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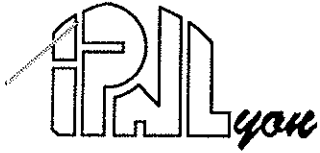
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Factorization properties of the diffraction dissociation of longitudinal photons

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Abstract

We develop the pQCD description of the diffraction dissociation (DD) of longitudinal photons. We demonstrate that the longitudinal diffractive structure function does not factor into the flux of pomerons and the partonic structure function of the pomeron, thus defying the usually assumed Regge factorization. In contrast to DD of the transverse photons, DD of the longitudinal photons is strongly peaked at $\beta = 1$. We comment on duality properties of DD in deep inelastic scattering.

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The longitudinal structure function $F_L(x, Q^2)$ is a fundamental quantity in deep inelastic scattering (DIS). In fixed target experiments, the measurement of $F_L(x, Q^2)$ requires the comparison of DIS at different energies of the lepton, at HERA one needs to vary either the lepton or proton beam energy or both. The exceptional case is diffraction dissociation (DD) of photons, which can be viewed as DIS on pomerons radiated by protons. Here the longitudinal diffractive structure function $F_L^D(x_{\mathbb{P}}, \beta, Q^2)$ could readily be measured varying the energy and/or $x_{\mathbb{P}}$ of the target pomeron at a fixed energy of the electron and proton beams. The experimental measurement of $F_L^D(x_{\mathbb{P}}, \beta, Q^2)$ could shed much light on the microscopic QCD structure of the pomeron; the corresponding experimental data from HERA will become available soon and thus the pQCD evaluation of $F_L^D(x_{\mathbb{P}}, \beta, Q^2)$ is one of the topical issues in the theory of the QCD pomeron.

The subject of the present communication is the derivation of the pQCD relationship between the mass spectrum in DD of longitudinal photons and the gluon structure function of the proton. We derive the relevant pQCD factorization scale and establish the pattern of breaking of the Regge factorization. Our results do clearly demonstrate that treating the pomeron as a hadronic state endowed with a well-defined flux in the proton and a partonic structure function is illegitimate.

We discuss the diffraction dissociation (DD) of (virtual) photons $\gamma^* + p \rightarrow X + p'$ into states X of mass M (large rapidity gap (LRG) events) and calculate the diffractive structure function defined by

$$(M^2 + Q^2) \frac{d\sigma^D(\gamma^* \rightarrow X)}{dt dM^2} \Big|_{t=0} = \frac{\sigma_{tot}(pp) 4\pi^2 \alpha_{em}}{16\pi Q^2} \{ F_T^D(x_{\mathbb{P}}, \beta, Q^2) + \varepsilon_L F_L^D(x_{\mathbb{P}}, \beta, Q^2) \}. \quad (1)$$

Here Q^2 is the virtuality of the photon, W and M are c.m.s. energy in the photon-proton and photon-pomeron collision, $\beta = Q^2/(Q^2 + M^2)$ is the Bjorken variable for the lepton-pomeron DIS, $x_{\mathbb{P}} = (Q^2 + M^2)/(Q^2 + W^2) = x/\beta$ is interpreted as the fraction of the momentum of the proton carried away by the pomeron, ε_L is the longitudinal polarization of the photon and α_{em} is the fine structure constant. The dimensional normalization factor $\sigma_{tot}(pp) = 40 \text{ mb}$ follows from the standard Regge theory convention [1, 2].

An assumption which is often made is that the pomeron can be treated as a hadronic state and $F_{T,L}^D(x_{\mathbb{P}}, \beta, Q^2)$ can be factored into the partonic structure function of the pomeron

$F_{2\mathbb{P}}(\beta, Q^2)$ and the flux of pomerons $\phi_{\mathbb{P}}(x_{\mathbb{P}})/x_{\mathbb{P}}$ in the proton [3, 4]:

$$F_{T,L}^D(x_{\mathbb{P}}, \beta, Q^2) = \phi_{\mathbb{P}}(x_{\mathbb{P}}) F_{T,L}^{\mathbb{P}}(\beta, Q^2). \quad (2)$$

This Ingelman-Schlein-Regge factorization, which has never been derived from the QCD analysis, involves a set of very strong assumptions on the diffractive cross section $d\sigma^D$: i) the $x_{\mathbb{P}}$ dependence is reabsorbed entirely in the Q^2, β and flavour independent pomeron flux function $\phi_{\mathbb{P}}(x_{\mathbb{P}})$, ii) the β, Q^2 and flavour dependence of $d\sigma_D$ are contained entirely in the structure function of the pomeron, iii) the ratio $R^D = F_L^D(x_{\mathbb{P}}, \beta, Q^2)/F_T^D(x_{\mathbb{P}}, \beta, Q^2)$ does not depend on $x_{\mathbb{P}}$. The purpose of the present communication is the pQCD derivation of $F_{T,L}^D(x_{\mathbb{P}}, \beta, Q^2)$ and the demonstration that none of the above properties i) to iii) holds in the pQCD.

Different aspects of the non-factorization in DD have already been discussed in [1, 2, 5, 6, 7]; the non-factorizable colour dipole approach to DD [1, 6] is well known to provide a very good quantitative description of the HERA data on LRG events [9, 10].

We start with diffraction excitation of photons into $q\bar{q}$ pairs, which dominates at large β and can be associated with DIS on the "valence" $q\bar{q}$ component of the photon. The formalism necessary for our purposes has been set up in [1, 5, 7]. The relevant pQCD diagrams for the colour singlet exchange in the t -channel are shown in Fig. 1. The mass of the diffractively excited state X is given by $M^2 = (m_f^2 + k^2)/z(1-z)$, where m_f is the quark mass, \vec{k} is the transverse momentum of the quark with respect to the γ^* -pomeron collision axis and z is the fraction of light-cone momentum of the photon carried by the (anti)quark. Other useful kinematical variables are $\varepsilon^2 = z(1-z)Q^2 + m_f^2$ and

$$q^2 = k^2 + \varepsilon^2 = (k^2 + m_f^2) \frac{M^2 + Q^2}{M^2} \quad (3)$$

After the standard leading $\log\kappa^2$ resummation, the cross sections of the forward ($t=0$) DD of longitudinal photons takes the compact form [1, 5]

$$\left. \frac{d\sigma_L}{dM^2 dk^2 dt} \right|_{t=0} = \frac{\pi^2}{6} e_f^2 \alpha_{em} \alpha_S^2(q^2) \cdot \frac{Q^2 (m_f^2 + k^2)^3}{M^7 \cos\theta \sqrt{M^2 - 4m_f^2}} \Phi_2^2. \quad (4)$$

Here e_f is the quark charge in units of the electron charge, θ is the quark production angle

with respect to the γ^* -pomeron collision axis,

$$\Phi_2 = \int \frac{d\kappa^2}{\kappa^4} f(x_{\mathbb{P}}, \kappa^2) \left[\frac{1}{\sqrt{a^2 - b^2}} - \frac{1}{k^2 + \epsilon^2} \right], \quad (5)$$

$a = \epsilon^2 + k^2 + \kappa^2$, $b = 2k\kappa$ and $f(x_{\mathbb{P}}, \kappa^2) = \partial G(x_{\mathbb{P}}, \kappa^2)/\partial \log \kappa^2$ is the unintegrated gluon structure function of the target proton. Following the analysis [5, 7] one can easily verify that after factoring out $(k^2 - \epsilon^2)/(k^2 + m_f^2)^3$ in (5), one will be left with the logarithmic κ^2 integration with q^2 being the upper limit of integration. Consequently, q^2 emerges as the pQCD factorization scale (it has already been used as such in the running strong coupling $\alpha_S(q^2)$ in (4)) and to the leading $\log q^2$,

$$\Phi_2 = \frac{M^4[(k^2 + m_f^2)(M^2 - Q^2) - 2m_f^2 M^2]}{(Q^2 + M^2)^3(k^2 + m_f^2)^3} G(x_{\mathbb{P}}, q^2). \quad (6)$$

Notice a zero of the $d\sigma_L$ at $(k^2 + m_f^2)(M^2 - Q^2) = 2m_f^2 M^2$. For light flavours, $d\sigma_L$ vanishes at $M^2 = Q^2$. Substituting (6) into (5), one readily finds

$$\left. \frac{d\sigma_L}{dM^2 dk^2 dt} \right|_{t=0} = \frac{\pi^2}{6} e_f^2 \alpha_{em} \alpha_S^2(q^2) G^2(x, q^2) \frac{Q^2 M [(k^2 + m_f^2)(M^2 - Q^2) - 2m_f^2 M^2]^2}{\cos \theta (Q^2 + M^2)^6 (k^2 + m_f^2)^3 \sqrt{M^2 - 4m_f^2}}. \quad (7)$$

Notice that in the DIS limit of $Q^2 \gg m_f^2$, the k^2 and M^2 dependences in (7) do factor, which leads to the simple β dependence $F_L^D \propto (1 - 2\beta)^2 \beta^3$. The r.h.s. of Eq. (7) decreases with k^2 only as k^{-2} and one has the logarithmic integration $\int^{\frac{1}{4}M^2} dk^2/(k^2 + m_f^2)$. The Jacobian peak singularity and the scaling violations in $G(x, q^2)$ further enhance the contribution from large $k^2 \sim \frac{1}{4}M^2 - m_f^2$. Consequently, the relevant pQCD factorization scale equals

$$q^2 \approx \frac{1}{4\beta} Q^2. \quad (8)$$

The k^2 integration produces a logarithmic factor of the form $\log(M^2/4m_f^2) = \log(Q^2(1 - \beta)/4\beta m_f^2)$. This factor has only a marginal effect on the β and Q^2 dependence. At asymptotically large $Q^2 \gg 4m_f^2$ the flavour symmetry is restored, but for Q^2 of practical interest there is a substantial suppression of the charm cross section, similar to the suppression of the charm structure function of the proton [8]. Suppressing this factor, to a logarithmic accuracy,

$$F_L^D(x_{\mathbb{P}}, \beta, Q^2) = e_f^2 \frac{2\pi}{3\sigma_{tot}(pp)Q^2} (1 - 2\beta)^2 \beta^3 \cdot \alpha_S^2\left(\frac{Q^2}{4\beta}\right) \cdot G^2\left(\frac{Q^2}{4\beta}, x_{\mathbb{P}}\right)^2, \quad (9)$$

which concludes the derivation of the longitudinal diffractive structure function. Because of the zero at $\beta = \frac{1}{2}$ and the β^3 dependence, the $F_{T,L}^D(x_{\mathbb{P}}, \beta, Q^2)$ is strongly peaked at $\beta = 1$, so that one can put, with a good accuracy, $\beta = 1$ in the factorization scale.

The salient features of the $F_L^D(x_{\mathbb{P}}, \beta, Q^2)$ are clearly seen from Eq. (9). First, it is short distance dominated and is exactly calculable in the realm of pQCD [11, 1]. Second, it has the higher twist dependence $\propto 1/Q^2$, a result known since [11]. Third, the β and $x_{\mathbb{P}}$ dependences do factorize. Fourth, the $x_{\mathbb{P}}$ - and Q^2 -dependences are inextricably entangled, the Regge factorization (2) breaks down and neither the concept of a Q^2 independent flux of pomerons nor the one of a pomeron structure function (which absorbs all the Q^2 dependence) do make sense. Regretfully, these concepts have become customary in the analysis and presentation of the experimental data. The above pQCD derivation shows unequivocally that if one wants to keep the pomeron structure function language, then one can do so only at the expense of modifying the Eq. 2 to allow for the Q^2 dependent pomeron flux function:

$$\phi_{\mathbb{P}}^L(x_{\mathbb{P}}, Q^2) = \left(\frac{G(x_{\mathbb{P}}, Q^2/4)}{G(x_0, Q^2/4)} \right)^2. \quad (10)$$

Here the normalization is $\phi_{\mathbb{P}}^L(x_0 = 0.03) = 1$ for every Q^2 [6]. Notice that the so defined $\phi_{\mathbb{P}}^L(x_0, Q^2)$ is flavour independent, in contrast to the diffraction dissociation of transverse photons where excitation of each and every new flavour entails the brand new pomeron flux function [7]. Then, with all the above reservations, one can define the longitudinal structure function of the pomeron,

$$F_{2\mathbb{P}}^L(\beta) = \Sigma_f e_f^2 \frac{A_f}{Q^2} (1 - 2\beta)^2 \beta^3, \quad (11)$$

where the normalization factors A_f are, to a first approximation, Q^2 independent.

For the numerical evaluation of the longitudinal cross section and the normalization factors A_f in (11) we rely upon the colour dipole gBFKL formalism [11, 1, 2]. The M^2 and/or β integrated DD cross section equals

$$\left. \frac{d\sigma_L^D}{dt} \right|_{t=0} = \int dM^2 \left. \frac{d\sigma_L^D}{dt dM^2} \right|_{t=0} = \frac{1}{16\pi} \int_0^1 dz \int d^2\vec{\rho} |\Psi_{\gamma^*}^L(Q^2, z, r)|^2 \sigma^2(x, r) \quad (12)$$

where

$$|\Psi_{\gamma^*}^L(Q^2, z, r)|^2 = \frac{6\alpha_{em}}{(2\pi)^2} \sum_i^{N_f} e_f^2 4Q^2 z^2 (1-z)^2 K_0(\varepsilon r)^2 \quad (13)$$

gives the colour dipole distribution in longitudinal photons [11] and $\sigma(x, r)$ is the colour dipole cross section from ref. [12, 13]. The resulting $\phi_{\mathbb{P}}^L(x_{\mathbb{P}}, Q^2)$ can conveniently be parameterized (for $1 \lesssim Q^2 \lesssim 100 \text{GeV}^2$) as

$$\phi_{\mathbb{P}}^L(x_{\mathbb{P}}, Q^2) = \left(\frac{x_0}{x_{\mathbb{P}}}\right)^{[a+d \log(Q^2/10)+f \log^2(Q^2/10)]} \cdot \left[\frac{x_{\mathbb{P}} + c}{x_0 + c}\right]^{[b+e \log(Q^2/10)]} \quad (14)$$

where Q^2 is in GeV^2 , $a = 0.456$, $b = 0.678$, $c = 0.012$, $d = 0.112$, $e = 0.078$ and $f = 0.01$.

The generalized flux function $\phi_{\mathbb{P}}^L(x_{\mathbb{P}}, Q^2)$ is flavour independent, the flavour dependent normalizations A_f in (11) can be determined equating the cross sections given by Eq. (12) and the M^2 and/or β -integrated Eq. (9). In the interesting range of $Q^2 \lesssim 100 \text{GeV}^2$, the result is : $A_{ud} = 0.82$, $A_s = 0.61$ and $A_c = 0.05$. Notice that $A_{ud} \approx A_s$, for the charm the flavour symmetry is strongly broken. This order of magnitude estimate for A_c is sufficient for evaluations of the numerically small charm cross section.

A very different situation occurs for the triple pomeron region of $\beta \ll 1$, which is dominated by DD into $q\bar{q}g\dots$ states. The factorization properties of DD in this region of β can clearly be seen from diffractive excitation of the $q\bar{q}g$ states of the photon, which gives the driving term of $F_{T,L}^D(x_{\mathbb{P}}, \beta, Q^2)$ at $\beta \ll 1$. In [2] it has been shown that DD in DIS is dominated by configurations in which the transverse separation ρ of the gluon from the $q\bar{q}$ pair is much larger than the $q\bar{q}$ separation r . Then, the DD cross section can be factored as

$$(Q^2 + M^2) \left. \frac{d\sigma_{T,L}^D}{dt dM^2} \right|_{t=0} \simeq \int dz d^2\bar{r} |\Psi_{\gamma^*}^{T,L}(Q^2, z, r)|^2 \cdot \frac{16\pi^2}{27} \cdot \alpha_S(r)r^2 \\ \times \frac{1}{2\pi^4} \cdot \left(\frac{9}{8}\right)^3 \cdot \int d\rho^2 \left[\frac{\sigma(x_{\mathbb{P}}, \rho)}{\rho^2}\right]^2 \mathcal{F}(\rho), \quad (15)$$

where $\mathcal{F}(\rho)$ provides an infrared cutoff at distances ρ exceeding the propagation radius for perturbative gluons, for a detailed discussion see [2, 6, 14]. The crucial point is that the $x_{\mathbb{P}}$ dependence in (15) decouples from the Q^2 and β dependence and is universal for the $d\sigma_T^D$ and $d\sigma_L^D$ (as well as flavour independent) as soon as $Q^2 \gtrsim 3 \text{GeV}^2$ [14], it is given by the pomeron flux function $f_{\mathbb{P}}(x_{\mathbb{P}})$ calculated in [6]. Eq. (15) gives the driving term of the leading- $\log\frac{1}{\beta}$ expansion of the DD cross section. It can be argued that at least to the leading- $\log\frac{1}{\beta}$, the diffractive structure function has the conventional GLDAP evolution

properties, the corresponding analysis [2] needs not be repeated here. The structure of the r integrations in (15) is only marginally different from that in the DIS structure function at small x . The detailed calculation of the ratio $R^{DIS} = \sigma_L/\sigma_T$ for DIS has been performed in [11, 12], the major finding is that $R^{DIS} \approx 0.2$ with a very weak Q^2 and x dependence. Consequently, we expect a close similarity of R^D at small β to R^{DIS} at small x . In Fig. 2 we present our results for R^D at $\beta \ll 1$ as a function of Q^2 , as it was anticipated it exhibits very weak Q^2 dependence, with the exception of the excitation of open charm, where the standard threshold behaviour $\propto Q^2/(Q^2 + 4m_c^2)$ is clearly seen.

For a numerical estimate of the longitudinal diffractive structure function we will use in the following: for the valence part, Eq. (11) with the flux parametrization Eq. (14) (which reproduces the exact result to a $\approx 10\%$ accuracy in the $(x_{\mathbb{P}}, Q^2)$ region relevant at HERA, $10^{-4} < x_{\mathbb{P}} < 0.03$, $1 \text{ GeV}^2 \leq Q^2 \leq 100 \text{ GeV}^2$); for the sea component the results for F_T^D from Ref. [6, 7] assuming $R^D = 0.2$ constant in the whole region.

In Fig. 3 we show how the transverse and longitudinal diffractive structure functions $F_{T,L}^D(x_{\mathbb{P}}, \beta, Q^2)$ evolve with Q^2 and $x_{\mathbb{P}}$. The Q^2 evolution of F_T^D is marginal, we show it for $Q^2 = 10 \text{ GeV}^2$. At small $\beta \ll 1$, the longitudinal contribution is small, $R^D \approx 0.2$, and both the F_T^D and F_L^D have identical $x_{\mathbb{P}}$ dependence. As $F_T^D(x_{\mathbb{P}}, \beta, Q^2)$ vanishes for $\beta \rightarrow 1$, at $\beta \gtrsim 0.9$ the diffractive structure function is entirely dominated by the $F_L^D(x_{\mathbb{P}}, \beta, Q^2)$. At fixed $x_{\mathbb{P}}$, the longitudinal structure function decreases with Q^2 , however at small $x_{\mathbb{P}}$ the higher twist behaviour $F_L^D \propto 1/Q^2$ is to a large extent compensated by the scaling violations in the gluon structure function. For this reason, F_L^D remains non-negligible even for Q^2 as large as $Q^2 \sim 100 \text{ GeV}^2$. Notice also the steeper $x_{\mathbb{P}}$ -dependence of $F_L^D(x_{\mathbb{P}}, \beta, Q^2)$ at large β as compared to the $x_{\mathbb{P}}$ dependence of F_T . In the typical kinematics of the HERA experiments $\epsilon_L \approx 1$ and the measured diffractive structure function roughly corresponds to $F_2^D = F_T^D + F_L^D$. In the range $x_{\mathbb{P}} = [10^{-3}, 10^{-2}]$ of the present HERA experiments, the $x_{\mathbb{P}}$ dependence of F_2^D can be parametrized by the law $\propto x_{\mathbb{P}}^{-\delta}$ to a $\approx 20\%$ accuracy. The so estimated exponent δ is shown in Fig. 4 for $Q^2 = 100 \text{ GeV}^2$ and $Q^2 = 10 \text{ GeV}^2$. The exponent δ rises towards $\beta \rightarrow 1$, takes a minimal value at moderately small β , then rises again towards small β . The approximation $\propto x_{\mathbb{P}}^{-\delta}$ for the $x_{\mathbb{P}}$ dependence is rather

crude; the value of the exponent δ depends on the range of $x_{\mathbb{P}}$, the explicit form of the $x_{\mathbb{P}}$ dependence is shown in [6, 7]. The values of δ evaluated for the range $x_{\mathbb{P}} = [10^{-3}, 3 \cdot 10^{-2}]$ are uniformly lower by $\approx 0.03-0.04$ than those shown in Fig. 3, however the form of the β dependence of the exponent δ is fully preserved. Because of the partial compensation, which has been described previously, of the higher twist behaviour of F_L by the scaling violations in the generalized flux (10), (14), the rise of the exponent δ towards $\beta \rightarrow 1$ persists at all the Q^2 and is quite relevant.

Finally, we wish to comment on the Bloom-Gilman-Drell-Yan-West duality-type relationship between the diffraction dissociation into the $q\bar{q}$ continuum at $\beta \rightarrow 1$ and the exclusive diffractive production of vector mesons:

$$\int_0^{M_V^2} dM^2 \frac{d\sigma_{T,L}^D}{dM^2} \propto \frac{1}{Q^2} \int_{\beta_0}^1 d\beta F_{T,L}^D(x, \beta, Q^2) \propto \sigma(\gamma_{T,L}^* N \rightarrow V_{T,L} N) \quad (16)$$

In the l.h.s. of (16), the integration goes over the resonance mass range $M^2 \sim M_V^2$ and/or over the large- β domain $1 - \beta \lesssim 1 - \beta_0 \sim \frac{M_V^2}{Q^2}$. In this domain, $x_{\mathbb{P}}$ coincides with the Bjorken variable x . The pQCD description of exclusive production of vector mesons has been developed in [15]. The longitudinal photons produce longitudinally polarized vector mesons, in [15] it was shown that $\sigma(\gamma_L^* N \rightarrow V_L N) \propto Q^{-6} G^2(x, \tau Q^2)$, the factor $\tau \sim 0.1 - 0.2$ in the factorization scale was derived in [13]. One recovers precisely the same x and Q^2 dependence after the integration of the mass spectrum (7) over $M^2 \lesssim M_V^2$, both the higher twist behaviour and the flat β dependence at $\beta \rightarrow 1$ in (9, 11) are crucial for this consistency. Similar consistency with the duality is found for the transverse photons. Namely, here the result of ref. [15, 13] for the exclusive cross section is $\sigma(\gamma_T^* N \rightarrow V_T N) \propto Q^{-8} G^2(x, \tau Q^2)$. The limiting behaviour of the mass spectrum for DD of transverse photons at $\beta \rightarrow 1$ has been derived in our previous paper [7],

$$F_T^D(x_{\mathbb{P}}, \beta, Q^2) \propto (1 - \beta)^2 G^2(x, q^2), \quad (17)$$

where the factorization scale q^2 equals

$$q^2 \sim \frac{m_f^2}{1 - \beta} = m_f^2 \left(1 + \frac{Q^2}{M^2}\right). \quad (18)$$

In the exclusive limit, $q^2 \propto Q^2$ and Eqs. (17),(18) entail the identical x and Q^2 dependence of the l.h.s. and r.h.s. of Eq. (16). Notice a remarkable conspiracy of breaking of the Regge

factorization in the longitudinal and transverse cross sections and of the β dependence of the pQCD factorization scales (8) and (18), which is crucial for the duality relationship between the exclusive vector meson production and diffraction dissociation to hold for both the longitudinal and the transverse photons.

Summary and conclusions. The presented QCD derivation of the mass spectrum for diffraction dissociation of longitudinal photons completes the analysis of the breaking of Regge factorization in diffractive deep inelastic scattering. The present results, together with those of our previous works [1, 6, 7] do unequivocally demonstrate that the Ingelman-Schlein-Regge factorization (2) is not born out by the QCD analysis of diffraction dissociation. The predicted breaking of the Regge factorization is strong and we look forward to the higher statistics data from HERA. Testing our predictions for the longitudinal diffractive structure function will be feasible in the near future, because the diffraction dissociation of photons is the unique process in which one can readily separate the longitudinal and transverse structure functions without varying the electron and proton beam energies.

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Figure captions

Fig.1 - One of the 16 Feynman diagrams for diffraction excitation of the qq state of the photon.

Fig.2 - R^D at $\beta \ll 1$ for the light quark (solid curve), the strange (dashed curve) and the charm (dot-dashed curve) components.

Fig.3 - $F_T^D(x_{\mathbf{P}}, \beta, Q^2)$ and $F_L^D(x_{\mathbf{P}}, \beta, Q^2)$ versus β at $x_{\mathbf{P}} = 0.03$ and $x_{\mathbf{P}} = 0.0003$. $F_T^D(x_{\mathbf{P}}, \beta, Q^2)$ is shown at $Q^2 = 10 \text{ GeV}^2$ (solid curve), while $F_L^D(x_{\mathbf{P}}, \beta, Q^2)$ is reported for $Q^2 = 10 \text{ GeV}^2$ (dot-dashed curve), $Q^2 = 50 \text{ GeV}^2$ (dashed curve) and $Q^2 = 100 \text{ GeV}^2$ (dotted curve).

Fig.4 - The exponent δ of the $x_{\mathbf{P}}$ dependence of the observed diffractive structure function $F^D = F_T^D + F_L^D$ at $Q^2 = 10 \text{ GeV}^2$ (dashed curve) and at $Q^2 = 100 \text{ GeV}^2$ (solid curve).

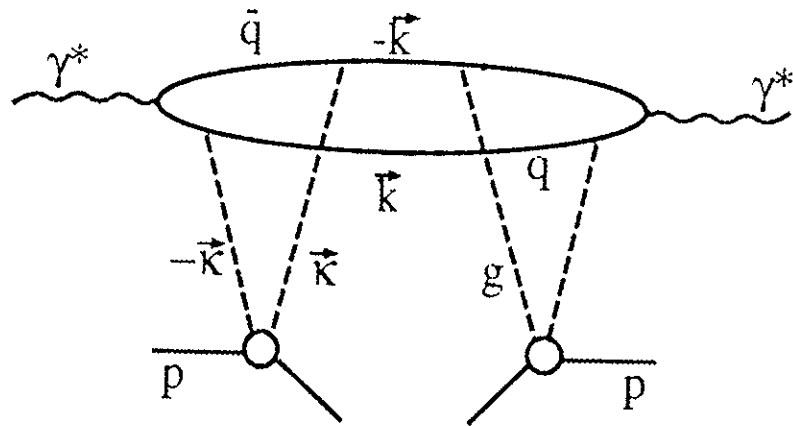


fig. 1

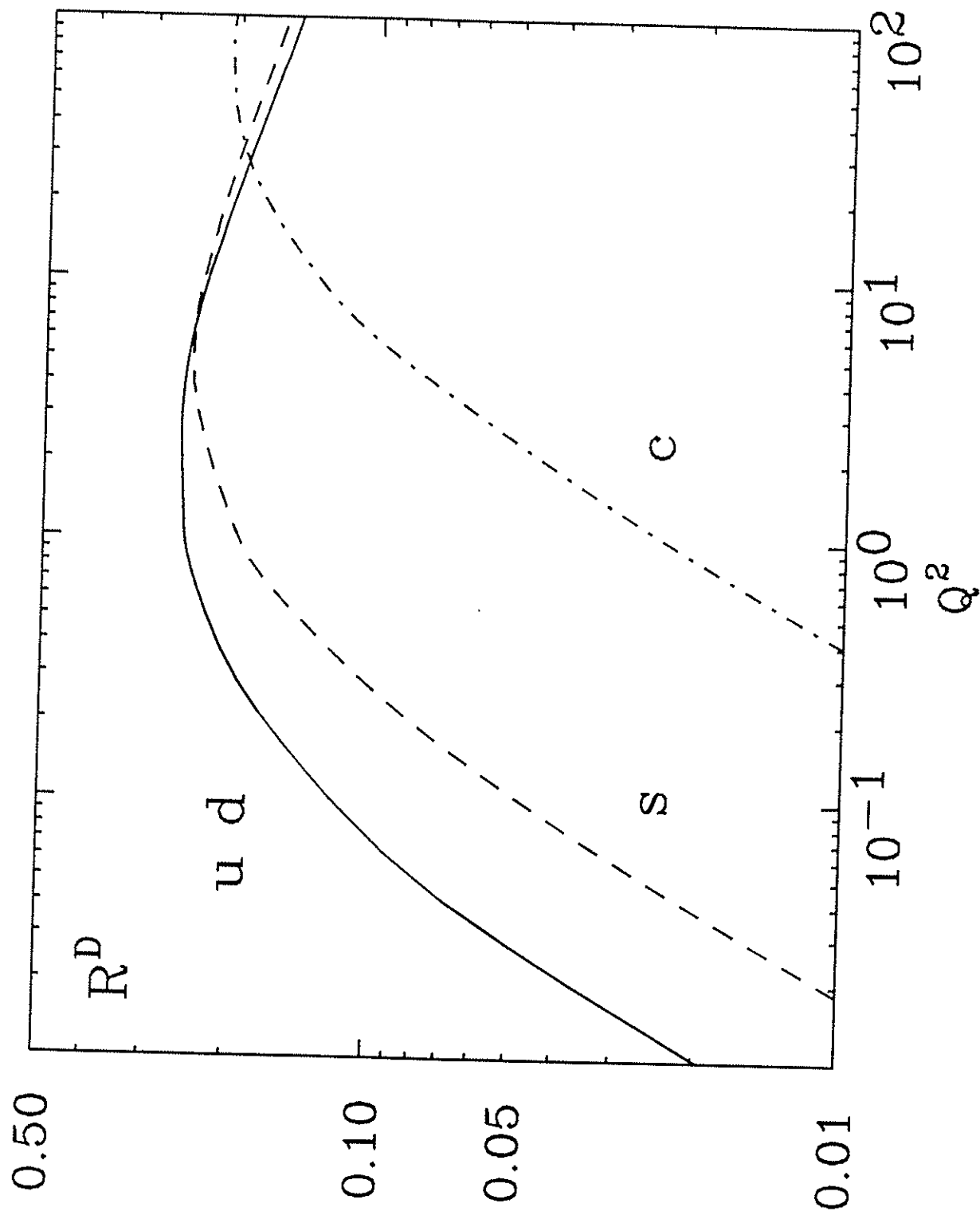


fig. 2

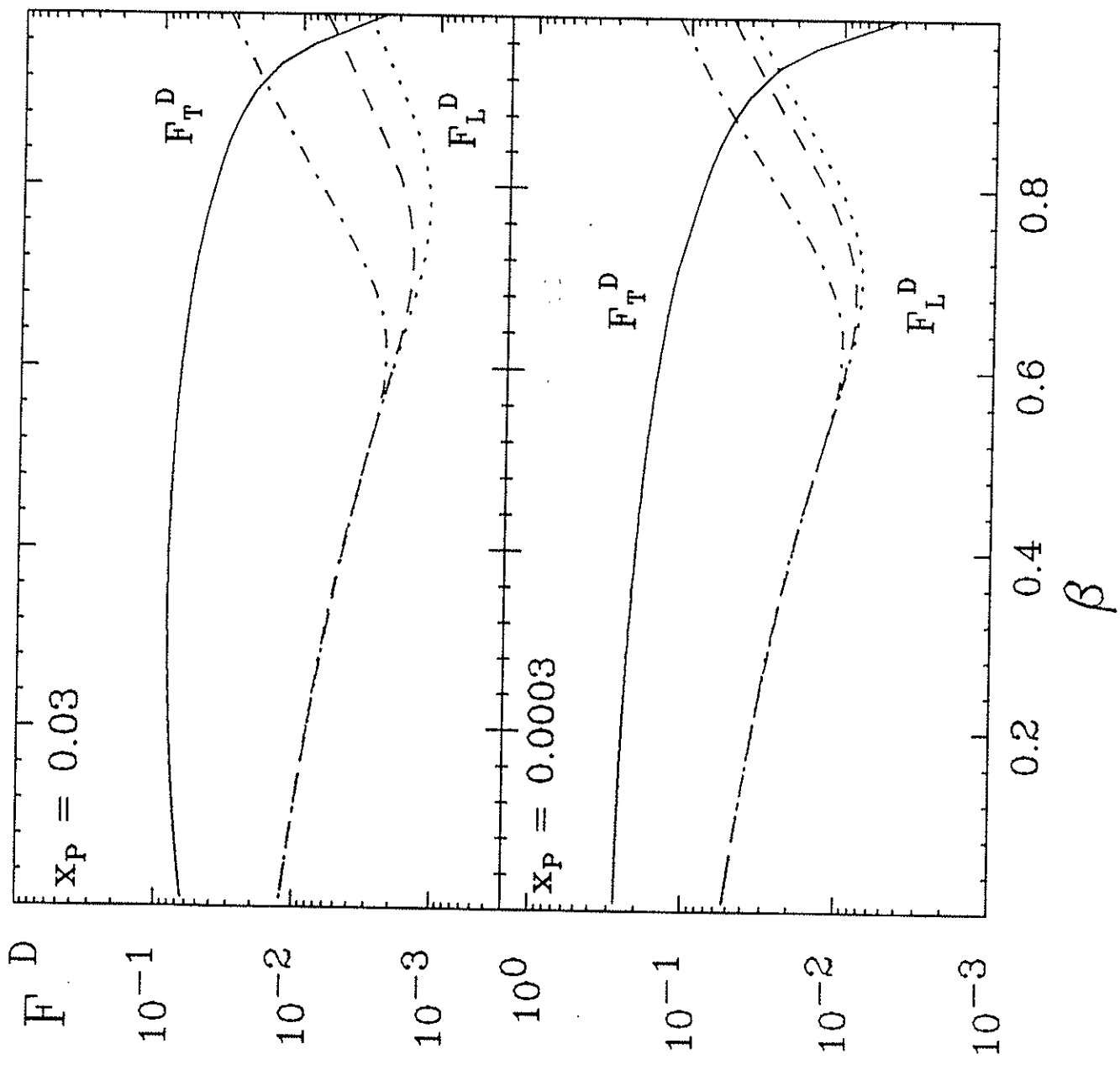


fig. 3

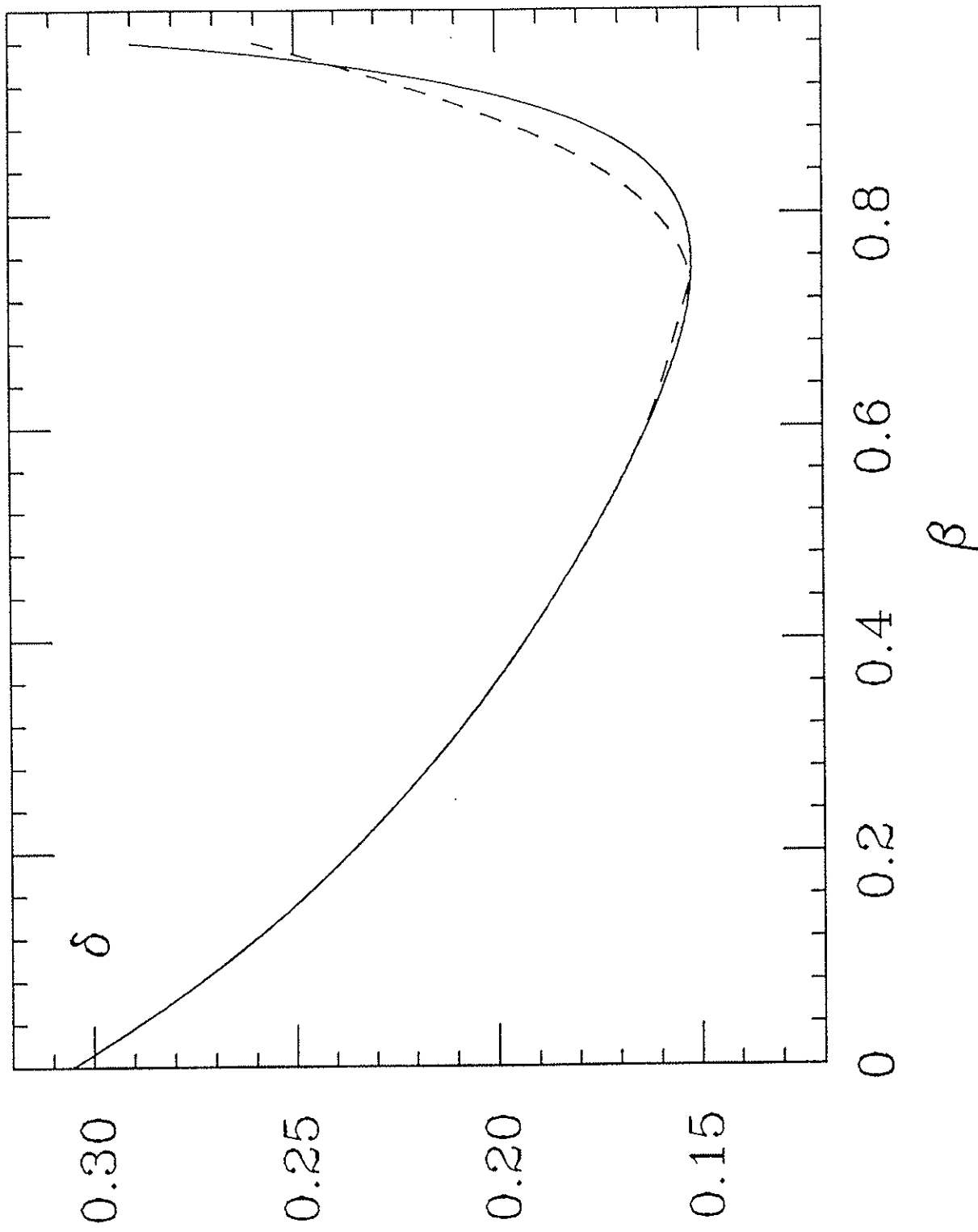


fig. 4