Charm as a Key to Diffractive Processes

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Abstract
The diffractive production of open charm in deep-inelastic scattering is studied in the semiclassical approach which has been proposed recently. In this approach, the leading order process contains a charm quark pair and an additional gluon in the diffractive final state. The $p_{\perp}$-spectrum and the diffractive mass distribution are evaluated and compared with predictions based on perturbative two-gluon exchange calculations for charm quark pair production. It is shown that the $p_{\perp}$-spectrum provides a clear test of the underlying partonic process whereas the diffractive mass distribution reflects the non-perturbative mechanism of colour neutralization.
The precise measurements of the diffractive structure function at small $x$ in the experiments at HERA \cite{1} provide a challenge to the theoretical description of deep-inelastic scattering based on QCD. The theoretical interest, as well as the difficulty, lies in the interplay between ‘soft’ and ‘hard’ processes which is the characteristic feature of diffractive deep-inelastic scattering. To disentangle both aspects of ‘hard diffraction’ the study of less inclusive processes will be decisive. First results on the diffractive production of open charm have already been reported by the H1 collaboration \cite{2}.

Diffractive charm production is theoretically interesting since one might expect the additional hard scale, provided by the charm mass, to ensure the applicability of perturbation theory. Indeed, several authors have considered perturbative two-gluon exchange as a mechanism for the excitation of open charm in diffraction \cite{3,4}.

The present paper investigates charm production in the framework of the semiclassical approach to diffraction, which treats the proton as a classical colour field \cite{5,6}. Working in the proton rest frame, the high energy scattering of a partonic fluctuation of the virtual photon is calculated. The leading process is the production of a $c\bar{c}g$-final state, with the relatively soft gluon being responsible for the large hadronic cross section of the process. In our analysis we concentrate on the comparison with the perturbative two-gluon model, with a $c\bar{c}$-final state, where the colour singlet exchange is kept hard by the large charm mass\cite{1}. This quantitative comparison assumes that higher order corrections to the two-gluon exchange are small for the considered observables.

As we shall see, the $p_t$-distribution of the produced charm jets is a clear probe of the underlying partonic process. The mass spectrum, on the other hand, reflects the additional soft partons necessarily present in the final state. These observables of diffractive open charm production can be used to understand the importance of soft and hard processes in inclusive diffraction.

**Cross sections for diffractive charm production**

Recall first the qualitative picture of diffraction in the semiclassical approach \cite{5,6}: for massless quarks the leading order process is the production of a colour neutral $q\bar{q}$-pair off the proton colour field. In this process the leading twist contribution comes entirely from the aligned jet region. This can be understood intuitively, since a highly virtual photon can split into two nearly on-shell quarks only if they share the photon’s longitudinal momentum in a highly asymmetric way.

In the case of heavy quarks, the aligned jet configuration is suppressed by the quark mass. The reason is the large off-shellness that is necessarily present if a virtual photon splits into two massive particles. As a result, the production of $c\bar{c}g$-final states, with the gluon being relatively soft, becomes the leading process.

One of the two contributing diagrams is displayed in Fig. \cite{1}. The corresponding cross sections, calculated in \cite{6} for massless quarks, generalize straightforwardly to the massive

\footnote{For a general discussion and references to other approaches see \cite{7}.}
case. The results can be given in the following form,

\[
\frac{d\sigma_L}{d\alpha dp_\perp^2 d\alpha' dk'_\perp^2} = \frac{e_c^2 \alpha' \alpha_s}{16\pi^2} \frac{Q^2 p_\perp^2}{\alpha(1-\alpha)^2 N^4} f(\alpha' N^2, k'_\perp),
\]

\[
\frac{d\sigma_T}{d\alpha dp_\perp^2 d\alpha' dk'_\perp^2} = \frac{e_c^2 \alpha' \alpha_s}{128\pi^2} \alpha' \left\{ \left[ \frac{2}{\alpha^2 + (1-\alpha)^2} \right] \left[ \frac{p_\perp^4 + a^4}{\alpha N^4} + 2p_\perp^2 m_c^2 \right] \right\} f(\alpha' N^2, k'_\perp),
\]

with

\[
f(\alpha' N^2, k'_\perp) = \int_{x_\perp} \int d^2 k_\perp \left( \frac{2}{\alpha' N^2} \right)^2 \left( \delta^{ij} + \frac{2k_i^j k_j^i}{\alpha' N^2} \right) \text{tr} \tilde{W}_{x_\perp}(k'_\perp - k_\perp)^2,
\]

and

\[
N^2 = Q^2 + \frac{p_\perp^2 + m_c^2}{\alpha(1-\alpha)} , \quad a^2 = \alpha(1-\alpha) Q^2 + m_c^2.
\]

Here \(\alpha = p_0/q_0\) and \(\alpha' = k'_0/q_0\) are the longitudinal momentum fractions carried by quark and gluon, \(Q^2 = -q^2\) is the photon virtuality. All transverse momenta are defined relative to the \(\gamma^*_\text{direction.}\)

![Figure 1: Open charm production off the proton colour field.](image)

The \(\alpha'\)- and \(k'_\perp\)-distributions are completely non-perturbative. They are governed by the Fourier transform of the colour field dependent function

\[
W_{x_\perp}(y_\perp) = \mathcal{A}(U^\dagger(x_\perp + y_\perp)U(x_\perp)) - 1 ,
\]

which is built from non-Abelian eikonal factors in the adjoint representation. These matrices correspond to light-like paths penetrating the colour field at transverse positions \(x_\perp\) and \(x_\perp + y_\perp\). To leading order in \(1/m_c\) the two quarks have a small transverse separation and act like a colour octet. The two eikonal factors arise from the propagation of this effective octet and the gluon.

The integration over the transverse momentum of the produced quarks is dominated by the hard scales \(Q^2\) and \(m_c^2\), implying \(l_\perp \approx -p_\perp\) for most of the events. Only the integration region where \(\alpha'\) and \(k'_\perp\) are small, in which the gluon tests large transverse distances in the proton, gives rise to a non-power-suppressed contribution to the total cross section. In this region the factor \(\alpha' N^2 + k_\perp^2\) in the denominator of Eq. (3), which comes from the propagator of the gluon in the wavefunction of the virtual photon, becomes small, \(\sim \Lambda^2\), where \(\Lambda\) is a typical hadronic scale. In the case where one of the charm quarks is soft the equivalent factor is prevented from becoming small by the quark mass,
one may then expand in the transverse momentum lost by the soft quark as it travels through the proton and the final result is power-suppressed, i.e. $O(\Lambda^2/m_c^2)$.

The cross sections of Eqs. (1) and (2) can be expressed as the convolution of the cross section of an ordinary partonic process, namely boson-gluon fusion, with a diffractive gluon density, which is given in terms of $\text{tr} W_{x_{1\perp}}^A$. Some of the predictions discussed in the following analysis are not sensitive to the particular form of the diffractive parton density; in this case we refer to our approach as ‘diffractive parton model’. Other predictions are sensitive to the eikonal approximation, which is used to treat the soft interaction with the proton, in this case we use ‘eikonal model’.

The following numerical analysis will compare our approach with two-gluon exchange calculations. The leading order cross sections for these, corresponding to the production of a $c\bar{c}$-pair, are available in the literature [3, 4]. Using our kinematical variables the results of [4] read

$$\frac{d\sigma_L}{d\alpha dp_{\perp}^2} = \frac{2e^2_\alpha e_m \alpha_s^2 \pi^2 [\xi G(\xi)]^2 C}{3(a^2 + p_{\perp}^2)^6} \left[ \alpha(1 - \alpha)^2 Q^2(a^2 - p_{\perp}^2)^2 \right],$$  

$$\frac{d\sigma_T}{d\alpha dp_{\perp}^2} = \frac{e^2_\alpha e_m \alpha_s^2 \pi^2 [\xi G(\xi)]^2 C}{6(a^2 + p_{\perp}^2)^6} \left[ 4(\alpha^2 + (1 - \alpha)^2)p_{\perp}^2 a^4 + m_c^2(a^2 - p_{\perp}^2)^2 \right],$$

where $\xi$ (or $x_{IP}$) is the longitudinal momentum fraction lost by the elastically scattered proton. We have simplified the formulae by neglecting the scale dependence of $\alpha_s$ and of the gluon density and by restricting ourselves to small transverse momenta of the exchanged gluons ($\ll a^2 + p_{\perp}^2$). At small $t$ the cross section is proportional to the square of the gluon density $G(\xi)$. The factor $C$ parameterizes the required extrapolation from $t \approx 0$ to the integrated cross section,

$$C = \left( \int \frac{d\sigma}{dt} dt \right) / \left( \left. \frac{d\sigma}{dt} \right|_{t=0} \right) \sim \Lambda^2.$$  

**Transverse momentum distribution and total cross sections**

The cross section formulae of the last section can be used to predict the $p_{\perp}$-distributions of the produced charm jets in each model. However, in order to make a comparison, the integrations over the kinematical parameters of the final state gluon, $\alpha'$ and $k_{\perp}'$, in Eqs. (1) and (2) need to be performed.

Consider the longitudinal case, Eq. (1). Introducing the definition of $W_{x_{1\perp}}^A$ in terms of a Fourier transformation explicitly, the $k_{\perp}'$-integration becomes trivial. The resulting $\delta$-function ensures that both factors $W_{x_{1\perp}}^A$ are evaluated at the same position $y_{\perp}$. Furthermore, it is convenient to replace the $\alpha'$-integration by a $u$-integration, with $u^2 = \alpha' y_{\perp}^2 N^2$. The resulting formula is

$$\frac{d\sigma_L}{d\alpha dp_{\perp}^2} = \frac{e^2_\alpha e_m \alpha_s}{2\pi} \frac{Q^2 p_{\perp}^2}{[\alpha(1 - \alpha)]^2 N^8} \int_{y_{\perp}} \int du u^3 g(u) \int_{x_{1\perp}} \left| \frac{\text{tr} W_{x_{1\perp}}^A(y_{\perp})}{y_{\perp}^4} \right|^2,$$  

(9)
where the dimensionless function

\[ g(u) = \left| \int \frac{d^2 b_\perp}{(2\pi)^2} \frac{(\delta^{ij} + 2b_i^\perp b_j^\perp) e^{iu_\perp b_\perp}}{1 + b_\perp^2} \right|^2, \quad u^2 \equiv u_\perp^2, \]  

(10)

has been introduced. Observing that

\[ g(u) \approx \frac{2}{\pi^2 u^4} \quad \text{at} \quad u \ll 1, \]  

(11)

the \( u \)-integration is found to be logarithmically divergent. Since the eikonal approximation will certainly fail for a gluon energy of \( O(\Lambda) \), it is natural to introduce the cutoff

\[ \alpha'_\text{min} = \Lambda/q_0. \]

Assuming a soft external field the \( y_\perp \)-integration in Eq. (9) is dominated by the region \( y_\perp^2 \sim 1/\Lambda^2 \), resulting in \( u^2_\text{min} \sim x = x_{Bj} \). Introducing the dimensionless constant

\[ h_A = \int y_\perp \int_{x_\perp} \left| \text{tr} W_{x_\perp}(y_\perp) \right|^2 y_\perp^4, \]  

(12)

the leading-ln(1/x)-contribution can now be extracted from Eq. (9). The transverse cross section follows completely analogously, and one obtains

\[ \frac{d\sigma_L}{d\alpha dp_\perp^2} = e_c^2 \alpha_\text{em} \alpha_s \ln(1/x) h_A \frac{\ln(1/x) h_A}{2\pi^3(a^2 + p_\perp^2)^4} \left[ \alpha (1 - \alpha) \right]^2 Q^2 p_\perp^2, \]  

(13)

\[ \frac{d\sigma_T}{d\alpha dp_\perp^2} = e_c^2 \alpha_\text{em} \alpha_s \ln(1/x) h_A \frac{\ln(1/x) h_A}{16\pi^3(a^2 + p_\perp^2)^4} \left[ \alpha^2 + (1-\alpha)^2 \right] \left[ p_\perp^4 + a^4 \right] + 2 p_\perp^2 m_c^2. \]  

(14)

These equations can be directly compared to the corresponding two-gluon results, Eqs. (6) and (7). In both models the \( \alpha \)-integration can be performed analytically.

The qualitative differences between the two models are particularly pronounced in the integrated cross section with a lower cut on transverse momentum of the charm quarks. Since most events are expected to have small \( y = Q^2/(sx) \) (cf. [1]), it is sufficient to consider \( \sigma = \sigma_T + \sigma_L \),

\[ \sigma(p_\perp^2, \min) = \int_{p_\perp^2, \min}^\infty dp_\perp^2 \frac{d\sigma}{dp_\perp^2}. \]  

(15)

Fig. 2 shows the dependence of the corresponding event fraction on the lower cut. We use \( m_c = 1.5 \text{ GeV} \) in our numerical analysis here and below. The diffractive parton model predicts a much stronger high-\( p_\perp \) tail of the distribution. For example, at \( Q^2 = 10 \text{ GeV}^2 \), more than 40\% of the events have \( p_\perp^2 > 5 \text{ GeV}^2 \), compared to only 7\% in the two-gluon model.

The physical reason for this difference between the two models is easily understood. As discussed above, the hard process in our model corresponds to boson-gluon fusion in the Breit frame. Therefore, the transverse momentum is logarithmically distributed between \( m_c^2 \) and \( Q^2 \). By contrast, the two-gluon process couples the small \( c\bar{c} \)-dipole directly to the hadron, resulting in a power suppression of small size configurations. The cross section comes entirely from the softest possible region, defined by the scale \( m_c^2 \). No logarithmic tail of higher transverse momenta appears.
Note that our analysis does not include $\alpha_s$-corrections for the two-gluon exchange result. Such corrections have been estimated in [4] and found to be important for diffractive masses much larger than $Q^2$, due to final state gluon radiation. This can affect the above $p_{\perp}$-distribution of the two-gluon exchange calculation.

For both models the complete $\alpha$- and $p_{\perp}$-integrated cross sections can be calculated in terms of elementary functions. Fig. 3 shows the $Q^2$-dependence of the resulting ratio $R_D^C = \sigma_L/\sigma_T$.

For our mainly qualitative discussion it is sufficient to display the $Q^2$- and $m_c^2$-dependence of the total cross sections in the limit $m_c^2 \ll Q^2$. In the diffractive parton model one finds

$$\sigma_L \sim \frac{1}{Q^2}, \quad \sigma_T \sim \frac{1}{Q^2} \left( \ln(Q/m_c) - \frac{1}{4} \right),$$

while the two-gluon model results read

$$\sigma_L \sim \frac{\Lambda^2}{Q^4} \left( \ln(Q/m_c) - \frac{3}{4} \right), \quad \sigma_T \sim \frac{\Lambda^2}{Q^2 m_c^2}.$$

The main qualitative difference is the suppression of the transverse cross section by $\Lambda^2/m_c^2$ in the two-gluon model. Furthermore, for very large $Q^2$, $R_D^D \sim m_c^2/Q^2$ in the two-gluon model whereas in the diffractive parton model $R_D^C \sim 1/\ln(Q/m_c)$.

**Charm anti-charm mass spectrum**

The different particle content of the diffractive final states, predicted by the leading order contributions of the two models, should be clearly visible in the resulting mass spectrum.
spectrum. In the two-gluon model the reconstructed diffractive mass should be peaked around the mass of the charm anti-charm pair which is given by

$$M^2 = M_{cc}^2 = \frac{p_{\perp}^2 + m_c^2}{\alpha(1-\alpha)}. \quad (18)$$

In contrast, for the leading order graphs in the eikonal model, which also contain a gluon in the final state, the corresponding diffractive mass is given by

$$M^2 = M_{cc}^2 + \frac{k_{\perp}^2}{\alpha'}. \quad (19)$$

To quantify the above distinction between the two models, the $M_{cc}$-distribution shall be calculated in the eikonal model for fixed $M^2$. By changing variables from $\alpha'$ and $p_{\perp}^2$ to $M^2$ and $M_{cc}^2$ in Eq. (2) it is possible to derive the differential cross section

$$\frac{d^2\sigma_T}{dM^2 dM_{cc}^2} = \frac{e_r^2 \alpha_{em} \alpha_s}{128\pi^3} \frac{s(b)}{b(Q^2 + M_{cc}^2)^5} \int_{\alpha_{\text{min}}}^{1-\alpha_{\text{min}}} \frac{d\alpha}{[\alpha(1-\alpha)]^3} \times$$

$$\left[ \left( \alpha^2 + (1-\alpha)^2 \right) \left\{ (\alpha(1-\alpha)M_{cc}^2 - m_c^2)^2 + a^4 \right\} + 2m_c^2(\alpha(1-\alpha)M_{cc}^2 - m_c^2) \right]$$

$$= \frac{e_r^2 \alpha_{em} \alpha_s}{64\pi^3} \frac{M_{cc}^4 s(b)}{b(Q^2 + M_{cc}^2)^5} \times$$

$$\left[ 2 \left\{ 1 + \frac{Q^4}{M_{cc}^4} + \delta \left( 1 - \frac{\delta}{2} \right) \right\} \text{Arctanh} \left( \sqrt{1-\delta} \right) - \left\{ \frac{(M_{cc}^2 - Q^2)^2}{M_{cc}^4} + \delta \right\} \sqrt{1-\delta} \right]$$

where $\alpha_{\text{min}} = (1 - \sqrt{1-\delta})/2$ with $\delta = 4m_c^2/M_{cc}^2$ and $b = (M^2 - M_{cc}^2)/(Q^2 + M_{cc}^2) = k_{\perp}^2/\alpha'N^2$. The dimensionless function $s(b)$ contains all the non-perturbative information.
and is closely related to $f$ of Eq. (2),

$$s(b) = \int d^2 k'_\perp (k'^2_\perp)^2 \int_{x_\perp} \left| \int \frac{d^2 k_\perp}{(2\pi)^2} \left( \delta^{ij} + \frac{2k^i_\perp k^j_\perp}{k^2_\perp} b \right) \frac{\text{tr} \tilde{W}_{x_\perp}^A (k'_\perp - k_\perp)}{k'^2_\perp + b k^2_\perp} \right|^2. \hspace{1cm} (21)$$

To arrive at a quantitative prediction the function $s(b)$ has to be analyzed in more detail. From Eq. (2) and the assumption of a smooth proton colour field it can be concluded that $\text{tr} \tilde{W}_{x_\perp}^A (y_\perp)$ is a smooth localized function of $y_\perp$, which vanishes together with its first derivative at $y_\perp = 0$. This results in the limiting behaviour $s(b) \sim b$ as $b \to 0$ and $s(b) \sim \text{const.}$ as $b \to \infty$. Since no large ratios are involved in the definition of $s$, it is natural to try the ansatz

$$s(b) = \frac{b}{(C_s + b)}, \hspace{1cm} (22)$$

where $C_s$ is a constant of $\mathcal{O}(1)$. Fig. 4 shows a typical spectrum in $M^2_{c\bar{c}}$ for the eikonal model and illustrates that the exact value of $C_s$ is not important. The two-gluon models are represented by a strip at $M^2_{c\bar{c}} = M^2$. Although the width of the strip is expected to be of the order of the hadronic scale we have shown a larger width in the figure to illustrate the experimental uncertainty of the diffractive mass measurement. The normalization is such that the area under each curve is unity.

A comparison of the diffractive mass, $M^2$, and the charm anti-charm pair mass, $M^2_{c\bar{c}}$, for diffractive charm events at HERA should, in principle, determine which of the mechanisms is responsible for the bulk of the events.
Ratio of charm in diffraction

Finally, we estimate the fraction of charmed events in the total diffractive cross section. Since this will involve strong assumptions about the soft part of the eikonal model, the results of this section are less reliable than the previous part of the paper.

The leading order diffractive cross section, determined by the production of a pair of light quarks in an aligned-jet type configuration, reads \[ \frac{d\sigma_T}{d\alpha'dk_{\perp}^2} = \frac{\alpha_{em}}{3\pi} \left( \sum e_q^2 \right) \int_{x_{\perp}} \left| \int \frac{d^2 k_{\perp}}{(2\pi)^2} \frac{k_{\perp} \text{tr} W_{\perp}^T (k'_{\perp} - k_{\perp})}{\alpha' Q^2 + k_{\perp}^2} \right|^2, \] where \( k'_{\perp} \) and \( \alpha' \) are transverse momentum and longitudinal momentum fraction of the outgoing soft quark and the sum is over the light quarks \( u, d, s \). The function \( W_{\perp}^T \) is defined as in Eq. (5), but with the colour matrices in the fundamental representation.

The corresponding total cross section, calculated as above, is found to be

\[ \sigma_T^{qq} = \frac{4\alpha_{em}}{9\pi^2 Q^2} \left( \sum e_q^2 \right) h_F, \] where the constant \( h_F \) is defined analogously to Eq. (12).

This has to be compared to the charm cross section, obtained by integrating Eq. (14). To leading \( \ln(1/x) \ln(Q^2) \)-accuracy the result reads

\[ \sigma_T^{c\bar{c}g} = \frac{e_c^2 \alpha_{em} \alpha_s \ln(1/x)}{6\pi^3 Q^2} \ln(Q/m_c) h_A. \]

The appearance of the above large logarithms suggests that an analogous \( qqg \)-cross section for light quarks will form an important correction of Eq. (24). However, this cross section is plagued by an infrared divergence in the region where all three particles have small \( p_{\perp} \). The leading-log contribution calculated in [6] can also be obtained by replacing the heavy quark mass in Eq. (25) with an infrared cutoff equal to \( \Lambda_{QCD} \),

\[ \sigma_T^{qqg} = \frac{\alpha_s \ln(1/x)}{6\pi^3 Q^2} \left( \sum e_q^2 \right) \ln(Q/\Lambda_{QCD}) h_A. \]

Neglecting the small longitudinal contribution (cf. Fig. [3]) the ratio \( r \) of charm in diffraction can now be given,

\[ r = \frac{\sigma_T^{c\bar{c}g}}{\sigma_T^{qq} + \sigma_T^{qqg} + \sigma_T^{c\bar{c}g}}. \]

A numerical evaluation within the present model requires the ratio \( h_A/h_F \). To estimate this ratio it is convenient to introduce the dimensionless functions

\[ j_R(y_2^2 \Lambda^2) = \Lambda^2 \int_{x_{\perp}} \left| \text{tr} W_{\perp}^R (y_{\perp}) \right|^2, \]
where $R \in \{A, F\}$ labels the representation. For smooth colour fields of finite transverse extension Eq. (29) implies $j_R(z) \approx a_R^2 z^2$ for $z \ll 1$ and $j_R(z) \approx b_R^2$ for $z \gg 1$, with two unknown constants $a_R$ and $b_R$.

In the region of small $z$ the matrices $U$ and $U^\dagger$ in the definition of $W$ can be expanded in powers of the gauge field $G_{\mu}$. The first term contributing to $\text{tr} W$ is of order $z^2 G^2$, and standard formulae of representation theory [9] give $a_A/a_F = 2N$ for the group $SU(N)$. In the region of large $z$, corresponding to large $y$ in Eq. (5), it is natural to assume that the matrices $U$ and $U^\dagger$ are not correlated. Note that in Eq. (4) an averaging procedure over all colour field configurations of the proton is understood [6]. In the case of strong fields the simplest assumption is that $U$ and $U^\dagger$ are uniformly distributed over $SU(N)$, in which case the first term of Eq. (5) vanishes under the above average. The second term is simply proportional to the dimension of the representation, so that $b_A/b_F = (N^2 - 1)/N$.

Although the constants $h_R$ are given by the simple formula

$$ h_R = \pi \int \frac{dz}{z^2} j_R(z), \quad (29) $$

knowing the limiting behaviour of $j_A(z)$ and $j_F(z)$ does not yet imply a knowledge of $h_A/h_F$. The problem is that the form of the interpolating function influences, in general, the integral in Eq. (29). Assuming, for simplicity, that the functional form does not depend on the representation, i.e.

$$ j_A(z) = \left( b_A/b_F \right)^2 j_F(z) \left( a_F b_A \right)/(a_A b_F), \quad (30) $$

the integrals for $R = A$ and $R = F$ are found to be related by a multiplicative factor, $h_A/h_F = (a_A b_A)/(a_F b_F) = 2(N^2 - 1)$.

Specifying $N = 3$ and using $x = 10^{-3}$, $Q^2 = 36 \text{ GeV}^2$, $\alpha_s(Q^2) = 0.22$, and $\Lambda_{QCD} = 200 \text{ MeV}$, the formally leading cross section in the massless quark case is found to be small, $\sigma_{q\bar{q}}^{q\bar{q}}/\sigma_{qq}^{qq} \approx 0.1$. Therefore the diffractive production of massless quarks is dominated by the $q\bar{q}g$ component of the $\gamma^* \text{ wavefunction}$. It is then natural to expect the ratio of charm in diffraction to be similar to inclusive deep-inelastic scattering (cf. [10]). Indeed, explicit evaluation of Eq. (27), with the parameters given above, results in a charm fraction of $r \approx 0.2$.

Conclusions

The $p_\perp$-spectrum and the diffractive mass distribution for the production of open charm have been evaluated in the semiclassical approach. The resulting cross sections can be expressed as the convolution of a cross section of an ordinary partonic process, namely boson-gluon fusion, with a diffractive parton density. In this way the physical content of the results becomes most transparent.

The $p_\perp$-spectrum provides a clear test of the underlying hard partonic process. It is found to have sizeable contributions from the region $p_\perp^2 \sim Q^2$, which are absent at leading order in the two-gluon model.
The diffractive mass distribution, on the other hand, reflects the soft interaction of the proton, which is treated in the eikonal approximation. The presence of the final state gluon yields a production cross section not suppressed by $\Lambda^2/m_c^2$ and a large ratio of charm in diffraction, comparable with the corresponding ratio in inclusive deep-inelastic scattering. It also allows a large difference between the diffractive mass $M^2$ and the invariant mass of the charm jets $M_{c\bar{c}}^2$. In contrast, both masses are equal, up to hadronization effects, for pure $c\bar{c}$-final states.

In summary, several observables have been identified which can discriminate between soft and hard mechanisms for colour singlet exchange in diffractive charm production. Diffraction with high-$p_\perp$ jets is kinematically very similar, since the hard scale $m_c$ is merely replaced by $p_\perp$. Therefore, we expect that the study of diffractive charm production will help to clarify the relative importance of soft and hard contributions to diffraction in general.

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