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PREDICTIONS FOR THE FORWARD CONE IN DIFFRACTIVE DEEP INELASTIC SCATTERING

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We calculate the diffraction slope B_D for diffractive deep inelastic scattering. We find a counterintuitive rise of B_D from exclusive diffractive excitation of vector mesons to excitation of continuum states with $M^2 \sim Q^2$. For the small-mass continuum we predict a rapid variation of B_D with M^2 on the scale m_V^2 and a sharp drop of B_D for a small-mass continuum above the vector meson excitation.

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The diffraction slope is one of the principal observables which measures the impact parameter structure of diffractive scattering. The comissioning of the leading proton spectrometer (LPS) of the ZEUS detector at HERA [1] gave a long awaited access to the transverse momentum transfer Δ and the diffraction slope $B_D = -\partial \log\left\{d\sigma_D/d\Delta^2)\right\}/\partial\Delta^2$ in diffractive deep inelastic scattering (DIS) $ep \to e'p'X$. The special interest in diffraction slope for diffractive DIS stems from the fact that besides the mass M of the excited state X there emerges a new large scale: the virtual photon's mass $\sqrt{Q^2}$. The principal issue is what B_D depends from: M^2, Q^2, the mass m_V of the ground-state vector meson in the corresponding flavour channel and/or the diffractive scaling variable $\beta = Q^2/(Q^2+M^2)$ (hereafter Q^2, x and $x_{\mathbf{P}} = x/\beta$ are the standard diffractive DIS variables).

This is a highly nontrivial issue because at fixed β diffraction proceeds into the high-mass continuum states X with $M^2 = Q^2(1-\beta)/\beta \gg m_V^2$. Our experience with diffraction of hadrons and/or real photons can be summarized as follows. For any two-body diffractive scattering $ac \to bd$, an essentially model-independent decomposition holds, $B_D = \Delta B_{ab} + \Delta B_{cd} + \Delta B_{int}$, where ΔB_{ij} comes from the size of the ij transition vertex and the relatively small ΔB_{int} comes from the interaction range proper [2, 3]. The values of ΔB_{ij} depend strongly on the excitation energy in the $i \to j$ transition, $\Delta M^2 = m_j^2 - m_i^2$. In elastic scattering, $i = j$, one finds $\Delta B_{ii} \approx 1/3R_i^2 \sim$ 4-6 GeV^{-2}, where R_i^2 is the mean squared hadronic radius, and typically $B_{el} \sim 10$ GeV^{-2}. The similar estimate $\Delta B_{ij} \approx 1/3R_i^2, 1/3R_j^2$ holds for diffraction into low-mass continuum states, $\Delta M^2 \lesssim m_N^2$, and diffraction into low-mass continuum and elastic scattering fall into the broad category of *exclusive* diffraction for which $B_D \sim B_{el}$. However, for excitation of high-mass continuum, $\Delta M^2 \gtrsim m_N^2$, often referred to as the triple-pomeron (3$\mathbf{P}$) and/or genuine *inclusive* region of diffraction, the size of the diffracting particle no longer contributes to the diffractive slope and $B_D = B_{3\mathbf{P}} = \Delta B_{pp} + \Delta B_{int} \sim 1/2B_{el} \approx 6$ GeV^{-2}. The above slope $B_{3\mathbf{P}}$ is about universal for all the diffracting beams and excited states X [3]). Furthermore, in the double high-mass diffraction $hp \to XY$, when $M_{X,Y} \gg m_N$, one is left with very small $B_D \sim \Delta B_{int} \sim$1.5-2 GeV^{-2} ([2, 4] and references therein). In real photoproduction the excitation scale is definitely set by the ground-state vector meson mass m_V. Perhaps the most dramatic example of this distinction between *exclusive* and *inclusive* diffraction is a drastic change of the diffraction slope from elastic, $pA \to pA$, to quasielastic, $pA \to p'A^*$, scattering of protons on heavy nuclei [5].

Another well understood diffractive process is elastic production of vector mesons $\gamma^* p \to p'V$. In this case the transverse size $\gamma^* \to V$ transition vertex, the so-called scanning radius

$$r_S = 6/\sqrt{Q^2 + m_V^2} \tag{1}$$

decreases with Q^2 (and m_V^2). This is a basis of the prediction [6] of $\Delta B_{\gamma^* V} \propto r_S^2$ and of the decrease of the diffraction slope B_V down to $B_V \approx B_{3\mathbb{P}}$ at very large Q^2, which is in good agreement with the experiment [7].

In this paper we report predictions for the Q^2, M^2 and flavour dependence of the diffraction slope for inclusive diffractive DIS. We demonstrate that in striking contrast to B_V for exclusive diffraction into vector mesons which exhibits strong dependence on Q^2, the diffraction slope B_D for inclusive diffractive DIS is a scaling function of β. The most paradoxical prediction is that in contrast to real photon and hadronic diffraction, in diffractive DIS B_D rises with the excited mass M reaching $B_D \sim B_{el}$ at $M^2 \sim Q^2$. Arguably, such an unusual behaviour of B_D derives from the scaling scanning radius r_S for diffraction excitation of continuum $q\bar{q}$ states [8],

$$r_S^2 \sim \frac{9}{m_f^2}(1-\beta)\,, \tag{2}$$

which rises towards small β, so that $\Delta B_{\gamma^* X} \propto r_S^2$ does not depend on Q^2 and rises substantially from $\beta \approx 1$ to $\beta \sim 1/2$. Earlier such a large, Q^2-independent $\Delta B_{\gamma^* X}$ has been conjectured in [9] and in the present communication we quantify this property of the diffraction slope by a direct calculation. Furthermore, we predict a substantial drop of B_D below $B_{3\mathbb{P}}$ for excitation of the small-mass continuum.

Finally, for very large excited masses, $M^2 \gg Q^2$, i.e., $\beta \ll 1$, even for the $q\bar{q}$ excitation one recovers the inclusive regime of small $\Delta B_{\gamma^* X}$ and B_D decreases back to $B_D \sim B_{3\mathbb{P}}$. This triple-pomeron limit of $\beta \ll 1$ is dominated by excitation of the $q\bar{q}g$ and higher Fock states of the photon, though, which is the genuinely *inclusive* process and by the same token as for hadronic diffraction one can argue [9] that B_D must not depend on Q^2 and that $B_D \approx B_{3\mathbb{P}}$. This anticipation has been confirmed by the first data from the ZEUS LPS: $B_D = 7.2 \pm 1.1(stat.)^{+07}_{-0.9}(syst.)$ GeV^{-2} for diffractive DIS ($5 < Q^2 <20$ GeV^{-2}) [1] and $B_D = 6.8 \pm 0.9(stat.)^{+1.2}_{-1.1}(syst.)$GeV^{-2} in real photoproduction ($Q^2 = 0$) [10].

We focus on diffractive excitation of the $q\bar{q}$ Fock states of the photon, which is known to dominate at $\beta \gtrsim 0.1$ [11]. The sample Feynman diagram for this process is shown in Fig.1, in which we show also all the relevant momenta. We base our analysis on the formalism [12], which we generalize to the non-forward case $\mathbf{\Delta} \neq 0$[1)].

If z and $(1-z)$ are fractions of the (lightcone) momentum of the photon carried by the quark and antiquark, respectively and $\mathbf{k}$ is the relative transverse momentum in the $q\bar{q}$ pair, then $M^2 = (m_f^2 + k^2)/z(1-z)$. The quark and antiquark are produced with the

[1)] The first calculation of the diffraction slope for the M^2 integrated cross section is found in [13], the preliminary results from the present study have been reported elsewhere [14].

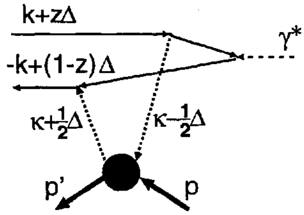

Fig.1. One of the four Feynman diagrams for diffraction excitation of the $q\bar{q}$ final state via QCD two-gluon pomeron exchange

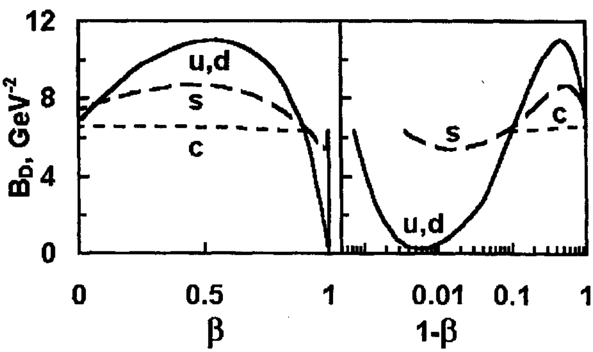

Fig.2. Our predictions for the β and flavour dependence of the diffraction slope B_D in diffractive DIS of transverse photons at $Q^2 = 100$ GeV^2

transverse momenta $\mathbf{k} + z\mathbf{\Delta}$ and $-\mathbf{k} + (1-z)\mathbf{\Delta}$ with respect to the $\gamma^* p$ collision axis. We focus on the transverse diffractive structure function (SF). To the leading $\log 1/x_{\mathbf{P}}$, for excitation of quarks of mass m_f and electric charge e_f,

$$F_T^{D(4)}(\mathbf{\Delta}^2, x_{\mathbf{P}}, \beta, Q^2) =$$

$$= \frac{8\pi e_f^2}{3\sigma_{tot}(pp)} \int \frac{d^2\mathbf{k}}{2\pi} \frac{(k^2+m_f^2)\beta}{(1-\beta)^2 J} \alpha_S^2(\overline{Q}^2) \left\{ [1 - 2z(1-z)] \mathbf{\Phi}_1^2 + m_f^2 \mathbf{\Phi}_2^2 \right\}, \quad (3)$$

where $J = \sqrt{1 - 4(k^2+m_f^2)/M^2}$, $\alpha_S^2(\overline{Q}^2)$ is the strong coupling, evaluated at the QCD hardness scale $\overline{Q}^2$ to be specified below, and $f(x_{\mathbf{P}}, \boldsymbol{\kappa}, \mathbf{\Delta})$ is the gluon density matrix [6, 15]. In the calculation of diffractive helicity amplitudes $\mathbf{\Phi}_1, \mathbf{\Phi}_2$ it is convenient to introduce

$$\psi(z, \mathbf{k}) = \frac{1}{\mathbf{k}^2 + m_q^2 + z(1-z)Q^2}, \quad \mathbf{\Psi}(z,k) = \mathbf{k}\psi(z,\mathbf{k}), \quad (4)$$

in terms of which

$$\mathbf{\Phi}_i = \int \frac{d^2\boldsymbol{\kappa}}{2\pi\kappa^4} f(x_{\mathbf{P}}, \boldsymbol{\kappa}, \mathbf{\Delta}_\perp)\, \phi_i \quad (5)$$

where

$$\phi_1 = \mathbf{\Psi}(z, \mathbf{r}+\boldsymbol{\kappa}) + \mathbf{\Psi}(z, \mathbf{r}-\boldsymbol{\kappa}) - \mathbf{\Psi}(z, \mathbf{r}+\frac{1}{2}\mathbf{\Delta}) - \mathbf{\Psi}(z, \mathbf{r}-\frac{1}{2}\mathbf{\Delta}), \quad (6)$$

$$\phi_2 = \psi(z, \mathbf{r}+\boldsymbol{\kappa}) + \psi(z, \mathbf{r}-\boldsymbol{\kappa}) - \psi(z, \mathbf{r}+\frac{1}{2}\mathbf{\Delta}) - \psi(z, \mathbf{r}-\frac{1}{2}\mathbf{\Delta}), \quad (7)$$

$$\mathbf{r} = \mathbf{k} - \frac{1}{2}(1-2z)\,\mathbf{\Delta}. \quad (8)$$

For small $\mathbf{\Delta}$ within the diffraction cone

$$\mathcal{F}(x, \boldsymbol{\kappa}, \mathbf{\Delta}) = \frac{\partial G(x,\kappa^2)}{\partial \log \kappa^2} \exp(-\frac{1}{2} B_{3\mathbf{P}} \mathbf{\Delta}^2) \quad (9)$$

where $\partial G/\partial \log \kappa^2$ is the conventional unintegrated gluon structure function [6]. The dependence of $\mathcal{F}(x, \kappa, \Delta)$ on $\Delta\kappa$ corresponds to the subleading BFKL singularities [15] and can be neglected at small $x_{\mathbf{I\!P}}$. The diffraction slope $B_{3\mathbf{I\!P}}$ in (9) is a nonperturbative quantity, it comes for the most part form the hadronic size of the proton, modulo to a slow Regge growth one can take $B_{3\mathbf{I\!P}} \sim 6$ GeV^{-2} [6].

In the present analysis we are mostly concerned with the β, Q^2 and flavour dependence of $\Delta B_{\gamma^* X}$ which comes from the $\mathbf{\Delta}$ dependence of ϕ_1 and ϕ_2, for our purposes it is sufficient to evaluate $\mathbf{\Phi}_1^2, \mathbf{\Phi}_2^2$ to an accuracy $\mathbf{\Delta}^2$. The calculation of amplitudes $\mathbf{\Phi}_1, \mathbf{\Phi}_2$ has been discussed to great detail in [8, 12, 16] and need not be repeated here. We simply cite the results starting with excitation of heavy quark-antiquark pair, when the fully perturbative quantum chromodynamics (pQCD) analytic calculation is possible:

$$F_T^{D(4)}(t, x_{\mathbf{I\!P}}, \beta, Q^2) = \frac{2\pi e_f^2}{9\sigma_{tot}(pp)} \frac{\beta(1-\beta)^2}{m_f^2} \Big[(3 + 4\beta + 8\beta^2) +$$

$$+ \frac{\Delta^2}{m_f^2}\frac{1}{10}\left(5 - 16\beta - 7\beta^2 - 78\beta^3 + 126\beta^4\right)\Big] \left[\alpha_s(\overline{Q}^2) G(x_{\mathbf{I\!P}}, \overline{Q}^2)\right]^2 \exp\left(-B_{\mathbf{I\!P}}\mathbf{\Delta}^2\right) \quad (10)$$

where the pQCD hardness scale equals

$$\overline{Q}^2 \approx m_f^2(1 + \frac{Q^2}{M^2}) = \frac{m_f^2}{1-\beta}. \quad (11)$$

The result (10) holds for the large-mass continuum, $M^2 \gg 4m_f^2$. As it has been shown in [8], the typical transverse size in the $\gamma^* \to q\bar{q}$ transition vertex is $1/\overline{Q}$, see eq. (2). For excitation of heavy flavors and/or for light flavours at $1 - \beta \ll 1$ the hardness scale $\overline{Q}^2$ is large and one is in the legitimate pQCD domain.

Consequently, the contribution to the diffraction slope from the $\gamma^* X$ excitation vertex equals

$$\Delta B_{\gamma^* X} = \frac{1}{m_f^2} \frac{16\beta + 7\beta^2 + 78\beta^3 - 126\beta^4 - 5}{10(3 + 4\beta + 8\beta^2)} \quad (12)$$

which is a rigorous pQCD result for heavy flavours. Evidently, it is a scaling function of β which does not depend on Q^2, which nicely correlates with the scanning radius being a function of β only. It rises from $\beta \sim 1$ to $\beta \sim 1/2$ and decreases in the inclusive limit of $\beta \to 0$. It diminishes the diffraction slope at $\beta \sim 1$, which can be attributed to the s-channel helicity nonconserving spin-flip transitions.

One can readily evaluate $\Delta B_{\gamma^* X}$ for the both terms $\propto \mathbf{\Phi}_1^2$ and $m_f^2 \Phi_2^2$, we only comment here that for the both terms the β dependence of $\Delta B_{\gamma^* X}$ is very similar to that given by eq. (12). Even for heavy flavours, the contribution to $F_T^{D(4)}$ from $m_f^2 \Phi_2^2$ is a numerically small correction to the dominant contribution from $\propto \mathbf{\Phi}_1^2$. This correction is even smaller for lighter flavours. As it has been discussed in [16], the scale μ_G of variation of the unintegrated gluon density in the soft-to-hard transition region becomes more important than the mass m_f of light quarks. For this reason, for light flavour excitation the contribution from $m_f^2 \Phi_2^2$ will be suppressed $\propto m_f^2/\mu_G^2$. Furthermore, the scale for $\Delta B_{\gamma^* X}$ will be set by $1/\mu_G^2$ rather than by $1/m_f^2$. One of the consequences is that the change of $\Delta B_{\gamma^* X}$ from strange to up/down quarks is much weaker than $\propto 1/m_f^2$, see Fig. 2 where we show our numerical results.

Although for light flavours the magnitude of $\Delta B_{\gamma^* X}$ is no longer pQCD calculable, the behaviour of the unintegrated gluon density in the soft-to-hard transition region is reasonably well tested from earlier calculations [16] of the diffractive SF $F_T^{D(4)}$ which agree with the experiment, and also from the small-Q^2 behaviour of the proton structure function [17]. The emergence of this second scale has only a marginal impact on the β-dependence of B_D what we here are concerned about. We checked that variations of B_D calculated using different soft-to-hard interpolations of the gluon structure function as described in [16] do not exceed $\sim$ 1 GeV^{-2} with the retention of the form of the β dependence of B_D.

In contrast to the scaling β dependence of B_D for finite β, for diffractive DIS into near-threshold small masses, $M^2 \sim m_V^2 \sim 4m_f^2$, i.e., for $1-\beta \propto M^2/Q^2 \ll 1$, we predict a strong M^2 dependence of the diffraction slope. The near-threshold region belongs to the pQCD domain even for light flavour excitation, because here the QCD hardness scale is large, $\overline{Q}^2 \approx 1/4(Q^2+m_V^2)$(for finite Q^2 and/or heavy flavours one must bear in mind the kinematical threshold $\beta \leq \beta_{th} = Q^2/(Q^2+4m_f^2) < 1$). The plane wave description of final states holds for the quark-antiquark relative velocity $v \gtrsim \alpha_S(\overline{Q}^2)$. In this case the small-v^2 expansion of diffractive SF is

$$F_T^{D(4)}(t, x_{\mathbf{IP}}, v, Q^2) = \frac{128\pi e_f^2}{3\sigma_{tot}(pp)} \frac{m_f^2}{Q^4} v \left[1 + \frac{\Delta^2}{6m_f^2} v^2\right] \left[\alpha_s(\overline{Q}^2) G(x_{\mathbf{IP}}, \overline{Q}^2)\right]^2 \exp\left(-B_{\mathbf{IP}} \mathbf{\Delta}^2\right). \tag{13}$$

The principal effect is that the diffraction slope *decreases* with the increase of v^2 and/or M^2:

$$\Delta B_{\gamma^* X} = -v^2/6m_f^2 . \tag{14}$$

Here for heavy flavours $v^2 = 1 - 4m_f^2/M^2$, for light flavours it only makes sense to speak of the continuum above the ground-state $1S$ vector mesons (ρ^0, ω, ϕ^0) and v^2 must be understood as $v^2 \sim 1 - m_V^2/M^2$. Consequently, for the small-mass continuum we predict very rapid variations of the diffraction slope B_D, see fig. 2, and here the relevant mass scale is m_V^2. The principal point is that B_D drops substantially, we leave open the scenario in which B_D becomes negative valued, i.e., there will be a forward dip, in a certain range of masses.

In the spirit of duality for diffractive DIS [18], diffraction excitation of the small-mass continuum above the $1S$ ground state vector meson is dual to production of radial excitations of vector mesons. Then, our finding of the near-threshold decrease of the diffraction slope with rising M^2 correlates nicely with the prediction that for the $V'(2S)$ states the diffraction slope is substantially smaller than for the ground state vector mesons $V(1S)$, which follows from the node effect [6]. The near-threshold drop of B_D is smaller for heavy flavours, in a nice conformity with the weaker node effect in diffractive production of heavy quarkonia.

The similar analysis can be repeated for the longitudinal diffractive structure function. Although it is of higher twist, it dominates diffractive DIS at $\beta \gtrsim 0.9$ [18, 16]. As far as the diffraction slope is concerned, the QCD hardness scale for diffraction excitation of longitudinal photons is large, $\overline{Q}^2 \approx \frac{1}{4\beta} Q^2$, the corresponding scanning radius is small and we expect $B_D \approx B_{3\mathbf{IP}}$.

We conclude with a somewhat academic observation on a sum rule for the M^2 integrated cross section of diffractive excitation of heavy $q\bar{q}$ pairs by transverse photons.

Namely, if one neglects the β dependence of the QCD hardness scale $\overline{Q}^2$ in (10), then one readily finds that for the M^2-integrated diffractive cross section $\Delta B_{\gamma^* X} = 0$ and $B_D = B_{3\mathbf{P}}$. Indeed, a closer inspection of the calculation of the M^2 integrated cross section shows that to the accuracy Δ^2 the dependence on $\mathbf{\Delta}$ can be eliminated by the change of the integration variable $d^2\mathbf{k} \to d^2\mathbf{r}$. One can trace the origin of this sum rule to a QCD gauge invariance properties of (5), it serves as a useful cross check of correspon-ding polynomial coefficients. This sum rule is of little practical value, though, because for the dominant excitation of light flavours $\overline{Q}^2$ is small, in the soft-to-hard transition region of a strong variation of the gluon structure function $G(x_{\mathbf{P}}, \overline{Q}^2)$ and the above outlined derivation is not applicable.

To summarize, we presented predictions from the standard two-gluon pomeron exchange mechanism for the forward cone in diffractive DIS. For the high-mass continuum excitation we predict that the diffractive slope B_D is a scaling function of β which has a counterintuitive rise from small masses to $M^2 \sim Q^2$, which has no analogue in diffraction of real photons and/or hadrons. For the small-mass continuum we predict a rapid variation of B_D with M^2 on the scale m_V^2 and a sharp drop of B_D for a small-mass continuum above the vector meson excitation. These predictions can be tested at HERA.

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