THE FORWARD CONE AND L/T SEPARATION IN DIFFRACTIVE DIS

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LPS provides access to new fundamental observables: the diffraction cone and azimuthal asymmetries. Diffraction cone has a unique rise of $B_T$ from the exclusive limit to excitation of continuum $M^2 \approx Q^2$ which is in striking contrast to experience with real photoproduction and hadronic diffraction. Azimuthal asymmetry is large and pQCD calculable at large $\beta$ and can be measured with LPS. It allows testing of the pQCD prediction of $L/T \gg 1$.

1 Helicity components of diffractive DIS

The detection of leading protons $p'$ from diffractive DIS $ep \rightarrow e'p'X$ gives access to several new observables: the diffraction slope $B_D$ which quantifies the impact parameter properties of diffractive DIS and new helicity structure functions. In this talk we report predictions for the nontrivial $\beta$-dependence of the diffraction slope and suggest a new method for measuring $R_D = d\sigma_D^L/d\sigma_D^T$ for diffractive DIS based on the azimuthal correlation of the $(e,e')$ and $(p,p')$ scattering planes.

The differential cross-section of the diffractive process $ep \rightarrow e'p'X$ reads

$$Q^2 y \frac{d\sigma}{dQ^2 dy dM^2 dp'^2} \frac{d\phi}{d\phi} = \frac{\alpha_{em}}{2\pi} \left[ \frac{1}{2} \left( 1 - y + \frac{y^2}{2} \right) \frac{d\sigma_D^T}{dM^2 dp'^2} + (1 - y) \frac{d\sigma_D^L}{dM^2 dp'^2} + (1 - y) \frac{d\sigma_D^T}{dM^2 dp'^2} \cdot \cos (2\phi) + (2 - y) \sqrt{1 - y} \frac{d\sigma_D^L}{dM^2 dp'^2} \cdot \cos (\phi) \right],$$

(1)

where $p_\perp$ is the $(p,p')$ momentum transfer, $\phi$ is the azimuthal angle between $(e,e')$ and $(p,p')$ scattering planes.

We focus on the $q\bar{q}$ excitation which dominates at large $\beta$. Following the technique developed in [3], we find (for the kinematical variables see Fig. 1)

$$\frac{d\sigma_D^L}{dM^2 dp'^2} = \frac{\alpha_{em}}{24\pi^2} \sum_f Z_f^2 \int d^2k \frac{1 - J^2}{4J} \alpha_5^2 [h_i(z_+) + h_i(z_-)],$$

(2)

where $z_\pm = \frac{1}{2}(1 \pm J)$, $J = \sqrt{1 - 4\frac{\mu^2 + m_f^2}{M^2}}$, $h_T = [1 - 2z(1 - z)]\Phi_T^2 + m_f^2 \Phi_T^2$, $h_L \cdot \cos (\phi) = 2z(1 - z)(1 - 2z)Q(\Phi_1^2)\Phi_2$, $h_L = 4z^2(1 - z)^2Q^2\Phi_2^2$, $h_T \cdot \cos (2\phi) = 2z(1 - z)[\Phi_1^2 - 2(\Phi_1^2)^2]$ and $m_f$ is the quark mass.
The helicity amplitudes $\Phi_1$, $\Phi_2$ equal

$$\Phi_j = \int d^2 \vec{k} f(x_{\text{IP}}, \vec{k}, \vec{p}_\perp) \left[ D_j(\vec{r}+\vec{k}) + D_j(\vec{r}-\vec{k}) - D_j(\vec{r}+\vec{p}_\perp/2) - D_j(\vec{r}-\vec{p}_\perp/2) \right],$$

(3)

where $\vec{D}_1(\vec{r}) = \vec{r} \cdot \vec{D}_2(\vec{r}) = \vec{r}/(\vec{r}^2 + \epsilon^2)$, $\vec{r} = \vec{k} - (1/2 - z)\vec{p}_\perp$, and $f(x_{\text{IP}}, \vec{k}, \vec{p}_\perp)$ is the off-forward unintegrated gluon density. Following (3), (4), (5) we find

$$\Phi_j \propto \int \bar{Q}^2 d^2 \vec{k} f(x_{\text{IP}}, \bar{Q}^2, p_\perp^2) = G(x_{\text{IP}}, \bar{Q}^2, p_\perp^2) \cdot (1 - \frac{1}{2} B_{3\text{IP}} p_\perp^2).$$

(4)

where the pQCD hardness scale equals

$$\bar{Q}^2 = (k^2 + m_f^2) \left( 1 + \frac{Q^2}{M^2} \right) = \frac{k^2 + m_f^2}{1 - \beta}.$$  

(5)

In Eq. (4) we parameterize the small-$p_\perp^2$ dependence by the diffraction slope $B_{3\text{IP}}$ which comes from the proton vertex and gluon propagation effects (see Fig. 1c). We emphasize that $B_{3\text{IP}}$ depends neither on $\beta$ nor flavor.

For excitation of heavy flavours we have the fully analytic results (for the discussion of twist-4 $F_T$ and $F_{TT'}$ see (4), (5))

$$F_T^{D(4)} = \frac{2\pi e_f^2}{9\sigma_{\text{tot}}^{pp}} \frac{\beta(1 - \beta)^2}{m_f^2} \left[ (1 - B_{3\text{IP}} p_\perp^2)(3 + 4\beta + 8\beta^2) 
+ p_\perp^2 \frac{1}{m_f^2} \frac{1}{10} (5 - 16\beta - 7\beta^2 - 78\beta^3 + 126\beta^4) \bar{G}_T^2 \right],$$

$$F_L^{D(4)} = \frac{2\pi e_f^2}{9\sigma_{\text{tot}}^{pp}} \frac{12\beta^3}{Q^2} \left[ (1 - B_{3\text{IP}} p_\perp^2)2(1 - 2\beta)^2 \bar{G}_L^2 \right].$$

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\[ F_{LT}^{(4)} = \frac{p_{\perp}^2}{Q} \cdot \frac{2\pi e_0^2 \beta^2 (1 - \beta)}{9\sigma_{\text{tot}}^2} \left[ (1 - \beta + 12\beta^2 + 672\beta^4) - \frac{m_f^2}{20} (1 - \beta) (1 - 7\beta + 23\beta^2 - 21\beta^3) G_T^2 \right] \]

where \( G_{T,L} = \alpha_s(Q^2_{T,L}) G(x_{pY}, Q^2_{T,L}) \) and \( Q^2_T = m_f^2, Q^2_L = \frac{Q^2}{4\beta^2} \). The principal point is that \( F_T \) and \( F_{LT} \) are dominated by the aligned jet configurations, \( k^2 \sim m_f^2 \), whereas \( F_L \) comes from the large \( k^2 \) jets, \( k^2 \sim M^2/4 \). In our calculations for light flavours we use the parameterization for the soft glue from 6.

2 Azimuthal asymmetry and \( L/T \) separation

In contrast to the inclusive DIS where \( R = \sigma_L/\sigma_T \ll 1 \), in diffractive DIS pQCD predicts \( R^D \gg 1 \) for \( \beta > 0.9 \), despite the fact that \( F_L \) is of higher twist. Because neither the proton nor electron energy will be changed at HERA, one must exploit the azimuthal asymmetry \( A_{LT}^{(4)} = F_{LT}^{(4)}/(F_T^{(4)} + F_L^{(4)}) \). The key observation is that both \( F_{LT} \) and \( F_T \) come from the aligned jet configurations and the \( LT/T \) ratio is model independent

\[ R_{LT/T}^{(4)} = \frac{F_{LT}^{(4)}}{F_T^{(4)}} = \frac{p_{\perp}}{Q} \cdot \frac{12\beta^3 (2 - 3\beta)}{(1 - \beta) (3 + 4\beta + 8\beta^2)}. \]

Consequently, the measurement of \( A_{LT}^{(4)} \) amounts to the measurement of \( R_{LT/T}^{(4)} \equiv F_{LT}^{(4)}/F_T^{(4)} = R_{LT/T}^{(4)} A_{LT}^{(4)} - 1 \). The predicted asymmetry is quite substantial in the interesting region of \( \beta \sim 0.9 \) (Fig. 2), can be measured with the ZEUS and H1 leading proton spectrometers and one can test the pQCD result \( R^D \gg 1 \) experimentally.

3 Peculiarities of the diffraction cone for diffractive DIS

The diffraction slope \( B_D \) is defined by the formula \( d\sigma(a p \rightarrow XY) \propto \exp(-B_D p_{\perp}^2) \). The experience with diffraction of hadrons and real photons can be summarized as follows. One can write down an essentially model-independent decomposition \( B_D = \Delta B_{aX} + \Delta B_{pY} + \Delta B_{\text{int}} \) in which \( \Delta B_{\text{int}} \) interaction (exchange) range, and \( \Delta B_{aX} \) and \( \Delta B_{pY} \) come from the (transverse) size of the \( aX \) and \( pY \) transition vertices. These contributions \( \Delta B_{aX,pY} \) depend strongly on the
excitation energy in the $i \to j$ transition $\Delta M^2 = m_j^2 - m_i^2$ and vanish for excitation of the continuum. $\Delta M^2 \gg 1 \div 2\text{GeV}^2$.

The experimental data on double diffraction dissociation $pp \to XY$ into high mass states $X, Y$ give $B_D \approx \Delta B_{int} \sim 1.5 - 2\text{GeV}^{-2}$. In single diffraction $ap \to Xp'$ into high mass states $B_D \approx \Delta B_{int} + \Delta B_{pp} \sim 6 - 7\text{GeV}^{-2}$ is independent of the projectile $a = p, \pi, K, \gamma$. It has been argued that in the triple-pomeron regime of diffractive DIS, $\beta \ll 1$, one must find $B_D \approx B_{3\mathbf{p}} \sim 6\text{GeV}^{-2}$, which has indeed been confirmed by the ZEUS collaboration.

Hereafter we focus on finite $\beta$, dominated by the $q\bar{q}$ excitation, $X = q\bar{q}$. For diffractive DIS at finite $\beta$ the excited mass is large, $M^2 = \frac{1-\beta}{\beta} Q^2 \gg m_i^2$, hence the continuum is excited and naively one would expect $\Delta B_{\gamma^*X} \approx 0$, and $B_D \approx B_{3\mathbf{p}}$ independently of $\beta$. Our principal finding is that this is not the case, $\Delta B_{\gamma^*X}$ is large and varies substantially with $\beta$.

Our results for the small-$p_T^2$ of diffractive structure functions are given by Eqs.(6). We focus on the transverse cross section which dominates at $\beta < 0.9$. The component $\Delta B_{\gamma^*X}$ comes from the term $\propto p_T^2 m_i^2$. These formulas are directly applicable for heavy flavour excitation. In the diffraction excitation of light flavours there is a sensitivity to the gluon structure function in the soft region, and the rate of variation of the gluon structure function in the soft region emerges as a scale instead of $\frac{1}{m_i^2}$. However, the qualitative features of the $\beta$ dependence do not change from heavy to light flavours. The numerical results are shown in Fig. 3.

The most striking prediction is the rise of $B_D$ when $\beta$ decreases from $\beta \sim 1$ to $\beta \sim \frac{1}{2}$. This rise can be related to the rise of the scanning radius discussed in Ref: $r_S^2 \sim \frac{1}{Q_T^2} \sim \frac{1-\beta}{m_i^2}$. Numerically, at $\beta \sim 1/2$ we have $\Delta B_{\gamma^*X} \sim \frac{1}{10 m_i^2}$.
Figure 3: Our predictions for the diffraction slope $B_D$ for the transverse structure function at $Q^2 = 100\text{GeV}^2$ and $x_{\text{IP}} = 0.001$.

In excitation of small masses, $4m_f^2 \ll M^2 \ll Q^2$, i.e., $\beta \to 1$, we predict a substantial drop of $B_D$, because here $\Delta B_{\gamma^*X} \sim -\frac{1}{4m_f}$. This is a legitimate pQCD domain because $\bar{Q}_T^2 = \frac{Q^2}{4}$ for $M^2 \sim 4m_f^2$. Very close to the threshold

$$F_T^{D(\gamma)}(p_T^2, x_{\text{IP}}, v, Q^2) = \frac{128\pi e_f^2}{3\sigma_{\text{tot}}^f} \frac{m_f^2}{Q^4} v^2 (1 - B_{3\text{IP}}^2) + \frac{p_T^2}{6m_f^2} v^2 |G_7^2|,$$

where $v = \sqrt{1 - \frac{4m_f^2}{M^2}}$. In the spirit of exclusive-inclusive duality, $B_D = B_{3\text{IP}}$ can be related to the diffraction slope for $V(1S)$ vector meson production, whereas the drop of $B_D$ for somewhat higher masses correlates nicely with the prediction $B_{V(2S)} \ll B_{V(1S)}$. The experimental observation of this large-$\beta$ drop of $B_D$ is not easy because of the masking effect of the longitudinal cross section for which $B_D \approx B_{3\text{IP}}$.

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