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# The Partonic Structure of the Generalized BFKL Pomeron : an Overview

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# THE PARTONIC STRUCTURE OF THE GENERALIZED BFKL POMERON: AN OVERVIEW

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## ABSTRACT

We report on QCD calculation of the cross section of diffractive DIS and of the partonic structure of the Pomeron in the generalized BFKL model. Our main result is that the Pomeron cannot be treated as a particle with uniquely defined structure function and flux from the proton. In fact different fluxes must be defined for the sea and the valence components. Furthermore, the fluxes for the charm and longitudinal valence component are different from the previous ones and cannot be defined as $\beta$ and $Q^2$ independent respectively.

Diffractive Deep Inelastic Scattering can be understood (if one assumes single pomeron exchange) as DIS on Pomerons radiated by protons. Many authors assumed[1)] that the Pomeron can be treated as a particle with a well-defined partonic structure and flux $f(x_{\mathbb{P}})/x_{\mathbb{P}}$ in the proton. Under this strong assumption, one is lead to the operational definition of the Pomeron structure function $F_2^{(\mathbb{P})}(\beta, Q^2)$:

$$(M^2+Q^2)\frac{d\sigma_D(\gamma^* \to X)}{dt\, dM^2}\bigg|_{t=0} = \frac{\sigma_{tot}^{pp}}{16\pi}\frac{4\pi^2\alpha}{Q^2} f_{\mathbb{P}}(x_{\mathbb{P}}) F_2^{(\mathbb{P})}(\beta, Q^2) \qquad (1)$$

(we use the Regge theory conventions). The factorization property (1) should be proved starting with a QCD treatment of diffractive DIS and it only makes sense if one can prove that the flux function $f_{\mathbb{P}}(x_{\mathbb{P}})$ can be defined so that the $Q^2$ dependence of the r.h.s. of (1) is concentrated in $F_2^{\mathbb{P}}(\beta, Q^2)$, which satisfies the conventional QCD evolution. We report here on the works[2)], based on the colour dipole approach to the BFKL Pomeron[3)], where we have found that the factorization (1) is broken.

The QCD process which corresponds to DIS on the valence $q\bar{q}$ pair is the diffraction excitation of the $q\bar{q}$ Fock state of the photon, which has the cross section

$$\frac{d\sigma_D}{dt}\bigg|_{t=0} = \frac{1}{16\pi}\int_0^1 dz \int d^2\vec{r}\, |\Psi_{\gamma^*}(Q^2, z, r)|^2 \sigma(x,r)^2 = \frac{4\pi^2\alpha}{Q^2}\int \frac{dr^2}{r^2}\, W(Q^2, r)\left[\frac{\sigma(x,r)}{r^2}\right]^2, \qquad (2)$$

where $\Psi_{\gamma^*}$ is the wave function of the $q\bar{q}$ pair of transverse size $\vec{r}$, $z$ is the fraction of the photon light-cone momentum carried by the quark, and $\sigma(x,r)$ is the colour

dipole cross section. The properties which allow reinterpretation of the cross section (2) as DIS on the valence $q\bar{q}$ state of the Pomeron are:
(i) The mass spectrum, which roughly follows[3)]

$$\left.\frac{d\sigma_D}{dM^2dt}\right|_{t=0} \propto \frac{M^2}{(Q^2+M^2)^3} = \frac{1}{Q^2}\cdot\frac{1}{Q^2+M^2}\cdot\beta(1-\beta)\,, \tag{3}$$

is $x$-independent to a good approximation.
(ii) At sufficiently high $Q^2$ the diffractive cross section satisfies the Bjorken scaling[3)].
(iii) $W(Q^2,r)$ is peaked at a large, $Q^2$-independent, hadronic size: this recalls the $Q^2$-independent large distance between the valence quark and antiquark in the pion, and we can analogously speak of DIS off the valence $q\bar{q}$ state of the Pomeron.

Using the colour dipole cross section[4)], which gives a good description of the HERA data on the proton structure, we find from (2) and (3) the form and absolute normalization (at $Q^2 = 10\text{GeV}^2$) of the valence $q\bar{q}$ structure function of the Pomeron $F_{val}^{(\mathbb{P})}(\beta) = 0.27\beta(1-\beta)$. The flux function $\phi_{\mathbb{P}}(x_{\mathbb{P}})$, defined subject to the normalization $\phi_{\mathbb{P}}(x_{\mathbb{P}}^0 = 0.03) = 1$, equals

$$\phi_{\mathbb{P}}(x_{\mathbb{P}}) = \frac{\int_0^1 dz \int d^2\vec{r}\,|\Psi_{\gamma^*}(Q^2_{\mathbb{P}},z,r)|^2\sigma(x_{\mathbb{P}},r)^2}{\int_0^1 dz \int d^2\vec{r}\,|\Psi_{\gamma^*}(Q^2_{\mathbb{P}},z,r)|^2\sigma(x^0_{\mathbb{P}},r)^2}\,. \tag{4}$$

The mass spectrum (3) for excitation of the $q\bar{q}$ state rapidly decreases at large $M^2 \gg Q^2$. The Born process for the $\propto 1/(M^2+Q^2)$ mass spectrum, typical of the so-called triple-Pomeron regime, is the diffraction excitation of the $q\bar{q}g$ Fock state of the photon. One obtains[3)] ($\vec{\rho}$ is the quark–gluon separation):

$$(Q^2+M^2)\left.\frac{d\sigma_D}{dtdM^2}\right|_{t=0} \simeq \int dz\, d^2\vec{r}\,|\Psi_{\gamma^*}(Q^2,z,r)|^2\cdot\frac{27}{32}\frac{\alpha_S(r)r^2}{2\pi^2}\int d\rho^2\left[\frac{\sigma(x_{\mathbb{P}},\rho)}{\rho^2}\right]^2\mathcal{F}(\mu_G\rho) \tag{5}$$

with an approximate factorization of $Q^2-$ and $x_{\mathbb{P}}-$ dependence ($\mathcal{F}(z) = z^2[K_1(z)^2 + zK_1(z)K_0(z) + \frac{z^2K_0(z)^2}{2}]$). This factorization property allows us to define the corresponding structure function and the Pomeron flux function. Furthermore, one can introduce the two-gluon wave function of the Pomeron[3)]

$$|\Psi_{\mathbb{P}}(\beta,\vec{r})|^2 = \frac{1}{f_{\mathbb{P}}(x_{\mathbb{P}})}\cdot\frac{81}{8\pi^4}\cdot\frac{1-\beta}{\beta}\cdot\frac{1}{\sigma_{tot}^{pp}}\left[\frac{\sigma(x_{\mathbb{P}},r)}{r^2}\right]^2\mathcal{F}(\mu_G r)\,, \tag{6}$$

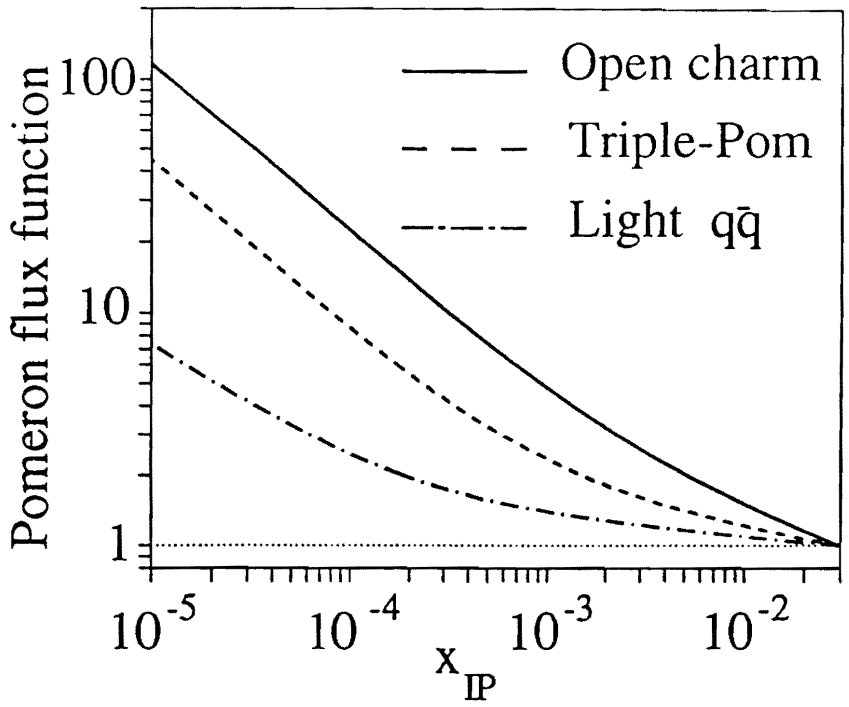
where $\vec{r}$ is the transverse separation of gluons in the Pomeron and $\beta$ is a fraction of the Pomeron's momentum carried by a gluon. The $1/\beta$ dependence is the usual soft-gluon behaviour, the factor $(1-\beta)$ is an educated guess. In the wave function (6), the $x_{\mathbb{P}}$-dependence approximately cancels out, and it gives the $x_{\mathbb{P}}$-independent gluon structure function of the Pomeron $G_{\mathbb{P}}(\beta) = \beta\int d^2\vec{r}\,|\Psi_{\mathbb{P}}(\beta,\vec{r})|^2 = 0.28(1-\beta)$.

Then, Eq. (5) gives the input sea structure function of the Pomeron for $\beta \ll 1$:

$$F_{sea}^{\mathbb{P}}(Q^2_{\mathbb{P}}) = \frac{81}{32\sigma_{tot}^{pp}}\frac{1}{f_{\mathbb{P}}(x_{\mathbb{P}})}\frac{Q^2_{\mathbb{P}}}{\pi\alpha}\int d\rho^2\left[\frac{\sigma(x_{\mathbb{P}},\rho)}{\rho^2}\right]^2\mathcal{F}(\mu_G\rho)\int_0^1 dz d^2\vec{r}\,|\Psi_{\gamma^*}(Q^2_{\mathbb{P}},r,z)|^2r^2\alpha_S(r). \tag{7}$$

As a starting approximation, we take $F_{sea}^{\mathbb{P}}(\beta,Q^2_{\mathbb{P}}) = C_{sea}(1-\beta)^2 = 0.063(1-\beta)^2$, with the normalization following from Eq. (7).

From the previous discussion, it follows that the two fluxes $\phi_{\mathbb{P}}(x_{\mathbb{P}})$ and $f_{\mathbb{P}}(x_{\mathbb{P}})$ are not identical (see the figure). The different $x_{\mathbb{P}}$-dependence of these two fluxes




- Our prediction for the pomeron flux function $\Phi^{(c\bar{c})}(x_{\mathbb{P}})$ for DD into open charm (the solid curve) in comparison with flux functions for excitation light flavours at $\beta \sim 0.5$ (the dot-dashed curve) and for the triple-pomeron region of $\beta \ll 1$.

is a dynamical prediction of this analysis, to be contrasted to conjectured forms of a universal flux of Pomerons[1]. Only at very large $\frac{1}{x_{\mathbb{P}}}$, well beyond the kinematical range of HERA, do the two fluxes have a similar $\frac{1}{x_{\mathbb{P}}}$-dependence.

Furthermore other factorization breaking effects are expected due to the heavy quarks[2a] and longitudinal[2b] valence contributions. In fact in both the cases the $\Psi_{\gamma^*}$ of eq. (4) differs from the light quarks one for transverse polarized $\gamma^*$.

An explicit calculation shows that the relevant factorization scale is $q^2 = (k^2 + m_f^2)\frac{M^2+Q^2}{M^2}$ where $k$ is the transverse momentum of quark with respect to the $\gamma^*$-pomeron collision axis and $m_f$ the quark mass.

The transverse cross section is dominated by small $k^2$ and one finds that for the charm component the flux is $\phi_{\mathbb{P}}^{c\bar{c}} = \left[\frac{G(x_{\mathbb{P}},q^2)}{G(x_{\mathbb{P}}^0,q^2)}\right]^2$, which acquires a $\beta$ dependence through the factorization scale $q^2 = \frac{m_c^2}{(1-\beta)}$. The numerical evaluation shows that it is steeper than the light quarks one (see the figure).

For what concerns the longitudinal cross section, the calculation shows that it is dominated by large $k^2$, leading to the factorization scale $q_L^2 = \frac{Q^2}{4\beta}$. This requires to introduce a $Q^2$ dependent flux $\phi_{\mathbb{P}}^L = \left[\frac{G(x_{\mathbb{P}},q_L^2)}{G(x_{\mathbb{P}}^0,q_L^2)}\right]^2$. Incidentally, this last flux, in contrast to the transverse photons case, is flavour-independent.

The longitudinal Pomeron structure function has the form $F_{2,\mathbb{P}}^L = \Sigma_f e_f^2 \frac{A_f}{Q^2}(1-2\beta)^2 \cdot \beta^3$ ($A_{ud} = 0.82\text{GeV}^2$, $A_s = 0.61\text{GeV}^2$ and $A_c = 0.05\text{GeV}^2$) and is dominant for large $\beta$, even at $Q^2 \approx 100 GeV^2$.

Quite a different situation occurs for the sea flux defined by eq. (7), which is universal for all the flavours and $\gamma^*$-polarizations.

Finally, further factorization breaking effects are expected when one considers charge currents diffractive DIS[6].

Albeit these factorization breaking effects are still too elusive to be detected in the available experimental data, they should be observed when more precise data will appear.

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