Direct and Resolved Pomeron in Rapidity Gap Cross Sections\textsuperscript{* ++}

H.-G. Kohrs
II. Institut für Theoretische Physik\textsuperscript{‡}, Universität Hamburg
and
Deutsches Elektronen–Synchrotron DESY, Hamburg, Germany

Abstract
We investigate the effect of a direct pomeron coupling to quarks on inclusive jet production in DIS and photoproduction. The direct pomeron coupling generates a point-like contribution to the diffractive part of the structure function $F_2$, which is analysed on the basis of the latest H1 and ZEUS data. Our model assumptions for the pomeron structure are consistent with the measured data.

1. Introduction

In diffractive production of hadronic final states in $ep$ scattering, the proton stays intact or becomes a low mass state. Between the direction of the proton remnant, which goes down the beam pipe, and the produced hadronic system there is no colour flow, which allows of the possibility to observe large gaps in rapidity between these directions. The experiments H1 and ZEUS at HERA have measured the portion of diffractive events to be $\approx 10\%$ of all events – not only in photoproduction ($Q^2 < 0.01\text{ GeV}^2$) but also in deep inelastic scattering (DIS) ($Q^2 > 10\text{ GeV}^2$).

This paper is organized as follows. First, we describe our ansatz to analyse diffractive $ep$ scattering. In section 3, we consider the pomeron structure function and the diffractive part of $F_2$. We find consistency with the latest HERA data. Finally, we analyse jet cross sections in diffractive inclusive photoproduction and in DIS.

2. Model for diffractive jet production

There exist various phenomenological models to describe the above mentioned diffractive nonperturbative QCD phenomena quantitatively. We follow a widely spread assumption, where the proton splits off a colourless object called pomeron ($IP$), which has the quantum numbers of the vacuum. Then, if factorization holds, the proton vertex can be parametrized by a $IP$–flux factor that depends on $t = (p - p')^2$, the momentum transfer to the pomeron, and $x_P$, the fraction of proton energy, that it carries away.

In fact, this has been done in the past by various authors, who fixed their parameters with the help of $p\bar{p}$ scattering data. Inspired by Regge phenomenology, Berger et al. found for the pomeron flux

$$f_{IP}(t, x_P) = \frac{\alpha^2}{16\pi} x_P^{1-2\alpha(t)}$$

with the Regge trajectory $\alpha = \alpha_0 + \alpha' t$ and the residue...
function \( \beta_{g/p}(t) = \beta_{g/p}(0)e^{\text{ht}}, \) where \( \alpha_0 = 1 + \epsilon, \epsilon = 0.085, \alpha' = 0.25 \text{GeV}^{-2}, \beta_{g/p}(0) = 58.74 \text{GeV}^{-2} \) and \( b_0 = 4.0 \text{GeV}^{-2}. \)

This definition of the factorization in a pomeron flux factor differs by a factor of \( \frac{\pi}{2} \) from the definition of Dounachie and Landshoff. In addition, they included the Dirac elastic form factor of the proton, which is given with the parameters \( \alpha = 24 \text{GeV}^{-2} \) and \( \beta = 4 \text{GeV}^{-2} \).

We get \( p_{\text{n}} \) and ZEUS data confirm this.

For our purposes, the momentum transfer \( t = (p - p')^2 \) to the pomeron has to be integrated out, since the proton remnant is (still) not tagged. We get the Dirac elastic form factor of the proton, which is given with the parameters \( \alpha = 24 \text{GeV}^{-2} \) and \( \beta = 4 \text{GeV}^{-2} \). Here, the factor of 1/2 comes in because of normalization to one proton vertex.

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The analysis of the authors suggests a behaviour of the diffractive parton distributions between \( (1 - \beta)^0 \) and \( (1 - \beta)^2 \) for \( \beta \rightarrow 1 \).

The parameter \( \tau \) describes the unknown ratio of the gluon to quark content of the pomeron. To get a first insight into the structure of our model, we restrict the number of free parameters and choose a value \( \tau = 3 \), which represents simple gluon dominance in the pomeron. We carry out the usual DGLAP–evolution to get the right \( Q^2 \) dependence of these functions.

Several groups considered the possibility of a direct pomeron coupling to quarks. A direct pomeron coupling corresponds to a \( \delta \)–function term in the pomeron structure function and produces a leading–twist behaviour in the \( p_T \) spectrum. Our purpose is to find criteria that allow us to see a direct pomeron coupling in the data.

As a consequence, similarly to \( \gamma \gamma \) scattering, the \( \gamma p^2 \rightarrow q\bar{q}^* \) cross section also contributes to the pomeron structure function. Here, we assume a direct vector coupling of the pomeron to the quarks with coupling strength \( c \). This is not really justified with respect to the \( C \) parity. However, the \( Q^2 \) dependence, that is \( \sim \log Q^2 \) at low \( x \), is only weakly dependent on the spin structure. This can be seen, for instance, if one replaces the vector coupling by a scalar one. To remove the collinear singularity, we introduce the regulator quark masses \( m_q \) and obtain for the point–like (pl) part

\[
\beta_{G_{q/p}}(\beta, Q^2) = \beta_{G_{q/p}}(1) = \beta_{G_{q/p}}(0)e^{\text{ht}} - \frac{\beta_0}{1 - \beta} \ln \frac{1 + v}{1 - v} + \left[ \beta^2 + (1 - \beta)^2 + \frac{4m_q^2}{Q^2} \right] \ln \frac{1 + v}{1 - v}.
\]
3. The diffractive contribution to $F_2$

To leading order in $\alpha_s$, only the quark distributions of the pomeron enter into the deep-inelastic $P$ structure function $F_2^P(\beta, Q^2)$, which is

$$F_2^P(\beta, Q^2) = \sum_q c_q^2 \beta \left[ G_{q/P}(\beta, Q^2) + G_{\bar{q}/P}(\beta, Q^2) \right] + 2 G_{pl}^P(\beta, Q^2).$$

The comparison with preliminary 1993 H1 data [1] and ZEUS data [2] is shown in figure 3 and figure 4.

If factorization holds, which is favoured by the experiments for a large range of $\beta$ and $Q^2$ values, the data points are proportional to $F_2^P(\beta, Q^2)$, i.e., $F_2(\beta, Q^2) = k F_2^P(\beta, Q^2)$ with a constant $k$ that is determined by the $P$-flux factor:

$$k = \frac{1}{v} \int_{x_{P}^{\text{min}}}^{x_{P}^{\text{max}}} dx_P \int_{t_{1}}^{t_{2}} dt f_{P/P}(x_P, t).$$

The absolute normalization due to $k$ is very sensitive to the integration bounds, $x_{P}^{\text{min}}$ and $x_{P}^{\text{max}}$, which are taken from the respective experiments (see figure 3 and figure 4). Unfortunately, no experimental errors on them have been published yet. A small reduction of the integration interval would improve our normalization to the data substantially.

We concentrate therefore on the discussion of the shapes of the data in comparison to our model. We find that our model fits the shape of the data curves well. Especially, the $Q^2$-evolution is fitted better for a combined ansatz, i.e., quarks in the pomeron and pl part (solid curves in figure 3), than for the DGLAP evolved quark distributions (short dashed line in figure 3) or pl part alone. An alternative possibility to fit the data has been represented in [2].

Finally, we fold the pomeron structure function [1] with the pomeron flux factor in eq. (1) to get the diffractive part of the proton deep-inelastic structure function $F_2^D(x, Q^2)$. The relation is

$$F_2^D(x, Q^2) = \int_{x_{P}^{\text{min}}}^{x_{P}^{\text{max}}} dx_P \int_{t_{1}}^{t_{2}} dt f_{P/P}(x_P, t) F_2^P(x/x_P, Q^2).$$

The upper bound $x_0 = 0.01$ is an experimental choice. In contrast to the analysis of the pomeron structure function, the variable Bjorken $x$ is now fixed (instead of $\beta$). In figure 5 we compare with 1993 H1 data [20].

Again, we emphasize the consistency of the choice $c = 1$ ($r = 3$) for the direct pomeron coupling in our model and the data. We use this value in the calculation of the diffractive jet cross sections.
4. Jets in diffractive inclusive photoproduction and DIS

The calculation of the differential jet cross section $d^2\sigma/dydp_T^2$ of the process depicted in figure 1 is straightforward. Here, $p_T$ and $y$ are the transverse momentum and rapidity of one outgoing jet. In the case of photoproduction, we perform the evaluation in the $ep$ laboratory system where the rapidity is positive for jets travelling in the proton direction. We have then the usual factorization of the photon flux factor at the electron vertex [16]. As is well known in photoproduction, the photon is resolved or couples directly to the final-state quarks. For the photon particle density functions, we take the parametrizations of GRV [17].

A more detailed discussion can be found in [18]. However, here we have included the improved quark distributions of the pomeron due to the performed $Q^2$-evolution with $Q = p_T$, the transverse momentum of the considered jet. Further, we use in our analysis more actual data of the run in 1993.

The differential inclusive one-jet cross section is obtained by integrating out all kinematic variables over the allowed ranges without regard to the rapidity of the second jet, while for the two-jet cross section, we demand explicitly that the second jet does not enter the cone that is set up by the first jet around the outgoing proton direction.

The results are shown in figure 6. Note the large rapidity gap in the forward direction between the diffractive and nondiffractive parts. In our analysis this is controlled by the $x_{B}^{max} = 0.01$ cut. This
The rapidity distribution of the a) one- and b) two-jet cross sections for fixed transverse momentum $p_T = 5$ GeV in the $ep$ laboratory system. Here, $y$ is defined to be positive for jets travelling in the proton direction. For comparison, the nondiffractive cross section obtained with CTEQ parametrizations of the proton structure functions is also shown (solid line). Since the 1993 event rates (data points) of H1 [21] and ZEUS [22] are not normalized to the luminosity, we can compare only the shapes.

The need of a direct pomeron coupling becomes clear, if one inspects the slope of the $p_T$ spectra in figure 7. As expected, with the direct coupling, the $p_T$ spectrum does not fall off so strongly compared to the resolved pomeron contribution ($c=0$) and is favoured by the shape of the data.

In DIS, the photon is always direct. The jet cross section has been calculated for our model in the $\gamma^*p$ c.m.s. Jets with positive rapidity are travelling now in the photon direction. With $p$, the four-momentum of the proton, $k$ the four-momentum of the electron, and $q$, the momentum transfer to the photon, we define the usual kinematic variables $s = (p+k)^2$, $Q^2 = -q^2$, $W^2 = (p+q)^2$, $x = Q^2/(2pq)$, $y_e = pq/pk$, $z = p_T e^y/W$. Finally, we set $\xi = x + p_T^2/(yz(1-z)s)$, which is the fraction of energy delivered from the proton to the subprocess (see figure 8).

A reduction of $r$ would increase the quark content in the pomeron due to the sum rule, eq. (8). But this would be accompanied by a reduction of the coupling $c$ to satisfy the bounds coming from the analysis of the diffractive part of $F_2$ in section 3.

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The inclusive one–jet cross section is then given by

\[ \frac{d^2\sigma}{dy dp_T} = \int_{a+b}^1 dy_e \int_{1-y_e}^{-b/y_e} dx \frac{d^4\sigma}{dx dy_e dy dp_T} \]  \hspace{1cm} (14) \]

where the kinematic bounds follow from the requirements \( Q^2 \geq Q_{\text{min}}^2 \) and \( W^2 \geq W_{\text{min}}^2 \): \( a = Q_{\text{min}}^2/s \), \( b = \max((2p_T \cosh y)^2, W_{\text{min}}^2)/s \).

In the resolved pomeron case, we have

\[ \frac{d^4\sigma}{dx dy_e dy dp_T} = \sum_{bi} \int_{\xi}^{x_{\text{max}}} \frac{dx_{\gamma}}{x_{\gamma}} G_{\gamma/p}(x_{\gamma}) G_{b/p}(\xi/x_{\gamma}, Q^2) \]

\[ \frac{2a s \alpha^2 Q_i^4}{\xi y g_{\gamma}^2 (1-z) s^2} \left\{ [1 + (1 - y_e)^2] \hat{h}_u + 2 (1 - y_e) \hat{h}_l \right\}. \]  \hspace{1cm} (15) \]

The contributions of the direct coupling is

\[ \frac{d^4\sigma}{dx dy_e dy dp_T} = \sum_{i} G_{\gamma/p}(x_{\gamma}) \]

\[ \frac{2c_i^2 \alpha^2 Q_i^2}{4\pi x y g_{\gamma}^2 (1-z) s^2} \]

\[ 6 \left\{ 1 + (1 - y_e)^2 \right\} \hat{h}_u + 2 (1 - y_e) \hat{h}_l \right\}. \]  \hspace{1cm} (17) \]

The functions \( \hat{h}_u = \frac{1}{3} (\hat{h}_g + \hat{h}_l) \) and \( \hat{h}_l \) depend on the Mandelstam variables \( s + i + \hat{u} = -Q^2 \) of the subprocesses. For \( \gamma q \rightarrow g q \), we have

\[ \hat{h}_g = \frac{4}{3} \left( -\frac{i}{s} - \frac{\hat{u}}{t} + \frac{2Q^2 \hat{u}}{st} \right), \]  \hspace{1cm} (18) \]

\[ \hat{h}_l = \frac{4}{3} \left( -2Q^2 t \right) \]  \hspace{1cm} (19) \]

while for \( \gamma q \rightarrow q \bar{q} \), we find

\[ \hat{h}_g = \frac{1}{2} \left( -\frac{i}{u} - \frac{\hat{u}}{t} - \frac{2Q^2 s}{tu} \right), \]  \hspace{1cm} (20) \]

\[ \hat{h}_l = \frac{1}{2} \left( \frac{4Q^2 s}{(Q^2 + s)^2} \right). \]  \hspace{1cm} (21) \]

Like in the photoproduction case, absolute experimental data for the rapidity distribution or \( p_T \) spectrum are not yet available to us. In figure 8, we compare the shape of the \( p_T \) spectrum with 1993 ZEUS data. The experimental conditions were \( Q^2 > 10 GeV^2 \), \( W > 140 GeV \) and \( 0.04 < y_e < 0.95 \). No special case of our model is favored, since the slopes of direct and resolved contribution are identical, but the shape can be approximately reproduced.

In our analysis, we did not consider particle production or hadronization effects etc. The Monte–Carlo–programs POMPYT by Bruni and Ingelman [18] or RAPGAP by Jung [19] have been developed for a wide class of pomeron models and allow, together with other programs, the study of event characteristics. They are widely used by the HERA collaborations to interpret the large rapidity gap data.

5. Conclusion

In summary, we have studied the effect of a direct pomeron coupling on diffractive jet production at HERA. The concept of pomeron structure functions with DGLAP \( Q^2 \)-evolution has to be enlarged if the pomeron has a direct coupling to quarks. We have included the additional point–like part in the analysis of diffractive \( F_2 \) data and find consistency for the assumption of a direct pomeron coupling to quarks. Some evidence for a direct coupling has been found in the \( p_T \) spectrum of photoproduction. However,
our analysis is model dependent and second, except for the discussion of $F_2^P$ and $F_2^D$, we have compared only the shapes and not the normalizations of the photoproduction and DIS cross sections.

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