HELICITY NON-CONSERVATION IN HIGH ENERGY EXCLUSIVE SCATTERING

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Abstract

We study a new mechanism for hadronic helicity flip in high energy hard exclusive reactions. Its fundamental feature is the breaking of rotational symmetry of the hard collision by a scattering plane in processes involving independent quark scattering. An important role is played by "chirally-odd" light-cone valence wave functions which carry non-zero orbital angular momentum and yet are leading in the high energy limit. There is no substantial suppression of the helicity violating process compared to helicity conserving ones as the momentum transfer Q^2 is increased over the experimentally accessible region $1\,GeV^2 < Q^2 < 100\,GeV^2$.

1 Introduction

The theory of hard elastic scattering in Quantum Chromodynamics (QCD) has evolved considerably over many years of work. A well-known procedure using the"quark-counting" diagrams has been given by LePage and Brodsky[1]. A consequence, and direct test, of the factorization defining this mechanism is the hadron helicity conservation rule:

$$\lambda_A + \lambda_B = \lambda_C + \lambda_D \ , \tag{1}$$

where the λ_j's are the helicities of the participating hadrons in the reaction $A+B \to C+D$. In general, the sum of the helicities of hadrons going into a reaction equals the sum going out, when one uses the quark-counting factorization. The fact that this rule is badly violated in almost every case tested leads one to suspect that another power behaved process causing helicity flip is present. In fact the "independent scattering" subprocess, introduced by Landshoff, is actually the leading process at very high energies.

We have shown[2] that the independent scattering mechanism predicts high-energy helicity *non-conservation*. Adopting a transverse position space formalism introduced by Botts and Sterman[3], we show that the details rest on non-perturbative wave functions that should be *measured* rather than calculated. These wave functions measure non-zero orbital angular momentum not taken into account by short distance expansions. We argue that the novel factorization properties of independent scattering processes cannot practically be reduced to the same ingredients used in the quark counting scattering. In any case, it is not necessary to flip a quark helicity: the new mechanism proceeds unimpeded in the limit of arbitrarily small quark mass and perfect chiral symmetry in the hard scattering.

2 Wave functions

In general, quark wave functions themselves are not particularly restricted in orbital angular momentum content, even in the high energy limit. For example, in the pseudoscalar meson case there are four wave functions allowed by parity symmetry,

$$\begin{aligned} \mathcal{P}_{\alpha\beta}(x, b_T; p) &= \int \frac{db^-}{2\pi}\, e^{ixp^+b^-} < 0|q_\alpha(b)\, \bar{q}_\beta(0)|\pi(p)>, \\ &= \{\mathcal{P}_{0\pi}\gamma_5 \not{p} + \mathcal{P}'_{0\pi}\gamma_5 + \mathcal{P}_{1\pi}\gamma_5[\not{p}, \not{b}_T] + \mathcal{P}'_{1\pi}\gamma_5\not{b}_T\}_{\alpha\beta} \end{aligned} \tag{2}$$

where $\mathcal{P}_\pi$'s are functions of the light-cone momentum fraction x of the valence quark and of the quark spatial separation b_T. The $\mathcal{P}_{1\pi}$-term scales with the same power of the "big" momentum p^+ as the $\mathcal{P}_{0\pi}$-term, which is an $SO(2)$ s-wave. Since the $\mathcal{P}_{1\pi}$-term has a b_T factor, which can be written in terms of $b_{T,x} \pm i b_{T,y}$, this term carries one unit of orbital angular momentum. In terms of power counting, then, the $m = 0$ and $m \neq 0$ amplitudes can be equally large.

In the case of vector mesons, isolating the dominant high-energy tensors which contain one power of the large scale p^+, we get for a longitudinally polarized ρ

$$\mathcal{P}(x, b_T; p, h=0) = \left(\mathcal{P}'_0 \not{p} + \tilde{\mathcal{P}}'_1 [\not{p}, \not{b}_T]\right), \tag{3}$$

and for a transversely polarized one

$$\mathcal{P}(x, b_T; p, |h| = 1) = \left(\mathcal{P}_0\,[\not\varepsilon_h, \not p] + \mathcal{P}_1\,\varepsilon_h.b_T\,\not p + \tilde{\mathcal{P}}_1\,[\not\varepsilon_h, \not p]\,\not b_T + \mathcal{P}_2\,\varepsilon_h.b_T\,[\not p, \not b_T]\right). \tag{4}$$

Let us stress that wave functions are not objects to be derived in perturbation theory, but instead represent the long-time non-perturbative evolution proceeding inside a hadron. The non-perturbative Hamiltonian of QCD does not conserve spin and orbital angular momentum separately, but instead generates mixing between orbital and spin angular momentum. Thus *if* a non-zero orbital angular momentum component somehow enters the hard scattering - and this is a crucial point - *then* the long-time evolution before or after the scattering can convert this angular momentum into the observed hadron spin. Because it is not necessary to flip a quark spin in the hard interaction, such a mechanism is totally consistent with the impulse approximation of perturbative QCD.

The challenge in high energy hadron scattering is therefore to find those large Q^2 processes in which non-zero orbital angular momentum enters, or in other words, to find those which are not "round". It turns out that in any treatment relevant to current energies the independent scattering process is not "round" but instead "flat", showing an extreme dependence on the scattering plane.

3 Q^2 Dependence of Helicity Non-conservation

Here we are concerned with the leading order description of helicity violating terms. Thus, we will consider $\mathcal{P}_{0\pi}$-type and $\mathcal{P}_{1\pi}$-type amplitudes on an equal footing. A crucial step is to elaborate a factorized form for the scattering amplitude, regarding radiative corrections. While the factorization of Ref. 1) does not apply, it is trivial to generalize the results of Ref. 3) to the case of the helicity-violating Dirac projections. A leading approximation to the soft region rearranges these corrections to obtain the following expression:

$$A(s,t) = \frac{\sqrt{2}Q}{2\pi|\sin\theta|}\int_0^1 dx\, t(\{xQ\})\, t'(\{(1-x)Q\}) \int_{-1/\Lambda}^{+1/\Lambda} db\, U(x,b,Q) \prod_{i=1}^{4} \mathcal{P}^{(S)}(x,b;p_i), \tag{5}$$

where the hard amplitude t and t' are evaluated at respective scales xQ and $(1-x)Q$ which are assumed to be large ($Q = \sqrt{s/2}$). Large logarithmic corrections to the process, with the coexistence of the two scales Q and $1/b$, are resummed in U, in such a way that $\mathcal{P}^{(S)}$ is a soft object, i.e. the non perturbative object necessary to connect short and long range physics. To evaluate the integral defined in Eq. (5), we approximate the Sudakov factor by its dominant expression at large Q

$$U(x,b,Q) \approx \exp - c \ln\frac{xQ}{\Lambda}\left[\ln\left(-\frac{\ln xQ/\Lambda}{\ln b\Lambda}\right) - 1 + \frac{\ln b\Lambda}{\ln xQ/\Lambda}\right] + (x \leftrightarrow 1-x), \tag{6}$$

with $c = 4\,\frac{4}{3}\,\frac{2}{11 - 2n_f/3} = 32/27$ for $n_f = 3$.

The helicity conserving hard amplitude is

$$a_0 = \left(\frac{\pi}{6}\right)^4 \frac{256g^4}{x^2\bar{x}^2s^2}\frac{s^4(s^2-3tu)+t^2u^2(s^2-tu)}{s^2t^2u^2}; \tag{7}$$

there are no b^1 or b^3 terms, due to the odd number of γ matrices. The first potentially helicity violating term is the amplitude a_2, containing a b^2 factor, which is found to be

$$a_2 = b^2\left(\frac{\pi}{6}\right)^4 \frac{2048g^4}{x^2\bar{x}^2s^2}\frac{s^4(s^2-3tu)-t^3u^3}{s^2t^2u^2}. \tag{8}$$

Inserting the hard parts a_i in Eq.(5), we get a value $A_i(Q,\theta)$ and perform the ratio of amplitudes $R_2 = A_2/A_0$. Results for our computation are displayed in Fig. 1. Normalizing arbitrarily the ratio to 1 at $\sqrt{s} = 2$GeV and $\theta = 90°$, we observe that R_2 decreases by a factor between 7 and 10 from $\sqrt{s} = 2$GeV to 20GeV. This is not much suppression. We have shown[2)] that various effects are likely to wash out even this suppression; this is represented by the shaded area in Fig.1. The angular dependence of R_2 is simply given by the ratio a_2/a_0.

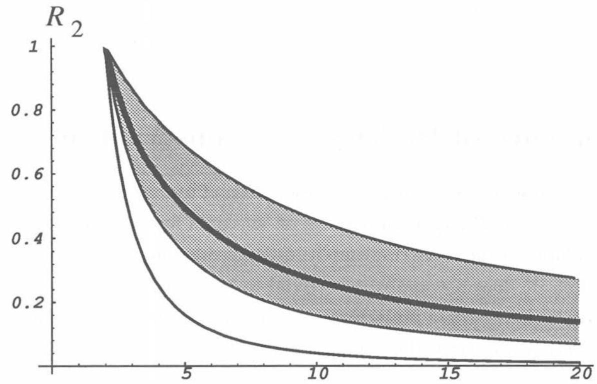

FIG. 1: The energy dependence of the R_2 ratio (thick line) and its naive $1/Q^2$ behavior (thin line). The shaded area indicates theoretical uncertainty. The ratio is normalized to unity at 2 GeV.

We are now able to consider an helicity violating process, namely $\pi\pi \to \rho_R\rho_R$. It is easy to verify the vanishing of the hard amplitude using the s-wave components of the external mesons. The first non zero term is a b^2 hard amplitude M_2 and the computation leads to

$$\begin{aligned} M_2(\pi\pi \to \rho_R\rho_R) &= \left(\frac{\pi}{6}\right)^4 \frac{128g^4}{x^2\bar{x}^2t^2u^2}b^2\left\{\frac{16(3s^2-7tu)}{3}\mathcal{P}_{1\pi}^2\mathcal{P}_{0\rho}^2 - \frac{t^3u^3}{s^4}\mathcal{P}_{0\pi}^2\mathcal{P}_{1\rho}^2\right. \\ &+ 8\frac{t^3u^3}{s^4}\mathcal{P}_{0\pi}^2\mathcal{P}_{0\rho}\mathcal{P}_{2\rho} - 16(s^2-3tu)\mathcal{P}_{0\pi}\mathcal{P}_{1\pi}\mathcal{P}_{0\rho}\mathcal{P}_{3\rho} \\ &+ \left. 4\frac{t^3u^3}{s^4}\mathcal{P}_{0\pi}^2\mathcal{P}_{1\rho}\mathcal{P}_{3\rho} + 4(s^2-3tu+\frac{t^2u^2}{s^2}-2\frac{t^3u^3}{s^4})\mathcal{P}_{0\pi}^2\mathcal{P}_{3\rho}^2\right\}. \end{aligned} \tag{9}$$

This amplitude M_2 has to be supplemented by U (Eq. (6)) and integrated over b and x.

Although this combination involves several unknown objects, notice that the angular dependence varies from one component to another. Therefore, it may be possible to analyze the contribution to helicity violation processes from different wave functions and use this information to deduce properties of the wave functions.

The energy dependence of this double helicity violating process at accessible energies can be derived as discussed above. The naive $1/Q^2$ factor is replaced by a milder suppression, due to the details of the independent scattering mechanism supplemented by Sudakov effects. At very large energies a $Q^{-1.10}$ ratio is obtained.

Including baryons is a necessary but quite intricate further step. The helicity density matrix of the ρ meson produced in $\pi p \to \rho p$ is a nice measure of helicity violating components. Experimental data[4] yield $\rho_{1-1} = 0.32 \pm 0.10$, at $s = 20.8\mathrm{GeV}^2$, $\theta_{CM} = 90°$, for the non-diagonal helicity violating matrix element. Without entering a detailed phenomenological analysis, we may use our results via the following line of reasoning. Assuming that the presence of the third valence quark, which is not subject to a third independent scattering, does not much alter the results, one views ρ_{1-1} as coming from the interference of an helicity conserving amplitude like $\pi\pi \to \rho_L \rho_R$ with a double helicity flip amplitude like $\pi\pi \to \rho_R \rho_R$. We then get an energy dependence for this matrix element similar to the one of R_2 in Fig.1.

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