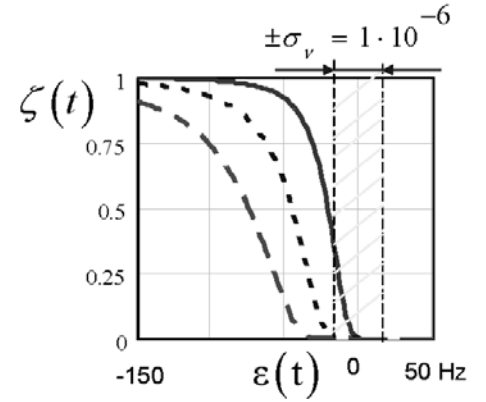


Figure 4: Resonant depolarization

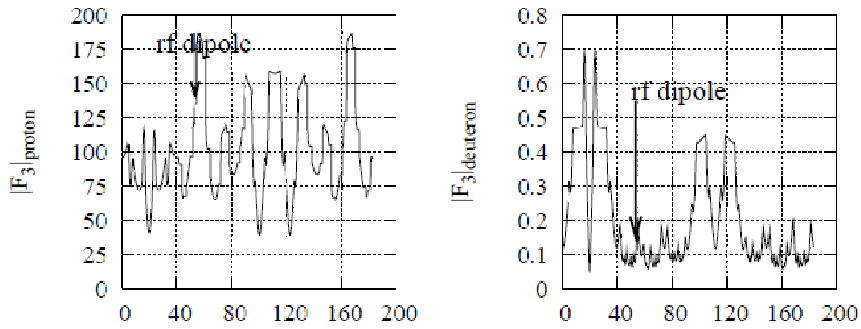


Figure 5: $|F_3|$ along the COSY orbit for protons and deuterons.

Till now, RF spin flip has been studied for protons and deuterons at the IUCF storage ring (see, for instance [17]) and at the synchrotron COSY. [18]. At first, the authors pointed out spin flip by RF-fields was going, as a rule, with very low polarization losses ($\ll 10^{-2}$ per pass). At both machines the experiments were carried out with RF-solenoids and RF-dipoles and the authors claimed to find “unexpected discrepancies” between the measured values and “the theoretical expressions for the spin flip resonance widths”. However, recent analysis of the COSY data, based on the formalism of the spin response functions, has found an amazing accordance of the experiment and the theory as for both kinds of the flippers, for both types of particles. [8] The main reason for the “unexpected discrepancies” is explained by Fig.5, which presents results of ASPIRRIN calculations of $|F_3|$ values along the COSY orbit for protons and deuterons, that are distinguished in more than 1000 times in the point of the RF flipper.

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RELATED PROBLEMS

REGULARIZATION OF SOURCE OF THE KERR-NEWMAN ELECTRON BY HIGGS FIELD

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Abstract

Regular superconducting solution for interior of the Kerr-Newman (KN) spinning particle is obtained. For parameters of electron it represents a highly oblated rotating bubble formed by Higgs field which expels the electromagnetic (em) field and currents from interior of the bubble to its domain wall boundary. The external em and gravitational fields correspond exactly to Kerr-Newman solution, while interior of the bubble is flat and forms a ‘false’ vacuum with zero energy density. Vortex of em field forms a quantum Wilson loop on the edge of the rotating disk-like bubble.

1.Introduction. Kerr-Newman (KN) solution for a charged and rotating Black-hole has $g = 2$ as that of the Dirac electron and paid attention as a classical model of electron coupled with gravity [1–7]. For parameters of electron (in the units $G = C = \hbar = 1$) $J = 1/2$, $m \sim 10^{-22}$, $e^2 \sim 137^{-1}$ the black-hole horizons disappear and the Kerr singular ring of the Compton radius $a = \hbar/2m \sim 10^{22}$ turns out to be naked. This ring forms the gate to a negative sheet of the Kerr geometry, significance of which is the old trouble for interpretation of the source of Kerr geometry. The attempts by Brill and Cohen to match the Kerr exterior with a rotating spherical shell and flat interior [8] had not lead to a consistent solution, and Israel suggested to truncate the negative sheet, replacing it by a thin (rotating) disk spanned by the Kerr ring [2]. This model led to a consistent source, however, it turned out to be build from a very exotic superluminal matter, besides, the Kerr singular ring appeared to be naked at the edge of the disk. López [4] transformed this model to an ellipsoidal rotating shell covering the Kerr ring. By means of especial choice of the boundary of the bubble, $r = r_0 = e^2/2m$ (r is the Kerr oblate spheroidal coordinate), he matched continuously the Kerr exterior with the flat interior, and therefore, he obtained a consistent KN source without the KN singularity. However the López model, like the other models, was not able to explain the origin of Poincaré stress, and a negative pressure was introduces by López phenomenologically. Meanwhile, the necessary tangential stresses appear naturally in the domain wall field models, in particular, in the models based on the Higgs field [5]. It suggested that a consistent field descriptions of this problem should contain, along with the KN Einstein-Maxwell sector, the sector of Higgs fields and sector of interaction between the em field (coupled to gravity) and Higgs fields corresponding to superconducting properties of the KN source. Up to our knowledge, despite several attempts and partial results, no one has been able so far to obtain a consistent field model of the KN source related with Higgs field. In this paper we present, apparently, first solution of this sort which is consistent in the limit of thin domain wall boundary of the bubble. The considered KN source represents a generalization of the López model and incorporates the results obtained earlier in [9, 5]. Although we consider here only the case of thin domain wall, generalization to the case of a finite thickness is also possible by methods described earlier in [5]. The described in [5] KN source was a bag formed by

a potential $V(r)$ interpolating between the external ‘true’ vacuum $V^{(ext)} = 0$ and a ‘false’ (superconducting) vacuum $V^{(in)} = 0$ inside the bag. It was shown that the corresponding Higgs sector may be described by $U(1) \times \tilde{U}(1)$ Witten field model with the given by Morris in [11] superpotential. It was shown also in [5, 9] that consistency of the Einstein-Maxwell sector with such a phase transition may be perfectly performed in the KS formalism with use of the Gürses and Gürsey ansatz [10]. However, consistent solution for interaction between the em and Higgs fields was not obtained in [5], and in this talk we improve this deficiency.

2. *Phase transition in gravitational sector.* Following [5, 9], for external region we use the exact KN solution in the Kerr-Schild (KS) form of metric

$$g_{\mu\nu} = \eta_{\mu\nu} + 2Hk_\mu k_\nu, \quad (1)$$

where $\eta^{\mu\nu}$ is metric of the auxiliary Minkowski background in Cartesian coordinates $x^\mu = (t, x, y, z)$. Electromagnetic (em) KN field is given by vector potential

$$A_{KN}^\mu = Re \frac{e}{r + ia \cos \theta} k^\mu, \quad (2)$$

where $k^\mu(x^\mu)$ is the null vector field which is tangent to a vortex field of null geodesic lines, the Kerr principal null congruence (PNC). For the KN solution function H has the form

$$H = \frac{mr - e^2/2}{r^2 + a^2 \cos^2 \theta}. \quad (3)$$

The used by Kerr especial spherical oblate coordinates r, θ, ϕ_K , are related with the Cartesian coordinates as follows

$x + iy = (r + ia)e^{i\phi_K} \sin \theta$, $z = r \cos \theta$. Vector field k^μ is represented in the form [1]

$$k_\mu dx^\mu = dr - dt - a \sin^2 \theta d\phi_K. \quad (4)$$

For the metric *inside of the oblate bubble* we use the KS ansatz (1) in the form suggested by Gürses and Gürsey [10] with function

$$H = \frac{f(r)}{r^2 + a^2 \cos^2 \theta}.$$

We set for the interior $f(r) = f_{int} = \alpha r^4$, which provides suppressing of the Kerr singularity and a constant curvature of the internal space-time.

For exterior, $r > r_0$, we use $f(r) = f_{KN} = mr - e^2/2$ corresponding to KN solution. Therefore, $f(r)$ describes a phase transition of the KS metric from ‘true’ to ‘false’ vacuum (see fig.1). We assume that the zone of phase transition, $r \approx r_0$, is very thin and metric is continuous there, so the point of intersection, $f_{int}(r_0) = f_{KN}(r_0)$, determines position of domain wall r_0 graphically and yields the ‘balance matter equation’ [5, 9],

$$m = m_{em}(r_0) + m_{mat}(r_0), \quad m_{em}(r_0) = \frac{e^2}{2r_0}, \quad (5)$$

which determines r_0 analytically. Interior has constant curvature, $\alpha = 8\pi\Lambda/6$. For parameters of electron $a = J/m = 1/2m \gg r_0 = e^2/2m$ and the axis ratio of the ellipsoidal bubble is $r_0/a = e^2 \sim 137^{-1}$, so the bubble has the form of highly oblated disk.

3. *Brief summary of the Higgs sector* [5].

The corresponding phase transition is provided by Higgs model with two complex Higgs field Φ and Σ , two related gauge fields A^μ and B^μ , and one auxiliary real field Z . This is given by Morris [11] generalization of the $U(1) \times \tilde{U}(1)$ field model used by Witten for superconducting strings [12]. The potential $V(r) = \sum_i |\partial_i W|^2$, where $\partial_1 = \partial_\Phi$, $\partial_2 = \partial_Z$, $\partial_3 = \partial_\Sigma$, is determined by superpotential $W = \lambda Z(\Sigma\bar{\Sigma} - \eta^2) + (cZ + m)\Phi\bar{\Phi}$, where c , m , η , λ are real constants. It forms a domain wall interpolating between the internal ('false') and external ('true') vacua. The vacuum states obey the conditions $\partial_i W = 0$ which yield $V = 0$ i) for 'false' vacuum ($r < r_0$): $Z = -m/c$; $\Sigma = 0$; $|\Phi| = \eta\sqrt{\lambda/c}$, as well as ii) for 'true' vacuum ($r > r_0$): $Z = 0$; $\Phi = 0$; $\Sigma = \eta$.

4. *Interaction between the KN and Higgs fields.* We set $B^\mu = 0$ and use only the gauge fields A^μ which interacts with the Higgs field $\Phi(x) = |\Phi(x)|e^{i\chi(x)}$ having a nonzero vev inside of the bubble, $|\Phi(x)|_{r < r_0} = \Phi_0$, and vanishing outside. The related part of the field model is a copy of the Nielsen-Olesen (NO) field model for a vortex string in superconducting media [13], however, there is principal difference in topology, since we describe superconducting interior of the bubble, contrary to superconducting external media of the NO model. The Lagrangian $\mathcal{L}_I = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}(\mathcal{D}_\mu\Phi)(\mathcal{D}^\mu\Phi)^*$, where $\mathcal{D}_\mu = \nabla_\mu + ieA_\mu$, leads to the equations $\nabla_\nu\nabla^\nu A_\mu = I_\mu = \frac{1}{2}e|\Phi|^2(\chi_{,\mu} + eA_\mu)$. In external region $|\Phi| = 0$ and the em field has to correspond to exact KN solution, $A^\mu \equiv A_{KN}^\mu$. Following López, we fix the boundary of bubble at $r_0 = r_e = e^2/2m$ (one half of the classical radius of electron), which yields flat interior, $\alpha = 0$, and allows us to use flat d'Alembertian and set $\mathcal{D}_\mu = \partial_\mu + ieA_\mu$ for $r < r_0$. Current has to be expelled from interior of the bubble to its boundary, and we set in interior $I_\mu = 0$, which yields

$$\square A_\mu = 0 = e|\Phi|^2(\chi_{,\mu} + eA_\mu^{(in)}). \quad (6)$$

The KN gauge field A_μ is given by (2) and (4). Inside of the bubble the specific Kerr angular coordinate ϕ_K turns out to be inconsistent with the simple angular coordinate of the Higgs field, $\phi = -i\ln[(x + iy)/\rho]$, $\rho = (x^2 + y^2)^{1/2}$. Using the relation between corresponding differentials $d\phi_K = \phi + \frac{adr}{r^2 + a^2}$, we transform vector potential on the boundary and inside of the bubble to the form

$$A_\mu dx^\mu = \frac{-er}{r^2 + a^2 \cos^2 \theta} [dt + a \sin^2 \theta d\phi] + \frac{2erdr}{(r^2 + a^2)} \quad (7)$$

which shows that radial component A_r is full differential. At the external side of the bubble $r = r_e + 0 = e^2/2m$, the value of potential is $A_\mu dx^\mu|_{r=r_e+0} = \frac{-er_e}{r_e^2 + a^2 \cos^2 \theta} [dt + a \sin^2 \theta d\phi] + \frac{2er_e dr}{(r_e^2 + a^2)}$. Note that the lines of KN vector-potential in equatorial plane are tangent to the Kerr singular ring, and by $\cos \theta = 0$, the potential takes the form

$$A_\mu^{(str)} dx^\mu = -\frac{2m}{e} [dt + ad\phi] + \frac{2er_e dr}{(r_e^2 + a^2)}. \quad (8)$$

In particular, the tangent component at the edge is $A_\phi^{(str)} = -2ma/e$. Since $J = ma$ we have $A_\phi^{(str)} = -2J/e$. Setting $J = \frac{n}{2}$, $n = 1, 2, \dots$ we find out that vector potential

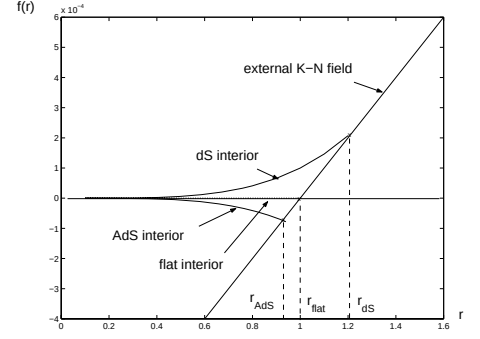


Figure 1: Matching of the metric of regular bubble interior with external KN field.

forms on the edge of bubble a closed quantized Wilson loop $S = \oint eA_\phi^{(str)} d\phi = -2\pi n$, which has to be matched with angular periodicity of the Higgs field $\Phi = \Phi_0 \exp(i\chi)$ and fixes its ϕ -dependence, $\Phi \sim \exp\{in\phi\}$. For the time-like component of A_μ inside of the bubble, the r.h.s. of (6) states $\chi_{,0} = -eA_0^{(in)}(r)$, which determines the $\chi_{,0}$ to be a constant corresponding to frequency of oscillations of the Higgs field, $\chi_{,0} = \omega = -eA_0^{(str)} = 2m$. Radial component of the KN field is a full differential, and being extended inside the bubble, it is compensated by Higgs field in agreement with the r.h.s. of (6). Therefore, the Higgs field acquires the form $\Phi(x) = \Phi_0 \exp\{i\omega t - i \ln(r^2 + a^2) + in\phi\}$.

For exclusion of the region of string-like loop at equator, $\cos\theta = 0$, the timelike and ϕ components of the gauge field have a chock crossing the boundary of bubble, which determines a distribution of circular currents over the bubble boundary.

5. *Consistency.* We find out that inside of the bubble $\mathcal{D}_\mu \Phi = i\Phi[\partial_\mu \chi + eA_\mu^{(in)}] \equiv 0$. Together with the result that $V^{(in)} = 0$, it leads to vanishing of the stress-energy tensor of matter inside of the bubble, $T_{\mu\nu}^{(int)} = (D_\mu \Phi)(\overline{D_\nu \Phi}) - \frac{1}{2}g_{\mu\nu}[(D_\lambda \Phi)(\overline{D^\lambda \Phi})]$, and provides the flatness of interior in agreement with our assumptions.

The obtained solution is regular and consistent in the limit of thin domain wall. It has important peculiarities: 1) the external KN field is spinning and gravitating, 2) the Higgs field is oscillating, 3) a quantum Wilson loop appears on the boundary of source.

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ABOUT SPIN PARTICLE SOLUTION IN BORN-INFELD NONLINEAR ELECTRODYNAMICS

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Abstract

The axisymmetric static solution of Born-Infeld nonlinear electrodynamics with ring singularity is investigated. This solution is considered as a static part of massive charged particle with spin and magnetic moment. The method for obtaining the appropriate approximate solution is proposed. An approximate solution is found. The values of spin, mass, and magnetic moment is obtained for this approximate solution.

The purpose of the present work is construction of the field model for massive charged elementary particle with spin and magnetic moment as a soliton solution of nonlinear electrodynamics.

In this approach the mass and the spin of the particle have fully field nature and must be calculated with integration of the appropriate densities over three-dimensional space.

The charge and the magnetic moment characterize the singularities and the behaviour of the fields at space infinity.

This theme was discussed in my articles (see [1–4]).

Here we considered the appropriate soliton solution with ring singularity in Born-Infeld nonlinear electrodynamics.

For inertial reference frames and in the region outside of field singularities, the equations are

$$\begin{aligned} \operatorname{div} \mathbf{B} &= 0, & \operatorname{div} \mathbf{D} &= 0, \\ \frac{\partial \mathbf{B}}{\partial t} + \operatorname{curl} \mathbf{E} &= 0, & \frac{\partial \mathbf{D}}{\partial t} - \operatorname{curl} \mathbf{H} &= 0, \end{aligned} \quad (1)$$

where

$$\begin{aligned} \mathbf{D} &= \frac{1}{\mathcal{L}} (\mathbf{E} + \chi^2 \mathcal{J} \mathbf{B}), & \mathbf{H} &= \frac{1}{\mathcal{L}} (\mathbf{B} - \chi^2 \mathcal{J} \mathbf{E}) \\ \mathcal{L} &\equiv \sqrt{|1 - \chi^2 \mathcal{I} - \chi^4 \mathcal{J}^2|}, & \mathcal{I} &= \mathbf{E} \cdot \mathbf{E} - \mathbf{B} \cdot \mathbf{B}, & \mathcal{J} &= \mathbf{E} \cdot \mathbf{B}. \end{aligned} \quad (2)$$

Here $\mathbf{E}$ and $\mathbf{H}$ are electrical and magnetic field strengths, $\mathbf{D}$ and $\mathbf{B}$ are electrical and magnetic field inductions.

Mass and spin is defined as three dimensional space integral from the appropriate densities:

$$m = \int_V \mathcal{E} dv, \quad s = \left| \int_V \mathbf{r} \times \mathcal{P} dv \right|, \quad (3)$$

where $\mathbf{P} = \frac{1}{4\pi} \mathbf{D} \times \mathbf{B}$ is the Poynting vector.

For Born-Infeld electrodynamics we have the following energy density:

$$\mathcal{E} = \frac{1}{4\pi\chi^2} \left(\sqrt{1 + \chi^2 (\mathbf{D}^2 + \mathbf{B}^2) + (4\pi)^2 \chi^4 \mathbf{P}^2} - 1 \right). \quad (4)$$

The behaviour of electrical and magnetic fields for particle solution at infinity is characterized by electrical charge and magnetic moment. For more details see my paper [3] In spherical coordinates the field components have the following form for $r \rightarrow \infty$:

$$\{D_r, D_\vartheta, D_\varphi\} \sim \{E_r, E_\vartheta, E_\varphi\} \sim \left\{ \frac{e}{r^2}, 0, 0 \right\}, \quad (5)$$

$$\{H_r, H_\vartheta, H_\varphi\} \sim \{B_r, B_\vartheta, B_\varphi\} \sim \left\{ \frac{2\mu \cos \vartheta}{r^3}, \frac{\mu \sin \vartheta}{r^3}, 0 \right\}. \quad (6)$$

The electrical charge and the magnetic moment characterize also the behaviour of fields near singularities.

Let us consider the toroid symmetry field configuration. We use the toroidal coordinate system $\{\xi \in [0, \infty], \eta \in (-\pi, \pi], \varphi \in (-\pi, \pi]\}$ which is obtained with rotation of bipolar coordinate system around the standing axis. It is convenient to use the new variable $s \equiv \frac{1}{\cosh \xi} \in [0, 1]$ instead of the variable ξ .

Let us consider the appropriate solution for linear electrodynamics ($\mathbf{D} = \mathbf{E}$, $\mathbf{H} = \mathbf{B}$).

The components of electromagnetic potential for the linear case have the following form:

$$A_0 = -\frac{e}{\rho_0 \sqrt{2} s} \sqrt{1 - s \cos \eta} P_{-\frac{1}{2}}^0(1/s), \quad (7)$$

$$A_\varphi = -\frac{\mu 2\sqrt{2}}{\rho_0^2 \sqrt{s}} \sqrt{1 - s \cos \eta} i P_{-\frac{1}{2}}^1(1/s). \quad (8)$$

where $P_n^m(x)$ is associated Legendre function.

The electrical and magnetic fields obtained from the potential $\{A_0, 0, 0, A_\varphi\}$ (7), (8) has the right behaviour at infinity ($r \rightarrow \infty$ in spherical coordinates) (5). The behaviour of the field near the singular ring ($s \rightarrow 0$) is the following:

$$\{E_\xi, E_\eta, E_\varphi\} = \{D_\xi, D_\eta, D_\varphi\} = \left\{ -\frac{e}{\pi \rho_0^2 s}, 0, 0 \right\} + o(s^{-1}), \quad (9)$$

$$\{B_\xi, B_\eta, B_\varphi\} = \{H_\xi, H_\eta, H_\varphi\} = \left\{ 0, -\frac{2\mu}{\pi \rho_0^3 s}, 0 \right\} + o(s^{-1}). \quad (10)$$

Let us search the appropriate solution for Born-Infeld electrodynamics in the following form:

$$A_0 = -\frac{e}{\rho_0 \sqrt{2}} \sqrt{1 - s \cos \eta} f_1(s, \eta), \quad (11)$$

$$A_\varphi = \frac{\mu \sqrt{2}}{4 \rho_0^2} \sqrt{1 - s^2} \sqrt{1 - s \cos \eta} f_2(s, \eta), \quad (12)$$

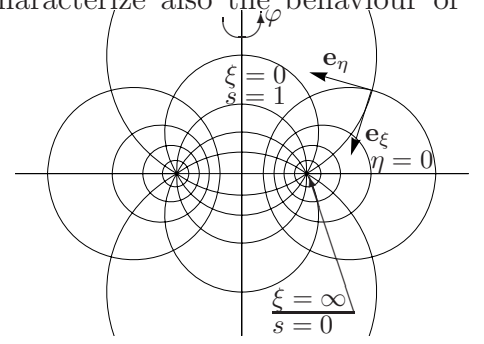


Figure 1: The section ($\varphi = 0 \cup \varphi = \pi$) of the toroidal coordinate system $\{\xi, \eta, \varphi\}$ ($\{s, \eta, \varphi\}$).

where the functions $f_i(s, \eta)$ must be found.

We take the condition of finiteness for the potentials (A_0, A_φ) and fields $(\mathbf{E}, \mathbf{B})$ as physically defensible.

We can seek the functions $f_i(s, \eta)$ in the form of formal double series, that is the power series in s and the Fourier series in η . But here we have the problem such that the Lagrangian $\mathcal{L}$ for an appropriate approximate solution can become nonanalytic function for sufficiently large values of variable s . To avoid the non-analyticity we must take the condition $(1 - \chi^2 \mathcal{I} - \chi^4 \mathcal{J}^2) \geq 0$. This condition can be satisfied with the help of exponential smoothing factors. Thus let the appropriate approximate solution be have the following form:

$$f_i(s, \eta) \approx f_i^{NK} = c_{i0} e^{-q_{i0} s^{N+1}} + \sum_{n=1}^N \sum_{k=0}^K c_{ink} e^{-q_{ijk} s^{N+1}} s^n \cos k\eta, \quad (13)$$

where the integers N and M characterize the order of the approximation, $q_{i0} \geq 0$, $q_{ink} \geq 0$.

Then we substitute the field functions $\{\mathbf{E}, \mathbf{B}\}$ obtained from the potentials (11), (12) with the functions (13) to equations

$$\text{div} \mathbf{D} = 0, \quad \text{curl} \mathbf{H} = 0 \quad (14)$$

with Born-Infeld relations $\mathbf{D}(\mathbf{E}, \mathbf{B})$ and $\mathbf{H}(\mathbf{E}, \mathbf{B})$ (2).

We have the fields near $s = 0$ in the following form:

$$\{E_\xi, E_\eta, B_\xi, B_\eta\} = \left\{ -\frac{e \sqrt{\rho_0^4 + \beta^2}}{\rho_0^2}, 0, 0, -\frac{e \beta}{\rho_0^2} \right\} + \mathcal{O}(s), \quad (15)$$

$$\{D_\xi, D_\eta, H_\xi, H_\eta\} = \left\{ -\frac{e}{\pi \rho_0^2 s}, 0, 0, -\frac{e \beta}{\pi \rho_0^2 s \sqrt{\rho_0^4 + \beta^2}} \right\} + \mathcal{O}(1), \quad (16)$$

where $\beta = \mu \left/ \frac{e \rho_0}{2} \right.$.

As we can see the fields $\{\mathbf{E}, \mathbf{B}\}$ near $s = 0$ is finite and the fields $\{\mathbf{D}, \mathbf{H}\}$ have the behaviour of the same type that in the linear case (9), (10).

The extraction of the coefficients for s^n and k -th Fourier harmonic in equations (14) allows to have the linear systems for obtaining the coefficients c_{ink} .

The coefficients c_{i0} must be obtained from the conditions

$$f_i(1, 0) = 1 \quad (17)$$

These relation satisfy the conditions (5) and (6) at space infinity ($r \rightarrow \infty$).

The approximate solution of the form (13) was found for $N = 3$ and $K = 3$.

The obtained approximate solution in appropriate units ($e = 1$, $r_0 = \sqrt{|e \chi|} = 1$) has the following free parameters: radius of the ring ρ_0 , magnetic moment μ , smoothing coefficients q_{i0} , q_{ijk} .

But the solution of Born-Infeld electrodynamics under investigation should not have continuous free parameters except for the electrical charge. Thus we must have a set of free parameters appropriate for the best approximation to solution.

The appropriate values of the free parameters was found by the direct numerical minimization of the action functional $\int (\mathcal{L} - 1) dv$ with the condition $(1 - \chi^2 \mathcal{I} - \chi^4 \mathcal{J}^2) \geq 0$.

The action with the approximate solution under consideration reach a local minimum for the following values:

$$\rho_0 \approx 0.95 r_0 \quad , \quad \mu \approx 0.80 \frac{e \rho_0}{2} \quad . \quad (18)$$

We have the following mass and spin for the obtained field configuration:

$$m \approx \frac{10}{\chi^2} \quad , \quad s \approx 0.1 \frac{e^2}{2\alpha} \quad , \quad (19)$$

where α is the fine structure constant and $e^2/(2\alpha) = \hbar/2$.

It should be noted that the value $\mu \sim e \rho_0/2$ can be considered as desirable for this model (see my article [5]).

Thus the obtained result is encouraging. But of course it must be considered as preliminary.

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SPINDYNAMICS

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Abstract

Geometrization of General Theory of Relativity and Quantum Mechanics leads to Gravidynamics and Spindynamics. The key points of spindynamics are spin symmetry (the fundamental realization of the concept of geometrical internal symmetry) and the spin field (space of defining representation of spin symmetry). The essence of spin is the bipolar structure of spin symmetry induced by the gravitational potential. The bipolar structure provides natural violation of spin symmetry and leads to the equations of spindynamics. Here we consider strong interactions and confinement as natural elements of spindynamics.

1. Introduction. The Standard Model offers to explain some experimental facts but it does not claim to find the origin of physical laws being one more attempt to describe Nature in a possibly most effective way. The nature and origin of new natural system of fundamental physical laws are defined by new concept of physical space, new concept of time and independence of any external or a priori conditions (principles of geometrization and self-organization of absolutely closed system) [1],[2],[3],[4]. Spindynamics involves all phenomena connected with spin and hence provides the new understanding of the artificial concepts of weak and strong isotopic spin and the Standard Model as a hole. The geometrical and physical nature of the strong interactions and confinement can be understood only in the framework of the nontrivial causal structure defining by the temporal field. The basic equation of temporal field has general solution and singular solution that in general is defined as geodesic distance [5]. The second solution is the starting point for understanding of the strong interactions and confinement on the basis of the new causal structure tightly connected with rotations.

2. Equations of Spindynamics. In Gravidynamics and Spindynamics the properties of time are defined by temporal field. The temporal field with respect to the coordinate system q^1, q^2, q^3, q^4 is denoted by $f(q) = f(q^1, q^2, q^3, q^4)$. The gradient of the temporal field (the stream of time) is the vector field $\mathbf{t}$ with the components

$$t^i = g^{ij} \frac{\partial f}{\partial q^j} = g^{ij} \partial_j f = g^{ij} t_j.$$

Equation of the temporal field is

$$D_{\mathbf{t}} f = g^{ij} \frac{\partial f}{\partial q^j} \frac{\partial f}{\partial q^j} = 1. \quad (1)$$

This equation means that intrinsic time of physical system flows equably without relation to anything external. We remind some definitions of the vector algebra and vector analysis

in the four-dimensional and general covariant form [1]. The operator rot is defined for the vector fields as follows:

$$(\text{rot}\mathbf{M})^i = e^{ijkl}t_j\partial_k M_l = \frac{1}{2}e^{ijkl}t_j(\partial_k M_l - \partial_l M_k),$$

where e_{ijkl} is the orientation of physical space ("element of volume"). It is easy to show that

$$(\mathbf{M}, \text{rot}\mathbf{N}) + \text{div}[\mathbf{MN}] = (\text{rot}\mathbf{M}, \mathbf{N}),$$

where

$$[\mathbf{MN}]^i = e^{ijkl}t_j M_k N_l$$

is a vector product of two vector fields $\mathbf{M}$ and $\mathbf{N}$, $\text{div}\mathbf{M} = \nabla_i M^i$. Thus, the operator rot is self-adjoint. We also mention that $(\text{grad}\varphi)_i = \Delta_i \varphi$, $\Delta_i = \nabla_i - t_i \nabla_{\mathbf{t}}$ and $\text{rot grad} = 0$, $\text{div rot} = 0$.

Equations of spindynamics include four scalar and four vector equations:

$$\begin{aligned} (\nabla_{\mathbf{t}} + \frac{1}{2}\varphi)\kappa &= \text{div}\mathbf{K} - m\mu \\ (\nabla_{\mathbf{t}} + \frac{1}{2}\varphi)\lambda &= \text{div}\mathbf{L} - m\nu \\ (\nabla_{\mathbf{t}} + \frac{1}{2}\varphi)\mu &= \text{div}\mathbf{M} + m\kappa \\ (\nabla_{\mathbf{t}} + \frac{1}{2}\varphi)\nu &= \text{div}\mathbf{N} + m\lambda \end{aligned} \quad (2)$$

$$\begin{aligned} (\nabla_{\mathbf{t}} + \frac{1}{2}\varphi)\mathbf{K} &= -\text{rot}\mathbf{L} + \text{grad}\kappa + m\mathbf{M} \\ (\nabla_{\mathbf{t}} + \frac{1}{2}\varphi)\mathbf{L} &= \text{rot}\mathbf{K} + \text{grad}\lambda + m\mathbf{N} \\ (\nabla_{\mathbf{t}} + \frac{1}{2}\varphi)\mathbf{M} &= \text{rot}\mathbf{N} + \text{grad}\mu - m\mathbf{K} \\ (\nabla_{\mathbf{t}} + \frac{1}{2}\varphi)\mathbf{N} &= -\text{rot}\mathbf{M} + \text{grad}\nu - m\mathbf{L}, \end{aligned} \quad (3)$$

where $\nabla_{\mathbf{t}} = t^i \nabla_i$, $\varphi = \nabla_i t^i$. From the first principles it follows that Equations (2) and (3) describe all phenomena connected with spin symmetry and spin.

3. Strong Interactions as subject of Spindynamics. We can consider (following the successive approximations method) that physical space is 4-dimensional Euclidean space E^4 with the Euclidean distance function. Evolution and causality in the physical space E^4 is defined by the scalar temporal field $f(q_1, q_2, q_3, q_4)$ which is a solution to the equation (1). In this case, Equation (1) has the general solution $f(q_1, q_2, q_3, q_4) = \mathbf{a} \cdot \mathbf{q} + a$, where $\mathbf{a} = (a_1, a_2, a_3, a_4)$ is a unit constant vector,

$$\mathbf{a} \cdot \mathbf{a} = 1,$$

and the singular solution

$$f(q_1, q_2, q_3, q_4) = \sqrt{(\mathbf{q} \cdot \mathbf{q})} = \sqrt{q_1^2 + q_2^2 + q_3^2 + q_4^2}.$$

The space cross-section of E^4 is defined by the field of time. For a given real number t , the space cross-section is defined by the equation $f(q^1, q^2, q^3, q^4) = t$. Thus, in the first case the space sections of E^4 are the Euclidean space E^3 and in the second one the space sections are the 3d spheres $\sqrt{\mathbf{q} \cdot \mathbf{q}} = t$. For the stream of time we have

$$t^i = t_i = \frac{\partial f}{\partial q_i} = \frac{q^i}{f}, \quad \partial_i t^i = \frac{3}{f}. \quad (4)$$

State of hadron matter is defined the operators M_ν , N_ν , D and the well-known homogenous harmonic polynomials

$$D_{mn}^j(q_1, q_2, q_3, q_4) = \sqrt{(j+m)!(j-m)!(j+n)!(j-n)!} \times \sum_k (-1)^{m+k} \frac{(q_1 + iq_2)^{m-n+k} (q_1 - iq_2)^k (q_3 + iq_4)^{j-m-k} (q_3 - iq_4)^{j+n-k}}{(m-n+k)!k!(j-m-k)!(j+n-k)!} \quad (5)$$

that form the full orthogonal system of functions on S^3 and are the eigenfunctions of these operators.

It is almost evident that the equations of spindynamics corresponding to the first causal structure describe the so-called electroweak interactions and in the second case these equations describe the strong interactions. Confinement actually means a new physical situation that arises together with new causal structure.

Now we write the equations of spindynamics (2) and (3) taking into account equation (4):

$$\begin{aligned} \frac{1}{f}(D + \frac{3}{2})\kappa &= \text{div } \mathbf{K} - m\mu \\ \frac{1}{f}(D + \frac{3}{2})\lambda &= \text{div } \mathbf{L} - m\nu \\ \frac{1}{f}(D + \frac{3}{2})\mu &= \text{div } \mathbf{M} + m\kappa \\ \frac{1}{f}(D + \frac{3}{2})\nu &= \text{div } \mathbf{N} + m\lambda \end{aligned} \quad (6)$$

$$\begin{aligned} \frac{1}{f}(D + \frac{3}{2})\mathbf{K} &= -\text{rot } \mathbf{L} + \text{grad } \kappa + m\mathbf{M} \\ \frac{1}{f}(D + \frac{3}{2})\mathbf{L} &= \text{rot } \mathbf{K} + \text{grad } \lambda + m\mathbf{N} \\ \frac{1}{f}(D + \frac{3}{2})\mathbf{M} &= \text{rot } \mathbf{N} + \text{grad } \mu - m\mathbf{K} \\ \frac{1}{f}(D + \frac{3}{2})\mathbf{N} &= -\text{rot } \mathbf{M} + \text{grad } \nu - m\mathbf{L}, \end{aligned} \quad (7)$$

where the operator D is defined as follows

$$D = q_1 \frac{\partial}{\partial q_1} + q_2 \frac{\partial}{\partial q_2} + q_3 \frac{\partial}{\partial q_3} + q_4 \frac{\partial}{\partial q_4}$$

From the first principles it is clear that equations of spindynamics (6) and (7) represent quantum mechanics of strongly interacting particles.

To complete the picture, we write the Maxwell equations in the form that corresponds to the second causal structure

$$\frac{1}{f}(D + 2)\mathbf{H} = -\text{rot } \mathbf{E}, \quad \frac{1}{f}(D + 2)\mathbf{E} = \text{rot } \mathbf{H}, \quad (8)$$

$$\mathbf{q} \cdot \mathbf{E} = \mathbf{q} \cdot \mathbf{H} = 0, \quad \text{div } \mathbf{E} = \text{div } \mathbf{H} = 0, \quad (9)$$

where

$$E_i = \frac{q_k}{f} F_{ik}, \quad H_i = \frac{q_k}{f} {}^* F_{ik}, \quad F_{ij} = \partial_i A_j - \partial_j A_i, \quad {}^* F_{ij} = \frac{1}{2} e_{ijkl} F^{kl}.$$

We recognize that if vector field $\mathbf{M}$ is a solution to the equation

$$\text{rot } \mathbf{M} = p \mathbf{M},$$

then

$$\mathbf{t} \cdot \mathbf{M} = 0, \quad \text{div } \mathbf{M} = 0.$$

It is clear that the self-adjoint operator rot have fundamental meaning but here we have only a possibility to find the eigenvalues of this operator. We set $P = f \text{rot}$ and for this operator polarization the following relations hold valid

$$P = -M_1^2 - M_2^2 - M_3^2 + N_1^2 + N_2^2 + N_3^2 = -M^2 + N^2, \quad P^2 = -2(M^2 + N^2).$$

This means that eigenvalues of the polarization operator equal $\pm p$, where $p = 2, 3, 4, \dots$.

4. Conclusion. In conclusion, we discuss the new situation with solution of spindynamics equations (6) and (7) and the Maxwell equations (8) and (9). From polynomials (5) we can construct the orthogonal system of vector fields on S^3 (potential and turbulent). The coefficients of expansion of the vector and scalar fields over these systems will be the functions of f . Using the properties of the harmonic polynomials in question we can derive from the equations of spindynamics and the Maxwell equations the systems of ordinary differential equations for these coefficients. It is natural to suppose that this system of equations will have physical solutions only for quite definite meanings of charge and mass. Thus, at first we have a possibility to find the solution of nonlinear system of equations that can represent the distinct hadrons and nucleus. It should be noted that a state of strongly interacting particle is characterized by the following orbital quantum numbers: $0, 1/2, 1, 3/2, \dots$.

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REMARK ON SPIN PRECESSION FORMULAE

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Abstract

It is common to represent spin as a 4-vector in relativistic formulations, but from a fundamental point this is incorrect : spin should be represented by a rank 2 antisymmetric tensor. Such a relativistic tensor is displayed, which for Dirac particles turns out to be the Foldy-Wouthuysen mean spin operator. The precession formula for such a rank 2 tensor is shown to be the same as that for a 4-vector.

As is well known from undergraduate physics, a vector $\mathbf{a}$, fixed in a rotating body, when viewed from an inertial frame has a time dependence

$$\frac{D\mathbf{a}}{Dt} = \frac{d\mathbf{a}}{dt} + \boldsymbol{\Omega} \wedge \mathbf{a}.$$

If spin is described by a 3-vector $\mathbf{S}$, the same formula holds for its rate of change in a rotating frame:

$$\frac{D\mathbf{S}}{Dt} = \frac{d\mathbf{S}}{dt} + \boldsymbol{\Omega} \wedge \mathbf{S}. \quad (1)$$

This is the *non-relativistic formula for spin precession*. To obtain a relativistic formula the usual procedure (see for example Weinberg [1]) is to replace the 3-vector $\mathbf{S}$ by a 4-vector $S_\mu = (S_0, \mathbf{S})$. With this formulation, the precession rate of a spinning object in a satellite in orbit round the Earth is given by

$$\boldsymbol{\Omega} = \boldsymbol{\Omega}_{deSitter} + \boldsymbol{\Omega}_{Lense-Thirring} = \frac{3GM}{2c^2 r^3} \mathbf{r} \times \mathbf{v} + \frac{GI}{c^2 r^3} \left\{ \frac{3(\boldsymbol{\omega} \cdot \mathbf{r})\mathbf{r}}{r^2} - \boldsymbol{\omega} \right\}.$$

and this formula for Lense-Thirring and de Sitter precession agrees well with observations.

It is useful to rewrite the formulae above in coordinate notation. If a body rotates about its z axis with angular velocity ω the coordinates proper to the body are

$$x' = \cos \omega t - y \sin \omega t, y' = x \sin \omega t + y \cos \omega t, z' = z, t' = t,$$

so the invariant interval in Minkowski spacetime is (with $\rho^2 = x^2 + y^2$)

$$ds^2 = \left(1 - \frac{\omega^2 \rho^2}{c^2}\right) c^2 dt^2 - dx^2 - dy^2 - dz^2 + 2\omega(ydxdt - xdydt).$$

Equivalently, with $ds^2 = g_{\mu\nu} dx^\mu dx^\nu$, the metric tensor is

$$g_{\mu\nu} = \begin{pmatrix} 1 - \frac{\omega^2 \rho^2}{c^2} & \omega y & -\omega x & 0 \\ \omega y & -1 & 0 & 0 \\ -\omega x & 0 & -1 & 0 \\ 0 & 0 & 0 & -1 \end{pmatrix}.$$

Equation (1) may then be written, using the usual formula for covariant differentiation

$$S_{\mu;\nu} = S_{\mu,\nu} - \Gamma_{\mu\nu}^{\lambda} S_{\lambda} \quad (2)$$

or

$$S_{i;0} = S_{i,0} - \Gamma_{i0}^{\lambda} S_{\lambda} = S_{i,0} - \Gamma_{i0}^k S_k, \quad (3)$$

where the connection coefficients $\Gamma_{\mu\nu}^{\lambda}$ are calculated from the above metric tensor. The relevant non-zero coefficients are $\Gamma_{10}^2 = \omega$, $\Gamma_{20}^1 = -\omega$, giving

$$\begin{aligned} S_{1;0} &= S_{1,0} - \Gamma_{10}^2 S_2 = S_{1,0} - \omega S_2, \\ S_{2;0} &= S_{2,0} + \omega S_1, \quad S_{3;0} = S_{3,0}, \end{aligned} \quad (4)$$

or (with $\omega = \omega \mathbf{n}$)

$$\frac{D\mathbf{S}}{Dt} = \frac{d\mathbf{S}}{dt} + \boldsymbol{\Omega} \wedge \mathbf{S}.$$

as in (1) above.

From a fundamental perspective, however, spin is *not* described by a 4-vector, but by a rank 2 antisymmetric tensor. Indeed, the total angular momentum $J_{\mu\nu}$ contains orbital and spin parts

$$J_{\mu\nu} = L_{\mu\nu} + S_{\mu\nu}$$

and as is well known $J_{\mu\nu}$ are the generators of the Lorentz group. The most profound analysis of spin, however, was made by Wigner [2] in 1939, who pointed out that to describe spin it is insufficient to consider the (homogeneous) Lorentz group. Instead the group of translations in space-time must be added, yielding the inhomogeneous Lorentz group, or Poincaré group. This has generators P_{μ} and $J_{\mu\nu}$, and spin is generated by *little group* of the Poincaré group. The structure of this group depends on the sign of $P^{\mu}P_{\mu}$. With metric $(+, -, -, -)$ massive particles have $P^{\mu}P_{\mu} = m^2$. In the nonrelativistic limit spin is described by the rotation group $SO(3)$ or $SU(2)$, and indeed this turns out to be true for all states with $P^{\mu}P_{\mu} > 0$:

$$P^{\mu}P_{\mu} > 0 \text{ (timelike states) : little group is } SU(2).$$

Wigner showed, however, that for other values $P^{\mu}P_{\mu}$ the little group has a different structure:

$$P^{\mu}P_{\mu} < 0 \text{ (spacelike states) : little group is } SO(2,1).$$

$$P^{\mu}P_{\mu} = 0 \text{ (null states) : little group is } E(2).$$

Thus, for example, the spin of spacelike (virtual) particles in Feynman diagrams is not properly described by the rotation group, but by the noncompact group $SO(2,1)$.

Wigner's analysis identified the group structure of the little groups but did not reveal the actual forms of the generators of these groups, leaving the question, what are these generators? The Casimir operators of the Poincaré group are $P^{\mu}P_{\mu}$ and $W^{\mu}W_{\mu}$, describing, in effect, the mass and spin of quantum states. Here W_{μ} is the Pauli-Lubanski vector

$$W_{\mu} = \frac{1}{2} \epsilon_{\mu\nu\kappa\lambda} J^{\nu\kappa} P^{\lambda}.$$

It is easy to see that for a (massive) particle at rest $(P_0, P_1, P_2, P_3) = (M, 0, 0, 0)$, $W_0 = 0$, and W_i generate the group $SU(2)$, as expected. Similarly for a spacelike state we may have $(P_0, P_1, P_2, P_3) = (0, P, 0, 0)$, and then W_0, W_1, W_2 are seen to generate $SO(2, 1)$, again as expected. These operators are, however, *non-covariant*. To obtain covariant operators first define [3]

$$W_{\mu\nu} = \frac{1}{M^2}[W_\mu, W_\nu],$$

which are seen to obey the commutation relation

$$[W_{\mu\nu}, W_{\kappa\lambda}] = \frac{1}{M^2}(\epsilon_{\mu\nu\kappa\rho}W_\lambda - \epsilon_{\mu\nu\lambda\rho}W_\kappa)P^\rho.$$

Now define also the *dual*

$$(W^D)^{\kappa\lambda} = \frac{1}{2}\epsilon^{\kappa\lambda\mu\nu}W_{\mu\nu}.$$

Then the linear combinations

$$X^{\mu\nu} = -i\{W^{\mu\nu} + i(W^D)^{\mu\nu}\}, \quad Y^{\mu\nu} = -i\{W^{\mu\nu} - i(W^D)^{\mu\nu}\} = i(X^D)^{\mu\nu}$$

are found to obey

$$[X_{\mu\nu}, X_{\kappa\lambda}] = -i(g_{\mu\kappa}X_{\nu\lambda} - g_{\mu\lambda}X_{\nu\kappa} + g_{\nu\lambda}X_{\mu\kappa} - g_{\nu\kappa}X_{\mu\lambda}),$$

and similarly for Y . These are the same as the commutation relations for $J_{\mu\nu}$. $X_{\mu\nu}$ and $Y_{\mu\nu}$ are therefore covariant spin operators. For Dirac particles they act as the spin operators for *left- and right-handed states*. The covariant spin operator may therefore be written as

$$Z_{\mu\nu} = P_L X_{\mu\nu} + P_R Y_{\mu\nu}. \quad (5)$$

Then, putting $X_i = \frac{1}{2}\epsilon_{ijk}X^{jk}$ and $Y_i = \frac{1}{2}\epsilon_{ijk}Y^{jk}$, the components of this rank 2 tensor which transform as a 3-vector may be written in the usual way as

$$Z_i = \frac{1}{2}(1 - \gamma_5)X_i + \frac{1}{2}(1 + \gamma_5)Y_i.$$

It then turns out, perhaps as an unexpected bonus, that Z_i is the Foldy-Wouthuysen *mean spin operator* [3, 4].

Having now identified a covariant relativistic spin operator with the required property that it is a rank 2 tensor, we may immediately write its covariant derivative as

$$Z_{\mu\nu;\lambda} = Z_{\mu\nu,\lambda} - \Gamma_{\mu\lambda}^\rho Z_{\rho\nu} - \Gamma_{\nu\lambda}^\rho Z_{\mu\rho}.$$

With the connection coefficients derived from the metric tensor above, these give

$$Z_{12;0} = Z_{12,0}; \quad Z_{23;0} = Z_{23,0} + \omega Z_{13}; \quad Z_{31;0} = Z_{31,0} - \omega Z_{32}$$

or

$$\frac{DZ_3}{dt} = \frac{dZ_3}{dt}; \quad \frac{DZ_1}{dt} = \frac{dZ_1}{dt} - \omega Z_2; \quad \frac{DZ_2}{dt} = \frac{dZ_2}{dt} + \omega Z_1,$$

as in equation (4) above. We see that the *same precession formula* is obtained when spin is represented as an antisymmetric rank 2 tensor, as when it is represented by a 4-vector.

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DYNAMICS OF SPIN IN NONSTATIC SPACETIMES

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Abstract

The quantum and classical dynamics of a particle with spin in the gravitational field of a rotating source is discussed. A relativistic equation describing the motion of classical spin in curved spacetimes is derived. The precession of the classical spin is in a perfect agreement with the motion of the quantum spin obtained from the Foldy-Wouthuysen approach for the Dirac particle in a curved spacetime. It is shown that the precession effect depends crucially on the choice of a tetrad.

In this paper, we consider two important problems. One aim is to generalize the methods of the Foldy-Wouthuysen transformations, that we previously used for the analysis of the spin in static gravitational fields, to the case of the stationary gravitational configurations. Another aim is to systematically investigate the dependence of the spin dynamics on the choice of a tetrad.

We denote world coordinate indices by Latin letters and reserve first letters of the Greek alphabet for tetrad indices. Spatial indices 1, 2, 3 are labeled by Latin letters from the beginning of the alphabet. The separate tetrad indices are distinguished by hats.

The line element for the Lense-Thirring (LT) metric [1] is given by

$$ds^2 = V^2 c^2 dt^2 - W^2 \delta_{ab} (dx^a - K^a c dt) (dx^b - K^b c dt), \quad (1)$$

with

$$V = \left(1 - \frac{GM}{2c^2 r}\right) \left(1 + \frac{GM}{2c^2 r}\right)^{-1}, \quad W = \left(1 + \frac{GM}{2c^2 r}\right)^2, \quad K^a = \frac{1}{c} \epsilon^{abc} \omega_b x_c. \quad (2)$$

The non-diagonal components of the metric (that reflect the rotation of the source) are described by the so-called Kerr field $\mathbf{K}$ that is given by Eq. (2) with

$$\boldsymbol{\omega} = \frac{2G}{c^2 r^3} \mathbf{J} = \left(0, \quad 0, \quad \frac{2GM a}{c r^3}\right), \quad (3)$$

where $\mathbf{J} = Mca\mathbf{e}_z$ is the total angular momentum of the source.

The *exact* metric of the flat spacetime seen by an accelerating and rotating observer also has form (1) where

$$V = 1 + \frac{\mathbf{a} \cdot \mathbf{r}}{c^2}, \quad W = 1, \quad K^a = -\frac{1}{c} (\boldsymbol{\omega} \times \mathbf{r})^a. \quad (4)$$

$\mathbf{a}$ describes acceleration of the observer and $\boldsymbol{\omega}$ is an angular velocity of a noninertial reference system.

The covariant Dirac equation for spin-1/2 particles and the orthonormal tetrad

$$e_i^{\hat{0}} = V \delta_i^0, \quad e_i^{\hat{a}} = W (\delta_i^a - K^a \delta_i^0), \quad a, b = 1, 2, 3, \quad (5)$$

can be transformed to the familiar Schrödinger form with the Hermitian Hamiltonian

$$\begin{aligned} \mathcal{H}' = & \beta mc^2 V + \frac{c}{2} [(\boldsymbol{\alpha} \cdot \mathbf{p}) \mathcal{F} + \mathcal{F}(\boldsymbol{\alpha} \cdot \mathbf{p})] \\ & + \frac{2G}{c^2 r^3} \mathbf{J} \cdot (\mathbf{r} \times \mathbf{p}) + \frac{\hbar G}{2c^2 r^3} \left[\frac{3(\mathbf{r} \cdot \mathbf{J})(\mathbf{r} \cdot \boldsymbol{\Sigma})}{r^2} - \mathbf{J} \cdot \boldsymbol{\Sigma} \right]. \end{aligned} \quad (6)$$

Dirac Hamiltonian (6) contains the first part describing the static gravitational field and the second one characterizing the contribution of rotation of the central body.

This FW transformation [2, 3] leads to the final Hamiltonian which is given by

$$\begin{aligned} \mathcal{H}_{FW} = & \mathcal{H}_{FW}^{(1)} + \mathcal{H}_{FW}^{(2)}, \quad \mathcal{H}_{FW}^{(2)} = \frac{2G}{c^2 r^3} \mathbf{J} \cdot \mathbf{l} + \frac{\hbar G}{2c^2 r^3} \left[\frac{3(\mathbf{r} \cdot \mathbf{J})(\mathbf{r} \cdot \boldsymbol{\Sigma})}{r^2} - \mathbf{J} \cdot \boldsymbol{\Sigma} \right] \\ & - \frac{3\hbar G}{8} \left\{ \frac{1}{\epsilon(\epsilon + mc^2)}, \left[\frac{2\{(\mathbf{J} \cdot \mathbf{l}), (\boldsymbol{\Sigma} \cdot \mathbf{l})\}}{r^5} + \frac{1}{2} \left\{ (\boldsymbol{\Sigma} \cdot (\mathbf{p} \times \mathbf{l}) - \boldsymbol{\Sigma} \cdot (\mathbf{l} \times \mathbf{p})), \frac{(\mathbf{r} \cdot \mathbf{J})}{r^5} \right\} \right. \right. \\ & \left. \left. + \left\{ \boldsymbol{\Sigma} \cdot (\mathbf{p} \times (\mathbf{p} \times \mathbf{J})), \frac{1}{r^3} \right\} \right] \right\} - \frac{3\hbar^2 c^2 G}{8} \left\{ (5p_r^2 - \mathbf{p}^2) \frac{2\epsilon^2 + \epsilon mc^2 + m^2 c^4}{\epsilon^4 (\epsilon + mc^2)^2}, \frac{(\mathbf{J} \cdot \mathbf{l})}{r^5} \right\}, \end{aligned} \quad (7)$$

where $\mathbf{l} = \mathbf{r} \times \mathbf{p}$ is an angular momentum operator, and the operator $p_r^2 = -\frac{\hbar^2}{r^2} \frac{\partial}{\partial r} \left(r^2 \frac{\partial}{\partial r} \right)$ is proportional to the radial part of the Laplace operator. The rotation-independent contribution $\mathcal{H}_{FW}^{(1)}$ has been calculated earlier [4].

The newly obtained operator of angular velocity of rotation of the spin in the static gravitational field is equal to

$$\begin{aligned} \boldsymbol{\Omega}^{(2)} = & \frac{G}{c^2 r^3} \left[\frac{3(\mathbf{r} \cdot \mathbf{J})\mathbf{r}}{r^2} - \mathbf{J} \right] - \frac{3G}{4} \left\{ \frac{1}{\epsilon(\epsilon + mc^2)}, \left[\frac{2\{\mathbf{l}, (\mathbf{J} \cdot \mathbf{l})\}}{r^5} \right. \right. \\ & \left. \left. + \frac{1}{2} \left\{ (\mathbf{p} \times \mathbf{l} - \mathbf{l} \times \mathbf{p}), \frac{(\mathbf{r} \cdot \mathbf{J})}{r^5} \right\} + \left\{ (\mathbf{p} \times (\mathbf{p} \times \mathbf{J})), \frac{1}{r^3} \right\} \right] \right\}. \end{aligned} \quad (8)$$

The semiclassical formula corresponding to Eq. (8) and describing the motion of average spin has the form

$$\boldsymbol{\Omega}^{(2)} = \frac{G}{c^2 r^3} \left[\frac{3(\mathbf{r} \cdot \mathbf{J})\mathbf{r}}{r^2} - \mathbf{J} \right] - \frac{3G}{r^5 \epsilon (\epsilon + mc^2)} [\mathbf{l}(\mathbf{l} \cdot \mathbf{J}) + (\mathbf{r} \cdot \mathbf{p})(\mathbf{p} \times (\mathbf{r} \times \mathbf{J}))]. \quad (9)$$

The presented quantum mechanical and semiclassical equations are principal new results.

Description of a spin requires the introduction of a tetrad. A choice of a tetrad means a selection of a local reference system of an observer.

There are infinitely many tetrads since a reference frame of an observer can obviously be constructed in infinitely many ways. In particular, from a given tetrad field e_i^α we can obtain a continuous family of tetrads by performing the Lorentz transformation $e_i'^\alpha = \Lambda^\alpha_\beta e_i^\beta$, where the elements of the Lorentz matrix $\Lambda^\alpha_\beta(x)$ are arbitrary functions of the spacetime coordinates. In practice, there are three most widely used gauges.

Schwinger gauge. Probably for the first time introduced independently by Schwinger [5] and Dirac [6] (and widely used in many works, including [7] and our current study), this choice demands that the tetrad matrix e_i^α and its inverse e_α^i satisfy $e_0^0 = 0$, $e_b^0 = 0$.

Landau-Lifshitz gauge (see, e.g., Ref. [8]) fixes the tetrad so that $e_0^a = 0$, $e_0^a = 0$.

Symmetric gauge. Using the Minkowski flat metric $g_{\alpha\beta} = \text{diag}(c^2, -1, -1, -1)$, we can move the anholonomic index down and define $e_{\alpha i} := g_{\alpha\beta} e_i^\beta$. Now assume that the resulting matrix is invariant under the transposition operation that can be symbolically written as $e_{\alpha i} = e_{i\alpha}$. Such a tetrad was used by Pomeransky and Khriplovich [9] and Dvornikov [10].

We choose the *Schwinger gauge* by specifying the coframe as (5). The other tetrads are obtained from our e_i^α with the help of the Lorentz transformation $e_i'^\alpha = \Lambda^\alpha_\beta e_i^\beta$, where

$$\Lambda^\alpha_\beta = \begin{pmatrix} \lambda & \lambda q_b/c \\ \lambda c q^a & \delta_b^a + (\lambda - 1) q^a q_b / q^2 \end{pmatrix}, \quad q^a = \xi \frac{W K^a}{V}, \quad \lambda = \frac{1}{\sqrt{1 - q^2}}. \quad (10)$$

The constant ξ conveniently parametrizes different choices of tetrads. Namely, for $\xi = 1/2$ the Lorentz matrix (10) transforms our tetrad to that of Pomeransky and Khriplovich, and for $\xi = 1$ we obtain the tetrad of Landau and Lifshitz.

The spin precession in the gravitational field of rotating object is given by

$$\boldsymbol{\Omega} = \frac{G}{c^2 r^3} \left[\frac{3\mathbf{r}(\mathbf{r} \cdot \mathbf{J})}{r^2} - \mathbf{J} \right] + \frac{\boldsymbol{\rho} \times \mathbf{v}}{c^2}, \quad (11)$$

where we denote

$$\boldsymbol{\rho} = \frac{2\gamma + 1}{\gamma + 1} \frac{GM}{r^3} \mathbf{r} + \frac{\gamma}{\gamma + 1} \frac{3G}{c^2 r^3} \left[\frac{\mathbf{r}}{r^2} (\mathbf{r} \cdot (\mathbf{J} \times \mathbf{v})) - \frac{2\xi}{3} \mathbf{J} \times \mathbf{v} + (2\xi - 1) \frac{(\mathbf{r} \cdot \mathbf{v})}{r^2} \mathbf{J} \times \mathbf{r} \right]. \quad (12)$$

Putting $\xi = 0$, thus specifying to the *Schwinger* tetrad, we find that the classical formula (11) perfectly reproduces the quantum result (9). If we choose another tetrad by putting $\xi = 1/2$ in (12), the equation (11) yields the result by Pomeransky and Khriplovich [9] and Dvornikov [10] which differs from Eq. (9).

For particles in a rotating frame, the angular velocity of spin precession is given by

$$\boldsymbol{\Omega} = -\boldsymbol{\omega} + \frac{\xi \gamma}{\gamma + 1} \frac{\mathbf{v} \times (\mathbf{v} \times \boldsymbol{\omega})}{c^2}. \quad (13)$$

The correct result for the Schwinger gauge ($\xi = 0$) was first obtained in [11]. An explanation of the dependence of the angular velocity of the spin precession on the gauge of a tetrad was presented recently [12] on the basis of the Thomas precession.

The Lense-Thirring effect or frame dragging is one of the most impressive predictions of the general relativity. This effect is currently analyzed in the Gravity Probe B experiment [13]. However, relativistic corrections to the LT effect are not observable in this experiment as well as in other experiments inside the solar system [14]. Nevertheless, it is necessary to take the relativistic corrections to the LT precession into account for the investigation of physical phenomena in the binary stars such as pulsar systems. In this case, both components of a system undergo a mutual LT precession about the total angular momentum $\mathbf{J}$. Since the spin precession effects are well observable [15], the use of

the results obtained in the present work may be helpful for the high-precision calculations of spin dynamics in the binaries.

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DSPIN-09 WORKSHOP SUMMARY

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Abstract

I will try to summarize several stimulating open questions in high energy spin physics, which were discussed during the five days of this workshop, showing also the striking progress recently achieved in this field.

1 Introduction

Once again this workshop has been very productive with a high density scientific program, since about 95 talks were presented. It has facilitated detailed discussions between theorists and experimentalists and, moreover since spin occurs in all particle processes, it is obvious that by ignoring this fundamental tool, we will miss an important part of the story. I will mention some substantial progress which have been achieved since DSPIN-07 and, whenever it is possible, to identify what we have learnt and what are the prospects. Although I will not cover technical subjects, I just want to mention some future projects, in particular E. Steffens, who reported about the status of the challenging PAX experiment in FAIR at GSI, supplemented by a bright theorist perspective by N. Nikolaev on polarized antiproton experiments. S. Nurushev gave a talk on the polarization program SPASCHARM at IHEP, in connection with the latest results from the PROZA experiment and various aspects of the preparation of the spin program at the Nuclotron in Dubna, were presented by Y. Gurchin, S. Piyadin and P. Kurilkin.

The opening talk of DSPIN-09 was given by A. Krisch who recalled us some unexpected large spin effects in pp elastic scattering obtained about twenty years ago, whose clear theoretical interpretation is not yet available (see Figs. 1a,b). These results provided the motivation to undertake a very successful Siberian Snake program, allowing to obtain high energy polarized proton beams, an essential element of the RHIC Spin program at BNL. Following the introduction, this summary talk contains several sections, first on the experimental side, with new results from COMPASS, HERMES, BELLE, JLab and RHIC and at the end a section devoted to several theoretical subjects.

2 New results from COMPASS

The COMPASS fixed-target experiment at the CERN SPS has produced over the last two years several interesting new results in different areas. R. Gazda presented an analysis of 2002-2007 data on the first moment of the structure function g_1 and semi-inclusive asymmetries, on proton and deuteron targets, which has led to a better determination of

the polarized valence quark and sea quark distributions with flavor separation. Concerning the important issue of the gluon polarization it has been extracted from high p_T hadron pairs, either in the quasi-real photoproduction ($Q^2 < 1\text{GeV}^2$) or in the DIS ($Q^2 > 1\text{GeV}^2$) regimes. L. Silva concluded that two independent analysis lead to compatible values, with a result consistent with zero, as shown in Fig. 2. One can also see in Fig. 2 that this small gluon polarization was confirmed by the method based on the measurement of open-charm asymmetries, as reported by K. Kurek, supplemented by a NLO QCD prediction. Longitudinal polarization of the Λ and $\bar{\Lambda}$ hyperons DIS events collected by COMPASS

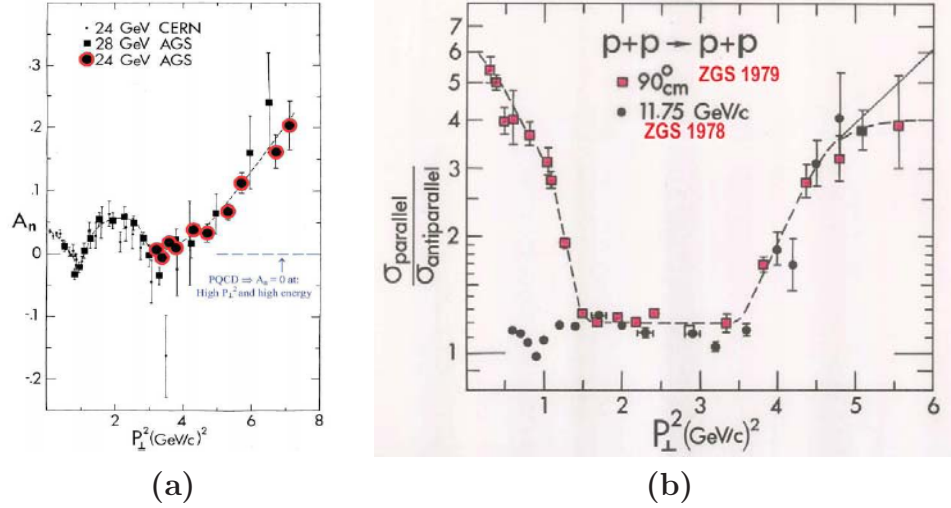


Figure 1: (a) The single-spin asymmetry A_n in pp elastic scattering as a function of p_T^2 (b) Ratio of pp differential cross sections for spins parallel or antiparallel along the normal to the scattering plane (Both taken from Krisch's talk).

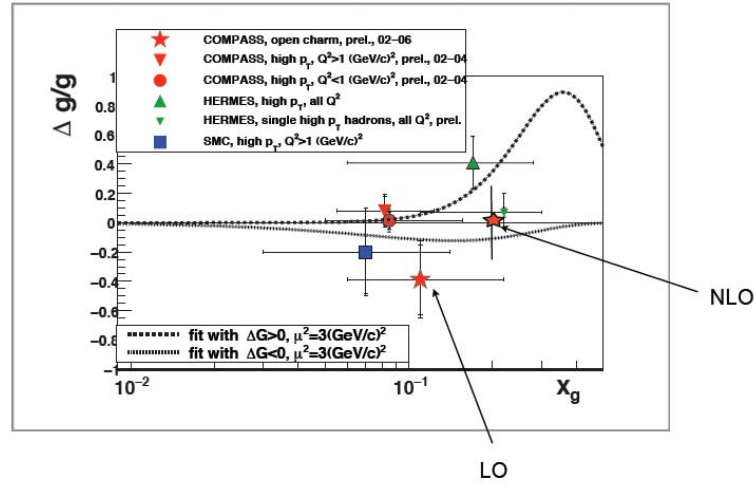


Figure 2: Comparison of the $\Delta G/G$ measurements from various experiments (Taken from Kurek's talk).

in 2003-2004, the world largest sample, was presented by V. Rapatskiy, with the goal to study Λ spin structure models. The results will be improved by adding a much larger events number from 2006-2007 data. The measurements of azimuthal asymmetries in semi-inclusive hadron production with a longitudinally polarized deuterium target, led

to some preliminary results reported by I. Savin, in various kinematical variables. The COMPASS experiment has put a serious effort in the determination of the transverse momentum dependent parton distributions in particular the Collins function, which gives access to the transversity distribution, and the Sivers function. Several new results were given in G. Sbrizzai's talk as shown in Figs. 3a,b and although the Collins asymmetries on proton are compatible with the predictions of the present picture, the agreement is marginal for the Sivers asymmetries. New data on transversely polarized proton will be taken in 2010 and the large increase in precision should, hopefully, solve this important problem.

The Generalized Parton Distributions (GPD) program proposed by COMPASS was covered by A. Sandacz. This project will explore intermediate x (0.01 - 0.1) and large Q^2 (8 - 12) GeV² and will be unique in this kinematical range before availability of new colliders. There is another future project by COMPASS, a complete program of Drell-Yan experiments for probing the hadron structure, with first data taking after 2012, which was presented by O. Denisov. Drell-Yan process is very rich because unpolarized angular distributions give access to the Boer-Mulders function, single spin-asymmetry allows to test the sign change from SIDIS of the Sivers function, a fundamental test of gauge theory and the double transverse spin asymmetry gives access to the transversity distributions.

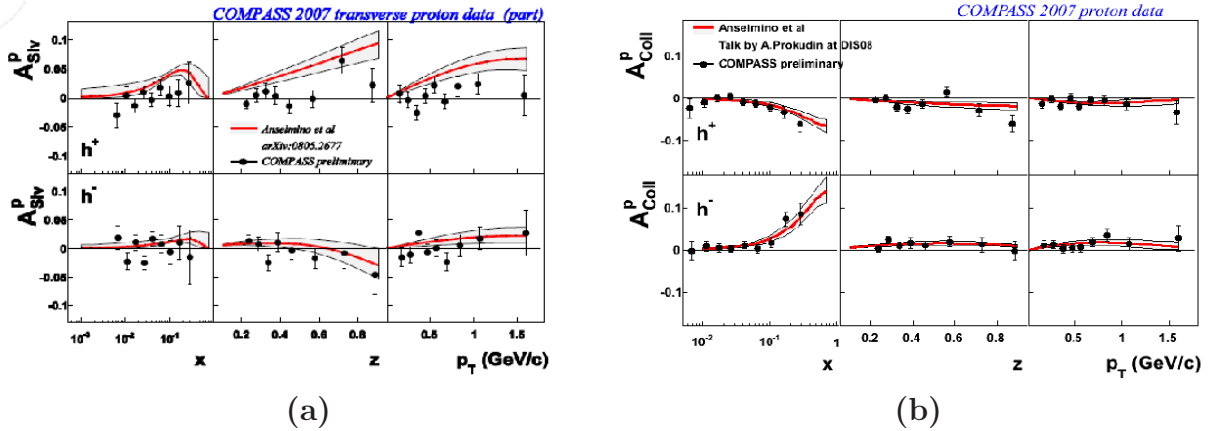


Figure 3: (a) COMPASS Sivers asymmetries on proton versus x , z and p_T . (b) COMPASS Collins asymmetries on proton versus x , z and p_T (Both taken from Sbrizzai's talk).

3 New results from HERMES

An overview of the HERMES experiment was given by V. Korotkov, including in the pure DIS sector, latest results on the $F_2^{p,d}$ and g_2^p structure functions and in the SIDIS sector, Collins asymmetries, Sivers asymmetries, azimuthal asymmetries in unpolarized scattering and strange quark distributions.

New results on Deeply Virtual Compton Scattering (DVCS) at HERMES were reported by A. Borissov, whose motivation is to get access to GPD. DVCS azimuthal asymmetries on proton provide a constraint on total angular momentum of valence quarks and allow the comparison with GPD model calculations. Along the same lines, S. Manayenkov gave a talk on a detailed study of exclusive electroproduction of ρ^0 , ϕ and ω at HERMES,

with a special emphasis on the extraction of the spin density matrix elements and tests of s-channel helicity conservation (SCHC). Violation of SCHC is observed in ρ^0 production, both on proton and deuteron, but no such signal is found for ϕ meson. Finally Λ physics at HERMES was covered by Y. Naryshkin, where they observed for Λ and $\bar{\Lambda}$, the longitudinal spin transfer from the beam, shown in Fig. 4, a very small longitudinal spin transfer from the target and a reliable transverse polarization.

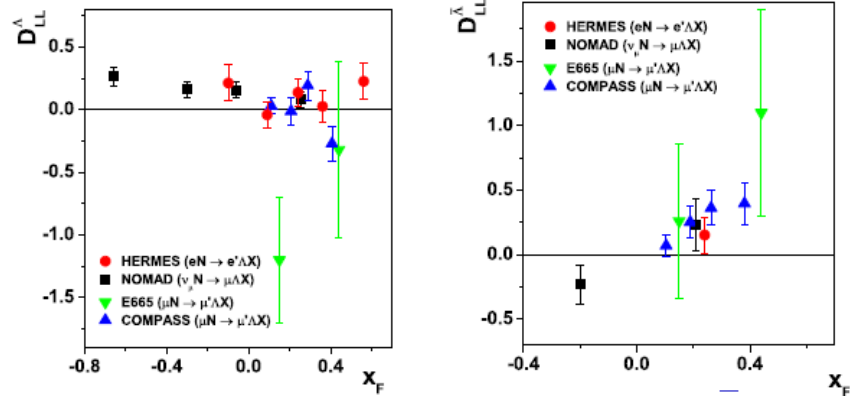


Figure 4: Spin transfer parameter D_{LL} versus x_F for Λ and $\bar{\Lambda}$ (Taken from Naryshkin's talk).

4 New results from BELLE and JLab

The measurement of quark transversity distributions can be done by different methods, using either pp collisions or SIDIS and e^+e^- collisions. A. Vossen recalled the first attempt to extract them, in a model dependent way, from BELLE data combined with Collins asymmetry in SIDIS. Then he described the forthcoming observation of an interference fragmentation function asymmetry in e^+e^- collisions at BELLE, another very promising method.

An introductory review talk of spin physics with CLAS in Hall B at JLab, was presented by Y. Prok, in particular some recent experimental results. The list of topics was covering, the status of $g_1^p(x, Q^2)$, including the resonance region, the large- x behavior of the A_1 asymmetry, dominated by valence quarks, the generalized Gerasimov-Drell-Hearn (GDH) sum rule, semi-inclusive processes, etc...The whole activity with CLAS was supplemented by several presentations. First, M. Mirazita with studies in SIDIS with longitudinal polarization and future transversely polarized target, in view of the determination of the transverse momentum dependent distributions. Second, C. Munos Camacho with a presentation of the GPD experimental program at JLab, measurements of DVCS/GPD both in Hall B and in Hall A, together with some exciting perspectives for the future with the 12GeV upgrade. Third, V. Drozdov told us that, concerning the CLAS EG4 experiment on the GDH sum rule in the low Q^2 region, the analysis is underway. Finally, E. Pasyuk who reported on very interesting results from CLAS on several polarization measurements, in exclusive photoproduction, after the new addition of a frozen spin target, with both longitudinal and transverse polarization.

A new measurement of the polarization transfer in ep elastic scattering, by the recoil polarization technique, was done in Hall C at JLab, allowing to extract the ratio of electric and magnetic form factors G_{Ep}/G_{Mp} at large Q^2 . The preliminary results were presented by V. Punjabi and are shown, with earlier results, in Fig. 5. They confirm the previous observation that this ratio decreases with increasing Q^2 and contradicts the expectation of a flat behavior, according to conventional wisdom. This measurement will be done up to $Q^2 = 15\text{GeV}^2$ with the 12GeV upgrade. Finally, the search for a 2γ contribution in

ep elastic scattering, a very closely related subject, was presented by Ch. Perdrisat, together with some preliminary results of the $G_{Ep}(2\gamma)$ experiment at JLab.

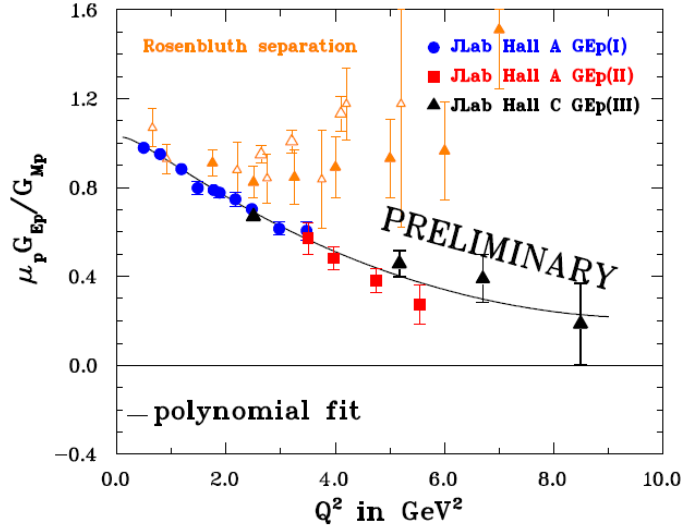


Figure 5: The proton elastic form factor ratio versus Q^2 (Taken from Punjabi's talk).

5 New results from RHIC

Brookhaven National Lab. operates since 2001 a polarized pp collider in RHIC, which will be running up to $\sqrt{s}=500\text{GeV}$, to perform a vast program of spin measurements. The energy is high enough to assume that NLO pQCD is applicable. The first talk due to D. Kawell was a presentation of the spin program at the PHENIX Collaboration, including new results on the gluon helicity distribution, from double helicity asymmetry measurements in π^0 , direct photon or heavy flavors production. A short run at the highest energy $\sqrt{s}=500\text{GeV}$ done in 2009 has given the first W signals and the parity-violating asymmetry A_L measurements, giving access to the quark and antiquark helicity distributions, are expected to be done in 2011. For transverse spin phenomena, new results were obtained and in Fig. 6 one displays a sizeable single-spin asymmetry in forward π^0 production. More data are certainly needed for a full understanding of the rise in x_F and the p_T dependence of A_N .

A review talk on polarimetry at RHIC, for both PHENIX and STAR, was presented by A. Bazilevsky, who told us that pp elastic scattering in the Coulomb Nuclear Interference (CNI) region is ideal for absolute polarimetry in a wide energy range. He also recalled that a large forward neutron single-spin asymmetry, for $x_F > 0$ was discovered in RHIC Run-2002, which is a very useful PHENIX local polarimetry. One can see in Fig. 7 the energy dependence of this asymmetry, very probably, a diffractive physics phenomena, whose theoretical interpretation is unfortunately not yet available.

Finally, we had two talks on STAR. First, L. Nogach presented new results on measurements of transverse spin effects in the forward region, for η^0 and π^0 production and discussed the future possibilities for similar measurements in jet production, Λ production and Drell-Yan production of dilepton pairs. Second, Q. Xu discussed the results on the

longitudinal spin transfer in Λ and $\bar{\Lambda}$ inclusive production at $\sqrt{s}=200\text{GeV}$, which is well described by pQCD.

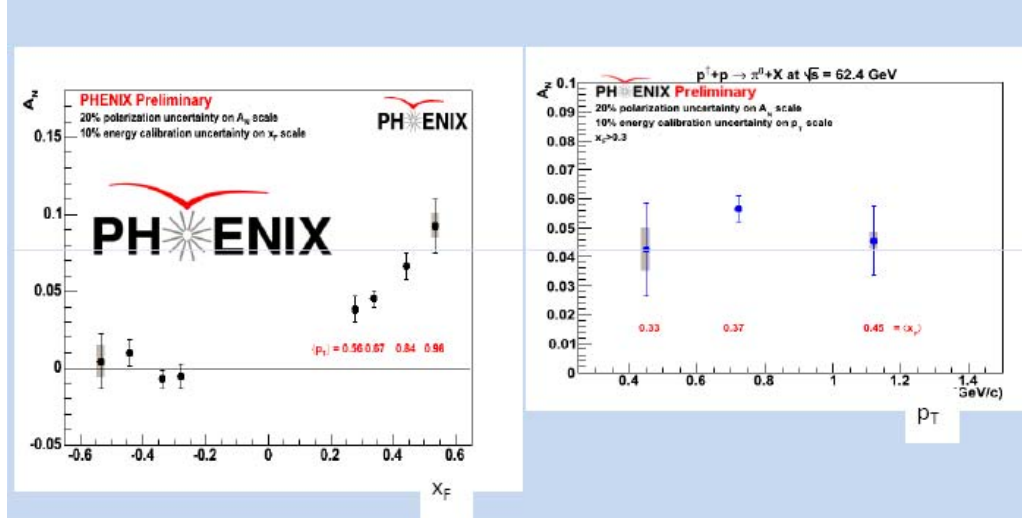


Figure 6: Single-spin asymmetry for π^0 inclusive production at $\sqrt{s} = 62.4 \text{ GeV}$ (Taken from Kawell's talk).

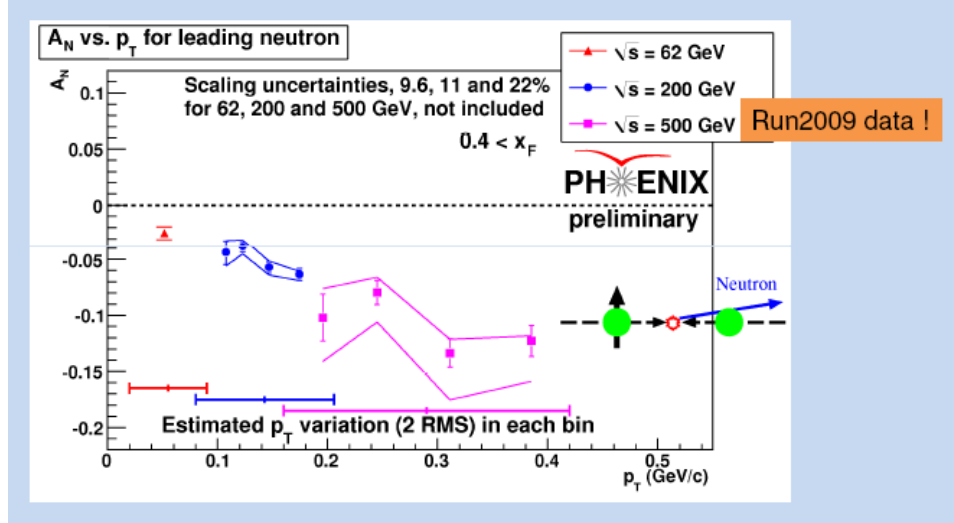


Figure 7: Single-spin asymmetry for forward neutron inclusive production at $\sqrt{s} = 62, 200, 500 \text{ GeV}$ (Taken from Bazilevsky's talk).

6 Theory

Concerning the subject of phenomenological studies of parton distributions functions (PDF), we heard a talk on some new developments in the quantum statistical approach of the unpolarized and polarized PDFs, stressing several very challenging points, in particular in the high x region. In his presentation A. Sidorov discussed different methods of QCD analysis of the polarized DIS data. He emphasized the importance of higher twist effects and showed that the correct determination of the PDFs depends crucially whether or not they are taken into account. In another contribution O. Shevchenko presented a

new NLO QCD parametrization of the polarized PDF, obtained by using all published results on DIS and SIDIS asymmetries. It generalizes an earlier parametrization released by COMPASS, based on only inclusive DIS data. New SIDIS data from COMPASS, which will be available in the near future, should allow to improve the quality of this parametrization. Another method based on evolution equations for truncated Mellin moments of parton densities has been presented by D. Kotlorz. It avoids uncertainties related to some poorly known x regions and may be a new tool for an accurate reconstruction of the PDFs. A relevant question concerning the kinematic regions in x and Q^2 , where DGLAP is legitimately applicable, in connection with the infrared dependence of the structure function g_1 , was discussed by B. Ermolaev. He recalled that the extrapolation of DGLAP into the small- x region is usually done by introducing in the parametrization of the PDFs, singularities which are theoretically groundless.

Ph. Ratcliffe first recalled us how painful it was for the QCD community to accept the existence of large single-spin asymmetries, before new QCD mechanisms were discovered. Then he discussed colour modification of factorization by relating the pQCD evolution of the Sivers function and the twist-three gluonic pole contribution.

A model independent determination of fragmentation functions (FF) was proposed by E. Christova, by considering cross-section differences. It does not provide the full information on the FFs, only part of it, but without any assumptions on the FFs and on the PDFs.

Some aspects of quark motion inside the nucleon have been discussed by P. Zavada. The Cahn effect, which is an important tool to measure the quark transverse motion, and the unintegrated unpolarized PDFs $f(x, \mathbf{k}_T)$, are studied in a covariant approach. In the framework of light-cone quark models, B. Pasquini presented a set of transverse momentum dependent (TMD) PDFs, allowing to derive some azimuthal spin asymmetries which were compared with available experimental data from CLAS, COMPASS and HERMES.

X. Artru described a very simple recursive fragmentation model with quark spin and its possible application to jet handedness and Collins effects for pions, either emitted directly or coming from vector meson decay. In his talk O. Teryaev discussed several issues related to the spin puzzle, the axial anomaly, the strangeness polarization and the role of heavy quarks polarization, suggesting a possible charm-strangeness universality and questioning when strange quarks can be heavy.

Cross sections and spin asymmetries in light vector meson leptonproduction were analyzed in the framework of generalized parton distributions by S.V. Goloskokov and the results are in good agreement with data from HERA, HERMES and COMPASS at various energies. Along the same lines, P. Kroll presented a talk on spin effects in hard exclusive meson electroproduction within the framework of the handbag approach, showing the failure of the leading-twist calculations.

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