Charged Current Diffractive Structure Functions

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Abstract

We present our study of the diffraction in charged current DIS. We analyse the perturbatively tractable excitation of heavy quarks, emphasizing the peculiarities of the Regge factorization breaking in excitation of open charm.

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The study of diffractive electromagnetic Deep Inelastic Scattering in QCD has rapidly advanced in last years ([1,10], for a recent review see [11]). Successful quantitative predictions

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for diffractive DIS are now available [4,8,11], the pQCD factorization scales for different components of the diffractive Structure Function have been derived [6,7,9], breaking of the diffractive factorization and of the DGLAP evolution for diffractive SF's have been understood and a large perturbative intrinsic charm in the pomeron has been established [4,6,7], the prediction of large perturbative higher twist contributions has been obtained [7,8].

The study of diffraction in charged current (CC) DIS, $ep \to \nu p' X$, will permit a further test of this approach. Rapidity gap events in CC DIS have already been observed at HERA [12] and with amassing more data on CC DIS a detailed comparison between the experiment and models for diffractive DIS will be possible [13]. Because of the parity non-conservation, in the CC case one has a larger variety of diffractive SF's compared to the neutral current electromagnetic (EM) case. To the lowest order in pQCD, CC diffractive DIS proceeds by the Cabibbo-favoured excitation of the $(u\bar{d})$ and $c\bar{s}$ dijet states. The unequal mass $(c\bar{s})$ is a particularly interesting testing ground where to investigate diffractive factorization breaking. A self-tagging property of charm jets gives better access to various diffractive structure functions, for instance, to $F_3^{D(3)}$.

The subject of this proceeding is the derivation of the above stated features of CC diffractive DIS and the discussion of its differences respect to EM diffractive DIS. For convenience we focus on the process $e^+p \to \bar{\nu} p' X$ already observed by ZEUS [12]. The discussion and results can easily be translated to the $e^-p \to \nu p' X$ process.

We start with necessary definitions. In diffractive CC e^+p scattering the experimentally measured quantity is the five-fold differential cross section $d\sigma^{(5)}(ep \to \nu p' X)/dQ^2 dx dM^2 dp_\perp^2 d\phi$. Here X is the diffractive state of mass M, p' is the secondary proton with the transverse momentum $\vec{p}_\perp$, $t = -\vec{p}_\perp^{\,2}$, ϕ is the angle between the (e, e') and (p, p') planes, $Q^2 = -q^2$ is the virtuality of the W^+ boson, $x, y, x_{\mathbf{IP}}$ and $\beta = x/x_{\mathbf{IP}}$ are the standard diffractive DIS variables.

The underlying subprocess is diffraction excitation of the W^+ boson, $W^+p \to p'X$. In the parity conserving EM DIS, the exchanged photon have either longitudinal (scalar), $s = \frac{1}{Q}(q_+n_+ - q_-n_-)$ or transverse, in the (e, e') plane, polarization t_μ (here $n_\pm$ are the usual light–cone vectors, $n_+^2 = n_-^2 = 0$, $n_+n_- = 1, q = q_+n_+ + q_-n_-$ and $q.s = q.t = 0$). In the parity-nonconserving CC DIS, the exchanged W^+ bosons have also the out-of-plane linear polarization $w_\mu = \epsilon_{\mu\nu\rho\sigma} t_\nu n^+_\rho n^-_\sigma$. We introduce also the usual transverse metric tensor $\delta^\perp_{\mu\nu} = \delta_{\mu\nu} + n^-_\mu n^+_\nu + n^+_\mu n^-_\nu = -t_\mu t_\nu - w_\mu w_\nu$. Then, the polarization state of the W^+ is described by the leptonic tensor

$$\frac{L_{\mu\nu}}{4} = \frac{2Q^2}{y^2}\Big[-\frac{1}{2}\delta^\perp_{\mu\nu}(1-y+\frac{1}{2}y^2) + \frac{1}{2}(1-y)(t_\mu t_\nu - w_\mu w_\nu) + (1-y)s_\mu s_\nu$$
$$+\frac{1}{2}(1-\frac{1}{2}y)\sqrt{1-y}(s_\mu t_\nu + s_\nu t_\mu) + \frac{i}{2}y(1-\frac{1}{2}y)(w_\mu t_\nu - w_\nu t_\mu) + \frac{i}{2}y\sqrt{1-y}(w_\mu s_\nu - s_\mu w_\nu)\Big] \quad (1)$$

which, upon contraction with the hadronic tensor leads to 6 different components for $d\sigma^{(3)}_i(W^+p \to p'X)/dM^2 dt d\phi$ ($i = T, L, TT', LT, 3$ and $LT(3)$):

$$\frac{y d\sigma^{(5)}(e^+p \to \bar\nu p'X)}{dQ^2 dy dM^2 dp^2_\perp d\phi} = \frac{G_F M^2_W Q^2}{4\sqrt{2}\pi^2(M^2_W+Q^2)^2}\Big\{(1-y+\frac{1}{2}y^2)\cdot d\sigma^{D(3)}_T - y(1-\frac{1}{2}y)\cdot d\sigma^{D(3)}_3$$
$$+(1-y)\cdot d\sigma^{D(3)}_L + (1-y)\cos 2\phi \cdot d\sigma^{D(3)}_{TT'}$$
$$+(1-\frac{1}{2}y)\sqrt{1-y}\cos\phi\cdot d\sigma^{D(3)}_{LT} - y\sqrt{1-y}\sin\phi\cdot d\sigma^{D(3)}_{LT(3)}\Big\}/dM^2 dp^2_\perp d\phi\,, \quad (2)$$

where G_F is the Fermi coupling, m_W is the mass of the W-boson. One can then define the dimensionless diffractive structure functions $F^{D(4)}_i$,

$$\frac{(Q^2+M^2)d\sigma^{(3)}_i(W^+p\to p'X)}{dM^2 dp^2_\perp} = \frac{\pi G_F M^2_W Q^2}{\sqrt{2}(Q^2+M^2_W)^2}\cdot\frac{\sigma^{pp}_{tot}}{16\pi}\cdot F^{D(4)}_i(p^2_\perp, x_{\mathbf{IP}}, \beta, Q^2)\,, \quad (3)$$

It is also useful to introduce the t–integrated SF's †

$$F^{D(3)}_i(x_{\mathbf{IP}}, \beta, Q^2) = \frac{\sigma^{pp}_{tot}}{16\pi}\int dp^2_\perp F^{D(4)}_i(p^2_\perp, x_{\mathbf{IP}}, \beta, Q^2)\,. \quad (4)$$

The diffractive SF's $F^{D(3)}_T, F^{D(3)}_L$ and $F^{D(3)}_3$ are counterparts of the familar $F_T = F_2 - F_L, F_L$ and F_3 for inclusive DIS of neutrinos, $F^{D(3)}_3$ and $F^{D(3)}_{LT(3)}$, are C- and P-odd and vanish in EM scattering. In the following we will not consider the azimuthal angle-dependent terms, but will limit our discussion to $F^{D(3)}_T, F^{D(3)}_L$ and $F^{D(3)}_3$.

Only relatively large $x \sim 10^{-2}$ are easily accessible in CC DIS [12,13]. As the experimental selection of diffractive events requires $x_{\mathbf{IP}}$ <0.05-0.1, the kinematical relation $\beta = x/x_{\mathbf{IP}}$ implies that the experimentally observed CC diffractive DIS will proceed at rather large β, dominated by the partonic subprocess $W^+p \to (u\bar d)p', (c\bar s)p'$.

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†The ZEUS/H1 definition [14] corresponds to $F^{D(3)}_2(H1/ZEUS) = \frac{1}{x_{\mathbf{IP}}}F^{D(3)}_2$, which blows up at $x_{\mathbf{IP}} \to 0$, and is much less convenient than our $F^{D(3)}_2$.

We focus on the $c\bar{s}$ excitation, analogous considerations apply to $u\bar{d}$. In what follows, z and $(1-z)$ are the fractions of the (light–cone) momentum of the W^+ carried by the charmed quark and strange antiquark respectively, $\vec{k}$ is the relative transverse momentum in the $q\bar{q}$ pair. The invariant mass of the di-jet final states equals

$$M^2 = \frac{k^2+\mu^2}{z(1-z)}\,, \tag{5}$$

where $\mu^2 = (1-z)m_c^2 + zm_s^2$. $m_c, (m_s)$ being the charmed (strange) quark mass. All SF's are calculable in terms of the same quark helicity changing and conserving amplitudes $\vec{\Phi}_1$ and Φ_2 introduced in [1,2]. Combining the formalism of [2] with the treatment of charm lepto-production in [16], we obtain (integration over the azimuthal orientation of $\vec{k}$ is understood, $\alpha_{cc} = \frac{G_F M_W^2}{2\pi\sqrt{2}}$)

$$\left.\frac{d\sigma^D_{L,T}}{dz dk^2 dt}\right|_{t=0} = \frac{\pi^2\alpha_{cc}}{3}\alpha_S^2(\bar{Q}^2)\left[A_{L,T}(z,m_s,m_c)\vec{\Phi}_1^2 + B_{L,T}(z,m_s,m_c)\Phi_2^2\right]\,, \tag{6}$$

(see later for a definition of $\bar{Q}^2$), where

$$A_T(z) = [1-2z(1-z)]\,, \tag{7}$$

$$B_T(z,m_s,m_c) = \left[m_c^2 - 2z(1-z)m_c^2 - z^2\Delta^2\right]\,, \tag{8}$$

$$A_3(z) = [2z-1]\,, \tag{9}$$

$$B_3(z,m_s,m_c) = \left[z^2 m_s^2 - (1-z)^2 m_c^2\right]\,, \tag{10}$$

$$A_L(z,m_s,m_c) = \frac{1}{Q^2}(m_s^2+m_c^2)\,, \tag{11}$$

$$B_L(z,m_s,m_c) = 4Q^2 z^2(1-z)^2 + 4z(1-z)\mu^2 + \frac{1}{Q^2}[\mu^4 + m_c^2 m_s^2] \tag{12}$$

with $\Delta^2 = m_c^2 - m_s^2$.

The amplitudes $\vec{\Phi}_1$ and Φ_2 were derived in [1,2]:

$$\vec{\Phi}_1 = \vec{k}\int\frac{d\kappa^2}{\kappa^4}f(x_{\mathbf{IP}},\kappa^2)\left[\frac{1}{k^2+\varepsilon^2} - \frac{1}{\sqrt{a^2-b^2}} + \frac{2\kappa^2}{a^2-b^2+a\sqrt{a^2-b^2}}\right]$$
$$\approx \frac{2\vec{k}\varepsilon^2}{(k^2+\varepsilon^2)^3}\int\frac{d\tau}{\tau}W_1(\omega,\tau)G(x_{\mathbf{IP}},\tau\bar{Q}^2)\,, \tag{13}$$

$$\Phi_2 = \int\frac{d\kappa^2}{\kappa^4}f(x_{\mathbf{IP}},\kappa^2)\left[\frac{1}{\sqrt{a^2-b^2}} - \frac{1}{k^2+\varepsilon^2}\right]$$
$$\approx \frac{k^2-\varepsilon^2}{(k^2+\varepsilon^2)^3}\int\frac{d\tau}{\tau}W_2(\omega,\tau)G(x_{\mathbf{IP}},\tau\bar{Q}^2)\,, \tag{14}$$

Here $f(x,\kappa^2) = \partial G(x,k^2)/\partial \log(\kappa^2)$ $(G(x,Q^2) = xg(x,Q^2))$ is the unintegrated gluon SF of the proton, $\varepsilon^2 = z(1-z)Q^2 + zm_s^2 + (1-z)m_c^2, a = \varepsilon^2 + k^2 + \kappa^2, b = 2k\kappa$ and

$$\bar{Q}^2 = \varepsilon^2 + k^2 = \frac{k^2 + zm_s^2 + (1-z)m_c^2}{1-\beta}\,, \tag{15}$$

emerges as the pQCD factorization scale. A second representation, in which

$$W_1(\omega,\tau) = \frac{(1+\omega^2)}{2\tau(1+\omega)}\left\{\left[1 - \frac{\Theta + \tau^2[1+(1-\omega)/\tau(1+\omega)]}{(1+w)\Theta^{3/2}}\right]\right.$$
$$\left. + \frac{4\tau^2(1+\tau)}{\Theta[\Theta^{1/2}+(1+\tau)]^2}\left[1 + \frac{1+\tau(1+\omega) - 2\omega\tau/(1+\tau)}{(1+\omega)\Theta^{1/2}}\right]\right\} \tag{16}$$

$$W_2(\omega,\tau) = \frac{(1+\omega)}{\tau(1-\omega)}\left[1 - \frac{\Theta + \tau^2[1+(1-\omega)/\tau(1+\omega)}{\Theta^{3/2}}\right] \tag{17}$$

where

$$\Theta = (1+\omega)^2\left[(1+\tau)^2 - \frac{4\omega\tau}{(1+\omega)}\right]. \tag{18}$$

$\tau = k^2/\bar{Q}^2$ and $\omega = k^2/\varepsilon^2$, is more convenient for the practical calculations based on available parameterization for the gluon SF of the proton $G(x,Q^2)$. Here the weight functions $W_i(\omega,\tau)$ have a narrow peak at $\tau \approx 1$ with the unit area under the peak, which gives the Leading LogQ^2 result [2,6]

$$\int \frac{d\tau}{\tau} W_i(\omega,\tau) G(x_{\mathbf{IP}}, \tau\bar{Q}^2) \approx G(x_{\mathbf{IP}}, \bar{Q}^2)\,, \tag{19}$$

valid for sufficiently large values of ω, which is equivalent to sufficiently large $\beta \gtrsim 0.1$ (of interest in the present study). For applications of the above formalism to the EM case see [6,7,9,10]. Also, despite the somewhat different appearance, the so-called soft color interaction model by Buchmüller et al. [17] is essentially identical to the above described picture of diffractive DIS (for more discussion on that see [11]).

At variance with the equal mass EM case, where $\bar{Q}^2 = (k^2 + m^2)/(1-\beta)$, now the factorization scale depends on z and then one expects different cross sections whether the charmed quark is produced in the forward (F) or the backward (B) hemisphere, with respect to the W momentum, in the rest frame of the diffractive state X. The two configurations differ by the value of the light-cone variable $z_{F,B} = \frac{1}{2}(1+\delta)\left[1 \pm \sqrt{1-4\frac{k^2+m_c^2}{M^2(1+\delta)^2}}\right]$, where we have introduced for brevity the variable $\delta = \frac{\Delta^2}{Q^2}\frac{\beta}{(1-\beta)}$. The pQCD scale is perturbatively large for large β even for light flavours, and for the charm component of the diffractive SF it is large for all β (see below).

Evaluating the light quark component of the diffractive SF at not so large β, one needs a model for the small-Q^2 behaviour of the gluon structure function $G(x, Q^2)$: in the following we will use the same form used in Ref. [8], which at large k^2 coincide with the GRV NLO parameterization [20].

For what concerns quark masses we assume $m_c = 1.5$ GeV, $m_s = 0.3$ GeV and $m_{u,d} = 150$ MeV. Variations of the charm mass by 10% have a little effect on the predicted SF, apart from shifting the threshold $\beta_c = Q^2/[Q^2 + (m_c + m_s)^2]$ (see below).

For evaluating $F_i^{D(3)}$ one needs to know the $p_\perp^2$ dependence of the diffractive cross section, which is usually parameterized as $d\sigma/dp_\perp^2 \propto \exp(-B_D p_\perp^2)$. As it was shown in [18,11], for heavy flavour excitation, for the perturbative transverse higher twist and longitudinal contributions $B_D \sim 6\text{GeV}^{-2}$, for light flavour contribution the diffraction slope B_D exhibits, at not so large β, a slight β-dependence, but for the purposes of this present exploratory study we can simply take $B_D(ud) = 9$ GeV^{-2}.

Various studies of diffractive DIS (as DIS off pomerons in the proton) have assumed diffractive factorization. The latter is not supported by QCD studies [4,6], and the present study of charm excitation in CC diffractive DIS offers more evidence to this effect. Nevertheless, wherever that would not lead to a grave confusion, we will speak of the perturbative intrinsic partons in the pomeron.

Separation of the pQCD subprocess of $W^+ \to c\bar{s}$ into the excitation of charm on the perturbative intrinsic strangeness in the pomeron and excitation of (anti)strangeness on the intrinsic (anti)charm is not unambiguous and must be taken with the grain of salt. In the naive parton model, in the former process charmed quark will carry the whole momentum of the W^+ and be produced with $z \approx 1$. In contrast, in the latter process, it is the strange antiquark which carries the whole momentum of W^+ and charmed quark is produced with $z \approx 0$, which suggests $z > \frac{1}{2}$ and $z < \frac{1}{2}$ as a compromise boundary between the two partonic subprocesses. However, the full fledged pQCD calculation leads to broad z distributions (for a related discussion of definition of the strangeness and charm density in $\nu N, \bar{\nu} N$ inclusive DIS see [16]). As a purely operational definition, we stick to a parton model decomposition $F_T^{D(3)} = F_{T(s)}^{D(3)} + F_{T(\bar{c})}^{D(3)}$ and $F_3^{D(3)} = F_{T(s)}^{D(3)} - F_{T(\bar{c})}^{D(3)}$, which is a basis for the results shown in Fig. 1.

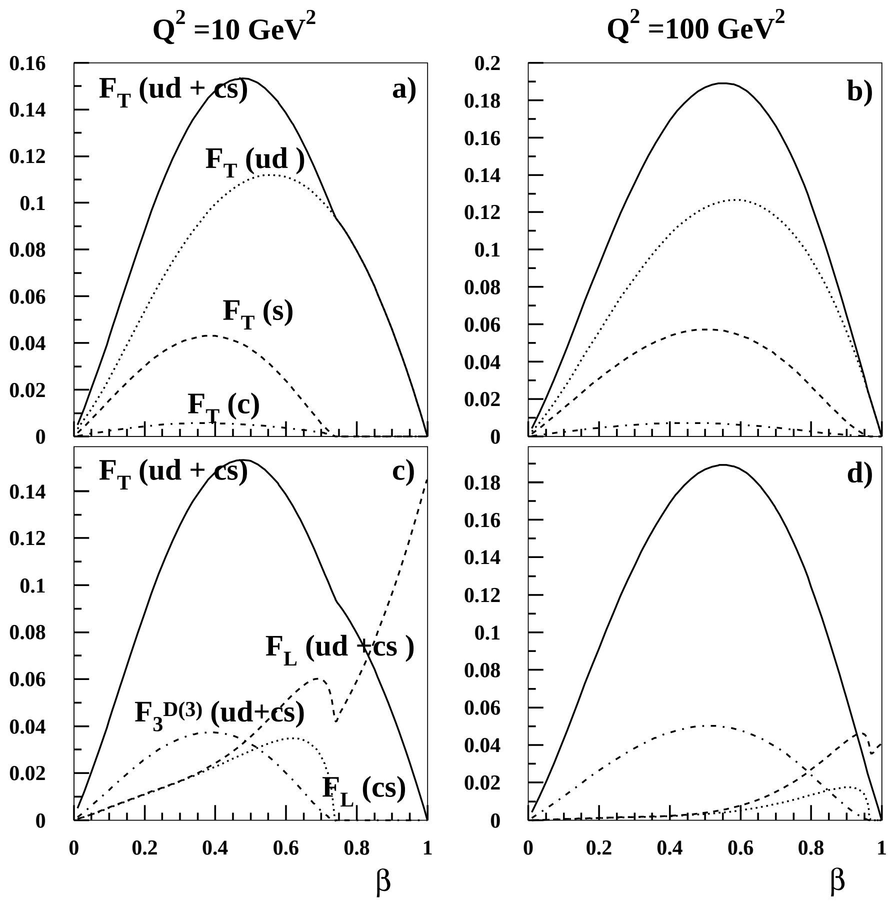

FIG. 1. The β dependence for $x_{\mathbb{P}} = 10^{-3}$: a) $F_T(u,d,s,c)$ [solid], $F_T(u,d)$ [dotted], $F_T(s)$ [dashed], $F_T(c)$ [dot–dashed] at $Q^2 = 10\text{GeV}^2$ b) the same as above for $Q^2 = 100\text{GeV}^2$ c) All flavours F_T [solid], F_3 [dot–dashed], F_L [dashed] and $F_L(cs)$ [dotted] at $Q^2 = 10\text{GeV}^2$ d) the same as above for $Q^2 = 100\text{GeV}^2$

With this definition, excitation of the charmed quark off the intrinsic strangeness, $F_{T(s)}^{D(3)}$, comes from terms $\propto z^2$ in (7),(8) and (9),(10). It is dominated by the forward production of charm w.r.t. the momentum of W^+ in the rest frame of the diffractive system X, but receives certain contribution also from $z < \frac{1}{2}$. Similarly, $F_{T(\bar{c})}^{D(3)}$ come from terms $\propto (1-z)^2$, is dominated by the forward production of strangeness (the backward production of charm), but still receives certain contribution from the forward charm production.

All the considerations of Ref. [7,8] for the longitudinal and transverse diffractive SF in electroproduction are fully applicable to the CC case at $Q^2 \gg m_c^2$. We consider first the backward charm, $z \ll 1$, for which Eq.s (5, 15) give $z \approx (k^2+\mu^2)/M^2$ and $\bar{Q}^2 \approx (k^2+m_c^2)/(1-\beta)$. Expanding the brackets of Eqs.(13, 14), in Eq.(6) and using the approximation (19) the k^2-integration in (6) gives dominant contributions to the transverse SF coming from the low-k^2 region but without entering deeply in the nonperturbative region for the heavy quark production. For $M^2 \gg m_c^2$ one finds for the low scales dominated contribution (including the LT and the first HT):

$$F_{T(\bar{c})}^{D(4)} \approx \frac{4\pi}{3\sigma_{tot}^{pp}} \frac{\beta(1-\beta)^2}{6m_c^2(1+\delta)} \left\{(3+4\beta+8\beta^2) + \frac{m_c^2}{Q^2}\frac{4\beta}{1-\beta}\right. \tag{20}$$
$$. \times \left.\left[\frac{5}{4}\frac{\Delta^2}{m_c^2}(1+8\beta^2) - (1-2\beta+4\beta^2)\right]\right\} \left[\alpha_S(\bar{Q}_L^2) G(x_{\mathbf{IP}}, \bar{Q}_L^2 \simeq \frac{m_c^2}{(1-\beta)})\right]^2$$

As in the EM case, the large k^2 ($k^2 \sim M^2/4$) dominated contributions come from the second term in (7). They are calculable in pQCD and the twist expansion starts with twist-4:

$$F_{T(\bar{c})}^{D(4)} = -\frac{16\pi}{\sigma_{tot}^{pp}} \frac{2}{3}\frac{\beta^2(1-\beta)}{Q^2(1+\delta)} \left(\beta^2 + 2\frac{\Delta^2}{Q^2}\frac{\beta^2(2\beta-1)}{(1-\beta)}\right) \left[\alpha_S(\bar{Q}_H^2) G(x_{\mathbf{IP}}, \bar{Q}_H^2 \simeq \frac{Q^2}{4\beta})\right]^2 \tag{21}$$

In (20,21) emerge additional higher-twist corrections $\propto (\Delta^2/Q^2)^n$, in a first approach we restrict ourselves to the genuine contributions and its first higher twist corrections.

The higher twist corrections to F_T^D receive thus both contributions from the low-scale region and the large scale $\bar{Q}_H^2$. The first term in Eq.(21) is substantially the same which has been discussed in Ref. [8] for EM current and it remains relevant even at relatively large value of Q^2 as the $1/Q^2$ factor is partially compensated by the growth of gluon SF $G(x_{\mathbf{IP}}, \bar{Q}_H^2)$.

Let us now consider the longitudinal cross section. The most important contribution comes from the term $z^2(1-z)^2Q^2$ in the B_L expansion (12), which is identical to that in the EM case. The k^2 integrated cross section is completely dominated by the short-distance contribution from high-k^2 jets, $k^2 \sim \frac{1}{4}M^2$. Upon the k^2 integration, to a logarithmic accuracy, we find the twist expansion of the longitudinal SF for the terms dominated by the large scale:

$$F_L^{D(4)} = \frac{16\pi}{\sigma_{tot}^{pp}} \frac{\beta^3\ (2\beta-1)}{3\ Q^2(1+\delta)} \left((2\beta-1) + \frac{\Delta^2}{Q^2} \frac{\beta(5-6\beta)}{(1-\beta)} \right) \left[\alpha_S(\bar{Q}_H^2) G(x_{\mathbf{I\!P}}, \bar{Q}_H^2) \right]^2 \qquad (22)$$

As the pQCD factorization scale, $\bar{Q}_H^2$, does not depend on flavours we predict a restoration of the flavours symmetry at asymptotically large Q^2. Again the scaling violation factor, $G^2(x_{\mathbf{I\!P}}, \bar{Q}^2)$, in (22) compensates to a large extent the higher twist factor $\frac{1}{Q^2}$ and the longitudinal SF remains large, and takes over F_T^D, in a broad range of Q^2 of the practical interest, see Fig. 1.

In (22), the leading twist-4 term is the same as for NC diffractive DIS. However, in the CC diffractive DIS, because non-conservation of weak current, extra higher twist contributions to $F_L^{D(3)}$ come from the expansion of B_L (always substantially dominated by the perturbative region).

Further terms (both twist-4 and higher), come from the term $\propto A_L$ in (6). They receive large contributions from the low k^2 region. In particular they assume a strong relevance for the charm–strange component where terms $\propto m_c^2/Q^2$ appear. We find for the A_L contribution:

$$F_L^{D(4)}[A_L] \approx \frac{4\pi}{9\ \sigma_{tot}^{pp}} \frac{(m_c^2 + m_s^2)}{m_c^2} \frac{\beta(1-\beta)^2}{Q^2(1+\delta)} (1 + 2\beta + 3\beta^2) \left[\alpha_S(\bar{Q}_L^2) G(x_{\mathbf{I\!P}}, \bar{Q}_L^2) \right]^2 \qquad (23)$$

Whereas the components coming from A_L is low scale dominated, it gives comparable contributions as the leading twist-4 in the small Q^2 region. The overall dependence of $F_L^{D(3)}(cs)$ on Q^2 and its decomposition in the A_L and B_L components are shown in Fig. 2.

Due to the symmetry $z, (1-z)$ of Eqs.(5, 7 - 12), one finds similar results subject to the replacement $m_c \to m_s$ for the forward production of charm (F_s^D) at $1 - z \lesssim m_s^2/m_c^2$.

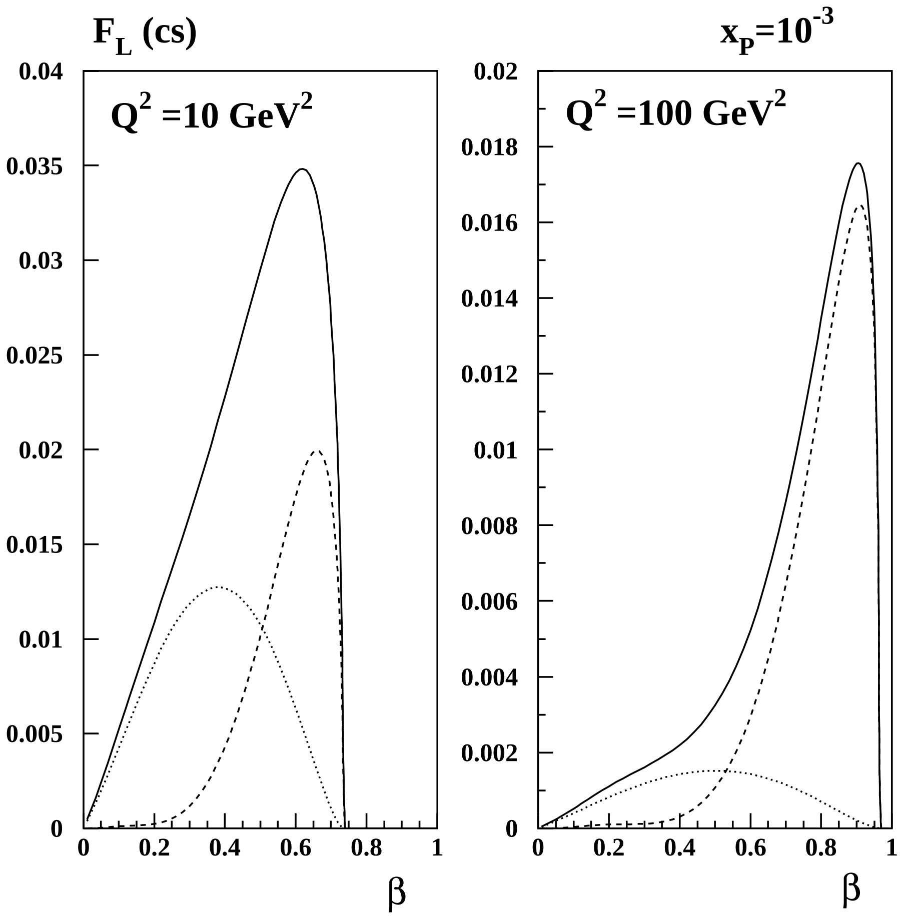

FIG. 2. $F_L^{D(3)}(cs)$ (solid) and A_L (Eq.23, dotted) and B_L (Eq.22, dot–dashed) components of F_L at Q^2=10,100 GeV^2.

A relevant point is that the pQCD scales $\bar{Q}^2(c)$ and $\bar{Q}(s)^2$ are different, both explicitly depend on β, and the $x_{\bf I\!P}$ and β dependences of $F_{T(i)}^{D(3)}$ are inextricably entangled. This gives another example where he Ingelman-Schlein factorization hypothesis, $F_2^{D(3)}(x_{\bf I\!P},\beta,Q^2) = f_{\bf I\!P}(x_{\bf I\!P})F_{2{\bf I\!P}}(\beta,Q^2)$, with process independent flux of pomerons in the proton $f_{\bf I\!P}(x_{\bf I\!P})$ and the $x_{\bf I\!P}$ independent pomeron SF $F_{2{\bf I\!P}}(\beta,Q^2)$, is not confirmed by pQCD calculation. The diffractive factorization breaking in CC diffractive DIS is especially severe, because for the same $c\bar{s}$ final state the pQCD factorization scale $\bar{Q}^2$ changes substantially from the forward to backward hemisphere: $\bar{Q}_s^2 \ll \bar{Q}_c^2$. Although the perturbative intrinsic charm component $F_{T(\bar{c})}^{D(3)}$ is suppressed by the mass of a heavy quark, it is still substantial and it is predicted to rise much steeper than the strange one as $x_{\bf I\!P} \to 0$. Furthermore, $F_{T(\bar{c})}^{D(3)}$ is truly of perturbative origin at all β, while $F_{T(s)}^{D(3)}$ has a non-negligible dependence to small scales up to $\beta \gtrsim 0.7$.

As an illustration of the diffractive factorization breaking, in Fig. 3 we show the effective exponent of the $x_{\mathbf{IP}}$ dependence

$$n_{eff} = 1 - \frac{\partial \log F_2^{D(3)}}{\partial \log x_{\mathbf{IP}}} \tag{24}$$

evaluated for $x_{\mathbf{IP}} = 3 \cdot 10^{-3}$.

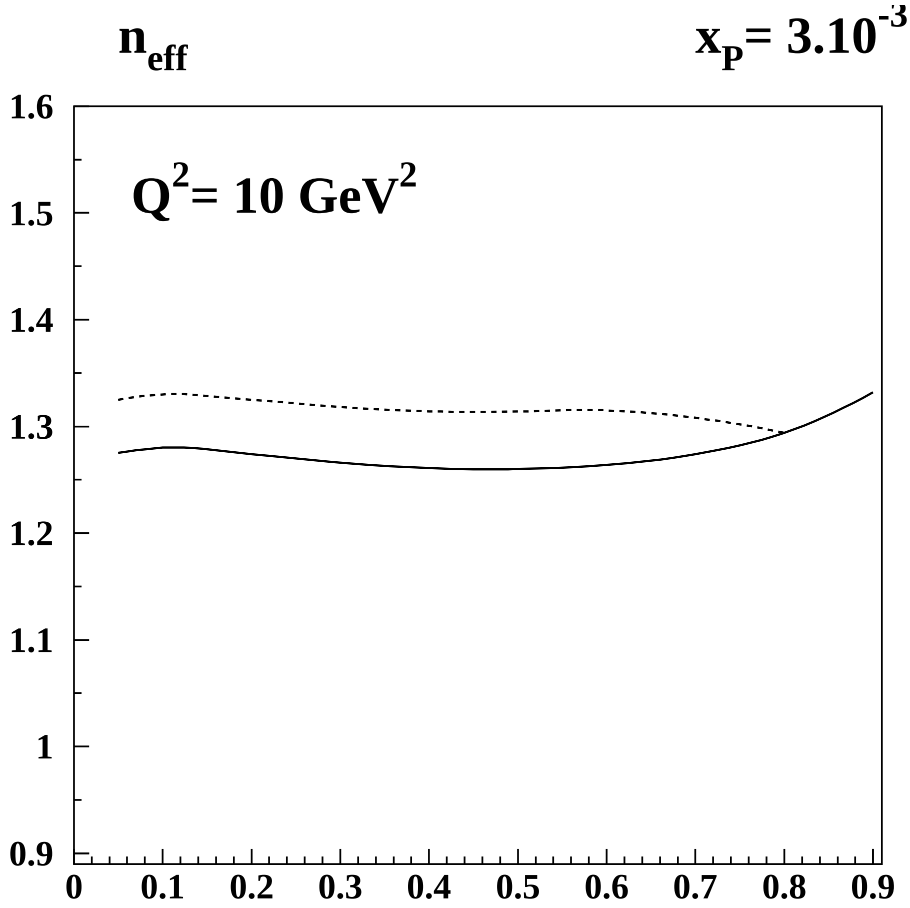

FIG. 3. The β dependence of n_eff for $x_{\mathbf{IP}} = 0.003$ and $Q^2 = 10\text{GeV}^2$. Solid line: ud component, dashed line: ud+cs

Evidently, at fixed β, the $c\bar{s}$ excitation is possible only for sufficiently large Q^2 such that $\beta_c > \beta$. For this reason, diffractive SF's exhibit strong threshold effects shown in Fig. 4, which are much stronger than in the NC case studied in [6,8]. Notice, that $F_3^{D(3)}$ vanishes below the $c\bar{s}$ threshold.

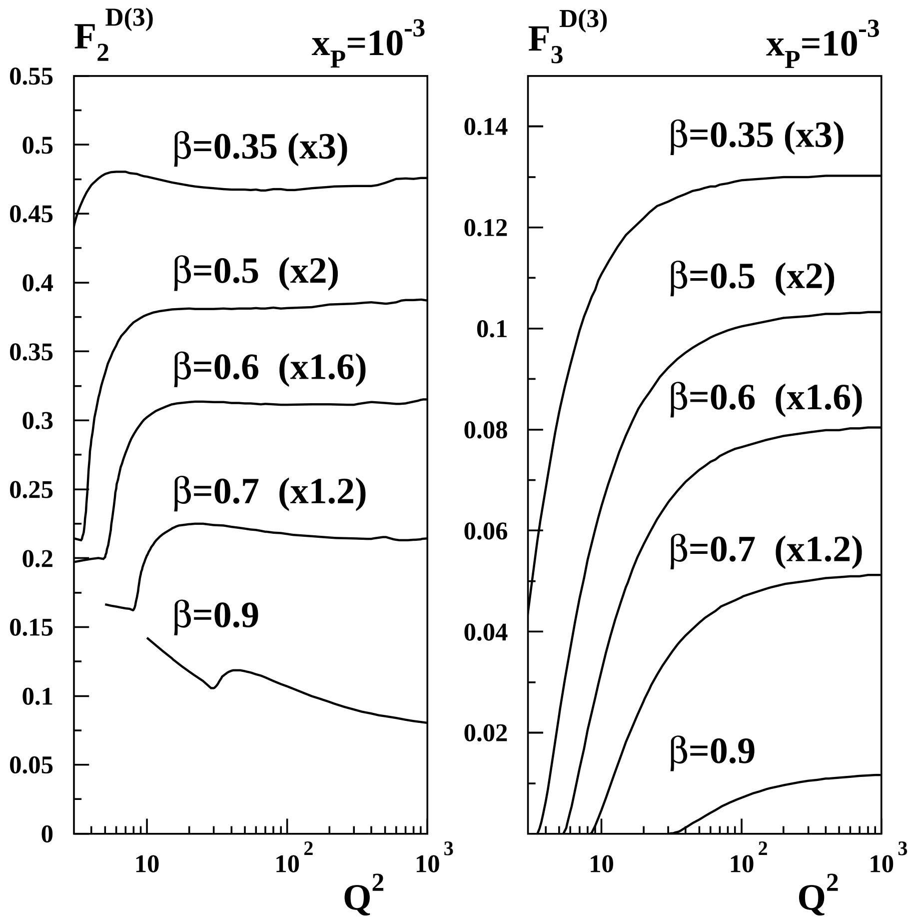

FIG. 4. Charm-strange threshold effect on Diffractive SF. First box: $F_2^{D(3)}$. Second box: $F_3^{D(3)}$.

We have also calculated the contribution to diffractive SF's from the top-(anti)bottom excitation. With $m_t = 180$ GeV and $m_b = 4.5$ GeV, our result is that, even for the very large $Q^2 = 10^4$ GeV^2, the $t\bar{b}$ contribution is about 50 times smaller than the $c\bar{s}$ one, making very difficult to observe this component at HERA.

Naively, one would expect $F_3^{D(3)} = 0$ for a quark-antiquark symmetric target as the Pomeron is. Indeed, because $A_3(z)$, and for equal mass case, $B_3(z)$ too, are antisymmetric about $z = \frac{1}{2}$, the contribution from $u\bar{d}$ excitation to $F_3^{D(3)}$ vanishes upon the integration over the u-jet production angles. In contrast to that, in the $c\bar{s}$ excitation there is a strong forward-backward asymmetry, and $F_3^{D(3)} = F_{T(s)}^{D(3)} - F_{T(\bar{c})}^{D(3)} \neq 0$. Our predictions for $F_3^{D(3)}$ are shown in Fig.s 1 and 4.

Up to now we have discussed the intermediate and large β regions, the small β one is, on the other hand, dominated by the so-called triple-pomeron component. In this case, diffraction proceeds via excitation of the soft gluon-containing $q\bar{q}g$ and higher Fock states of the photon. As it has been discussed to great detail in [3,4], at $\beta \ll 1$ and only at $\beta \ll 1$, and with certain reservations, one can apply the standard parton model treatment to diffractive DIS. For instance, the conventional fusion of virtual photons with the gluon from the two-gluon valence state of the pomeron becomes the driving term of diffractive DIS.

For this component, the results for the diffractive SF of light quarks coincide (once the opportune couplings of weak interaction are substituted to the EM ones) with those presented in Ref. [4]. However, for the charm–strange part, it must be considered that now a charm quark always appears together with a strange one, leading to a threshold ($Q^2_{cs} = 4GeV^2$), which is intermediate between the strange ($Q^2_{ss} = 1GeV^2$) and the charm ($Q^2_{cc} = 10GeV^2$) electromagnetic DIS thresholds in analogy to our discussion concerning the usual DIS [6] (Referring to the notation of [4] one finds $A_{cs} = 0.08$.)

We have presented the calculation of charged current diffractive structure functions in QCD and carried on a comparison with electromagnetic case. Both charged current and electromagnetic diffraction share the property of diffractive factorization breaking. For instance, we find different $x_{\mathbf{I\!P}}$ dependences of the intrinsic u, d, strangeness and charm composition of the pomeron. Futhermore, we predict, for CC, even a different $x_{\mathbf{I\!P}}$ dependence for the production of charm quark in the forward and backward direction.

Other new features of CC diffraction respect to EM one, are the emergence of substantial $F_3^{D(3)}$, and the large higher twist contributions to the longitudinal structure function. These predictions will be testable with the accumulation of a larger statistic on CC diffraction at HERA.

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