Large Rapidity Gap Events in DIS

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Abstract. Diffractive scattering in DIS is discussed in terms of the perturbative two-gluon model and numerical results for $F_2^D$ are presented.

Large Rapidity Gap Events can only be explained in conjunction with color zero or vacuum exchange, since gaps with color exchange are exponentially suppressed. The vacuum exchange is generally divided up into two groups, the subleading meson exchanges and the Pomeron which represents the leading exchange at very high energies or small x-Bjorken. In most cases the proton emerges intact from the reaction (Single Diffractive Dissociation) and proton dissociation (or Double Diffractive Dissociation) is only considered as background which is, however, not completely negligible.

The simplest representation of the Pomeron within QCD is a two gluon pair with zero net color charge. The full structure of the Pomeron is of course much more complicated and higher order contributions have to be added. Fortunately, the first correction, which is due to multi gluon exchange, does not require any recalculation as will be explained. The process depicted in Fig.1 consists of two basic contri-
butions, the production of a single quark-antiquark pair and the emission of an extra gluon. Higher order multi parton final states may be included by switching on QCD-evolution, in this talk, however, only the lower orders are considered.

A very convenient and intuitive approach is the wave function formalism which associates the parton content of the photon with a color dipole. The simple quark-antiquark dipole yields the following wave function:

\[
\Psi^\gamma_h(\alpha, k_t) = \begin{cases} 
\sqrt{2} \frac{\alpha}{k_t^2 + \alpha(1-\alpha)Q^2} & \text{for } \gamma = +1 \text{ and } h = +1 \\
\sqrt{2} \frac{\alpha}{k_t^2 + \alpha(1-\alpha)Q^2} & \text{for } \gamma = +1 \text{ and } h = -1 \\
\sqrt{2} \frac{\alpha}{k_t^2 + \alpha(1-\alpha)Q^2} & \text{for } \gamma = -1 \text{ and } h = +1 \\
\sqrt{2} \frac{\alpha}{k_t^2 + \alpha(1-\alpha)Q^2} & \text{for } \gamma = -1 \text{ and } h = -1 
\end{cases}
\] (1)

and

\[
\Psi^\delta_h(\alpha, k_t) = \frac{\alpha(1-\alpha)Q}{k_t^2 + \alpha(1-\alpha)Q^2} \quad \text{for } \gamma = 0 \text{ and } h = \pm 1
\] (2)

where \(\gamma\) denotes the photon and \(h\) the quark helicity. The variables correspond to the Sudakov-decomposition of the momentum \(\tilde{k}\) (see Fig.1): \(\tilde{k} = \alpha Q' + \frac{k_t^2}{\alpha s} p + k_t\).

Similar results can be found in [1–3].

In the case of gluon emission one finds that for large enough photon virtuality \(Q^2\) the quark-antiquark state (the shaded area in Fig.1.b) effectively acts as a gluon state. This leads to the notion of a single gluon dipole which in analogy to the quark-antiquark dipole previously is described by the following wave function:

\[
\Psi^{mn}(\alpha, k_t) = \frac{1}{\sqrt{\alpha(1-\alpha)Q^2}} \frac{k_t^2 \delta^{mn} - 2 k_t^m k_t^n}{k_t^2 + \alpha(1-\alpha)Q^2}.
\] (3)

This expression goes beyond earlier results in [1,2,4] which were restricted to the triple Regge limit (large diffractive mass) where \(\alpha(1-\alpha)Q^2\) is much smaller than \(k_t^2\). In this case the \(\delta^{mn}\)-term drops out when all diagrams are summed up. It, however, becomes important when the extension towards small masses is considered.

In both cases the wave function has to be folded with the unintegrated gluon structure function \(F\) (see Fig.1) which can be extracted from the inclusive \(F_2\)-data (for more details see [5]). The \(t\)-dependence (momentum transfer) was added by hand using the diffractive slope as measured by ZEUS [6]. It is interesting to note
that the semiclassical approach of ref. [7] obtains the same results as presented in eqs.(1-3). The identity

\[ W(k_\perp) = 2 \frac{\mathcal{F}(k_\perp)}{k_\perp^2} - 2 \delta^2(k_\perp) \int d^2k'_\perp \frac{\mathcal{F}(k'_\perp)}{k'^2_\perp} \] (4)

provides the link between the classical gluon field \( W \) in ref. [7] and the unintegrated structure function \( \mathcal{F} \). A perturbative expansion of the underlying eikonal factors shows that \( W \) can be interpreted as multi gluon exchange. It therefore represents a major improvement over the simple two-gluon model. This is of particular relevance, since the diffractive process is dominated by low \( k_t \), i.e. soft contributions.

In the following section the main numerical results for \( F_2^D(x_P, \beta, Q^2) \) as derived from eqs.(1-3) in combination with an appropriate parametrization for the unintegrated structure function \( \mathcal{F} \) are presented (\( \beta = Q^2/(M^2 + Q^2) \) and \( x_P = (M^2+Q^2)/(W^2+Q^2) \)). First we examine the \( x_P \)-distribution (Fig.2). The Pomeron intercept depends on the scale \( k_t^2/(1-\beta) \) which leads to the breaking of Regge-factorization. With increasing \( Q^2 \) and/or increasing \( \beta \) the \( x_P \)-distributions become steeper, i.e. the process harder. Since the scale is not directly coupled to \( Q^2 \) the process never becomes purely hard and the Pomeron-intercept lies only slightly above the value for the soft Pomeron.

The shape of the \( \beta \)-spectrum turns out to be rather flat when all three contributions associated with the longitudinal and transverse production of quarks (large and medium \( \beta \), \( F_b \) and \( F_a \) in Fig.3.a) and the emission of an extra gluon are combined. The latter fills the low and only the low \( \beta \)-region (\( F_c \) in Fig.3.a). The fact that this contribution is rapidly decreasing with increasing \( \beta \) is an important consequence of the \( k_t^2\delta^m \)-term in eq.(3). With increasing \( Q^2 \) the small \( \beta \) region starts
FIGURE 3. $\beta$-spectrum at $x_F = 5 \cdot 10^{-4}$.

rising (log($Q^2$)-correction) whereas the large $\beta$ region decreases which indicates the higher twist nature of the longitudinal part.

The $Q^2$-dependence is more clearly depicted in Fig.4.a. Due to the $Q^2$ cutoff in the phase space all graphs show in the beginning a rise with $Q^2$. Only at rather large $Q^2$ the asymptotic regime is approached and the higher twist contribution (dominant when $\beta = 0.9$) turns down while the purely leading twist part (dominant when $\beta = 0.5$) flattens out. This delay in the asymptotic behavior is model depen-

dent as can be seen in Fig.4.b. Only the transverse part of the quark-antiquark production which is supposed to give a flat leading twist contribution is presented here. The phenomenological Pomeron model as compared to the BFKL-Pomeron [8] reaches the asymptotic limit faster. The observed deviations from the leading twist behavior are in general related to higher twist corrections which turn out be non negligible even at rather large $Q^2$ (see also [9]).

Finally we examine the azimuthal distribution (Fig.5). The azimuthal angle is defined in the transverse plane with respect to the lepton. The interesting feature
is the peak at $\pi/2$ which is strongly enhanced when a hard final state is required, i.e. exclusive jets without a remnant (cut on $k_t$). But already the inclusive final state shows some oscillation (dashed line).

**FIGURE 5.** Azimuthal distribution.

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