SUMMARY TALK - HADRONIC INTERACTION SESSION

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Abstract: We review the recent work on calculating jets and large p_T events by perturbative QCD. Special emphasis is placed on the QED origins of the ideas. We also discuss prompt lepton pairs and baryonium.

Résumé: Nous rendons compte des calculs récents concernant les jets et les évènements à grand p_T dans le cadre du chromodynamique quantique (QCD) perturbatif. Les origines (électrodynamique quantique: QED) de ces idées sont mises en évidence. Nous discutons aussi de la physique des paires de leptons à grands moments transverses et de la physique du baryonium.

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Introduction

In the summary I shall focus on three subjects: calculating with QCD in perturbation theory, prompt lepton pairs, and baryonium. QCD perturbation theory will be treated in greatest detail because I feel it has been the key new theoretical development at the conference.

I. Calculating Jets and Large p_T Events by Perturbative QCD.

In spite of its beauty, QCD has been frustrating inasmuch as so disappointingly little can be derived about the observed hadrons - for example, one cannot prove that quarks and gluons bind to form p's and π's. The trouble is that while we hope eventually to use non-perturbative techniques, for example solitons or instantons, the only reliable calculational technique available at present is perturbation theory -the expansion in powers of the quark-gluon coupling strength g^2.

Calculations of the vacuum polarization indicate that $\alpha_k \equiv g_k^2/4\pi$, the "effective" or "running" coupling constant, varies with momentum transfer k in the manner depicted in Fig.1. At $k << 1$ GeV, one finds $\alpha_k \gtrsim 1$: perturbation theory clearly fails. At $k > 1$ GeV, the logarithmic falloff of α_k raises the hope that perturbation theory can be applied. But unfortunately,as in QED, the actual expansion parameter in most practical calculations turns out to be $\alpha_k \ell n\, k/m > 1$, where m is either the quark or the gluon mass, so the hope that perturbation theory will converge is frustrated.

Why does α_k commonly appear multiplied by $\ell n k/m$? A very general insight into the phenomenon, based on elementary quantum mechanics, was offered by Lee and Nauenberg[1)] in 1964. Consider the Hamiltonian

$$H = H_o + \alpha H_I \quad (1)$$

with

$$H_o \psi_n = E_n \psi_n . \quad (2)$$

In second order perturbation theory we have

$$\psi'_i = \psi_i + \alpha \sum_{n \neq i} \frac{(H_I)_{ni} \psi_n}{E_i - E_n} . \quad (3)$$

The key point is that large changes occur, even for small α, if the states are nearly degenerate.

Suppose, for example, that ψ_i represents the electron state and the ψ_n represent the continuum of $e+\gamma$ states. We let γ have

Nevertheless, factors such as $\alpha_k \ \ell n E/m_q \ \ell n E/m_H$, of order $\gtrsim 1$, remain. The phenomena of bremsstrahlung in QCD and QED seem so closely related that it is natural to try to use the very general Lee-Nauenberg analysis to locate experimental quantites which are free of these logarithms.

I shall discuss several cases which illustrate the main points of the recent work.

Case A is the reaction

$$e^+e^- \to \gamma \to \quad \text{hadrons} . \tag{13}$$

Old application: In $e^+e^- \to$ hadrons the γ is colorless and all hadronic final states, degenerate and nondegenerate, are summed over. Thus all $\ell n \ p/m_q$ and $\ell n \ p/m_{gluon}$ terms should cancel and the rate should be expansible in powers of α_k. And indeed, one finds by explicit calculation

$$R \equiv \frac{\sigma(e^+e^- \to \gamma \to \text{hadrons})}{\sigma(e^+e^- \to \gamma \to \mu^+ \mu^-)} = \sum_i Q_i^{\,2} \left(1 + \frac{\alpha_k}{\pi} \right) . \tag{14}$$

The fact that the second term does not contain a logarithm, and is thus only a correction of order 20 %, is of course crucial to the phenomenological use of R as an indicator of fundamental charges.

New application: While R is very important, the information on quark jets contained in perturbation theory has been lost by summing over final states. Sterman and Weinberg[3] sought to retain the information on jets while obtaining a convergent expansion in α_k at the same time. For this purpose they considered the cross-section

$$(e^+e^- \to \gamma \to \ q\bar{q} + q\bar{q} \text{ gluons} + \ldots)$$

for events with all but a fraction ε of the energy lying inside a pair of opposing cones of half-angle δ (Fig.3). As discussed in the photon case, the gluon bremsstrahlung is strongest when k and θ (the angle between the gluon line and the emitting q or $\bar{q}$ line) are small. These are the configurations where the $q\bar{q}$ gluon state is nearly degenerate with the $q\bar{q}$ state. To sum over all approximately degenerate states one must integrate not only over small k (the upper limit being characterized by ε rather than ΔE in the Sterman-Weinberg formalism) but also over $\theta \leq \delta$ (ie. over hard collinear gluons). The dominant correction term is thereby softened a small fictitous mass m_γ ; thus the continuum E_n begins at $E_o + m_\gamma$. In this example the sum in (3) becomes an integral of type $\int dE_n/(E_i - E_n)$ with lower limit $E_i + m_\gamma$ and some high energy cutoff M, and the change in state is

$$\psi'_i - \psi_i \sim \alpha \ln M/m_\gamma \; . \qquad (4)$$

The diagnosis that the logarithm is caused by nearly degenerate states immediately suggests the cure:
<u>To obtain physical quantities expansible in α_k rather than $\alpha_k \ln k/m$ one must sum over the nearly degenerate states</u>.
The logs resulting from mixing among the nearly degenerate states then cancel out. This is a reflection of what happens in the degenerate limit where $\ln m_\gamma$ becomes singular. The standard prescription for eliminating the singularity is to diagonalize H_I in the subspace of degenerate states, which of course can only be done if we include all the degenerate states together.

A famous example is the behavior of QED as $m_\gamma \to 0$. The rate for each of the individual processes

$$\begin{aligned} &ep \to ep \\ &ep \to ep\gamma \\ &\cdots\cdots\cdots \end{aligned} \qquad (5)$$

has logarithmic infrared divergences as $m_\gamma \to 0$. But when these degenerate final states are summed over, one finds that

$$\sum_n \sigma(ep \to ep + n\gamma)$$

with

$$\sum_{i=0}^{n} E_i \leq \Delta E \qquad (6)$$

is finite.

Another example is the logarithmic divergence of QED as $m_e \to 0$. Of course, this example is less famous because m_e is not really zero, and the effective expansion parameter $\alpha \ln k/m_e$ which occurs remains substantially less than one for the physical value of m_e. But in QCD we shall be interested in the analogous expansion parameter $\alpha_k \ln k/m_q$, which can exceed unity. The $\ln m_e$ divergence arises from the familiar property that bremsstrahlung from a fast-moving charge is emitted preferentially at small angles. For example in $ep \to ep\gamma$, the amplitude for radiation off the final charge line (Fig.2) is proportional to

$$A = \frac{2\,\varepsilon.p_2}{(k+p_2)^2 - m_e^2} \qquad (7)$$

where ε_μ is the photon polarization, and we treat the electron as spinless for simplicity. Using the on-shell kinematics $k^2=0$, $p_2^2=m_e^2$ we find

$$A = \frac{\varepsilon.\,p_2}{k.p_2} \;\cdots\cdots\cdots$$

$$= \frac{\varepsilon.\,p}{k(E_2 - p_2\cos\,\theta)} \;\cdots\cdots\cdots\cdots \qquad (8)$$

where θ is the angle between $\vec{p}_2$ and $\vec{k}$. At $p_2 >> m_e$, we have $E_2 \simeq p_2$ and (8) becomes

$$A \simeq \frac{\varepsilon\,.\,p_2}{kp_2\;(1-\cos\theta + \frac{m_e^2}{2p_2^2})} \;\cdots\cdots\cdots \qquad (9)$$

which exhibits clearly the preferential emission at small angles and the role of m_e in cutting it off. When the energy dominator becomes small the e and e γ states become nearly degenerate; this happens in (9) not only at $k\to 0$ but also at $\theta\to 0$. When $m_e=0$ the small angle behavior $(1-\cos\theta)^{-1} \sim \theta^2/2$ of the denominator is only partially compensated by the effect of transverse γ polarization in the numerator $[\varepsilon.p_2= -\vec{\varepsilon}\,.\vec{p}_2=-p_2\,\sin\theta \sim -p_2\theta]$; overall the amplitude is of order $1/k\,\theta$. Squaring the amplitude and integrating over phase space, we find that the key factors at small k and θ are

$$\int\frac{d^3k}{k_o}\quad\int d\cos\theta\;\left|\frac{1}{k\theta}\right|^2 \sim \int\frac{dk}{k}\;\int\frac{d\theta}{\theta} \qquad (10)$$

The integral over k, cut off by a fictitious photon mass, leads to $\ell n\; p_2/m_\gamma$; the integral over θ, cut off by m_e/p_2 as we have seen in Eq.(9), leads to $\ell n\; p_2/m_e$.

In a familiar case such as the calculation of $\sigma(ep\to ep)$ at high energy and large angle, the $O(\alpha)$ radiative corrections reduce the elastic rate by (approximately) a factor

$$[\,1 - \alpha c \,\ell n \frac{p}{m_e}\, \ell n \frac{p}{m_\gamma}\,] \qquad (11)$$

where c is of order 1. When we add $\sigma(ep \to ep\gamma)$ with $k \leq \Delta E$ as dictated by the experimental conditions, the overall reduction factor is softened to the finite value

$$[\,1 - \alpha c \ell n \frac{p}{m_e}\, \ell n \frac{p}{\Delta E}\,], \qquad (12)$$

still a substantial reduction because the numerous nearly-degenerate final states involving hard $(k > \Delta E)$ photons with small θ are not counted by the experiment.

In 1959 Kinoshita and Sirlin[2)] noted a case where the experimental conditions imply a sum over all the nearly-degenerate states-hard, nearly-collinear photons as well as soft photons -with the corresponding disappearance of $\ell n\, m_e$ as well as $\ell n\, m_\gamma$ factors. Specifically they calculated the order α radiative corrections to $\mu \to e\, \nu\bar{\nu}$. As usual the virtual-photon correction reduced the rate for $\mu \to e\, \nu\bar{\nu}$ by a factor of the form (11). Adding the rate for $\mu \to e\, \nu\bar{\nu}\gamma$ with $k \leq \Delta E$ removed the $\ell n\, m\gamma$ singularity, leaving the standard correction of the form (12). Finally, when the <u>total</u> rate was calculated including all hard as well as soft photons (i.e. by raising ΔE to its kinematics limit) the remaining $\ell n\, m_e$ singularity cancelled leaving a small correction of order α . To summarize : one finds a divergent $[O(\alpha \ell n\, p/m_e\, \ell n\, p/m\gamma)]$ change in the final state ($e\, \nu\bar{\nu}$ replaced by $e\, \nu\bar{\nu}\gamma$ and, eventually, multiphoton states), a substantial but finite $[\,O(\alpha\, \ell n\, p/m_e)]$ change in electron energy (depletion of large p_e increase in small p_e events as the hard collinear photons borrow energy) and a small $[O(\alpha)]$ change in overall rate.

With these classic results for QED in mind, it is quite easy to understand the recent proposals for QCD. Of course, QCD differs in certain respects, for example:
i) Both gluons and quarks are colored so both can radiate bremsstrahlung.
ii) In confined QCD, hadrons have no overall color, so there are no true infrared divergences -all integrals representing color radiation are cut off at a long -wavelength or low- frequency scale set by the hadron binding energy m_H.

from $\alpha_k \ln \sqrt{s}/m_q \ln \sqrt{s}/m_{gluon}$ to $\alpha_k \ln \delta \ln \varepsilon$. For sufficiently large δ and ε (eg. $\delta = 15°$ and $\varepsilon = 0.2$) the corrections are of order α_k with no large logarithms. Thus Sterman and Weinberg obtain calculable jets characterized by a cone with

$$p_T \sim p_L \delta \sim \frac{\sqrt{s}}{2} \delta \qquad (15)$$

rather than by the usual cylinder characterized by a constant p_{T_o}. If they had included in their jets only $p_T < P_{T_o}$, gluons would have been emitted copiously at larger p_T, ie. cylindrical jets would not contain most of the events and would be subject to large corrections of order $\alpha_k \ln \sqrt{s}/p_{T_o}$.

In summary: the perturbative QCD corrections reduce the $q\bar{q}$ final state by a factor of form $1 - \alpha_k \ln \sqrt{s}/m_q \ln \sqrt{s}/m_H$, ie. by close to 100 %, largely replacing it by a $q\bar{q}$ gluon state, but when the two nearly degenerate states are added the overall rate is changed by only $O(\alpha_k)$. The original back-to-back momenta of the q and $\bar{q}$ get spread over a distribution that peaks within opposing cones of order 15° (2-jet events) with a tail at larger angles (3-jet events).

In attempting to verify the Sterman-Weinberg proposal one encounters a complication: the non-perturbative conversion of quarks and gluons into hadrons. One assumes (without good theoretical justification) that this introduces a further contribution of order 350 Mev to the transverse momentum within a jet. At present energies ($\sqrt{s} \lesssim 8$ GeV) the perturbative contribution $p_T \propto \sqrt{s}$ is not sufficiently greater than the non perturbative contribution 350 Mev to verify that p_T is rising -especially since jets are not visible below about $\sqrt{s} = 4$ GeV. However, when $\sqrt{s}$ is increased to the 20-30 GeV range in the next generation of colliding beams, a clean test should be possible. The outcome will be crucial for the theory.

Case B is typified by the reaction

$$p + p \rightarrow p + p \qquad (16)$$

at large angles such as 90 °. A typical subprocess is $qq \rightarrow qq$ with gluon exchange. Evidently pp scattering is an exclusive process, with nearly-degenerate states not summed in either the initial or final state. Therefore, the Lee-Nauenberg type argument cannot be employed to justify the use of perturbation theory in this case.

Case C is typified by the reaction

$$e + p \rightarrow e + X. \qquad (17)$$

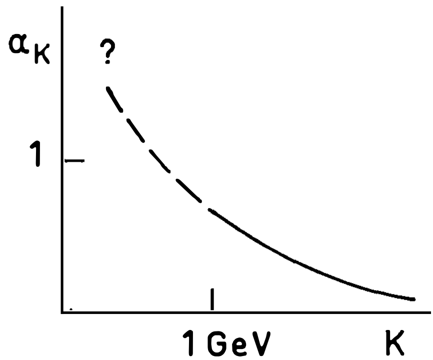


Fig. 1 : The dependence of $\alpha_k = g_k^2/4\pi$ on k in QCD.

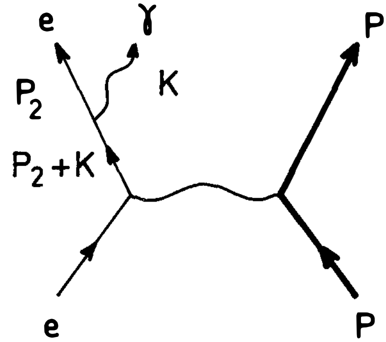


Fig. 2 : The lowest order Feynman diagram for radiation off the final electron line in ep → epγ.

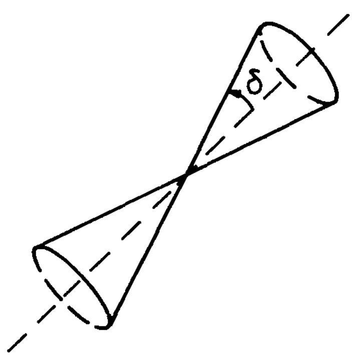


Fig. 3 : Two opposing cones of half-angle δ.

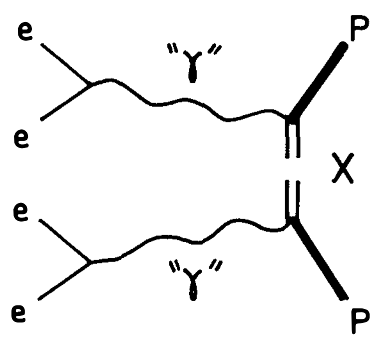


Fig. 4 : Diagramatic representation of the squared matrix element for ep → eX.

Old application: It is well known that the cross-section factors into the known ee"γ" vertex times the absorptive part of the forward Compton amplitude for a virtual photon (Fig.4).

$$\sigma(ep \to eX) \;\alpha\; \text{Im Amp}\,(\text{"}\gamma\text{"}p \to \text{"}\gamma\text{"}p;\ 0°) \tag{18}$$

This forward Compton amplitude is normally studied by means of the operator product expansion, but it is interesting to look at it from the present perspective. The forward amplitude represents an important case intermediate between examples A and B. From one point of view, although all final states are summed over, states nearly degenerate with the initial state are not summed; therefore we expect that the radiative corrections give large logarithms and a non-convergent perturbation series.

On the other hand, we are dealing with forward elastic scattering, and from QED we expect that if no acceleration of charge (color) occurs, there should be no radiation. So we have essentially a case of ∞ times zero, and closer investigation is required.

For QED the investigation was made by Kinoshita in a classic 1962 paper[4]. For QCD, while not everyone is convinced yet, analogous results have been found in low-order perturbation theory and, in leading log approximation, to all orders[5]. The Kinoshita result, for forward elastic amplitudes, is that large logarithms do occur, but only in the form of multiplicative factors associated with the charged (colored) external lines. These large, factoring logarithms refer to the dissociation of e into e+ γ, q into q + + gluon, etc.... In our case this implies

i) Everything about "γ"p → "γ"p can be calculated in perturbation theory except the state

$$|p\rangle = C_1\,|qqq\rangle + C_2\,|qqq\ \text{gluon}\rangle + C_3\,|qqqq\bar{q}\rangle \ldots \tag{19}$$

which must be treated by phenomenological "structure functions".

ii) Because of the factorization property, the structure function for $|p\rangle$ is independent of the process; thus once determined in "γ"p → "γ"p it can be applied to any process one desires.

iii) If one attempts to estimate the proton structure functions in perturbation theory via diagrams such as Fig.(5), one obtains contributions of order $\alpha_k \ln(Q^2/m_q^2)$. These are the "scaling violations".

iv) The qualitative trend of the scaling violations can be seen as follows: radiation of photons (gluons) occurs in response to acceleration of the charge (color) and grows with acceleration. Thus the gluon emission grows with the momentum transfer Q^2 received from

"γ". But the gluon takes energy from the original parton. Thus, with growing Q^2, the distribution of Feynman x in the proton structure function will shift to lower values -quarks with large x will be depleted and quarks with low x will be enhanced. This is the same trend found in work based on renormalization group calculations on "moments" of the distribution function[6].

New application: recently Politzer[7], Hinchliffe and Llewellyn-Smith[7], Sachrajda[8], and others have extended the method of Case C to a class of inclusive reactions which, unlike $e^- p \to e^- X$, could not be treated by previous QCD analyses. An example is the large p_T behavior of $pp \to \pi X$. To reduce this to the previous case I use the Mueller relation

$$\sigma(pp \to \pi X) \propto \mathrm{Im\,Amp}\,[\,\bar{\pi}\, pp \to X \to \bar{\pi}\, pp;\ 0°\,] \,. \tag{20}$$

Once again the initial state is unsummed and unaccelerated, so the Kinoshita analysis implies a factorization, with all $\alpha_k \ell nE/m_q$ terms absorbed into $|p\rangle$ structure functions and the $|\pi\rangle$ fragmentation function. The remaining effects involve only powers of α_k and are thus calculable.

Example I (Contogouris, Gaskell and Papadopoulos[9]; Field[10]). The behavior

$$d\sigma/dp_T(pp \to \pi X) \sim p_T^{-n} \tag{21}$$

expected at fixed $x_T = 2p_T/\sqrt{s}$ has posed a famous problem for QCD. Experimentally one has

$$n_{exp} \sim \begin{cases} 8.3, & p_T = 2\text{-}6 \text{ GeV} \\ 6.6, & p_T = 5\text{-}16 \text{ GeV}^{11)} \end{cases} \tag{22}$$

Theoretically, estimates in the lowest order [$O(g^4)$] using the scattering of valence quarks [Fig.(6)] yield

$$n_{th} = 4 \,. \tag{23}$$

In addition to Fig.(6) there are also $O(g^4)$ diagrams involving gluon constituents. These have recently been included[12]; they increase the magnitude at low x_T without changing the prediction $n_{th}=4$.

The new work[9,10] modifies the effective value of n_{th} by including the following effects:

i) The logarithmic variation in the "running" coupling constant g_k^2 (this occurs in the $O(g^6)$ and higher corrections to the rate).
ii) $\ell n\ p_T/m_q$ corrections in the structure and fragmentation functions (again these occur in the $O(g^6)$ and higher corrections to the rate).
iii) "Intrinsic p_T".
The result is a qualitative success; each of the new effects increases the effective value of n_{th}, and each increases n_{th} more at intermediate than at high p_T, so all effects act to reduce the discrepancy between theory and experiment.

For example, the logarithmic variation of the running coupling constant is essentially

$$g^2(p_T) \sim \frac{g^2(p_{T_O})}{1 + \frac{25}{24\pi^2}\, g^2(p_{T_O})\ \ell n\ p_T/p_{T_O}} \tag{24}$$

The factor $g^4(p_T)$ in the rate therefore falls with increasing p_T. In the present p_T range, a $\ell n\ p_T$ variation is approximately equivalent to $p_T^{1/4}$; thus $g^4(p_T) \sim p_T^{-1/2}$ and effect (i) contributes a shift $\Delta n_{th} \simeq 1/2$ in the power. Evidently this shift falls with p_T. The logarithmic scaling violations in each structure and fragmentation function contribute a similar shift in the effective power behavior. Finally, the assumption that the incoming parton distribution has an "intrinsic p_T spread" makes it easier to achieve total p_T on the order of 2 or 3 GeV, but of course has little effect on reactions with really large p_T.

Quantitatively these effects can add up to change the effective n_{th} from 4 to 8 at intermediate p_T, but the numerical result is sensitive to parameters. The contribution (i) from the running coupling constant is reliable but small. The contribution (ii) from the scaling violations is large but somewhat less reliable: since the corrections are large one should go beyond $O(g^6)$; this can be done by means of renormalization group analysis on the moments of the distribution but the coefficient of each moment is a parameter to be fit (from deep inelastic scattering in the case of the proton structure function). Finally the "intrinsic p_T" (iii) is the largest contribution of all, but is completely phenomenological. Thus it cannot be said that n=8 is predicted by the theory; rather the new theoretical developments appear to have converted a major

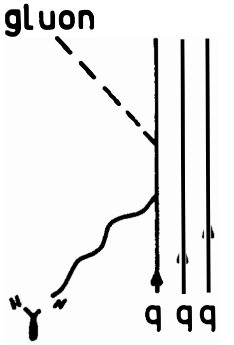


Fig. 5 : Diagram for a typical matrix element contributing to the scaling violation in the proton structure function.

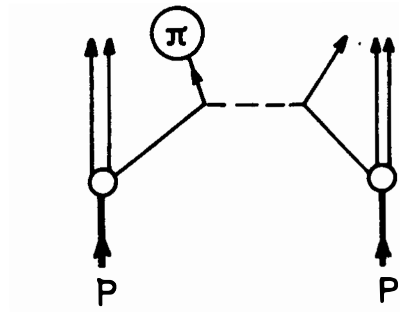


Fig. 6 : Diagram representing lowest order valence quark contribution to pp → π X.

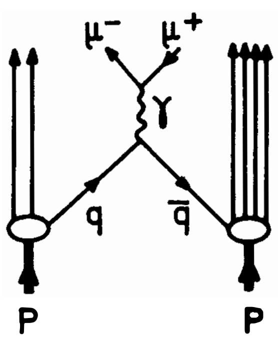


Fig. 7 : Drell-Yan diagram for $pp \rightarrow \mu^+\mu^- X$.

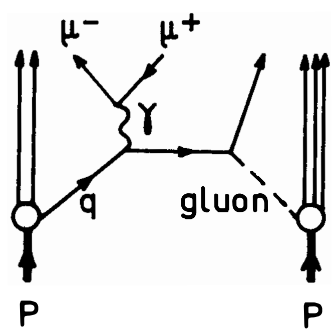


Fig. 8 : Order g^2 contribution to $pp \rightarrow \mu^+\mu^- X$ involving a gluon constituent of the proton.

discrepancy between theory and experiment into a more minor problem of understanding details.

Example II (work reported at this meeting by Petronzio[13] and Michael[14]) . The cross-section

$$d\sigma/dp_T(pp \rightarrow \text{"}\gamma\text{"} + X)$$
$$\qquad\qquad \llcorner\!\rightarrow \mu^+\mu^-$$

can be treated in the same way as $pp \rightarrow \pi X$[7,8]. The order g^o contribution is given by the Drell-Yan diagram (Fig.7). Even if the quark constituents are given a reasonable amount of intrinsic p_T, it is too small to explain the several per cent of events which have p_T= 2 to 5 GeV. The $O(g^2)$ corrections, for example Fig.(8), introduce gluon constituents and are capable of fitting the p_T=2 to 5GeV events. That brings us to the second major topic of the meeting.

II. Prompt Lepton Pairs.

When "prompt" leptons were first studied at large p_T, their origin was quite unclear. By now a fairly detailed picture has emerged. If we plot $d\sigma(pp \rightarrow \mu^+\mu^- X)/dM_{\mu\mu}$ versus $M_{\mu\mu}$ (Fig.9) we find three regions where different production mechanisms, each interesting in its own right, are at work:

Region I consists of the peaks at $M_{\mu\mu} \simeq$ 3 and 9 GeV. These peaks arise from basically new physics (charmonium and upsilon production).

Region II is the straight part of the curve at $M_{\mu\mu} >$ 3 GeV. It is interpreted in terms of hard constituent-constituent collisions such as the Drell-Yan mechanism (Fig.7). We note that the original reaction

$$p + p \rightarrow \mu^+\mu^- + X \qquad (25)$$

does not, by itself, provide a very incisive test of the $q\bar{q}\rightarrow\gamma\rightarrow\mu^+\mu^-$ subcollision of Drell-Yan; while the distribution of valence quarks in the proton is rather well known from $e^-p \rightarrow e^-X$, the distribution of "sea" $\bar{q}$'s in the proton is less well determined.

The more recently studied reaction

$$\pi + p \rightarrow \mu^+\mu^- + X \qquad (26)$$

reported on by Pilcher[15] and Romana[16] has the advantage that it can proceed by collision of a valence $\bar{q}$ from the pion with a valence q from the proton. Even without knowing the detailed distribution of partons within the pion one can say that

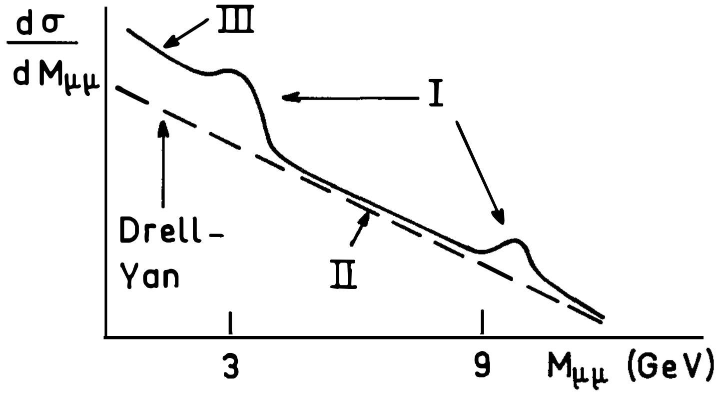


Fig. 9 : $d\sigma\ (pp \to \mu^+\mu^-\ X)/dM_{\mu\mu}$ experimental (solid line) and Drell-Yan (dotted line).

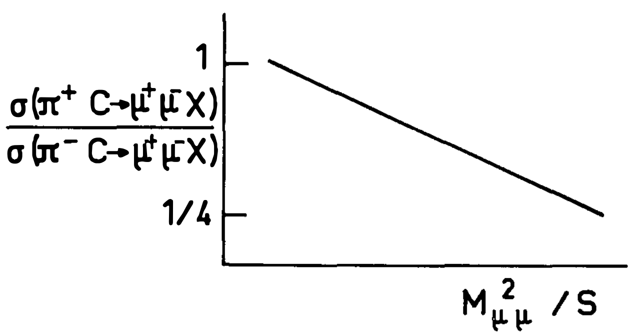


Fig. 10 :Schematic representation of the data for $\sigma(\pi^+C \to \mu^+\mu^-\ X)/\sigma(\pi^-C \to \mu^+\mu^-\ X)$.

$$\frac{\sigma(\pi^+ C \to \mu^+\mu^- X)}{\sigma(\pi^- C \to \mu^+\mu^- X)} \longrightarrow \frac{1}{4} \tag{27}$$

at large Feynman x (where valence constituents dominate), because the valence $\bar{q}$ in π^+ is $\bar{d}(Q^2_d = 1/9)$ whereas the valence $\bar{q}$ in π^- is $\bar{u}$ $(Q^2_u = 4/9)$. This prediction is nicely satisfied by the data (Fig.10). That the presence of a valence $\bar{q}$ in the pion truly favors $\sigma(\pi p \to \mu^+\mu^- X)$ over $\sigma(pp \to \mu^+\mu^- X)$ at large Feynman x is now strikingly verified by the data[16] (Fig.11) which shows the ratio of these two processes reaching 300 at large $M^2_{\mu\mu}/s$. In the near future it should be possible to extract the parton distribution within the pion from $\pi p \to \mu^+\mu^- X$ data.

When it becomes experimentally feasible, study of the reaction

$$\bar{p} + p \to \mu^+ + \mu^- + X \tag{28}$$

will also be interesting as emphasized by Lederman[17]. Here the $\bar{p}$ is the source of valence $\bar{q}$'s, with the same (relatively well-known) x-distribution as the valence q's in p. Thus, knowledge of the absolute rate for this reaction would provide one of our few clean tests of the color factor of 3.

The successful fits to the Drell-Yan model achieved in $pp \to \mu^+\mu^- X$ involve the rates integrated over p_T. The model fails to describe the small fraction of events in which the pair has large p_T. Here other processes must be at work, and we have described earlier in the talk how gluon constituents and QCD corrections can be used to explain the data. In the particular case of $pp \to \mu^+\mu^- X$ with its absence of valence $\bar{q}$'s it is believed that these corrections may be significant even at small p_T.

Region III refers to $M_{\mu\mu} < 3$ GeV. The $\mu^+\mu^-$ pairs are most numerous here, and most of the early events which called attention to the puzzle of large p_T leptons came from this region. Nevertheless, this region has been less productive of insights into the fundamental mechanisms involved. The Drell-Yan prediction is typically a factor of 10 below the data at $M_{\mu\mu} < 3$ GeV, and no simple quark picture or other comprehensive explanation has worked. It appears that here one is in the relatively low-momentum region where perturbative QCD breaks down and the quark dynamics becomes more complicated.

For example, M. Duong-Van[18] reported on a model which makes no mention of quarks, but is based on the subprocess $\pi\pi \to \gamma \to \mu^+\mu^-$

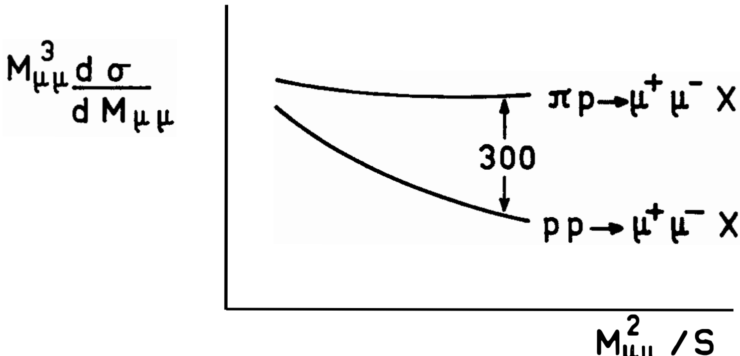


Fig. 11 : Schematic representation of the data for $M^3_{\mu\mu}\ d\sigma/dM_{\mu\mu}$ in $\pi p \to \mu^+\mu^-\ X$ and $pp \to \mu^+\mu^-\ X$.

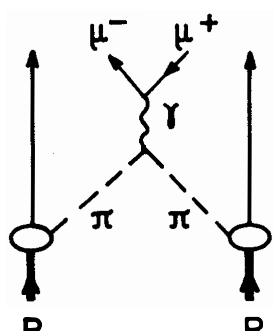


Fig. 12 : Diagram contribution of the subprocess $\pi\pi \to \gamma \to \mu^+\mu^-$ to $pp \to \mu^+\mu^-\ X$.

(Fig.12). Taking account of the pion form factor in the time-like region, this SLAC model fits the data well at the ρ peak and below, while falling too low at $M_{\mu\mu} > 1$ GeV (Fig.13). Expressed in quark language, this model would involve soft $q\bar{q}$ pairs and gluons in a complicated way.

An alternative explanation of pair production at $M_{\mu\mu} < 3$ GeV, based on quark bremsstrahlung, had been proposed by Farrar and Frautschi[19] and others.

In addition to the Dalitz μ-pair conversions of virtual photons it required real γ's in copious amounts [$\gamma/\pi > 10\%$ at large s, p_T]. The experiment of the Willis group, reported to us by Rehak[20], attacked this question by studying μ pairs down to $M_{\mu\mu}$ = a couple of hundred MeV. They were able to account for essentially all μ pairs without quark bremsstrahlung, and thus [using the tight connection between low-$M_{\mu\mu}$ Dalitz pairs and real γ's in a bremsstrahlung mechanism] estimate a limit

$$\gamma_{direct} / \pi^o < 1\ \%$$

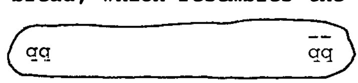
on real γ's in a range p_T=2 to 3 GeV, $\sqrt{s}$=52 GeV where a substantially higher value was expected in the quark bremsstrahlung model.

III. Baryonium, etc.

A third major theme of the meeting was baryonium. Strictly speaking the name "baryonium" refers to B=0 levels with small Γ_{meson}. Sometimes, but not always, they also have small Γ_{total}. Recently this subject has flowered forth experimentally[21] to the point where there are now on the order of 10 levels at $M > 2M_N$ and 5 levels at $M < 2M_N$ that are candidates for baryonium.

Baryonium states have long been expected on the basis of NN potentials suitably crossed to the $N\bar{N}$ channel[22]. They are required by the quark duality diagrams for $N\bar{N} \rightarrow N\bar{N}$[23]. More recently, they have been extensively treated by Johnson and Thorn, Jaffe, Chan and Høgaasen and others[24] in the MIT bag model, where the original baryonium states appear as just one example of a whole class of multiquark resonances.

I propose to call the Johnson-Thorn-Jaffe-Chan-Høgaasen theory of these multiquark resonances the "baguette model" in honor of the long thin French bread, which resembles the highly streched bags



used in the model. The length $r \sim \sqrt{s}$ of the baguette ensures that

$$\ell = |\vec{r} \times \vec{p}| \sim \sqrt{s} \times \sqrt{s} \sim s \tag{29}$$

lies as high as possible, near the leading Regge trajectory. The high orbital angular momentum plays the essential role of inhibiting the decay into mesons via recombination of the q's and $\bar{q}$'s at the ends of the baguette.

A crucial test of the bag model for multiquark configurations such as $qq\bar{q}\bar{q}$ is that exotic resonances are also predicted. Why complicate in this manner the highly successful $q\bar{q}$ model of mesons when it has predicted exactly the observed states up to now, and when no exotic state has ever been well authenticated ? Jaffe[24] and Høgaasen[24] have given us the answer: the quarks and gluons of QCD provide degrees of freedom that should express themselves in a richer spectrum of mesons than is provided by $q\bar{q}$ alone, and the semiphenomenological bag Lagrangian predicts a greatly expanded spectrum.

Let us review from another standpoint some of the reasons why multiquark levels are both expected and hard to see. I begin by reminding you of the dual-resonance model plot of J versus M^2 (Fig.14) where the degeneracy at a given M^2 increases rapidly as one proceeds from the leading to the daughter trajectories. The degeneracies of lower trajectories are so great that the overall level density $\rho(M)$ in this model increases as

$$\rho(M) \sim e^{bM} \tag{30}$$

with b of order m_π^{-1}.

Next consider the J versus M^2 plot from the point of view of the bag model (Fig.15). This model also predicts $\rho(M) \sim e^{bM}$. On the other hand, taking nonexotic mesons as an example, the density of $q\bar{q}$ states only rises as a power, $\rho(M) \sim M^p$. Such states are dominant only at low M or on the leading trajectory. As the mass is increased (or as we proceed down from the leading trajectory) successively more complicated states such as $qq\bar{q}\bar{q}$, $q\bar{q}$ gluon etc... are found, and it is the sum over all of these states which grows exponentially.

If we fix our attention on a particular set of quantum numbers J, J_z, B,S,I,Q..., the level density still grows as e^{bM}. The typical level width is $\Gamma \gtrsim m_\pi$ for most hadron resonances. Thus the spacing between levels rapidly becomes less than Γ , ie the levels overlap above a mass which is on the order of 2 GeV for low J,

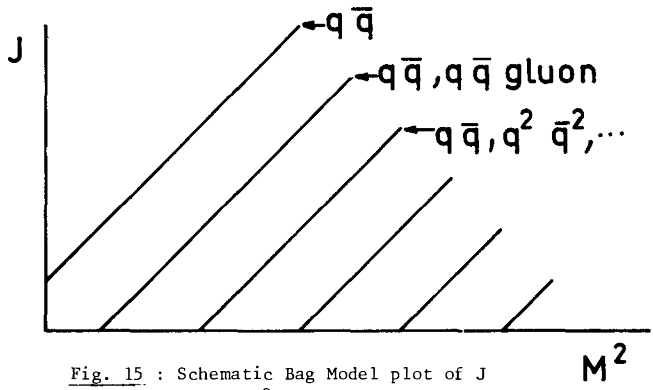


Fig. 15 : Schematic Bag Model plot of J versus M^2.

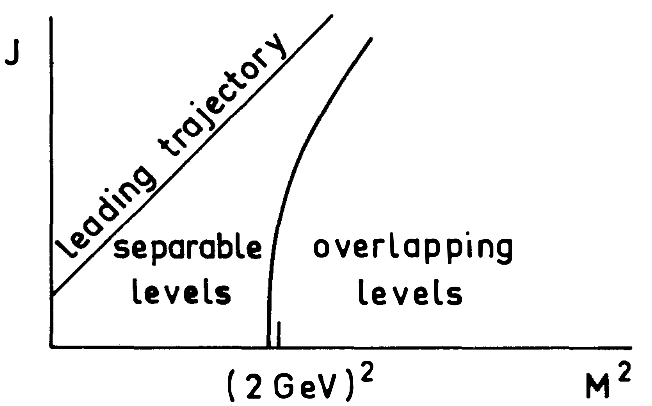


Fig. 16 : Schematic plot indicating where Γρ (M,J) is large enough to make individual levels (of a given J^P, B, S, I,...) overlap.

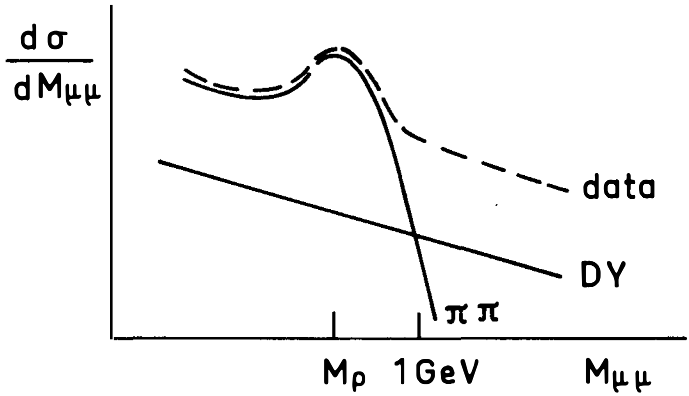


Fig. 13 : Schematic comparison of the data, Drell-Yan (DY) prediction, and $\pi\pi \to \gamma \to \mu^+\mu^-$ subprocess prediction for $pp \to \mu^+\mu^- X$.

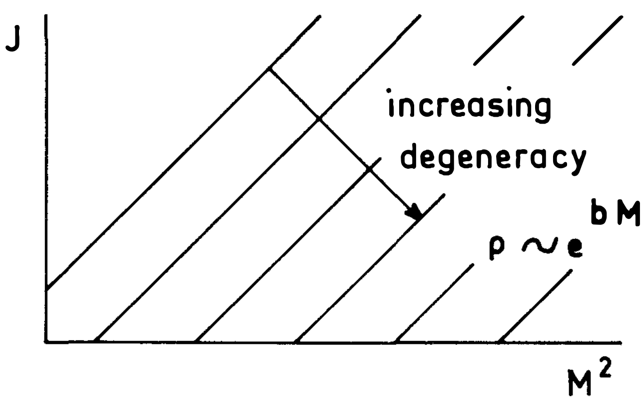


Fig. 14 : Dual Resonance Model plot of J versus M^2.

somewhat higher for higher J[25] (Fig.16). In the overlap region, ordinary levels do not stand out as individual resonance peaks and cannot be isolated even by phase shift analysis on a given reaction. Thus they are hard to see ; one is reduced to looking for statistical effects such as Ericson fluctuations[26]. The relevance for the bag model is that most multiquark levels are in the overlap region.

It follows that only rather special multiquark states have a good chance to stand out experimentally:

i) "Baguette" states near the top trajectory may have $\Gamma_{tot} \ll m_\pi$, or at least have small Γ_{meson}, as discussed earlier.

ii) In an exotic channel, the first couple of exotic resonances should not overlap even if they have rather high masses (as predicted by Jaffe[24]) and normal widths of order m_π.

But even in these favorable cases, the multiquark states usually couple rather weakly to "normal" states. Thus theorists should furnish not only lists of states, but also suggestions for favorable production and formation reactions.

Conclusion

If there was a common thread running through this meeting, it was a process of taking more seriously the gluons and associated quark pairs suggested by the full QCD dynamics, even though the solution of QCD remains as elusive as ever.

In conclusion I wish to thank the founder and organizer of the Moriond Conferences, Tran Thanh Van, for his efforts towards making this meeting so fruitful and pleasant.

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