Are charm and high-$p_\perp$ jets the keys to understanding diffraction in DIS ?*

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ABSTRACT

Following an introduction which explains some basic ideas about diffraction, the semiclassical approach to diffraction in deep inelastic scattering (DIS) is outlined. Some phenomenological tests of this picture, concerning the $p_\perp$ and mass spectra of open charm and jet production are given.

1. Introduction

To introduce the main topic of this talk, I would like, first of all, to recall some basic ideas concerning diffraction that it is useful to have at the back of one's mind when thinking about diffraction in high energy processes such as DIS at small x. I will then discuss predictions for diffractive events containing either high $p_\perp$-jets or charm quarks in the final state, for two-gluon exchange models and the semiclassical approach.

1.1. The optical model

Let us begin by reminding ourselves of the phenomena of diffraction in optics. Consider a broad beam of plane polarized light incident on a small piece of Polaroid, which has its axis of polarization at an angle, θ, to the polarization direction. The component of the light with its polarization parallel to this axis will be transmitted and that perpendicular to it will be absorbed. The transmitted wave, just behind the Polaroid now contains two components, with polarizations parallel and perpendicular to that of the incident wave. A new state, degenerate in energy to the initial one, has been "diffracted into existence" by the piece of Polaroid.

Good and Walker developed an optical model for diffraction in the high energy scattering of hadrons as early as 1960 [1]. They consider the possibility that a beam of hadrons, of type A and mass M, incident on a heavy nucleus at rest, produces two new particles, B and C, of combined rest mass M^*, leaving the nucleus in its ground state. The longitudinal kick given to the nucleus needs to be fairly small to avoid breaking it up : $q_{||} \leq m_\pi/A_N^{1/3}$, where m_π is the mass of the pion and A_N is the nucleon number of the nucleus. As long as hadron A has a three momentum, P, of

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at least

$$P_{th} = \frac{(M^2 - M^{*2})A_N^{1/3}}{m_\pi} \tag{1}$$

then the process is energetically possible. According to quantum mechanics the dressed nucleon $|\tilde{A}\rangle$ is a complicated superposition of 'bare' states, $|B_i\rangle$, which have the same quantum numbers as A. These, in turn, may be re-expanded in terms of a set of 'physical' dressed states $|\tilde{D}_j\rangle$,

$$|\tilde{A}\rangle = \sum_i b_{Ai}|B_i\rangle = \sum_{ij} b_{Ai}d_{ij}|\tilde{D}_j\rangle\,. \tag{2}$$

For very high beam energies $P \gg P_{th}$, the system A spends a very short time in the heavy nucleus on the timescale of its quantum fluctuations, we may consider these fluctuations as essentially frozen during the interaction.

In the nuclear medium A will have a different expansion, $|\tilde{A}\rangle = c_{Ak}|\tilde{C}_k\rangle$. Each component $|\tilde{C}_k\rangle$ constitutes an eigenstate of the interaction and may be absorbed differently by the nucleus. For the scattering matrix just behind the target we have

$$|S\rangle = |I\rangle - |T\rangle = (1-\bar{\eta})|\tilde{A}\rangle + \sum_i (\bar{\eta} - \eta_i)c_{Ai}|C_i\rangle. \tag{3}$$

Provided the eigenstates have different absorption coefficients ($\bar{\eta} \neq \eta_i$) then new physical states $|C_i\rangle$, "degenerate in mass" with A, may be diffracted into existence by the interaction with the nuclear target.

1.2. High energy hadron-hadron diffractive scattering, rapidity gaps, light cone co-ordinates and variables for diffractive DIS

Consider very high energy diffractive scattering of two hadrons: hadron A of momentum P_A, travelling in the $+z$ direction, loses only a very small amount of its momentum and remains intact ($P'_A \gtrsim 0.99 P_A$); hadron B, travelling in the opposite direction, breaks up. It is very useful to introduce *light cone co-ordinates* for the patrons in these hadrons:

$$\begin{aligned} q_i^+ &= q_i^0 + q_i^z = m_i^\perp(\cosh y_i + \sinh y_i) = m_i^\perp e^{y_i}\,, & (4)\\ q_i^- &= q_i^0 - q_i^z = m_i^\perp(\cosh y_i - \sinh y_i) = m_i^\perp e^{-y_i}\,, & (5)\\ q_i^\perp &= (q_i^x, q_i^y)\,, & (6)\\ m_i^\perp &= \sqrt{q_i^{\perp 2} + m_i^2}\,, & (7) \end{aligned}$$

where y_i is the rapidity of parton i and is approximately equal to its pseudorapidity η_i provided $m_i^2 \ll q_i^{\perp 2}$. A typical parton in hadron B has a large q^- component and a small q^+ component and the converse is true for a parton in hadron A. Momentum conservation

$$P_A + P_B = P'_A + \sum_{i\epsilon B'} q_i \tag{8}$$

then yields

$$\xi = \frac{P_A^+ - P_A'^+}{P_A^+} = \frac{\sum_{i\epsilon B'} q_i^+ - P_B^+}{P_A^+} \approx \frac{\sum_i m_i^\perp e^{-(y_A - y_i)}}{m_A} \tag{9}$$

for the momentum fraction, ξ, lost by hadron A. We can immediately see that for small $\xi \lesssim 0.01$ we will necessarily have a *rapidity gap*, $\eta_{\mathrm{gap}} \approx y_A - y_{i,max}$, in the final state and that we can get a reasonable experimental measure of ξ, even without measuring the final state hadron, by considering hadrons close to the gap and those with high $q_i^{\perp 2}$ which contribute most in the above sum (for a more detailed discussion see Appendix A of Collins *et al* [2]).

In diffractive DIS one has to integrate over phase space, d^3P', of the scattered proton. Integrating over the angle, one has two additional scattering variables t and ξ. The scattering variables usually used to describe the events are

$$x = \frac{Q^2}{Q^2 + W^2} \quad , \quad M^2 = (P - P' + q)^2 , \tag{10}$$

$$t = (P - P')^2 \quad , \quad \xi \approx \frac{M^2 + Q^2}{W^2 + Q^2} , \tag{11}$$

$$\beta = \frac{Q^2}{2q.(P - P')} = \frac{Q^2}{M^2 + Q^2 - t} \approx \frac{Q^2}{M^2 + Q^2} \approx \frac{x}{\xi} , \tag{12}$$

where W^2 is the centre of mass energy of the $\gamma^* P$-subprocess, q and $q^2 = -Q^2$ are the four momentum and virtuality of the photon, and M is the mass of the diffractive final state. In order that the proton does not break up, the momentum transfer in the interaction should not be too large: $t \lesssim 1\mathrm{GeV}^2$. We will be primarily concerned with events for which $M^2 \approx Q^2 \ll W^2$, i.e. with the diffractive production of mass states approximately degenerate with the virtuality of the photon.

1.3. Eigenstates of diffraction in DIS

An analogous process to the optical model discussed above is the production of an e^+e^- pair from a highly energetic virtual photon in a Coulomb field, such as that surrounding a heavy nucleus. The pair is a virtual fluctuation of the virtual photon and, provided the individual virtualities of the fermions are small, will require only a small 'kick' from the field to push them onto mass shell.

This notion forms the basis of the semiclassical approach to diffraction in DIS. Viewed from the proton's rest frame, DIS at very high energies (small x) corresponds to the scattering of 'frozen' parton configurations in the virtual photon from the colour field of the proton. Those configurations in which one of the partons carries only a small fraction of the photon's longitudinal momentum and has a low transverse momentum have a small 'off-shellness' and may be diffracted into existence by receiving a small momentum transfer from the proton. The presence of this 'wee' parton in the photon implies that the fluctuation develops a large transverse size by the time it reaches the proton.

To see this, consider the lowest order fluctuation of the photon into a $q\bar{q}$ -pair, of transverse momentum $p_\perp^2$. Momentum conservation at the photon vertex gives

$$\frac{\Delta_p^2}{1-\alpha}+\frac{\Delta_l^2}{\alpha}=-\left(Q^2+\frac{p_\perp^2+m_q^2}{\alpha(1-\alpha)}\right), \tag{13}$$

where l, Δ_l^2 and m_q are the four-momentum, virtuality and mass of the anti-quark; $\alpha = l_z/q_z \approx l_0/q_0$ is longitudinal momentum fraction of the photon carried by the anti-quark; Δ_p^2 is the quark virtuality. Hence, in order to produce partons that are close to mass shell we need the asymmetric configurations in which α or $1-\alpha \ll 1$, $p_\perp^2 \ll Q^2, m_q^2 = 0$. The uncertainty principle then gives the following estimate for the lifetime of these fluctuations

$$\Delta\tau \sim \frac{1}{\Delta E} \sim \frac{1}{m_p x_{bj}}\left(\frac{Q^2}{Q^2+p_\perp^2/\alpha(1-\alpha)}\right), \tag{14}$$

which is much longer, at small x_{bj}, than the typical timescales involved in the interaction of the partons with the proton, which are of the order of the inverse of the mass, m_p, of the proton.

Miettinen and Pumplin [3] were the first to recognize that frozen parton fluctuations, in general, correspond to the eigenstates of diffraction. A very nice discussion of many of the points made here can be found in Chapter 7 of the recent text book by Forshaw and Ross [4].

2. Two-gluon exchange models

The simplest QCD model one can think of for diffraction is the exchange of two gluons in a colour singlet in the t-channel. Several recent papers present predictions for diffractive production of charm [5] and high $p_\perp$-jets [6], for the pure $q\bar{q}$ final state, which have this mechanism of diffractive exchange in common. Unfortunately a complete $\mathcal{O}(\alpha_s)$ calculation, which would also include a gluon in the final state is not yet available. In terms of the diffractive structure function, these corrections are expected to be important when the diffractive mass is large compared to Q^2. The differential cross sections for transversely and longitudinally polarized photons can be written, in the double leading logarithmic approximation, in terms of the square of the gluon density, $G(\xi)$, as follows

$$\begin{aligned}\frac{d^2\sigma_L}{d\alpha dp_\perp^2} &= \frac{2\sum_q e_q^2\alpha_{em}\alpha_s^2\pi^2[\xi G(\xi,(p_\perp^2+m_q^2)/(1-\beta))]^2 C}{3(a^2+p_\perp^2)^6}[\alpha(1-\alpha)]^2 Q^2(a^2-p_\perp^2)^2\,, \\ \frac{d^2\sigma_T}{d\alpha dp_\perp^2} &= \frac{\sum_q e_q^2\alpha_{em}\alpha_s^2\pi^2[\xi G(\xi,(p_\perp^2+m_q^2)/(1-\beta))]^2 C}{6(a^2+p_\perp^2)^6}\end{aligned} \tag{15}$$

$$\begin{aligned} &\times \quad \left[4(\alpha^2+(1-\alpha)^2)p_\perp^2 a^4 + m_q^2(a^2-p_\perp^2)^2\right], && (16)\\ a^2 &= \alpha(1-\alpha)Q^2 + m_q^2 . && (17)\end{aligned}$$

The factor C parameterizes the required extrapolation from $t \approx 0$ to the integrated cross section, $C \approx \Lambda^2$ (where Λ is a typical hadronic scale).

3. The semiclassical approach

In this section predictions for the diffractive production of heavy flavours and high-$p_\perp$ jets are presented for $q\bar{q}$ and $q\bar{q}g$ final states, in the semiclassical approach [7]. Phenomenological implications of this approach have been investigated in recent preprints [8 9].

3.1. $q\bar{q}$ final states

The longitudinal and transverse differential cross sections for the leading order $q\bar{q}$ fluctuation are as follows

$$\frac{d\sigma_L}{d\alpha dp_\perp'^2} = \frac{4\Sigma_q e_q^2 \alpha_{em}}{3(2\pi)}[\alpha(1-\alpha)]^2 \int_{x_\perp} \left| \int \frac{d^2p_\perp}{(2\pi)^2} \frac{\mathrm{tr}\tilde{W}^{\mathcal{F}}_{x_\perp}(p_\perp'-p_\perp)}{a^2+p_\perp^2}\right|^2 , \qquad (18)$$

$$\frac{d\sigma_T}{d\alpha dp_\perp'^2} = \frac{\Sigma_q e_q^2 \alpha_{em}}{3(2\pi)}(\alpha^2+(1-\alpha)^2) f^{\mathcal{F}}(a^2, p_\perp'), \qquad (19)$$

$$f^{\mathcal{F}}(a^2,p_\perp') = \int_{x_\perp} \left| \int \frac{d^2p_\perp}{(2\pi)^2} \frac{p_\perp \mathrm{tr}\tilde{W}^{\mathcal{F}}_{x_\perp}(p_\perp'-p_\perp)}{a^2+p_\perp^2}\right|^2 , \qquad (20)$$

$$a^2 = \alpha(1-\alpha)Q^2 + m_q^2 . \qquad (21)$$

The eikonal factor

$$W^{\mathcal{F}}_{x_\perp}(y_\perp) = U^\dagger(x_\perp+y_\perp)U(x_\perp) - 1 \qquad (22)$$

is built from the non-Abelian eikonal factors U and $U^\dagger$ of the quark and antiquark whose light-like paths penetrate the colour field of the proton at transverse positions $x_\perp$ and $x_\perp + y_\perp$, respectively. The superscript $\mathcal{F}$ is used because the quarks are in the fundamental representation of the gauge group. The function $f^{\mathcal{F}}$ depends on the Fourier transform of the colour trace of the eikonal factor which depends in general on the transverse momentum lost by parton. Taking its trace projects onto the colour singlet configurations relevant for diffraction. For a soft colour field, large transverse momentum transfers are exponentially suppressed. The factors in the denominator and numerator in the integrand of Eq.(20) reflect the propagator and spin of this parton, respectively.

In the region $\alpha(1-\alpha)Q^2 > \Lambda^2$, we can expand the denominator of Eq.(20) in the momentum transfer lost by the quark and, from the properties of the eikonal function,

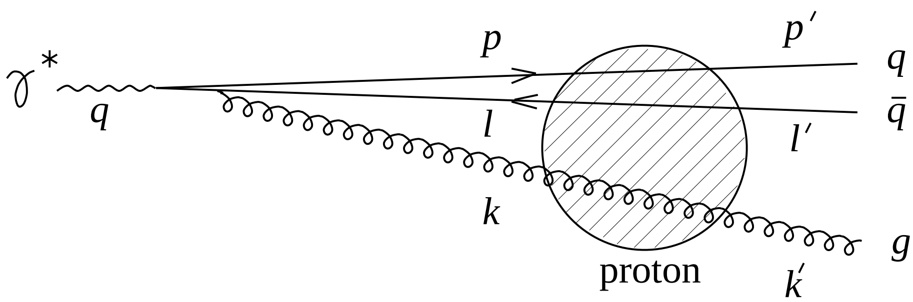

Fig. 1. Dominant partonic process for diffractive charm and high-$p_\perp$ jet production in the semiclassical approach.

show that there is no leading twist diffraction in this region. For $\alpha(1-\alpha)Q^2 < \Lambda^2$, for low $p_\perp$ quarks, i.e. the asymmetric configurations, this expansion is no longer possible, at least for massless quarks, since the denominator becomes small. As a result one gets a leading twist contribution to the transverse cross section in this region. The addition powers of $\alpha(1-\alpha)$ in Eq.(18) ensure that the contribution to the longitudinal cross section is higher twist.

In the case of a massive fermion, such as the charm quark, the expansion is always possible since the denominator is limited by the fermion mass and the cross section for diffractive charm production is suppressed by Λ^2/m_c^2. Similarly high-$p_\perp$ configurations are suppressed by $\Lambda^2/p_\perp^2$. All of this can be seen indirectly from size of the virtualities in Eq.(13).

The integration over t is implicit in Eq.(20), as reflected in the single integration over $x_\perp$. Reinstating the t-integration and expanding in $y_\perp$ (for high-$p_\perp$ configurations) we can relate the resulting $x_\perp$-integral to the gluon density and the cross section, $\sigma(y_\perp)$, for scattering a small dipole of size $y_\perp$ from the proton,

$$\int_{x_\perp} \mathrm{tr} W_{x_\perp}(y_\perp) = \frac{3\sigma(y_\perp^2)}{2} = \frac{\pi^2}{2}\alpha_s y_\perp^2 \xi G(\xi)\,. \tag{23}$$

The differential cross sections are then exactly those given in Eqs.(15,16) for the two-gluon exchange models. We see that the mechanism for the diffractive production of jets or open charm is essentially 'hard' and corresponds to the exchange of two gluons in the t-channel.

3.2. $q\bar{q}g$ final states

The $\mathcal{O}(\alpha_s)$ calculation with an additional gluon in the photon fluctuation proceeds along similar conceptual lines. Momentum conservation, for the fluctuation with a wee gluon and two high $p_\perp$-jets, now gives

$$\frac{\Delta_p^2}{1-\alpha} + \frac{\Delta_l^2}{\alpha} + \frac{\Delta_k^2}{\alpha'} = -\left(Q^2 + \frac{p_\perp^2 + m_q^2}{\alpha(1-\alpha)} + \frac{k_\perp^2}{\alpha'}\right), \tag{24}$$

for a gluon of transverse momentum $k_\perp$, virtuality Δ_k^2 and momentum fraction α'. Leading twist diffraction can now result in two cases: either the gluon is wee ($\alpha' \ll 1, k_\perp^2 \ll Q^2$) or the fermion is wee and massless (to avoid the mass suppression). The other two partons may emerge in a high transverse momentum configuration to produce jets. For the case with a wee gluon, shown in Fig.(1), we have

$$\frac{d\sigma_L}{d\alpha dp_\perp^2\, d\alpha' dk_\perp'^2} = \frac{\sum_q e_q^2 \alpha_{em}\alpha_s}{16\pi^2} \frac{\alpha' Q^2 p_\perp^2}{[\alpha(1-\alpha)]^2 N^4} f^{\mathcal{A}}(\alpha' N^2, k'_\perp)\,, \tag{25}$$

$$\frac{d\sigma_T}{d\alpha dp_\perp^2\, d\alpha' dk_\perp'^2} = \frac{\sum_q e_q^2 \alpha_{em}\alpha_s}{128\pi^2} \frac{\alpha' \left\{[\alpha^2 + (1-\alpha)^2]\,[p_\perp^4 + a^4] + 2p_\perp^2 m_q^2\right\}}{[\alpha(1-\alpha)]^4 N^4} f^{\mathcal{A}}(\alpha' N^2, k'_\perp)\,, \tag{26}$$

$$f^{\mathcal{A}}(\alpha' N^2, k'_\perp) = \int_{x_\perp} \left| \int \frac{d^2k_\perp}{(2\pi)^2} \left(\delta^{ij} + \frac{2k_\perp^i k_\perp^j}{\alpha' N^2}\right) \frac{\mathrm{tr}\tilde{W}^{\mathcal{A}}_{x_\perp}(k'_\perp - k_\perp)}{\alpha' N^2 + k_\perp^2} \right|^2 , \tag{27}$$

$$N^2 = Q^2 + \frac{p_\perp^2 + m_q^2}{\alpha(1-\alpha)}\,. \tag{28}$$

A heavy quark mass or high-$p_\perp$ ensures that the $q\bar{q}$ pair stays small in transverse space and behaves like a gluon terms of colour. The superscript $\mathcal{A}$ is used since we are effectively testing the proton's field with two "gluons". The factor in round brackets and the denominator factor in Eq.(27) reflect the spin and propagator of the wee gluon, respectively. Integration over the final state variables of the wee gluon in the leading $\ln(1/x)$ approximation gives

$$\frac{d\sigma_L}{d\alpha dp_\perp^2} = \frac{\sum_q e_q^2 \alpha_{em}\alpha_s \ln(1/x) h_{\mathcal{A}}}{2\pi^3 (a^2 + p_\perp^2)^4} [\alpha(1-\alpha)]^2 Q^2 p_\perp^2\,, \tag{29}$$

$$\frac{d\sigma_T}{d\alpha dp_\perp^2} = \frac{\sum_q e_q^2 \alpha_{em}\alpha_s \ln(1/x) h_{\mathcal{A}}}{16\pi^3 (a^2 + p_\perp^2)^4} \left[(\alpha^2 + (1-\alpha)^2)\,(p_\perp^4 + a^4) + 2p_\perp^2 m_q^2\right]\,, \tag{30}$$

$$h_{\mathcal{A}} = \int_{y_\perp}\int_{x_\perp} \frac{\left|\mathrm{tr}W^{\mathcal{A}}_{x_\perp}(y_\perp)\right|^2}{y_\perp^4}\,. \tag{31}$$

In the wee massless fermion case we have

$$\frac{d\sigma_L}{d\alpha dp_\perp^2\, d\alpha' dk_\perp'^2} = \frac{4\Sigma_q e_q^2 \alpha_{em}\alpha_s}{9\pi^2} \frac{Q^2}{[\alpha(1-\alpha)]^2 N^4} f^{\mathcal{F}}(\alpha' N^2, k'_\perp)\,, \tag{32}$$

$$\frac{d\sigma_T}{d\alpha dp_\perp^2\, d\alpha' dk_\perp'^2} = \frac{\Sigma_q e_q^2 \alpha_{em}\alpha_s}{9\pi^2 p_\perp^2 N^4} \left[N^4 - 2Q^2(N^2 + Q^2) + \frac{N^4 + Q^4}{\alpha(1-\alpha)}\right] f^{\mathcal{F}}(\alpha' N^2, k'_\perp)\,, \tag{33}$$

where the functional $f^{\mathcal{F}}$ is given in Eq.(20). Now $\alpha', k'_\perp$ refer to the final state variables of the slow quark (anti-quark) and $\alpha, p_\perp$ to the fast anti-quark (quark). Since the fast anti-quark and gluon stay very close together in transverse space, in terms of colour they behave like an anti-quark. Integration gives

$$\frac{d\sigma_L}{d\alpha dp_\perp^2} = \frac{16\Sigma_q e_q^2 \alpha_{em}\alpha_s}{27\pi^3} \frac{Q^2}{[\alpha(1-\alpha)]N^6} h_{\mathcal{F}} , \tag{34}$$

$$\frac{d\sigma_T}{d\alpha dp_\perp^2} = \frac{4\Sigma_q e_q^2 \alpha_{em}\alpha_s}{27\pi^3 N^6 p_\perp^2} \left[N^4 - 2Q^2(N^2+Q^2) + \frac{(N^4+Q^4)}{\alpha(1-\alpha)} \right] h_{\mathcal{F}} , \tag{35}$$

with $h_{\mathcal{F}}$ defined like $h_{\mathcal{A}}$.

Comparing Eqs.(29,30) with Eqs.(34,35) we see that the configurations with a wee gluon are enhanced by $\ln(1/x)$ at small x relative to those with a wee fermion. In addition there may be an additional large suppression of the slow fermion contributions due to colour factors ($h_{\mathcal{A}} \approx 16 h_{\mathcal{F}}$ [8]). As a result we claim that the configurations with a wee gluon dominate over those with a wee fermion in the small-x region relevant to diffraction and we shall ignore the latter from now on. In contrast to the $q\bar{q}$ case the diffractive production of high-$p_\perp$ jets and charm is now 'soft' and relies on the non-perturbative interaction of the wee gluon with the proton.

3.3. Boost to the Breit frame

We now consider what happens when one boosts from the Proton's rest frame to the Breit frame (in which q = (0,0,0,Q) and the proton is fast). As an example we will consider the $q\bar{q}g$ fluctuation with a wee gluon. The general relations for the boost along the z-axis, in light cone co-ordinates of Eqs.(4-7), are

$$a'^{+} = \gamma(1-\beta)a^{+} \quad ; \quad a'^{-} = \gamma(1+\beta)a^{-} . \tag{36}$$

For this particular boost

$$\gamma(1-\beta) = \frac{m_p x}{|Q|} = \frac{1}{\gamma(1+\beta)} . \tag{37}$$

The typical four momentum of the slow gluon before interaction is

$$k = (\frac{\Lambda^2}{m_p x}, -\Lambda x, \vec{k}^\perp) , \tag{38}$$

where, crucially, its '-' component is *negative*. The boost merely rescales the components to give

$$k' = (\frac{\Lambda^2}{|Q|}, -\frac{\Lambda |Q|}{m_p}, \vec{k}^\perp) . \tag{39}$$

This negative energy outgoing gluon may then be reinterpreted as a positive energy *incoming* parton in the proton [10]. In contrast, the final state wee gluon is forced to

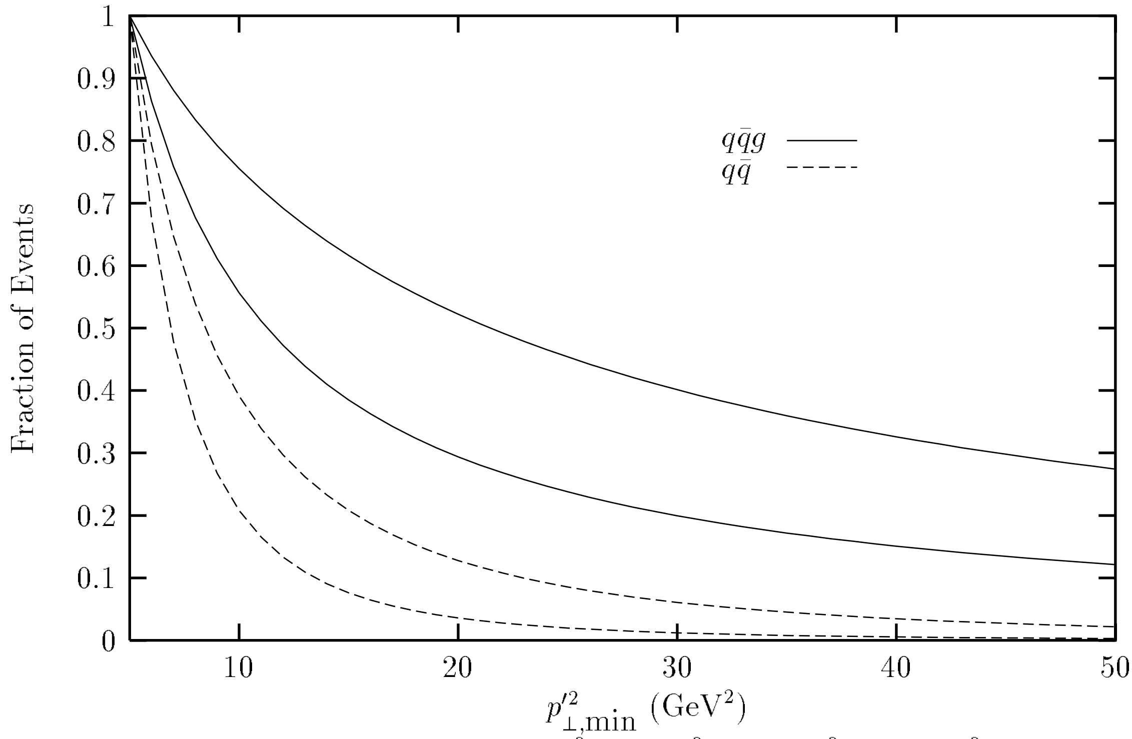

Fig. 2. The fraction of diffractive events with $p'^2_\perp$ above $p'^2_{\perp,\text{min}}$ for Q^2 of 10 GeV^2 and 100 GeV^2 (lower and upper curve in each pair).

have a positive '-' component because it is on shell and ends up in the final state. So, in the Breit frame, the process of Fig.(1) may be interpreted as boson-gluon fusion with an additional gluon in the final state.

4. Features of diffractive final states

We propose two phenomenological tests which reveal the nature of the underlying process and reflect the mechanism of colour neutralization which produces the different diffractive final states. Firstly, one may ask how many high-$p_\perp$ events survive above a given minimum $p^2_\perp$. To examine this we plot the quantity

$$\sigma(p^2_{\perp,\text{min}}) \;=\; \int_{p^2_{\perp,\text{min}}}^{\infty} dp^2_\perp \int_0^1 d\alpha \frac{d^2\sigma}{dp^2_\perp\, d\alpha} \tag{40}$$

in Fig.(2) for massless quarks. Each curve is normalized to its value at $p^2_{\perp,\text{min}} =$ 5 GeV^2. The figure clearly reveals a much harder spectrum for the $q\bar{q}g$ final states than for $q\bar{q}$. For diffractive charm production we may set $p^2_{\perp,\text{min}} = 0$ and get similar results [8].

We may also examine the expected mass spectra for the different final states. Let M_j be the invariant mass of the two-jet system in diffractive events. The measure-

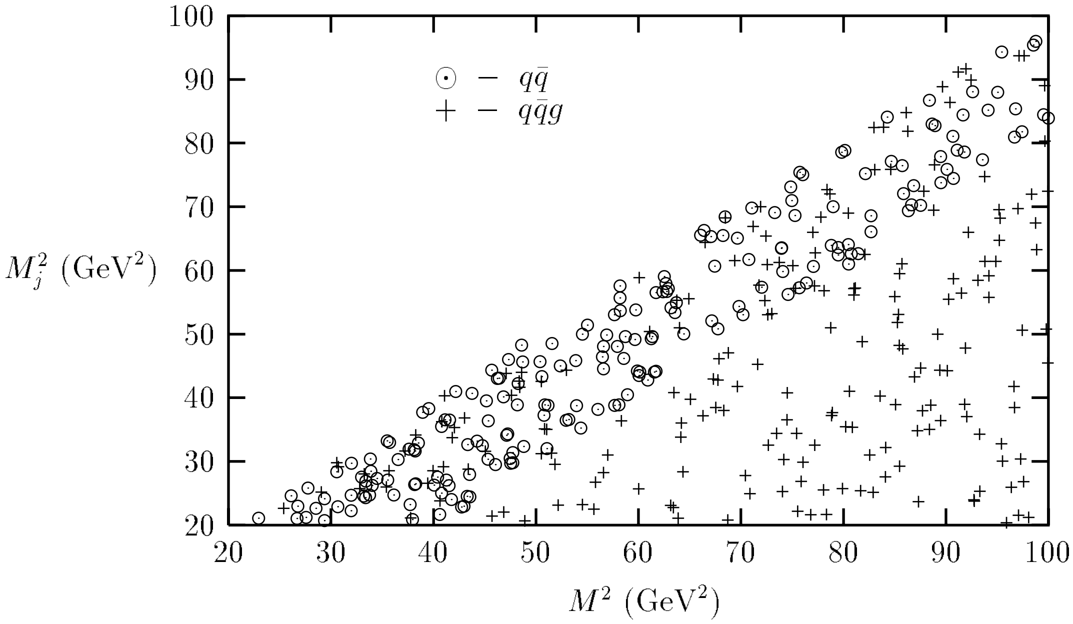

Fig. 3. Distributions in M^2 and M_j^2 of diffractive events originating from $q\bar{q}$ and $q\bar{q}g$ final states.

ment of this observable provides, in principle, a clean distinction between $q\bar{q}$ final states, where $M_j^2 = M^2$, and $q\bar{q}g$ final states, where $M_j^2 < M^2$. In practice, however, this requires the contribution of the wee gluon to the diffractive mass, which is responsible for the difference between M^2 and M_j^2, to be sufficiently large. To quantify the expectation within the semiclassical approach we consider the transverse photon contribution to the differential diffractive cross section $d\sigma/dM^2 dM_j^2$.

To illustrate the experimental implications for high-$p_\perp$ jets we have shown in Fig. (3) the positions in M^2 and M_j^2 of two sets of 200 events, scattered randomly according to the two above distributions. The δ-function of the $q\bar{q}$ case has been replaced by a uniform distribution of M_j^2 in a band $M^2 > M_j^2 > M^2 - 20$ GeV^2. This allows for hadronization effects and, more importantly, for a large experimental uncertainty of the mass of the jet system. The scatter plot clearly exhibits the distinctive features of the underlying partonic processes, even for this limited number of events. With sufficient statistics a determination of the relative weight of soft colour-singlet exchange, relevant for $q\bar{q}g$ final states, and hard colour-singlet exchange, relevant for $q\bar{q}$ final states, should be feasible.

Clearly this is also possible for diffractive charm. In Fig.(4) normalized mass spectra comparing the mass of the charm pair, $M_{c\bar{c}}$, with the total diffractive mass, M, are shown. Here the pure $c\bar{c}$ final state $M_{c\bar{c}} = M$, is represented by a block of width 10 GeV^2. The other curves in the figure represent the $c\bar{c}g$ final state of Fig.(1) Details about how the latter curves were arrived at and the meaning of C_s are given in the preprint [8].

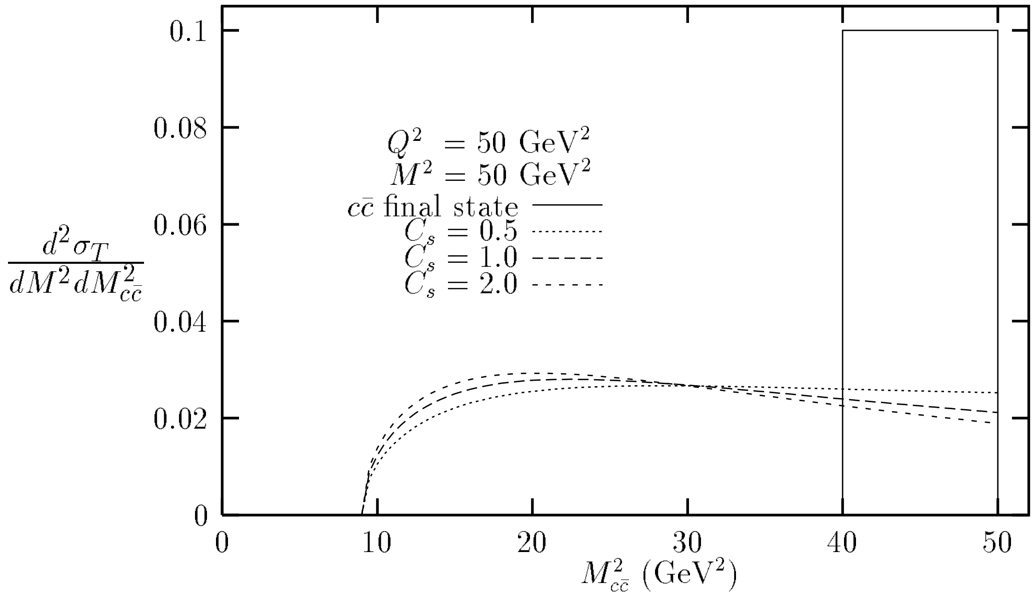

Fig. 4. Normalized mass spectra for the $c\bar{c}g$ and $c\bar{c}$ final states calculated in the semiclassical approach

First results on open charm production in diffraction for the 1994 running period were presented by the H1 collab., at Warsaw [11]. Both H1 and ZEUS have recently presented results on open charm and high-$p_\perp$ jets in diffractive DIS [12]. Given the increase in statistics of HERA for the 1996 running period, it is hoped that these phenomenological tests may be performed very soon.

An additional means of distinguishing the underlying diffractive mechanism is the energy dependence of the two processes. In the two gluon model a steeply rising gluon density, taken from a fit to F_2 for example, produces a rise in diffractive charm events that is twice as steep with energy. In contrast the semiclassical approach the energy dependence is flat, at least at this order, corresponding to a classical bremstrahl spectrum of gluons.

5. References

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