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PERTURBATIVE QCD AND JETS

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Abstract

A brief review of some of the recent progress in perturbative QCD is given.

Rencontre de Moriond 86 SPhT/86/077

1 . Introduction

In this talk I shall give a brief summary of some of the topics which have received considerable attention in studies of perturbative QCD over the last year. This is by no means an exhaustive account of what has been done, but rather a sampling of the progress and new understanding which has been achieved.

2 . Heavy Quark Production

I would like to separate the discussion of heavy quark production into two parts, (A) heavy quark production in gluon jets and (B) heavy quark production in hadronic collisions in general

2.A. Heavy Quarks in Gluon Jets

In leading orders of QCD ρ, the number of heavy quark pairs in a gluon jet of energy Q, is given by

$$\rho = \frac{1}{3\pi}\int_{4M^2}^{Q^2} \frac{dK^2}{K^2}\,\alpha(K)\left(1 + \frac{2M^2}{K^2}\right)\sqrt{\frac{1}{4} - \frac{M^2}{K^2}}\;\mathcal{N}_g(Q^2, K^2) \qquad (1)$$

where

$$\mathcal{N}_g = \left(\frac{\ln Q^2/\Lambda^2}{\ln K^2/\Lambda^2}\right)^a \frac{\exp\left\{\left[\frac{2C_A}{\pi b}\ln Q^2/\Lambda^2\right]^{1/2}\right\}}{\exp\left\{\left[\frac{2C_A}{\pi b}\ln K^2/\Lambda^2\right]^{1/2}\right\}} . \qquad (2)$$

M is the mass of the heavy quark and lowest order perturbation theory is obtained by setting $\mathcal{N}_g$, the gluon multiplicity at mass K, equal to 1.

One reason which makes a measurement of ρ especially interesting is that one can do a reliable estimate[1] of the non-perturbative corrections to (1). The leading non-perturbative corrections can be expressed in terms of $\langle\alpha F^2\rangle$, the gluon condensate introduced by Shifmann, Vainshtein and Zakharov[2]. $\langle\alpha F^2\rangle$ is one of the parameters which must be added to perturbative calculations to take into account the fact that the perturbation series is not Borel summable. Thus in the limit of zero mass light quarks one finds

$$\delta\rho = \frac{\langle\alpha F^2\rangle}{N^2-1}\,\frac{\alpha(M)}{M^4}\left(-\frac{1}{30}C_F + \frac{53}{3780}C_A\right) \qquad (3)$$

with N the number of colors and C_F and C_A the Casimir operators in the fundamental and adjoint representations respectively. (3) gives a very small correction to (1).

Light quark condensates can contribute an additional correction to ρ proportional to

$$\frac{m\langle\bar{\psi}\psi\rangle}{M^4} .$$

Though an exact calculation of light quark condensates to ρ has not been performed, they likely make an extremely small contribution also.

Thus Eq.(1) should give a very accurate prediction for the number of heavy quarks in a gluon jet with the main uncertainty coming from the values of Λ and M chosen. The M appearing in (1) is the current algebra quark mass so that a good measurement of ρ could give important information toward determining this bare mass. Of course, especially at present energies, charm production in gluon jets is the dominant flavor in such heavy quark production. I believe that it is very interesting and important to make a good measurement of ρ for this process.

2.B. *Heavy Quark Production in Hadronic Reactions*

Significant progress in understanding heavy quark production has been made in the last year or so. In particular the concepts of gluon-gluon fusion, flavor excitation and intrinsic heavy quarks have undergone a significant revision in the past year. At present the experimental situation is still unsettled with regard to charm production in hadronic collisions[3] though we may expect considerable clarification in the next year or so. At present there is no data on production of quarks heavier than charm in hadronic collisions.

There are three firm statements which can now be made about heavy quark production which were not completely apparent a year ago. (i) At leading order in $\alpha(M)$ the gluon-gluon fusion term, $gg \to Q\bar{Q}$, shown in Fig. 1, should give the dominant contribution to heavy quark production[4] as

$$\sigma = \int dX_1 dX_2 \; G(X_1, M^2) G(X_2, M^2) \sigma_{gg\to Q\bar{Q}} . \qquad (4)$$

$\sigma_{gg\to Q\bar{Q}}$ is the order α^2 contribution to $Q\bar{Q}$ production by gluons having momenta $X_1 p_1$ and $X_2 p_2$ respectively, while M is the mass of the heavy quark. The total cross section is proportional to $\alpha^2(M^2)/M^2$. (ii) Intrinsic heavy quark components in the wave function of the colliding hadrons give a contribution of size Λ^2/M^4 to the cross section[5] . (iii) The concept of flavor excitation is no longer operative[4].

The arguments for (i), recently given by Collins, Soper and Sterman are really quite straightforward. The essential idea is that heavy quark production involves momentum transfers at least of order M and so constitutes a hard process taking place over short times. The process is then essentially no different than massive μ-pair or Z-producton, except that in heavy flavor production gluons play the dominant role. Eq.(4) can be systematically improved in powers of $\alpha(M^2)$. We shall come back to the question of what energies are actually necessary to realize the dominant role played by the lowest orders of perturbation theory.

With respect to (ii) high energy hadrons undoubtedly have intrinsic heavy quark components in their wave functions with probability of size Λ^2/M^2. It requires a hard scattering to set free these heavy quarks, thus resulting in a production cross section of size Λ^2/M^4. Such a cross section is certainly negligable for quarks heavier than charm and is likely true for charm.

With regard to (iii) the idea of flavor excitation, illustrated in Fig. 2, was that a heavy quark pair might exist in the wave function of a fast hadron, however, in a perturbative rather than in an intrinsic (higher twist) sense. While undergoing a scattering with another hadron this heavy quark system would then be set free, generally in the fragmentation region of the hadron of which it had been a part. This view of large-X charm production is not incorrect, however, this component is already included in the gluon-gluon fusion contribution[4].

Despite the progress made in understanding heavy flavor production in hadronic collisions there remain many questions which are still unanswered or for which only partial answers exist.

- 1 - I have said that gluon-gluon fusion, as shwon in Fig. 1, dominates heavy flavor production at high energies. Consider now graphs of the type shown in Fig. 3. Indeed, such graphs are order $\alpha(M^2)$ smaller than the leading contribution, however, the fact that

$$\frac{\sigma_{gg\to gg}}{\sigma_{gg\to Q\bar{Q}}} \gtrsim 100 \qquad (5)$$

at order α^2 means that higher order contributions could be competitive with the leading term. Whether or not this is so is not known at present. It is very important to evaluate the $\alpha(M^2)$ corrections to gluon-gluon fusion before one can have any confidence in theoretical calculations of heavy flavor production.

- 2 - Recently Ellis and Leon[6] have done the order $\alpha(M^2)$ correction to gluon-gluon fusion *in the large-X region.* They find a very small correction, but in the process of their work they noticed that the charm cross section predicted from gluon-gluon fusion depends strongly on the choice of the charm mass, much more strongly than the dimensional factor $1/M^2$. I think the origin of this dependence is clear. To produce a charm-anti-charm system at mass $\mathcal{M}$ in a collision of center of mass energy $\sqrt{s}$ means that the X_1 and X_2 of the colliding gluons satisfy $X_1 X_2\, s = \mathcal{M}^2$. Especially at fixed target energies this requires X values which are not too small and the rapidly falling gluon distribution strongly favors $\mathcal{M}$ values as close to threshold as possible. Forcing $\mathcal{M}$ close to threshold can vitiate the use of perturbation theory. The application of perturbation theory requires that one be able to produce a quark-anti-quark pair in a mass region $2M \leq \mathcal{M} \leq 2CM$, with C a constant not close to 1, (C = 2 might be a reasonable choice,) without encountering a strong X-dependence of the gluon distributions in that range of $\mathcal{M}$. As an example take M = 1.5GeV then perturbation theory should be valid so long as one does not encounter strong gluon structure function dependence in the X-region up to $X \approx \frac{4M}{\sqrt{s}} = \frac{6}{\sqrt{s}}$, at least for perfectly central production. Now if the gluon distribution varies like $XG(X) \sim (1-X)^5$, perturbation theory should be valid when $\left(1 - \frac{6}{\sqrt{s}}\right)^{10}$ is close to 1. (Recall that the gluon distribution comes in quadratically in gluon-gluon fusion.) Such a criterion cannot be met at fixed target or even ISR energies. Thus we would expect a strong dependence on the choice of the charm mass and in addition, we might expect a strong energy dependence as one goes from fixed target and ISR energies. However, at collider energies the mass dependence should again become weak with the cross section depending on M^2 only as $1/M^2$. Here again, as in the case of heavy quarks coming from gluon jets considered in the previous section, a good measurement of the cross section for charm production when combined with a higher order calculation and a knowledge of the small-X gluon distribution could provide a determination of the bare charm mass.

3 . Small-X Physics

Interest in small-X physics was stimulated a few years ago when the reliability of the Altarelli-Parisi equation at SSC energies was questioned because of the rapid growth in gluon and sea-quark distributions predicted in the small-X region. As we shall see in a moment there are indeed important corrections to the Altarelli-Parisi equation at very small values of X, but such corrections do not modify predictions for new article production at SSC energies since these modifications are only effective at mass values of a few GeV.

To see that the gluon distribution must increase rapidly at small values of X one may take the Altarelli-Parisi equation as a starting point. However, the Altarelli-Parisi equation requires input X-distributions at some Q_0^2, so we assume that $XG(X,Q_0^2)$ is given. (Quark interactions are not very important in small-X evolutions and we shall only keep the purely gluonic part of the theory in our present discussion.) Then

$$XG(X,Q^2) = \int_X^1 \frac{dX'}{X'} K\left(X/X',Q^2,Q_0^2\right) X'G\left(X',Q_0^2\right) . \qquad (6)$$

now,

$$K\left(X/X',Q^2,Q_0^2\right) = \exp\left\{2\sqrt{\frac{C_A}{\pi b} \ln\left(\frac{\ln Q^2/\Lambda^2}{\ln Q_0^2/\Lambda^2}\right) \ln X'/X}\right\} \qquad (7)$$

with k a known and slowly varying function when X'/X is large. If $X'G\left(X',Q_0^2\right)$ does not grow strongly with X', the small-X behavior of $XG(X,Q^2)$ will be determined by the growth of K when X/X' is small. In general, however, we have no reason to expect $X'G\left(X',Q_0^2\right)$ to be slowly varying. Indeed it likely grows[7] something like $(X')^{-1/2}$ so long as X' is not too small, so that the small X behavior of $XG(X,Q^2)$ is determined both by the initial distribution and by K. Since the initial distribution is not calculable within perturbative QCD the small X behaviour of $XG(X,Q^2)$ cannot be accurately calculated. Nevertheless the growth of K for X' fixed as X becomes small gives a lower bound for the X dependence of $XG(X,Q^2)$ and we may conclude that $XG(X,Q^2)$ grows rapidly for small X, at least in that domain where the Altarelli-Parisi equation is valid.

Now I would like to review the physical picture[8] of this growth in $XG(X,Q^2)$ at small X. Suppose we consider a proton of momentum $\vec{p}$, in a frame where $\vec{p} = \vec{1}_z p$ and with p chosen very large. In this frame $XG(X,Q^2)$ represents the number of gluons in the proton, per unit rapidity interval, with longitudinal momentum centered about Xp and having transverse size $|\Delta \underline{b}| \approx 1/Q$. Now when $XG(X,Q^2) \gtrsim Q^2R^2$, with R the proton radius, the gluons within this unit of rapidity begin to spatially overlap in the longitudinally thin disc which they occupy. When X is further decreased we may expect the gluons in this rather dense system to scatter and annihilate with one another, eventually reaching a saturation limit as $X \to 0$. When the gluon density becomes large enough that gluon interactions and annihilations are no longer negligable, the Altarelli-Parisi equation ceases to be valid and (6) does not represent the actual gluon density[8].

In the domain of X and Q^2 where gluon interactions due to gluon crowding are small one can in fact give a modified Altarelli-Parisi equation which takes these interactions into account. In this case one has[8,9]

$$Q^2 \frac{\partial}{\partial Q^2} XG(X,Q^2) = \frac{\alpha C_A}{\pi} \int_{X'}^{1} \frac{dX'}{X'} \frac{X}{X'} \gamma(X/X') X'G(X',Q^2) - \frac{4\pi^3}{N^2-1} \left(\frac{\alpha C_A}{\pi}\right)^2 \frac{1}{Q^2} \int_X^1 \frac{dX'}{X'} (X')^2 G^{(2)}(X',Q^2) \tag{8}$$

as the modified equation. $G^{(2)}$ is a two gluon correlation. The second term in (8) represents gluon recombination or, equivalently, gluon shadowing. (8) should be a valid equation for small X so long as the second term is small compared to the first term. At X and Q^2 values for which both terms on the right-hand-side of (8) are comparable this equation can no longer be trusted as higher gluon correlations also become important.

Nevertheless one should be able to use (8) to get an indication of where the non-linear effects become important. We may estimate the regions where gluon saturation occurs by setting the two terms on the right-hand-side of (8) to be equal and by using $\gamma(X) \to 1$ the small-X limit of this anomalous dimension. Let us do this estimate first for a large nucleus. Our normalization is such that[9]

$$G^{(2)}(X,Q^2) = \frac{\left[G(X,Q^2\right]^2}{\frac{8}{9}\pi R^2} A \qquad (9)$$

where $G^{(2)}$ is the two gluon correlation in a large nucleus of size R and G is the protons gluon distribution. Anticipating $Q^2 \approx 1$ GeV^2 we set $\frac{\alpha C_A}{\pi} \approx \frac{1}{3}$ and obtain

$$XG(X,Q^2) \approx 15Q^2 A^{-1/3} \qquad (10)$$

as the boundary of saturation. For $A^{1/3} \sim 5$-6 it should be possible to achieve saturation of gluon densities in the region $X \lesssim 0.02$-0.03 with $Q^2 \approx 1$-2GeV^2. Such regions would be available in heavy ion collisions at RHIC and it is exactly these gluons which should eventually thermalize to give the central region of the quark-gluon plasma expected at such a collider.

An analogous estimate for proton-proton or proton-anti-proton collisions is a little more difficult because the proton is not really a loosely bound system. Nevertheless, in order to get an estimate of the Q^2 and X-values for which gluon saturation might occur in a high energy hadron-hadron collision we shall take as a crude model the picture in which a proton is a loosely bound system of three valence constituent quarks. Then, proceeding as in the nuclear case, one needs

$$XG(X,Q^2) \approx \frac{25}{\pi\alpha} Q^2 R^2 \qquad (11)$$

for saturation, with R now the proton radius. For $Q^2 \sim 1$GeV^2 and R = 1fm one needs extremely small values of X, probably near 10^{-4} in order to have a chance of reaching saturation. Clearly heavy ions appear more efficient in producing such dense systems of gluons, although we shall see in Section 4 that it is possible to produce a high gluon density at much lower energies and at much larger Q^2 over a severely restricted part of the proton.

The dense system of gluons which we have been talking about is part of the wave function of a high momentum heavy ion or proton. It will not be in either a kinetic or flavor equilibrium, although such an equilibrium may indeed take place after a collision. Are such high density systems of interest and if so, why are they of interest ? What I find fascinating about such high density gluon systems is the fact that one is producing very high chromodynamic field strengths. One can see the size of the potential and field strengths involved by the following argument : in light-cone gauge $XG(X,Q^2)$ represents the number of gluons of size $|\Delta \underline{b}| \approx 2/Q$ in a unit of rapidity centered about X. XG also represents the size of A_μ^2 due to quanta of this type. Now the linear and non linear terms in (8) balance when $XG \sim Q^2/\alpha$. This means that saturation occurs, when making field measurements coherently over a transverse size $r \approx 2/Q$ at values $\vec{E}^2 \sim \vec{B}^2 \sim \frac{1}{r^2}\frac{1}{\alpha}$. These are very large field strengths, of such size that the individual quantum description of the field is no longer appropriate. If there are interesting non perturbative effects in QCD, at distances small compared to $1/\Lambda$, they would likely show up most clearly in such high field configurations. In any case, the physics of the regime of high field strength gluon saturation is that of non perturbative QCD though a non perturbative regime where the normal condensates reflecting the non perturbative vaccum are irrelevant.

4 - Minijets and Related Topics

Minijets, jets of transverse momentum greater than, say, 5GeV are strongly produced at the CERN collider having an inclusive cross section of order 10mb at the highest collider energies. Because α is small at $p_\perp$ = 5GeV one has the possibility of explaining at least portions of this data in terms of perturbative QCD. To see that the problem of relating minijet data to perturbative QCD is in fact non trivial we begin by reminding the reader that the cross section for two jet production in the minijet regime is given by

$$\frac{d\sigma}{dy_1 dy_2 dp_\perp^2} = \frac{\pi}{s} G(X_1, p_\perp^2) G(X_2, p_\perp^2) \quad \hat{\sigma}_{gg \to gg} \tag{12}$$

where $\hat{\sigma}$ is the elementary cross section as given, for example, in Ref. 10. (We shall see later that (12) is probably not a reliable starting point but this is the usual formula and will suffice for our present purposes). For simplicity we have limited ourselved to gluon jets, which should dominate minijet production. Then the total inclusive production for a minijet having $p_\perp > M$ is given by

$$\sigma(M,s) = \int \frac{dX_1}{X_1} \frac{dX_2}{X_2} \frac{d\sigma}{dy_1 dy_2 dp_\perp^2} \tag{13}$$

Now, is it possible to calculate the large s behavior of $\sigma(M,s)$ for M fixed ? Clearly large s values require samll X values of $G(X,M^2)$ and we have seen in the last section that one is not in a position to do a reliable calculation of the small X behavior of G. Thus we feel that a reliable calculation of the energy dependence of minijet production is not possible purely within perturbative QCD. However, if accurate values of gluon distributions were available from deeply inelastic lepton measurements at small values of X it should be possible to predict σ, and we shall indicate how this can be done a little later in this section.

In the remander of this section I would like to describe a minijet measurement which can be carefully discussed within perturbative QCD. The discussion given below summarizes part of a piece of work I have done in collaboration with H. Navelet and which should be available in a more complete version soon. The measurement is most easily described in terms of the two jet inclusive cross section, $\frac{d\sigma}{dy_1 dy_2 d^2p_{1\perp} d^2p_{2\perp}}$. Suppose $y_2 > y_1$ with $y_1 - Y_{min} = \Delta_1$, with $Y_{max} - y_2 = \Delta_2$ where Y_{max} and Y_{min} are the maximum and minimum values of y_2 and y_1. Then define

$$\sigma(M,s) = \int d^2p_{1\perp} d^2p_{2\perp}\, \Theta\,(p_{1\perp}-M)\ \Theta\,(p_{2\perp}-M)\ \frac{d\sigma}{dy_1 dy_2 d^2p_{1\perp} d^2p_{2\perp}}\ , \quad (14)$$

where the variables Δ_1 and Δ_2 have been suppressed. We may use factorization to write

$$\sigma(M,s) = \left(\frac{\alpha C_A}{\pi}\right)^2 \frac{\pi^3}{2M^2}\, X_1\left(G\left(X_1, M^2\right) + \frac{4}{q}\, Q\left(X_1, M^2\right)\right) X_2\left(G\left(X_2, M^2\right) + \frac{4}{q}\, Q\left(X_2, M^2\right)\right) f \quad (15)$$

where $G + \frac{4}{q}Q$ is the usual gluon plus quark factor which appears in jet physics. In lowest order perturbation theory $f = 1$ while in a leading logarithmic approximation $f = f(\alpha y)$ with $y = ln\left(\frac{X_1 X_2 s}{M^2}\right)$. We shall be mainly concerned here with a discussion of the behavior of $\sigma(M,s)$, or f, in the leading logarithmic approximation. The α which appears in (15) and throughout the rest of this discussion is $\alpha(M)$ which is already small when M is in the range of 5-10GeV.

In the leading logarithmic approximation one finds[11,12]

$$f = \frac{1}{2\pi}\int_{-\infty}^{\infty} \frac{d\nu}{\nu^2 + 1/4}\, e^{\frac{2\alpha C_A}{\pi} y\, \chi(\nu)} \quad (16)$$

where $\chi(\nu) = -\gamma - \mathrm{Re}\ \psi(\frac{1}{2} + i\nu)$ with γ the Euler constant and ψ the logarithmic derivative of the Γ-function. For real values of ν, $-\chi(\nu)$ is monotonically decreasing function of $|\nu|$ away from $\nu = 0$. At large values of αy the asymptotic behavior of f is determined by the saddle point at $\nu = 0$ in (16). In fact we are interested in intermediate values of y since, as we shall see later, the leading logarithmic approximation cannot be reliable for too large values of y. In order to determine f in an intermediate region of y we have evaluated the perturbation expansion for f in some detail. For $y \gtrsim 2$ the leading term in the large y expansion give a good representation of f as

$$f \sim \frac{e^{\frac{\alpha C_A}{\pi} y\, 4 ln\, 2}}{\sqrt{\frac{\alpha C_A}{2} y\, \zeta(3,1/2)}}\ , \quad (17)$$

while for $y \lesssim 2$ perturbation theory in αy converges rapidly. Thus the effective rate of growth for $y \gtrsim 3$ of $\sigma(M,s)$ is close to $\sqrt{s}$.

The growth, in s, indicated by (17) is the growth due to the bare perturbative Pomeron in QCD. What is new about our discussion is that for a limited range in s this behavior should in fact dominate the cross section. Let us in fact determine the range in y over which the bare perturbative Pomeron can be expected to correctly give f. In general we may write f as

$$f = f_0(\alpha y) + \alpha f_1(\alpha y) + \alpha^2 f_2(\alpha y) + \ldots \qquad (18)$$

where, now, f_0 contains the leading logarithms. Now the leading logarithmic approximation breaks down when αf_1 and $\alpha^2 f_2$ etc are comparable to f_0. This happens when

$$y = \frac{2}{\frac{\alpha C_A}{\nu} 4 \ln 2} \ln(C/\alpha) \qquad (19)$$

with the constant in (19) undetermined at present. When y attains the value indicated in (19) the higher Pomeron exchanges become important and presumably slow the rate of growth indicated in (17) to a $\ln^2 s$ growth. The picture should be something like that shown in Fig. 4.

The cross section given by (15) is not small and is certainly in the 1-10mb region through the energy range of the CERN collider. The growth produced by (15) as one increases s is due to multijet production. That is, given the two jet trigger we have been discussing it is very likely that additional jets will be present, at least when y is not too small. The trigger for the two jet inclusive process picks out a parton-parton cross section which, in terms of the quantities appearing in (15), is given by

$$\left(\frac{\alpha C_A}{\pi}\right)^2 \frac{\pi^3}{2M^2} f \qquad (20)$$

What we have tried to show here is that the high energy behavior of this parton-parton cross section involves j-plane singularities, as usual, but that now one can justifiably use weak coupling techniques although ultimately, when (19) is attained this weak coupling regime will correspond to a non perturbative high field strength regime; a high field strength regime now, however, only over a small part of the proton.

5 . Classical Simulations in High Energy Reactions

Recently a number of vary ambitious schemes have been proposed for studying complete events in e^+e^- and in hadronic collisions[13-19]. The practical motivation is clear in that such information is needed in order to estimate backgrounds and uncertanties in a given detector. Theoritically, these programs are quite interesting because they attempts to represent a complicated quantum mechanical process in terms of classical time evolution. In fact at first sight it is not at all clear that Monte Carlo simulations of high energy events have anything but rough empirical validity. The situation in e^+e^- collisions is simpler so let me begin there.

5.A. Classical Simulations in e^+e^- Annihilation

The models developed to describe e^+e^- annihilation events are of types, branding models and string models. Both types of models are quite successful in explaining quite detailed properties of e^+e^- events. The models are, however quite different in spirit.

-1 - Branching models can be schematically viewed as shown in Fig. 5. Immediately after the virtual photon has converted into an energetic quark-anti-quark pair, the quark may emit gluons which then emit more gluons. The probability of a quark or gluon emitting another gluon is taken to be proportional to perturbative QCD matrix elements. When the quarks and gluons go below a mass Q_0 they are converted into hadrons by some non-perturbative hadronization model, usually a cluster or string model. The evolution of individual quanta between mass Q, the mass of the virtual photon and Q_0 is thus given by the branching note that the procedure is classical in that there is no addition of amplitudes and no interference in the branching. The Q^2-dependence of events is given only by the branching part of the model and does not depend on the method of hadronization.

How can one possibly hope to represent a complicated quantum system without taking interference into account ? That this is not completely impossible comes about because of two key facts. *(i)* In large regions of phase space there is effectively only one non-negligable Feynman graph. *(ii)* In another large region of phase space, where in fact many Feynman graphs interfere, interference is complete and the quantum amplitude is exactly zero. In the classical simulation one may simply supress these regions of phase space. It is a remarkable fact that the Marchesini-Webber branching algorithm gives results which agree with perturbative QCD in a rather systematic way. To make this statement more precise, define

$$\sigma_{n_1 n_2 \ldots n_r} = \frac{1}{\sigma} \int \frac{d\sigma}{dX_1 \ldots dX_r} X_r^{n_1} \ldots x_r^{n_r} \, dx_1 \ldots dx_r \tag{21}$$

Then

$$\frac{\sigma_{n_1 \ldots n_r}}{(n\gamma}= C_{n_1 \ldots n_r} \left(1 + d_{n_1 \ldots n_r} \sqrt{\alpha} + O(\alpha)\right) \tag{22}$$

with the constants C and d correctly given by the Marchesini-Webber algorithm.

Nevertheless a number of questions remain with respect to branching. *(i)* Is it possible to include, systematically, order α and higher corrections ? *(ii)* Perhaps even more importantly, is it possible to get the right Q^2

dependences at the constant and $\sqrt{\alpha}$ levels for observables which involve angular correlations ? (Note that all angle have been integrated out in the expressions in (21) and (22)). As we shall see a little later the Marchesini-Webber model has angular correlations only partially built in. At present it is not at all clear that any classical branching model can give angular correlations in a systematic way.

- 2 - String models were originally introduced as an attemps to follow the hadronic time evolution of a high energy collision by Preparata and his collaborators[20]. They have been most completely developed and successfully utilized by the Lund collaboration. In Fig. 6 I have illustrated the physical picture where a virtual proton converts into a quark-anti-quark pair which then separates and form a thin flux tube whose ends connect to the pair. After a certain amount of stretching the flux tube breaks when a new quark-anti-quark pair is created out of the vaccum. The process continues until the final strings become physical hadrons. The string itself does not have vibrational degrees of freedom, but it does have "kink" degrees of freedom corresponding to hard gluon emission.

The Lund model has been very successful phenomenologicaly. The relationship of the model to QCD is, however, not yet clear. In particular the early stages of evolution which should be governed by perturbation theory do not appear to have any connection with the Feynman diagrams one usually associates with perturbative QCD. At the moment the Lund model is very useful, however, to become a truly interesting model theoretically it is quite important that a firm connection with the fundamental theory be made.

5.B. *Monte Carlo Simulations in Hadronic Interactions*

The grandfather of those models whose object is to generate realistic events in a hadron-hadron collision is ISAJET. Until recently such models treated hard scattering events well as far as, say, jet cross sections and the evolution of jets in the final state are concerned. There was also a phenomenological treatment of the soft particles present in such events. The progress made recently has been a proper inclusion of initial state radiation off the lines involved in the hard scattering. The various elements are schematically illustrated in Fig. 7.

Nevertheless, such simulations are, perhaps, still at the stage of being an art rather than a science. For example it is not known whether there is important interference between initial state radiation and beam fragments or between either of these with final state radiation. A realistic way of putting in multiple minijets is far from clear. The problems here are very important from a practical point of view, but at the moment they seem very hard as far as keeping systematic contact with QCD. except for those features depending only on large transverse momentum.

6 - Coherence and the string effect

One of the important successes of the Lund model as opposed to independent fragmentation models was its prediction ofthe so-called "string-effect". The experimental situation is the following : consider a three jet event in a e^+e^- collision, viewed in the plane of three jets. Project the momenta of the associated particles onto that plane. Define the gluon jet to be the jet of smallest momentum. Then, experimentally, one observes an excess of associated particles between the gluon jet and aither of the quark jets and a depletion of associated particles betrween the two quark jets.The excess or depletion is, for example, as compared to a prediction using independent fragmentation of the three jets. In the string model[18] this comes about because the two fragmenting strings, each connecting one of the quarks to the gluon are moving away from the straight line connecting the quark and anti-quark.

Recently an incisive analysis of this effect has been given within QCD[21]. This analysis points out, at the same time, the great virtues of the Marchesini-Webber model and some of its limitations. To simplify the analysis let us work at the level of the leading term in the $1/N_c$ expansion[21,22]. Referring to Fig. 8, we may write the cross section for emission of a soft gluon of momentum k as

$$\frac{d\sigma}{d\Omega_k} = \frac{\alpha C_A}{4\pi^2}\frac{dk}{k}\sigma_{3Jet}\ W \tag{23}$$

with

$$W = \frac{1-\cos\theta_{1g}}{(1-\cos\theta_1)(1-\cos\theta_g)} + 1 \longrightarrow 2 \ . \tag{24}$$

Now imagine fixing θ and varying the azimuthal angle which the gluon, k, makes with the jet plane. W achieves a maximum when k is in the plane of the jets and between the gluon and quark-1. W is at a minimum in theplane of the jets and between the two quarks. This is the string effect.

Now one may write (24) as

$$2W = \left\{\frac{1}{1-\cos\theta_1} + \frac{\cos\theta_1-\cos\theta_{1g}}{(1-\cos\theta_1)(1-\cos\theta_g)}\right\} + \left\{\frac{1}{1-\cos\theta_g} + \frac{\cos\theta_g-\cos\theta_{1g}}{(1-\cos\theta_1)(1-\cos\theta_g)}\right\} + 1 \leftrightarrow 2 \ . \tag{25}$$

Suppose we average the first term in brakects azimuthally about line 1 and the second term azimuthally about the gluon, g, and similarly for those terms where $1 \leftrightarrow 2$. Then

$$2W = \frac{1}{1-\cos\theta_1}\,\Theta\,(\theta_{1g}-\theta_1) + \frac{1}{1-\cos\theta_2}\,\Theta\,(\theta_{2g}-\theta_2) + \frac{1}{1-\cos\theta_g}\,(\Theta(\theta_{1g}-\theta_g) + \Theta\,(\theta_{2g}-\theta_g)) \quad . \tag{26}$$

(26) is not exactly the same as (25) but it is a reasonable average approximation. (26) is the expression used in the Webber Monte Carlo and gives an excellent fit to the data.

Thus the Marchesini-Webber model has built in enough of the coherence of QCD to give a good fit even to an effect where coherence is of the utmost importance. This is very nice. However, the model does not have the leading order QCD effects exactly built in for such angular correlations and that is a shame. Whether or not it is in principle possible to exactly build in such angular correlation is not known at present.

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Fig. 1

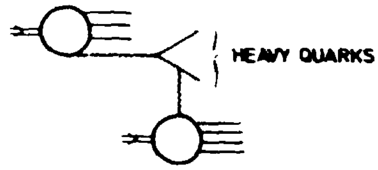

Fig. 2

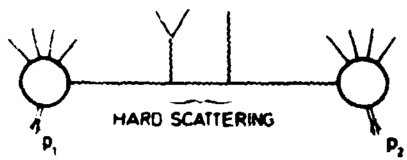

Fig. 3

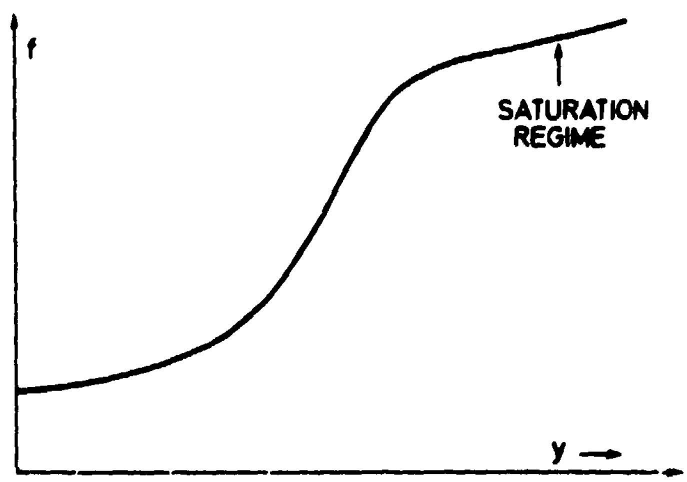

Fig. 4

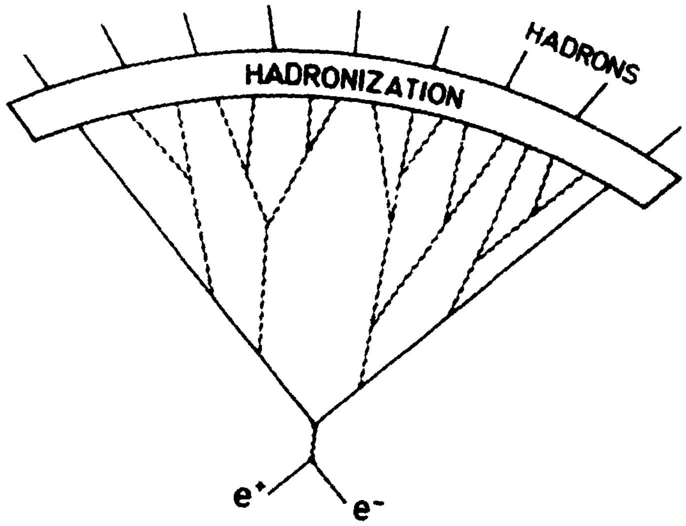

Fig. 5

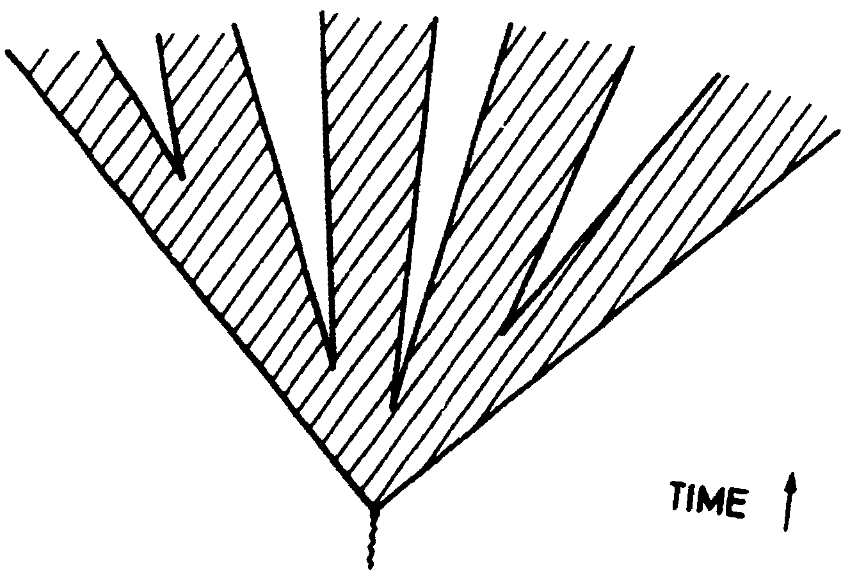

Fig. 6

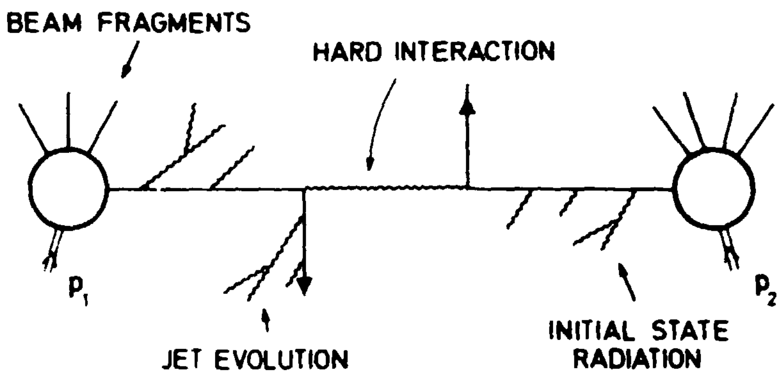
BEAM FRAGMENTS
HARD INTERACTION
P1
P2
JET EVOLUTION
INITIAL STATE
RADIATION

Fig. 7

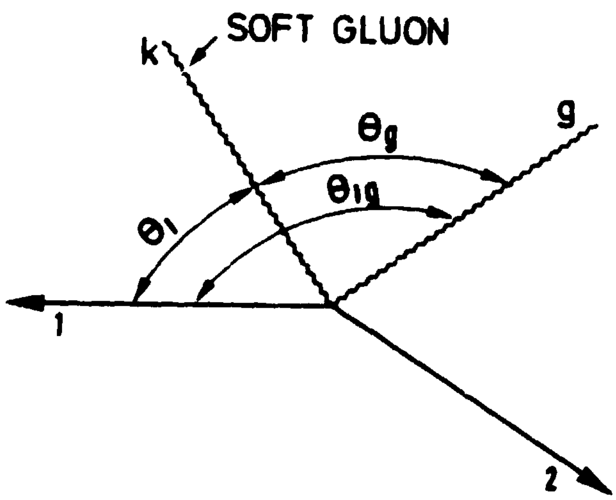
SOFT GLUON
k
g
θg
θ1g
θ1
1
2

Fig. 8