QCD and Diffraction in DIS

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Abstract

Coherence phenomena, and the non-universality of parton structure of the effective Pomeron are explained. New hard phenomena directly calculable in QCD such as diffractive electroproduction of states with \( M^2 \ll Q^2 \) as well as new options to measure the light-cone wave functions of various hadrons are considered. An analogue of Bjorken scaling is predicted for the diffractive electroproduction of \( \rho \) mesons at large momentum transfers and for the production of large rapidity gap events, as observed at HERA. A phenomenological QCD evolution equation is suggested to calculate the basic characteristics of the large rapidity gap events. The increase of parton densities at small \( x \) as well as new means to disentangle experimentally soft and hard physics are considered. We discuss constraints on the increase of deep inelastic amplitudes with \( Q^2 \) derived from unitarity of the S matrix for collisions of wave packets. New ways to probe QCD physics of hard processes at large longitudinal distances and to answer the long standing problems on the origin of the Pomeron are suggested. Unresolved problems and perspectives of small \( x \) physics are also outlined.

Résumé

Nous presentons une revue sur le rôle respectif de QCD dure et molle dans les réactions diffraction inclusives et exclusives.

1. Introduction

The aim of this talk is to outline QCD predictions for color coherence phenomena – a result of nontrivial interplay of hard and soft QCD physics specific for high energy processes (for more detailed discussion see [1]). Coherence phenomena provide an important link between the well understood physics of hard processes and the physics of soft processes which at present is mostly phenomenological. The soft/hard interplay is elaborated for the exclusive deep inelastic processes \( \gamma^* + N \rightarrow a + N \) for \( M_a^2 \ll Q^2 \) directly calculable in QCD. These processes provide new methods of investigating the structure of hadrons and the origin of the Pomeron and allow to search for new forms of hadronic matter in heavy ion collisions (for a review and references see [2]). The phenomenon of coherence reveals itself in high energy processes through a large probability of occurrence of diffractive processes and through their specific properties. Thus in this report we concentrate mostly on diffractive processes.

2. Interaction cross section for small size wave packet.

One of the striking QCD predictions for hard processes dominated by large longitudinal distances is that if a hadron is found in a small size configuration of partons

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it interacts with a target with a small cross section. The prediction which follows from the factorization theorem for hard processes in QCD is in variance with many phenomenological approaches based on pre-QCD ideas and on quark models of hadrons.

A sufficiently energetic wave packet with zero baryon and color charges localized in a small transverse volume in the impact parameter space can be described by a $q\bar{q}$ pair. This conclusion follows from asymptotic freedom in QCD which implies that the contribution of other components is suppressed by a power of the strong coupling constant $\alpha_s$ and/or a power of $Q^2$. A familiar example of such a wave packet is a highly longitudinal photon or a $\gamma^*$ in a $q\bar{q}$ state. Within the parton model the cross section for the interaction of such a photon with a target is suppressed by an additional power of $Q^2$. But at the same time the probability for a longitudinal photon to be in a large transverse size configuration (soft physics=parton model) is suppressed by a power of $Q^2$. These properties explain why reactions initiated by longitudinally polarized photons are best to search for new QCD phenomena.

The cross section for a high-energy interaction of a small size $q\bar{q}$ configuration off any target can be unambiguously calculated in QCD for low $x$ processes by applying the QCD factorization theorem. In the approximation when the leading $\alpha_s \ln \frac{Q^2}{\Lambda_{QCD}^2}$ term in $x$ terms are accounted for \[\frac{1}{3}\] the result is

$$\sigma(b^2) = \frac{\pi^2}{3} \left[ b^2 \alpha_s(Q^2) xG_T(x,Q^2) \right]_{x=Q^2/s, Q^2 \gg b^2} ,$$

where $b$ is the transverse distance between the quark $q$ and the antiquark $\bar{q}$ and $G_T(x, Q^2)$ is the gluon distribution in the target $T$ calculated within this approximation. In this equation the $Q^2$ evolution and the small $x$ physics are properly taken into account through the gluon distribution. To derive similar equation in the leading $\alpha_s \ln \frac{Q^2}{\Lambda_{QCD}^2}$ approximation one should account for all hard processes including diagrams where (anti)quarks in the box diagram with production of one hard gluon. The final result has the same form as eq. \[1\], but with $G_N(x, Q^2)$ calculated in the leading $\alpha_s \ln \frac{Q^2}{\Lambda_{QCD}^2}$ approximation. It also contains a small contribution due to sea quarks. Eq. \[1\] accounts for the contribution of quarks $Q$ whose masses satisfy the condition: $l_c = \frac{2\rho_0}{4m_Q^2 + Q^2} \gg r_N^2$. The estimate $Q^2 \approx \frac{1}{b^2}$ was obtained in \[1\] by numerical analysis of the $b$-space representation of the cross section of the longitudinally polarized photon, $\sigma_L$, and requiring that $G_T$ is conventional gluon distribution calculated in the leading $\alpha_s \ln \frac{Q^2}{\Lambda_{QCD}^2}$ approximation.

There is a certain similarity between equation \[3\] and the two gluon exchange model of F. Low \[5\] and S. Nussinov \[6\], as well as the constituent quark 2 gluon exchange model of J. Gunion and D. Soper \[8\]. The factor $b^2$ which is present in the QCD expression \[1\] for the cross section is also present in these models. The major qualitative distinction between the results of QCD calculations and expectations of the two gluon exchange models is that the nonperturbative QCD physics is accounted for in equation \[1\] through experimentally measured quantities - the gluon and the sea quark distributions. The latter are particularly relevant for the fast increase of the cross section at small $x$, for the increase of leading twist nuclear shadowing with decreasing $x$, for the seemingly slow decrease with $Q^2$ of higher twist processes. All those effects are characteristic for QCD as a nonabelian gauge quantum field theory which predicts an increase of parton densities in hadrons with $\frac{1}{b}$ in contrast to quantum mechanical models of hadrons. In QCD the inelastic cross section for the collision of a sufficiently energetic small size, colorless two gluon configuration off any target is \[1\]

$$\sigma(b^2) = \frac{3\pi^2}{4} \left[ b^2 \alpha_s(Q^2) xG_T(x,Q^2) \right]_{x=Q^2/s, Q^2 \gg \lambda/b^2} ,$$

where the parameter $\lambda$ is likely to be similar to the one present in the case of scattering of a $q\bar{q}$ pair off a target. The difference compared to equation \[1\] is in the factor $9/4$ which follows from the fact that gluons belong to the octet representation of the color group $SU(3)$, while quarks are color triplets.

3. Electroproduction of vector mesons in QCD.

One of the examples of a new kind of hard processes calculable in QCD is the coherent electroproduction of vector mesons off a target $T$,

$$\gamma^* + T \rightarrow V + T ,$$

where $V$ denotes any vector meson ($\rho, \omega, \phi, J/\Psi$) or its excited states.

The idea behind the calculation of hard diffractive processes is that when $l_c = \frac{1}{2m_N x}$ exceeds the diameter of the target, the virtual photon transforms into a hadron component well before reaching the target and the final vector meson $V$ is formed well past the target. The hadronic configuration of the final state is a result of a coherent superposition of all those hadronic fluctuations of the photon of mass $M$ that satisfy equation \[1\] $\frac{2l_c}{1 + M^2/Q^2} \gg r_N$. Thus, as in the more familiar leading twist deep inelastic processes, the calculation should take into account all possible
hadronic intermediate states satisfying this condition. The use of completeness over diffractively produced intermediate hadronic states allows to express the result in terms of distributions of bare of quarks and gluons as in the case of other hard processes. The matrix element of electroproduction of a vector meson \( A \) can be written as a convolution of the light cone wave function of the photon \( \psi^{\gamma \rightarrow |n|} \), the scattering amplitude for the hadron state \(|n\rangle\), \( A(nT) \), and the wave function of the vector meson \( \psi_V \)

\[
A = \psi^* \frac{\gamma}{|n|} A(nT) \otimes \psi_V .
\]

In the case of a longitudinally polarized photon with high \( Q^2 \) the intermediate state \(|n\rangle\) is a \( q\bar{q} \) pair. As was mentioned in the previous chapter section, it can be demonstrated by direct calculations that the contribution of higher Fock state components and soft physics are suppressed by a factor \( \frac{1}{Q^2} \) and/or powers of \( \alpha_s \). The proof of this result resembles the calculation of the total cross section for the deep inelastic scattering in QCD. The situation is qualitatively different in the case of other hard processes. The matrix element of the total cross section for the deep inelastic scattering in \( QCD \) is substituted by the gluon line. This leads to an extra term \( \alpha_s \ln \frac{Q^2}{\Lambda_{QCD}^2} \) as compared to 1, with \( \alpha_s \) the coupling constant.

It can be shown that under certain kinematic conditions this interaction of a \( q\bar{q} \) pair with the target is given by equation (1). In the leading order in \( \alpha_s \ln x \ln \frac{Q^2}{\Lambda_{QCD}^2} \) the leading Feynman diagrams for the process under consideration are a hard quark box diagram with two gluons attached to it and convoluted with the amplitude for the gluon scattering off a target.

One can consider the same process in the leading \( \alpha_s \ln \frac{Q^2}{\Lambda_{QCD}^2} \) approximation. In this case one has to include also the diagrams where one hard quark line is substituted by the gluon line. This leads to an extra term \( \alpha_s \ln \frac{Q^2}{\Lambda_{QCD}^2} \) in the calculation of the target interaction with a target is given by equation (1) [3, 4, 9].

Since Feynman diagrams are Lorentz invariant it is possible to calculate the box part of the amplitude in terms of the light-cone wave functions of the vector meson and the photon and to calculate the bottom part of the diagram in terms of the parton wave function of the proton. This mixed representation is different from the QCD improved parton model which only uses the light-cone wave function of the target.

The next step is to express this amplitude through the parton distributions in the target. The calculation of the imaginary part of the relevant Feynman diagram shows that the fractions of the target momentum carried by the exchanged gluons \( x_i \) and \( x_f \) are not equal,

\[
x_i - x_f = x, \quad \text{for } M^2_V \ll Q^2
\]

We neglect terms \( O(\frac{Q^2}{M_V^2}) \) as compared to 1, with \( l_t \) the transverse momentum of the exchanged gluons. Within the QCD leading logarithmic approximation

\[
\alpha_s \ln \frac{Q^2}{\Lambda_{QCD}^2} \sim 1
\]

or

\[
\alpha_s \ln x \ln \frac{Q^2}{\Lambda_{QCD}^2} \sim 1
\]

when terms \( \sim \alpha_s \) are neglected, the difference between \( x_i \) and \( x_f \) can be neglected and the amplitude of the \( q\bar{q} \) interaction with a target is given by equation (1) [3, 4, 9].
We are now able to calculate the cross section for the production of longitudinally polarized vector meson states when the momentum transferred to the target $t$ tends to zero, but $Q^2 \to \infty$.

$$\left| \frac{d\sigma^{L \to \gamma N}}{dt} \right|_{t=0} = \frac{12\pi^2 T_{V\to e^+e^-} m_V \alpha_s^2(Q) \eta_V^2}{\alpha_{EM} Q^0 N_c^2} \Gamma_V \int (xG_T(x,Q^2) + \frac{i\pi}{2} \frac{d}{d \ln x} G_T(x,Q^2))^2. \quad (11)$$

$\Gamma_{V\to e^+e^-}$ is the decay width of the vector meson into $e^+e^-$. The parameter $\eta_V$ is defined as

$$\eta_V = \frac{1}{2} \int \frac{dz}{z} \Phi_V(z), \quad (12)$$

where $\Phi_V$ is the light cone wave function of the vector meson. At large $Q^2$ equation (11) predicts a $Q^2$ dependence of the cross section which is substantially slower than $1/Q^6$ because the gluon densities at small $x$ increase with $Q^2$. Numerically, the factor $\alpha_s^2(Q^2)G^2(x,Q^2)$ in equation (11) is $\propto Q^6$ with $n \sim 1$. An additional $Q^2$ dependence of the cross section arises from the transverse momentum overlapping integral between the light-cone wave function of the $\gamma_L$ and that of the vector meson $\Psi$, expressed through the ratio $I_V(Q^2)$

$$I_V(Q^2) = \frac{\int_0^1 \frac{dz}{z(1-z)} \frac{d^2k_t}{Q^2 + (k_t^2 + \mu^2)^4} \psi_V(z,k_t)}{\int_0^1 \frac{dz}{z(1-z)} \frac{d^2k_t}{Q^2 + (k_t^2 + \mu^2)^4} \psi_V(z,k_t)}. \quad (13)$$

In ref. [8] it was assumed that $I_V(Q^2) = 1$ as for $Q^2 \to \infty$ the ratio $I_V(Q^2)$ tends to 1. But for moderate $Q^2$ this factor is significantly smaller than 1. For illustration we estimated $I_V(Q^2)$ for the following vector meson wave function: $\psi_V^{(1)}(z,k_T^2) = \frac{e^{-z(1-z)}}{(k_T^2 + \mu^2)^2}$. The momentum dependence of this wave function corresponds to a soft dependence on the impact parameter $b = \exp(-\mu b)$ in coordinate space. We choose the parameter $\mu$ so that $\langle k_T^2 \rangle^{1/2} < 0.3 \pm 0.6 \text{ GeV}/c$.

Our numerical studies show that the inclusion of the quark transverse momenta leads to several effects:

- Different $k_T$ dependence of $\psi_V$ leads to somewhat different $Q^2$ dependence of $I_V(Q^2)$. Thus measuring of $Q^2$ dependence of electroproduction of vector meson states may become an effective way of probing $k_T$-dependence of the light-cone $q\bar{q}$ wave function of vector mesons.

- The $Q^2$ dependence of $I_V$ for production of vector mesons build of light quarks $u,d,s$ should be very similar.

- For electroproduction of charmonium states where $\mu_c \sim \mu_{4\pi}$ the asymptotic formula should be only valid for extremely large $Q^2$.

The NMC data [10] and the HERA data [11] on diffractive electroproduction of $\rho$ mesons are consistent with several predictions of equation (11):

- a fast increase with energy of the cross section for electroproduction of vector mesons (proportional to $x^{-\delta}$ for $Q^2 = 10$ GeV$^2$) (figure 1 [8])

- the dominance of the longitudinal polarization $\frac{d\sigma}{d\cos \theta} \propto Q^2$.

- the absolute magnitude of the cross section within the uncertainties of the gluon densities and of the $k_T$ dependence of the wave functions (figure 1)

- the $Q^2$ dependence of the cross section for $Q^2 \sim 10$ GeV$^2$ which can be parameterized as $Q^{-n}$ with $n \sim 4$. The difference of $n$ from the asymptotic value of 6 is due to the $Q^2$ dependence of $\alpha_s^2(Q^2)G_N^2(x,Q^2)$ and of $I_V^2$ which are equally important in this $Q^2$ range.

We discussed above (see also section 8) that the perturbative regime should dominate in the production of transversely polarized vector mesons as well, though at higher $Q^2$. This may be manifested in the $x$-dependence of the ratio $\frac{d\sigma}{d\cos \theta}$ for fixed $Q^2$. At intermediate $Q^2 \sim 10$ GeV$^2$ where hard physics already dominates in $\sigma_L$, $\sigma_T$ may still be dominated by soft nonperturbative contributions. For these $Q^2$ the ratio should increase with decreasing $x \sim x^2G_N^2(x,Q^2)$. At sufficiently large $Q^2$ where hard physics dominates for both $\sigma_L$ and $\sigma_T$ the ratio would not depend on $x$.

The $t$ dependence of the cross section is given by the square of the two gluon form factor of the nucleon $G_{2n}(t)$. Practically no $t$ dependence should be present in the block of $\gamma^* + g$ interaction for $-t \ll Q^2$. Thus the $t$ dependence should be universal for all hard diffractive processes.

Experimentally the data on diffractive production of $\rho$ mesons for $Q^2 \geq 5$ GeV$^2$ [12], on photoproduction of $J/\Psi$ mesons [13] and even on neutrino production of $D^*$ mesons [14] show a universal $t$ behavior corresponding to $G_{2n}(t) = \exp(Bt)$ with $B \approx 4 \pm 5$ GeV$^{-2}$. A certain weak increase of $B$ is expected with increasing incident energy due to the so called Gribov diffusion [15], but this effect is expected to be much smaller than for soft processes. However in the limit $Q^2 = \text{const}$ and $s \to \infty$ it is natural to expect an

† This fast increase with decreasing $x$ is absent in the non-perturbative two–gluon exchange model of Donnachie and Landshoff [12] which leads to a cross section rising as $\sim x^{-0.14}$ at $t = 0$ and to a much weaker increase of the cross section integrated over $t$. 
onset of a soft regime, which is characterized both by a slowing down of the increase of the cross section with increasing \(s\) and by a faster increase of the slope \(B\) with \(s\),

\[
\frac{\partial \ln B}{\partial \ln s} \bigg|_{s=\infty, Q^2=\text{const}} \approx \alpha_s^2 \ln \frac{s}{x_f} \approx 0.25 \text{GeV}^{-2}.
\]  

(14)

For further discussion see section 10.

We want to point out that for \(M_X^2 < Q^2\), the effect of QCD radiation is small. This is because bremsstrahlung corrections due to radiation of hard quarks and gluons are controlled by the parameter \(\alpha_s \ln \frac{1}{x_f}\) which is small since in the reaction considered here \(x_i \sim x_f\). This argument can be put on a formal ground within the double logarithmic approximation where only terms \(\sim \alpha_s \ln \frac{1}{x_f} \ln \frac{Q^2}{\Lambda^2}\) are taken into account. One can consider a more traditional approximation where terms \(\sim \alpha_s \ln \frac{Q^2}{\Lambda^2}\) are taken into account but terms \(\sim \alpha_s\) are neglected. Within these approximations it is legitimate to neglect the contribution of the longitudinal momentum as compared to the transverse one. This is a special property of small \(x\) physics. Thus the difference between \(x_i\) and \(x_f\) leads to an insignificant correction.

Formula (11) correctly accounts for nonperturbative physics and for the diffusion to large transverse distances characteristic for Feynman diagrams, because in contrast to the naive applications of the BFKL Pomeron the diffusion of small size configurations to large transverse size is not neglected.

Electroproduction of \(J/\Psi\) mesons has been calculated in the whole \(Q^2\) range in [10] within the leading \(\alpha_s \ln x\) approximation of QCD for the interaction with a target and the nonrelativistic charmonium quark model for \(J/\Psi\)-meson wave function. If we would apply (eq. (11)) at \(Q^2 = 0\) the result of ref. [10] coincides with the nonrelativistic limit of our result if \(t_V\) is assumed to be equal 1. At the same time the inclusion of the transverse momentum distribution of \(c\) quarks in the \(J/\Psi\) wave function significantly suppresses the cross section of the diffractive electroproduction of \(J/\Psi\) mesons for \(Q^2 \leq m_{j/\Psi}^2\). In particular, account of the quark Fermi motion within the model of ref. [10] using realistic charmonium models lead to suppression of photoproduction cross section by a factor 4 \(\div 8\) depending on the charmonium model (see discussion in ref. [10]). Remember that transverse distances essential in the photoproduction of the \(J/\Psi\) meson are \(\sim \frac{1}{m_c}\) which are comparable to the average interquark distance in the \(J/\Psi\) wave function. Since the energy dependence of diffractive photoproduction of \(J/\Psi\) is consistent with pQCD prediction of [10] the disagreement with the absolute prediction may indicate an important role for the interaction with interquark potential. Note that in the limit where it is possible to justify the application of PQCD (eq. (11)) \((m_{j/\Psi}^2 \ll Q^2)\) it is necessary to use distribution of bare \(c\) quarks within \(J/\Psi\) meson instead of charmonium wave functions to account for the screening of color fields of \(c\) quarks.

Another interesting process which can be calculated using the technique discussed above is the production of vector mesons in the process \(\gamma_p^* + p \rightarrow V + X\) in the triple Reggeon limit when \(-t \geq \text{few GeV}^2\) and \(-t \ll Q^2\). In this kinematical domain the dominant contribution is due to the scattering of the two gluons off a parton of the target \(g + g + \text{parton} \rightarrow \text{parton}\). To avoid the uncertainties related to the vector meson wave function it is convenient to normalize the cross section of this process to that of the exclusive vector meson production at \(t = 0\) [17, 1]

\[
\frac{d\sigma_{\gamma_p^* + p \rightarrow V + X}}{dt} \bigg|_{t=0} = \frac{9}{8\pi} \alpha_s^2 \left| \ln \frac{Q^2}{k^2} \right|^2 \times \int \frac{[G_p(y', k^2) + \frac{32}{81} S_p(y', k^2)] dy'}{[x G_p(x, Q^2)]^2},
\]

(15)

where \(S_p\) is the density of charged partons in the proton, \(\nu = 2m_N q_0, x = Q^2/\nu, k^2 = -t, y = -t/(2(q_0 - p_{vo}) m_N + pv_{vo}\) the energy of the vector meson and all variables are defined in the nucleon rest frame.

It follows from equation (15) that the cross section of the process \(\gamma_p^* + p \rightarrow V + X\) should decrease very weakly with \(t\) and therefore it is expected to be relatively large at \(-t \sim \text{few GeV}^2\). Similarly to the approach taken in [18, 19] one can easily improve equation (15) to account for leading \(\alpha_s \ln x\) terms. Equation (15) is a particular case of the suggestion (and of the formulae) presented in reference [18], that semi-exclusive large \(t\) diffractive dissociation of a projectile accompanied by target fragmentation can be expressed through the parton distributions of the target. The advantage of the process considered here as compared to the general case is the possibility to prove the dominance of hard PQCD physics for a longitudinally polarized photon as the projectile and the lack of \(t\) dependence in the vertex \(\gamma^* + g \rightarrow g + V\). These advantages allow to calculate the cross section without free parameters.

Production of transversely polarized vector mesons by real or virtual photons in the double diffractive process \(\gamma^* + p \rightarrow V + X\) has been calculated recently within the approximation of the BFKL Pomeron in [20]. The calculation was performed in the triple Reggeon limit for large \(t\) but \(s \gg -t\). Contrary to reactions initiated by longitudinally polarized photons this calculation is model dependent; the end point
nonperturbative contribution to the vertex \( \gamma^*_T + g \rightarrow g + V \), and therefore to the whole amplitude, leads to a contribution which is not under theoretical control. This problem is familiar to the theoretical discussions of high \( Q^2 \) behavior of electromagnetic form factors of hadrons.

Perturbative QCD predicts also approximate restoration of \( SU(3) \) symmetry in the production of vector mesons at large \( Q^2 \) and significant enhancement of the production of \( J/\Psi \) meson as compared to \( SU(4) \):

\[
\rho^0 : \omega : \phi : J/\Psi = 9 : 1 : (2 \ast 1.0) : (8 \ast 1.5) .
\]

This prediction is valid for \( Q^2 \gg m_V^2 \) only. Pre-asymptotic effects are important in the large \( Q^2 \) range. They significantly suppress the cross section for production of charmonium states (see above discussion). Thus the value of the \( J/\Psi/\rho \) ratio would be significantly below the value given by eq. (16) up to very large \( Q^2 \). For example the suppression factor is \( \sim 1/2 \) for \( Q^2 \sim 100 \text{GeV}^2 \). At the same time it is likely to change very little the predictions for \( \rho, \omega, \phi \)-meson production, since the masses of these hadrons are quite close and their \( q\bar{q} \) components should be very similar.

At very large \( Q^2 \) the \( q\bar{q} \) wave functions of all mesons converge to a universal asymptotic wave function with \( \eta_V = 3 \). In this limit further enhancement of the heavy resonance production is expected

\[
\rho^0 : \omega : \phi : J/\Psi = 9 : 1 : (2 \ast 1.2) : (8 \ast 3.4) .
\]

It is important to investigate these ratios separately for the production of longitudinally polarized vector mesons where hard physics dominates and for transversely polarized vector mesons where the interplay of soft and hard physics is more important.

Equation (16) is applicable also for the production of excited vector meson states with masses \( m_V \) satisfying the condition that \( m_V^2 \ll Q^2 \). In this limit it predicts comparable production of excited and ground states. There are no estimates of \( \eta_V \) for these states but it is generally believed that for \( \rho', \omega' \) and \( \phi' \) it is close to the asymptotic value, and as a rough estimate, we will assume that \( \eta_V = \eta_V' \). Using the information on the decay widths from the Review of Particle Properties [21] we find that

\[
\begin{align*}
\rho(1450) : \rho^0 & \approx \omega(1420) : \omega \approx 0.3 \\
\rho(1700) : \rho^0 & \approx \omega(1600) : \omega \approx 1.0 \\
\phi(1680) : \phi & \approx 0.6, \quad \Psi' : J/\Psi \approx 0.5 .
\end{align*}
\]

In view of substantial uncertainties in the experimental widths of most of the excited states and substantial uncertainties in the values of \( \eta_V' \), and the ratio \( \frac{\Gamma_V}{\Gamma_{\eta_V}} \) these numbers can be considered as good to about a factor of 2. The case of \( \Psi' \) where \( \Gamma_V \) is well known is less ambiguous. In this case estimates using charmonium models indicate a significant suppression as compared to the asymptotic estimate up to \( Q^2 \sim 20 \text{GeV}^2 \) where this suppression is \( \sim 0.5 \). In spite of these uncertainties it is clear that a substantial production of excited resonance states is expected at large \( Q^2 \) at HERA. A measurement of these reactions may help to understand better the dynamics of the diffractive production as well as the light-cone minimal Fock state wave functions of the excited states. It would allow also to look for the second missing excited \( \phi \) state which is likely to have a mass of about 1900 MeV to follow the pattern of the \( \rho, \omega, J/\Psi \) families.

The predicted relative yield of the excited states induced by virtual photons is expected to be higher than for real photons. Another interesting QCD effect is that the ratio of the cross section for the diffractive production of excited and ground states of vector mesons should increase with decreasing \( x \) and increase \( Q^2 \). This is because the energy denominator \( \frac{m_V^2 + Q^2}{(1-x)} \) relevant for the transition \( V \rightarrow q\bar{q} \) (with no additional partons) should be large and positive. Thus the heavier the excited state, the larger Fermi momenta should be important. Thus the gluon distributions should enter at larger virtualities in the case of \( V' \) production.

To summarize, the investigation of exclusive diffractive processes appears as the most effective method to measure the minimal Fock \( q\bar{q} \) component of the wave functions of vector mesons and the light-cone wave functions of any small mass hadron system having angular momentum 1. This would be very helpful in expanding methods of lattice QCD into the domain of high energy processes.

4. Electroproduction of photons.

The diffractive process \( \gamma^* + p \rightarrow \gamma + p \) offers another interesting possibility to investigate the interplay between soft and hard physics and to measure the gluon distribution in the proton. We shall consider the forward scattering in which case only the transverse polarization of the projectile photon contributes to the cross section. This follows from helicity conservation. In this process, in contrast to reactions initiated by longitudinally polarized highly virtual photons, soft (nonperturbative) QCD physics is not suppressed. As a result, theoretical predictions are more limited. Within QCD one can calculate unambiguously only the derivative of the amplitude over \( \ln \frac{Q^2}{m_V^2} \) but not the amplitude itself. However for sufficiently small \( x \) and large \( Q^2 \), when \( \alpha_s(Q^2) \ln \frac{Q^2}{m_V^2} \ln x \) is large, PQCD predicts the asymptotic behavior of the whole
amplitude.

It is convenient to decompose the forward scattering amplitude for the process $\gamma^* + p \rightarrow \gamma + p$ into invariant structure functions in a way similar to the case of deep inelastic electron-nucleon scattering. Introducing the invariant structure function $H(x, Q^2)$, an analogue of $F_1(x, Q^2)$ familiar from deep inelastic electron scattering off a proton, we have

$$\frac{d\sigma}{dt} \bigg|_{t=0} = \pi \alpha_s^2 \frac{H(x, Q^2)^2}{s^2} .$$

(19)

When $Q^2$ is sufficiently large, QCD allows to calculate the $Q^2$ evolution of the amplitude in terms of the parton distributions in the target. As in the case of deep inelastic processes it is convenient to decompose $H(x, Q^2)$ in terms of photon scattering off flavors of type $i$

$$H(x, Q^2) = \sum_i c_i^2 h_i(x, Q^2) ,$$

(20)

where the sum runs over the different flavors $i$ with electric charge $e_i$. It is easy to deduce the differential equation for $h_i$, the analogue of the evolution equation for the parton distributions.

$$\frac{dh_i(x, Q^2)}{d \ln Q^2} = \frac{\alpha_s(Q^2)}{2\pi} \int \frac{d z}{z} \left[ P_{gq} \left(\frac{x}{z}\right) G_p(z, Q^2) + P_{qq} \left(\frac{x}{z}\right) q_i(z, Q^2) \right] \left[ 1 + \frac{x}{z} \left(1 - \frac{x}{z}\right) \right] + O(\alpha_s^2) .$$

(21)

Here $P_{gq}$ and $P_{qq}$ are the splitting functions of the GLDAP evolution equation. The factor $1 + \frac{x}{z} \left(1 - \frac{x}{z}\right)$ takes into account the difference of the virtualities of the initial and final photon. The solution of this equation is

$$h_i(x, Q^2) = h_i(x, Q_0^2) + \frac{\alpha_s(Q_0^2)}{2\pi} \ln \frac{Q^2}{Q_0^2} \int \frac{1}{x} \frac{d x}{z} \left[ P_{gq} \left(\frac{x}{z}\right) G_p(z, Q_1^2) + P_{qq} \left(\frac{x}{z}\right) q_i(z, Q_1^2) \right] \left[ 1 + \frac{x}{z} \left(1 - \frac{x}{z}\right) \right] + O(\alpha_s^2) .$$

(22)

Usually it is assumed that the soft components of the parton distributions increase at small $x$ more slowly than the hard ones. If this is the case, at sufficiently small $x$, in the leading $\alpha_s \ln x$ approximation, the first term in equation (22) can be neglected. As a result one can obtain the asymptotic formula for the whole $H(x, Q^2)$ and not only for its derivative.

Similarly to the case of electroproduction of photons it is not difficult to generalize the $Q^2$ evolution equation to the amplitude for the diffractive production of transversely polarized vector mesons. One of the consequences of this evolution equation is that, at asymptotically large $Q^2$ and small $x$, the production cross section has the same dependence on the atomic number of a target as in the case of longitudinally polarized vector mesons.

5. Coherent Pomeron.

It is interesting to consider high-energy hard processes in the diffractive regime with the requirement that there is a large rapidity gap between the diffractive system containing the high $p_t$ jets and the target which can remain either in the ground state or convert to a system of hadrons. In PQCD such a process can be described as an exchange of a hard gluon accompanied by a system of extra gluons which together form a color neutral state. It was predicted that such processes should occur in leading twist. (Note that in reference it was stated that this process should rather be a higher twist effect. This statement was due to some specific assumptions about the properties of the triple Pomeron vertex).

The simplest example is in the triple Reggeon limit the production of high $p_t$ jets in a process like

$$h + p \rightarrow jet_1 + jet_2 + X + p$$

(23)

where the final state proton carries practically the whole momentum of the initial proton. The initial particle can be any particle including a virtual photon. To probe the new PQCD hard physics the idea is to select a final proton with a large transverse momentum $k_t$. One can demonstrate that this selection tends to compress initial and final protons in small configurations at the moment of collision. In this case the use of the PQCD two gluon exchange or two–gluon ladder diagrams becomes legitimate. A nontrivial property of these processes is a strong asymmetry between the fractions of the target momentum carried by the two gluons (the contribution of the symmetric configurations is a higher twist effect with the scale determined by the invariant mass of the produced two jets). Thus one expects gluon bremsstrahlung to play a certain role. However since the proton is in a configuration of a size $\sim 1/k_t$ this radiation is suppressed by the small coupling constant: $\sim \alpha_s(k_t^2) \ln(\frac{Q^2}{k_t^2})$. When $k_t$ tends to 0 this radiation may suppress significantly the probability of occurrence of events with large rapidity gaps.

The prediction is that such a process appears as a leading twist effect.

$$\frac{d\sigma}{dp_t^2} \sim \frac{1}{p_t^1} .$$

(24)

This prediction is in an apparent contradiction with a naive application of the factorization theorem in QCD which states that the sum of the diagrams with such soft gluon exchanges cancels in the inclusive cross section.
However in reaction \[23\] we selected a certain final state with a white nucleon hence the usual proof of the factorization theorem does not hold anymore — there is no cancelation between absorption and radiation of soft gluons \[24\]. This conclusion was checked in a simple QED model with scalar quarks \[20\].

It was suggested by Ingelman and Schlein \[27\] to consider scattering off the Pomeron as if the Pomeron were an ordinary particle and to define parton distributions in the effective Pomeron. In this language the mechanism of hard interaction in diffraction discussed above would contribute to the parton distribution in the Pomeron a term proportional to

\[
\delta(1 - x) \quad \text{or} \quad \frac{1}{(1 - x)}. \tag{25}
\]

This term corresponds to an interaction in which the Pomeron acts as a whole. Hence the term coherent Pomeron. In this kinematical configuration the two jets carry practically all the longitudinal momentum of the Pomeron. The extra gluon bremsstrahlung discussed previously renders the \(x\) dependence somewhat less singular at \(x \to 1\) but the peak should be concentrated at large \(x\) \[25, 24, 26\]. There are no other known mechanisms generating a peak at large \(x\). The recent UA(8) data \[28\] on the reaction \(p + \bar{p} \to jet_1 + jet_2 + X + p\), with the proton transverse momentum in the range \(2 \text{ GeV}^2 \geq k_1^2 \geq 1 \text{ GeV}^2\), seem to indicate that a significant fraction of the two jet events corresponds to the \(x \sim 1\) kinematics. It is thus possible that the coherent Pomeron contributes significantly to the observed cross section \[2\].

The prediction is that the contribution of the coherent Pomeron to diffractive electroproduction of dijets at \(p_t^2 \gg Q^2\) should be suppressed by an additional power of \(Q^2\)

\[
\frac{d\sigma^{\gamma^*p \to 2jet_x + X + p}}{dp_t^2} \sim \frac{1}{p_t^4} \frac{1}{Q^2}
\]
as compared to

\[
\frac{d\sigma^{\gamma^*p \to 2jet_x + X}}{dp_t^2} \sim \frac{1}{p_t^4}
\]

for other hard processes originating from the hard structure of the virtual photon.

The complicated nature of the effective Pomeron should manifested itself in several ways in hard diffraction \[18, 24\].

(i) There should be a significant suppression of the\[†\] coherent Pomeron mechanism at small \(t\) due to screening (absorptive) effects since at small \(t\) the nucleon interacts in an average configuration. This suppression should be larger for \(pp\) scattering than for \(\gamma p\) scattering since absorptive corrections increase with the increase of the total cross section (for \(\gamma p\) interaction the VDM effective total cross section at HERA energies is \(\leq 30\) mb).

(ii) Due to the contribution of soft physics, the effective Pomeron structure function as determined from the low \(t\) diffractive processes should be softer than for large \(t\) diffraction.

Therefore it would be very important to compare hard diffractive processes induced by different projectiles and to look for deviations from the predictions based on the simplest assumption that the Pomeron has an universal parton distribution \[31\].

6. Forward electroproduction of jets.

Forward diffractive photo and electroproduction of high \(p_t\) jets off a nucleon target (in the photon fragmentation region) \(\gamma^* + N \to jet_1 + jet_2 + N\) is another promising process to investigate the interplay of soft and hard physics. We shall confine our discussion to the kinematical region

\[
\frac{-\langle r_N^2 \rangle t_{\text{min}}}{3} \approx \left( \frac{Q^2 + M_{\bar{q}q}^2}{2q_0} \right)^2 \left( \frac{r_N^2}{3} \right) \ll 1, \tag{26}
\]

where

\[
M_{\bar{q}q}^2 = \frac{(m_q^2 + p_t^2)}{z(1 - z)} \tag{27}
\]
is the square of the invariant mass of the produced \(q\bar{q}\) system, \(m_q\) is the mass of quarks and \(z\) is the fraction of photon momentum carried by the \(q\) or \(\bar{q}\). In this regime the coherence of the produced hadron states allows to express the amplitude through the gluon distribution in the target.

An interesting effect occurs in the photoproduction. The contribution of a single Feynman diagram with the 2 gluon exchange in the \(t\) channel contains terms \(R_1 \approx \frac{p_t}{p_t + M_{\bar{q}q}}\) and \(R_2 \approx \frac{m_q}{p_t}\). Here is the mass of a bare quark, \(M\) can be calculated through \(m\) in pQCD but in general accounts for the nonperturbative physics. We omit constants and \(\sigma\) matrixes in this dimensional estimate and restrict ourselves to the contribution of large \(p_t\) only. A cancelation occurs when the sum of diagrams is considered. It accounts for the fact that the sum of diagrams describes the scattering of a colorless dipole.

Naively we should expect that after cancelation \(R_1\) term should become \(R_1 \approx \frac{p_t}{(p_t + M_{\bar{q}q})^2}\). But in reality it becomes \(R_1 \approx \frac{M_{\bar{q}q}^2}{(p_t + M_{\bar{q}q})^2} \approx \frac{1}{p_t^2}\). \(R_2\) term after cancelation in the sum of diagrams becomes \(R_2 \approx \frac{m_q}{p_t}\).
Thus cross section of forward photoproduction of $q\bar{q}$ pair $d\sigma/dtdp^2$ contains terms: $\frac{M^2}{p_t^2}$, $\frac{M^4}{p_t^4}$ and $\frac{M^2}{p_t^2}$.

Since mass of light quark is small it is reasonable to put it 0. It is not legitimate to put $M = 0$. So expected asymptotical behavior is $\frac{M^4}{p_t^4}$. Thus photoproduction of charm should dominate hard diffractive photoproduction processes for $p_t \geq m_c$. [30]

Photoproduction of high $p_t$ jets originating from the fragmentation of light flavors is predominantly due to next to leading order processes in $\alpha_s$.

The diffractive electroproduction of dijets seems to be the dominant process in the region of $M^2_{q\bar{q}} \leq Q^2$, while in the region $M^2_{q\bar{q}} > Q^2$ exclusive dijet production is one of many competing processes contributing to the diffractive sector like radiation of gluons from quark and gluon lines.

In the approximation when only leading $\alpha_s$ in $\frac{Q^2}{\Lambda_{QCD}^2}$ or leading $\alpha_s$ in $x \ln \frac{Q^2}{\Lambda_{QCD}^2}$ terms are kept, the off mass shell effects in the amplitude for the $q\bar{q}$ interaction with a target are unimportant. Therefore the total cross section of diffractive electroproduction of jets by longitudinally polarized photons can be calculated by applying the optical theorem for the elastic $q\bar{q}$ scattering off a nucleon target and equation (1) for the total cross section of $q\bar{q}$ scattering off a nucleon:

$$\sigma(\gamma^*\gamma^* + N \rightarrow jet_1 + jet_2 + N) = \frac{1}{16\pi B} \int \psi^2_{\gamma^*\gamma^*}(z, b) \cdot (\sigma(b^2)) dzd^2b$$

(28)

Here B is the slope of the two gluon form factor discussed in section 4 and $\psi_{\gamma^*\gamma^*}(z, b)$ is the wave function of the longitudinally polarized photon. Essentially the same equation is valid for the production by transversely polarized virtual photons of two jets which share equally the momentum of the projectile photon.

In Ref.[32] it has been assumed that diffractive production of jets off a proton is dominated by hard physics and that soft physics is unimportant. The formulae obtained under this assumption resembles equation (28) but with the gluon distribution in a target calculated within the leading $\alpha_s \ln \frac{1}{x}$ approximation. In view of the nontrivial interplay of soft and hard physics of large longitudinal distances this approach is difficult to justify in QCD. To visualize this point let us consider the effect of nuclear shadowing in diffractive electroproduction of jets. If the assumption that hard PQCD dominates at each stage of the interaction were correct, nuclear shadowing should be numerically small and suppressed by a power of $Q^2$. On the contrary, in QCD at sufficiently small $x$ and fixed $Q^2$ nuclear shadowing is expected to be substantial and universal for all hard processes. This conclusion is supported by current data on nuclear shadowing in deep inelastic processes.

Dijet production has been also considered in the constituent quark model of the proton [33, 34]. In this approach the cross section for diffraction is expressed through a convolution of the quark distribution in the virtual photon, the distribution of constituent quarks in the proton and their interaction cross section. A later generalization of this model [34] includes the gluon field of constituent quarks. In QCD though, hard processes have to be expressed in terms of bare partons and not constituent ones. This is due to the use of completeness of the intermediate hadronic states in hard processes.

Equation (28) implies that in this higher twist effect the contribution of large $b$, that is of the nonperturbative QCD, is enhanced as compared to the large $b$ contribution to the total cross section. This result has been anticipated in the pre-QCD times [35] and has been confirmed in QCD [36]. A similar conclusion has been reached in the constituent quark model [34] approach which however ignores characteristic for QCD increase of parton distributions at small $b$. In QCD the hard contribution may become dominant only at rather small $x$ and large $Q^2$. A similar conclusion has been reached for the cross section of diffractive processes, calculated in the approximation of the BFKL Pomeron [37] in the triple Reggeon region when the mass of the produced hadronic system is sufficiently large $M^2 \gg Q^2$.

Note that PQCD diagrams which were found to dominate in the large mass diffraction [37] are different from those expected from the naive application of the BFKL Pomeron [32, 34] and lead to different formulæ.

To calculate this process within the more conventional leading $\alpha_s \ln \frac{1}{x}$ approximation it is necessary to realize that in the kinematical region where $M^2_{q\bar{q}} \sim Q^2$ the fractions of nucleon momentum carried by the exchanged gluons are strongly different, $x_{\text{hard}} \simeq 2x$ but $x_{\text{soft}} \ll x$. This is qualitatively different from the case of the vector meson production considered in section 3 in which the two values of $x$ of the gluons were comparable. This is because in the case of dijet production the masses of the intermediate states are approximately equal to the mass of the final state. As a result of the asymmetry of the two $x$ values the overlap integral between the parton wave functions of the initial and final protons cannot be expressed directly through the gluon distribution in the target. However at sufficiently small $x$ and large $Q^2$, when the parameter $\frac{Q^2}{x} \ln x \ln \frac{Q^2}{x} \sim 1$, electroproduction of high $p_t$ dijets can be expressed through the gluon distribution in a target but in a more complex way.

In this particular case the factorization theorem can be applied after the first two hard rungs attached to the photon line, which have to be calculated exactly. The lower part of the diagram can be then expressed through the gluon distribution in the target since the asymmetry
between the gluons becomes unimportant in the softer blob. The proof is the same as for the vector meson electroproduction. The cross section is proportional to
\begin{equation}
\frac{d\sigma}{dt}(\rho_{g\gamma^* N} \to q\bar{q}) \sim \left( \frac{\alpha_s(Q^2)}{Q^2} \right)^2 G_N(\tilde{x}, Q^2) \tilde{x}, \end{equation}
where \(\tilde{x}\) is the average of the quark distribution functions in the photon jet amplitude, \(\tilde{x} \gg x\), and \(A_{\gamma^* + gg\to jet_1 + jet_2}\) is the hard scattering amplitude (which includes at least 2 hard rungs) calculated in PQCD.

One of the nontrivial predictions of QCD is that the decomposition of the cross section for a longitudinally polarized photon in powers of \(Q^2\) becomes inefficient at small \(x\). This is because additional powers of \(1/Q^2\) are compensated to a large extent by the increase with \(Q^2\) of \(\left[ \alpha_s(Q^2)xG(x, Q^2) \right]^2 \sim \frac{Q^2}{x}\) (see equations (29), (30)). Thus the prediction of QCD is that electroproduction of hadron states with \(M_N^2 \ll Q^2\) by longitudinally polarized photons, formally a higher twist effect, should in practice depend on \(Q^2\) rather mildly. The contribution of such higher twist effects to the total cross section for diffractive processes may be considerable, as high as 30–40%. One of the observed channels, the electroproduction of \(\rho\) mesons, constitutes probably up to 10% of the total cross section for diffractive processes. So far a detailed quantitative analysis of this important issue is missing. On the experimental side, it would be extremely important to separate the longitudinal and transverse contributions to diffraction.

7. Can diffractive cross sections raise forever?

We have demonstrated above that cross sections of hard diffractive processes are related to cross section of interaction of small color dipole with the target which increase fast with incident energy. However such increase cannot be sustained forever. The simplest way to obtain an upper limit for the range of energies where such increases should stop is to consider here the scattering of a small object, a \(q\bar{q}\) pair, from a large object, a nucleon. If only hard physics was relevant for the increase of parton distributions at small \(x\), the radius of a nucleon should not increase (small Gribov diffusion). Under this assumption the unitarity limit corresponds to a black nucleon. In this case the inelastic cross section cannot exceed the geometrical size of the nucleon
\begin{equation}
\sigma(g\bar{q}N) = \frac{\pi^2}{3} b^2 \alpha_s(1/b^2)xG_N(x, b^2) < \pi r_N^2. \quad (30)
\end{equation}

To find the value of \(r_N\) in eq. (30) we use the optical theorem to calculate the elastic cross section for a \(q\bar{q}\) pair scattering off a nucleon,
\begin{equation}
\sigma_{el} = \frac{\sigma_{tot}}{16\pi B} \quad (31)
\end{equation}
where \(B\) is the slope of the elastic amplitude (cf. discussion in section 4). It follows from Eqs. (30), (31) and condition that \(\sigma_{inel} + \sigma_{el} = \sigma_{tot}\) that the unitarity limit is achieved when the elastic cross section is equal to the inelastic cross section \(\sigma_{el} \leq \sigma_{inel}\). Based on this we find \(r_N^2 = 4B \approx 16 \text{ GeV}^{-2} \approx (0.8 \text{ fm})^2\) is the radius of a nucleon. It follows from the above equations that practically the same estimate is obtained from the assumption that \(\frac{\sigma_{el}}{\sigma_{tot}} \sim (0.3 - 0.5)\).

Applying these inequality for the cases of \(\sigma_L\) and \(\rho\)-meson production we find \(B\) that unitarity limit is reached for \(x_{\sigma_L}(Q^2 = 5\text{ GeV}^2) \sim 3 \times 10^{-5}\), \(x_{\sigma_L}(Q^2 = 10\text{ GeV}^2) \sim 6 \times 10^{-6}\), \(x_{\rho}(Q^2 = 5\text{ GeV}^2) \sim 3 \times 10^{-4}\), \(x_{\rho}(Q^2 = 10\text{ GeV}^2) \sim 2 \times 10^{-4}\).

The use of the amplitude for \(q\bar{q}\) pair scattering off a nucleon to deduce the limit allows to account accurately for nonperturbative QCD effects through the unitarity condition for such an amplitude. On the other hand if the increase of parton distributions is related to soft physics as well then the cross section may be allowed to increase up to smaller \(x\) values.

The black disc limit for \(\sigma_{\gamma N}\) has been discussed earlier (for a review and references see [8, 9]). The difference compared to previous attempts is that we deduce the QCD formulae for the cross section of a \(q\bar{q}\) pair scattering off a hadron target. For this cross section the geometrical limit including numerical coefficients unambiguously follows from unitarity of the \(S\)-matrix, that is the geometry of the collision. As a result we obtain an inequality which contains no free parameters. Recently a quantitative estimate of the saturation limit was obtained [10] by considering the GLR model of nonlinear effects in the parton evolution and requiring that the nonlinear term should be smaller than the linear term. The constraint obtained for \(xG_p(x, Q^2)\) is numerically much less restrictive compared to our result. Even a more stringent restriction follows for the interaction of a colorless gluon pair off a nucleon from the requirement that the inelastic cross section for the scattering of a small size gluon pair should not exceed the elastic one
\begin{equation}
\sigma(ggN) = \frac{3\pi^2}{4} b^2 \alpha_s(1/b^2)xG_N(x, b^2) < \pi r_N^2. \quad (32)
\end{equation}

For \(b = 0.25 \text{ fm}\) the geometrical limit is achieved for \(x \sim 10^{-3}\).

We want to point out that the black disc limit implies a restriction on the limiting behavior of the cross sections for hard processes but does not allow to
calculate it. The dynamical mechanism responsible for slowing down of the increase of parton distributions so that they satisfy equations (33) is not clear. In particular the triple Pomeron mechanism for shadowing suggested in [14] does not lead to large effects at HERA energies especially if one assumes a homogeneous transverse density of gluons [13].

The theoretical analysis performed in this section does not allow to deduce restrictions on the limiting behavior of parton distributions in a hadron. Beyond the evolution equation approximation and/or leading \( \alpha_s \ln x \ln Q^2 \) over \( \alpha_s^2 QCD \) approximation the restriction on the cross sections of deep inelastic processes cannot be simply expressed in terms of parton distributions in a hadron target.

We want to draw attention to the fact that nonperturbative QCD effects play an important role in the contribution of higher twist effects to \( \sigma_L(\gamma^* p) \). This is evident from the impact parameter representation of the contribution to \( \sigma_L(\gamma^* + p) \) of \( n \) consecutive rescatterings of small transverse size \( q \bar{q} \) pairs. This contribution is proportional to

\[
Q^2 \int |\psi_{\perp L}(z, b^2)|^2 dz d^2b \left[ \alpha_s(1/b^2)b^2 x \sigma_G(x, b) \right]^n.
\]

The inspection of this integral shows that for large \( n \geq 3 \), \( b \) which dominates under the integral does not decrease with increasing \( Q^2 \) for \( x \sim 10^{-3} \div 10^{-4} \). We use as estimate \( \alpha_s x \sigma_G(x, Q^2) \propto \sqrt{Q} \) which follows from the evolution equation for small \( x \). (This QCD effect is absent in the applications [34] of the constituent quark model). Thus if higher twist effects were really important in small \( x \) physics, it would imply that the small \( x \) physics is the outcome of an interplay of hard (small \( b \)) and soft (large \( b \)) QCD. To illustrate this point let us consider the cross section of diffractive electroproduction of hadrons with masses \( M^2 \sim Q^2 \) by transversely polarized photons. Applying the same ideas as in the case of longitudinally polarized photons we would obtain a similar expression as given by equation [23]. The important difference is that the wave function of a transversely polarized photon is singular for \( z \to 0 \) or 1. As a result the contribution of large impact parameters \( b \) in the wave function of the photon should give the dominant contribution to the integral in a wide kinematical range of \( x \) and \( Q^2 \). This has been understood long ago – see discussion in sections 8-9. A similar conclusion has been achieved recently [34] within the constituent quark model. (Note however that this model ignores the increase of gluon distribution with \( Q \) typical for QCD and therefore overestimates the nonperturbative QCD contribution). Thus such type of diffractive processes should depend on energy in a way similar to the usual soft hadron processes.

A good example of the consequences of the interplay of small \( b \) and large \( b \) physics is that in electroproduction of small mass states the unitarity limit may become apparent at larger \( x \) than in the case of the total cross section of deep inelastic processes.

8. Diffraction in DIS at intermediate \( Q^2 \)

It has been understood long ago that the production of almost on mass shell quarks by virtual photons should give a significant contribution to the total cross section for deep inelastic scattering at small \( x \) [14]. One of the predictions of this approach (which is essentially the parton model approximation) is a large cross section for diffractive processes. The QCD \( Q^2 \) evolution does not change this physical picture radically. The only expected modification of the picture is the appearance of a number of hard jets in the current fragmentation region [23] typical for including \( \alpha_s \ln Q^2 \) terms. It is often stated that the dominance of the BFKL Pomeron in diffractive processes predicts the dominance of final states consisting of hard jets [23, 10]. However this prediction is not robust since the analysis of Feynman diagrams for hard processes in QCD finds strong diffusion effects into the region of small transverse momenta of partons (see [13] and references therein). Recent HERA data [13] seem to support the picture with a dominance of events with small \( k_t \). Thus it seems worthwhile to investigate the role of nonperturbative QCD physics in diffractive processes.

The interaction of a virtual photon with a target at intermediate \( Q^2 \) and small \( x \), when gluon radiation is negligible, can be considered as a transformation of \( \gamma^* \) into a \( q \bar{q} \) pair which subsequently interacts with the target. In this case an important role is played by the quark configurations in which the virtuality of the quark interacting with the target is small,

\[
k_{qt} \sim k_{t0}, \quad \alpha_q = \frac{m_q^2 + k_{qt}^2}{Q^2}. \tag{33}
\]

Here \( \alpha_q \) denotes the light-cone fraction of the photon momentum carried by the slower quark and \( k_{t0} \) is an average transverse momentum of partons in the hadron wave function, typically \( k_{t0} \sim 0.3 - 0.4 \) GeV.

In the language of non-covariant perturbation theory the \( q \bar{q} \) configurations described by (33) correspond to an intermediate state of mass \( m^2 \sim Q^2 \) and of transverse size \( \sim \frac{1}{k_{t0}} \geq 0.5 \) fm. These configurations constitute a tiny fraction \( \sim \frac{k_{qt}^2}{Q^2} \) of the phase volume kinematically allowed for the \( q \bar{q} \) pair. However the interaction in this case is strong – similar to the interaction of ordinary hadrons, since the virtuality of the slower quark is small and the transverse area occupied by color is large. The contribution of these configurations leads to Bjorken scaling since the total cross section is proportional to
and in the parton model only these configurations may contribute to the cross section. Hence Bjorken has assumed [4] that all other configurations are not important in the interaction though the underlying dynamics of such a suppression was not clear at that time [35]. Accounting for the \( \frac{1}{q^2} \) factor in the Gribov dispersion representation allowed him to reconcile this dispersion representation with scaling. He suggested to refer to these configurations as aligned jets since both quarks have small transverse momenta relative to the photon momentum direction. In further discussions we will refer to this approach as that of the Aligned Jet Model (AJM). Note that in terms of the Feynman fusion diagram the aligned jet contribution arises only for transversely polarized virtual photons. This is because the vertex for the transition \( \gamma^*_L \rightarrow q\bar{q} \) is singular \( \sim \frac{1}{t} \) when the fraction of the photon momentum \( t \) carried by the slowest quark (antiquark) tends to 0. For the case of a longitudinally polarized photon the naive aligned jet approximation produces results qualitatively different from expectations in QCD where the contribution of symmetric jets dominates. This is because in QCD the dominant contribution to the \( \gamma^*_L \)-nucleon cross section arises from the region of large \( k_{qt} \sim \frac{Q^2}{4 \pi} \).

In QCD the interaction of quarks with large relative transverse momenta with a target is suppressed but not negligible. The suppression mechanism is due to color screening since \( q\bar{q} \) configurations with large \( k_t \) correspond, in the coordinate space, to configurations of small transverse size, \( b \sim \frac{1}{k_t} \) for which equation (1) is applicable. It is easy to check that the contribution of large \( k_t \) also gives a scaling contribution to the cross section. The practical question then is which of the two contributions dominates at intermediate \( Q^2 = Q_0^2 \approx 4 \text{ GeV}^2 \), above which one can use the QCD evolution equations. To make a numerical estimate we assume that the \( q\bar{q} \) configurations with \( k_{qt} \leq k_{t0} \), in which color is distributed over a transverse area similar to the one occupied by color in mesons, interact with a cross section similar to that of the pion. A comparison with experimental data for \( F_{2p}(x \sim 0.01, Q_0^2) \) indicates that at least half of the cross section is due to soft, low \( k_t \) interactions [34, 35]. A crucial check is provided by applying the same reasoning to scattering off nuclei in which the interaction of the soft component should be shadowed with an intensity comparable to that of pion-nucleus interaction. Indeed the current deep inelastic data on shadowing for \( F_{2A}(x, Q^2) \) are in reasonable agreement with calculations based on the soft mechanism of nuclear shadowing [34, 35].

Similarly to the case of hadron-nucleon and hadron-nucleus interactions, the interaction of \( \gamma^* \) in a soft hadron component naturally leads to diffractive phenomena. Application of the Gribov representation with a cutoff on the \( k_t \) of the aligned jets in the integral leads to a diffractive mass spectrum for the transversely polarized virtual photon

\[
\frac{d\sigma}{dM^2} \propto \frac{1}{(M^2 + Q^2)^2}.
\]

The two major differences compared to the hadronic case are that elastic scattering is substituted by production of states with \( M^2 \approx Q^2 \) and that the contribution of configurations of small spatial size is larger for \( \gamma^*_L \).

If the aligned jet configurations were dominant, the fraction of cross section of deep inelastic \( \gamma^* N \) scattering due to single diffractive processes would be

\[
R_{\text{AJM}} \approx \frac{\sigma_{\text{diff}}}{\sigma_{\text{tot}}} \approx 0.25. \tag{35}
\]

Our numerical estimates indicate that for \( Q^2 \sim Q_0^2 \) and \( x \sim 10^{-2} \) the AJM contributes about \( \eta \sim 60 - 70\% \) of the total cross section. So we expect that in this \( Q^2 \) range the probability for diffraction is

\[
R_{\text{single diff}} = \eta R_{\text{AJM}} \approx 15\%. \tag{36}
\]

This probability is actually related in a rather direct way to the amount of shadowing in interactions with nuclei in the same kinematic regime, so it is quite well determined by the nuclear shadowing data.

To estimate the probability of events with large rapidity gaps one has to add the processes of diffractive dissociation of the nucleon and double diffraction dissociation, leading to an estimate

\[
P_{\text{gap}} = (1.3 - 1.5) R_{\text{single diff}} \sim 0.2. \tag{37}
\]

This is rather close to the observed gap survival probability for photoproduction processes [47].

The characteristic features of the AJM contribution which can be checked experimentally are the charge and flavor correlations between the fastest and the slowest diffractively produced hadrons which should be similar to those in \( e^+e^- \rightarrow \text{hadrons at } M^2 \approx Q^2 \).

Another important feature of the soft contribution which distinguishes it from the contribution of hard processes is the \( t \) dependence of the cross section for \( M^2 \leq Q^2 \). Since the size of the configurations is comparable to that of the pion one may expect that the \( t \) slope of the cross section, \( B \), should be similar to that of the pion-nucleon interaction, i.e. \( B \geq 10 \text{ GeV}^{-2} \) which is much softer than for hard processes where we expect \( B \approx 4 \text{ GeV}^{-2} \) (see discussion in section 3). The large value of the slope for the soft component is also natural in the parton type logic where only slow partons interact. It is easy to check that for \( -t \gg k_{t0}^2 \sim 0.1 \text{ GeV}^2 \) the mass of the produced hadron system is
larger than the mass of the intermediate state by factor $\frac{1}{\sqrt{\frac{Q}{Q_0}}}$.
Thus for large $t$ the production of masses $M \leq Q$ is suppressed. Therefore the study of the $t$-dependence of diffraction can be used to disentangle the contribution of soft and hard mechanisms.

This discussion indicates also that the contribution of non-diagonal transitions \( M^{2n} \rightarrow M'^{2n} \) leads to a weaker decrease of the differential cross section with $M^2$ than given by equation (34). Besides at large $M^2 \sim few$ $Q^2$ one expects an onset of the dominance of the triple Pomeron mechanism which corresponds to

$$\frac{d\sigma}{dM^2} \sim \frac{1}{Q^2 M^2}. \quad (38)$$

9. \( Q^2 \) evolution of the soft contribution in diffraction.

The major difference between the parton model and QCD is the existence in QCD of a significant high $p_t$ tail in the parton wave functions of the virtual photon and the proton. This is the source of the violation of Bjorken scaling observed at small $x$. It is thus necessary to modify the aligned jet model to account for the hard QCD physics.

It is in general difficult to obtain with significant probability a rapidity gap in hard processes in perturbative physics. Confinement of quarks and gluons means that a gap in rapidity is filled by gluon radiation in PQCD and subsequently by hadrons [49]. It is possible to produce diffraction in perturbative QCD but the price is a suppression by powers of the coupling constant $\alpha_s$ and/or powers of $Q^2$. In first approximation in calculating diffraction in deep inelastic processes at small $x$ we will thus neglect diffraction in PQCD. In the following analysis, for the description of large rapidity gap events, we shall use the QCD modification of the AJM model suggested in [86] as well as the suggestion of Yu.Dokshitzer [64] to add to the conventional evolution equation the assumption of local duality in rapidity space between quark-gluon and hadron degrees of freedom.

In the course of the following considerations it will be convenient to switch to the Breit frame. In this frame the photon has momentum $(0, -2xP)$ and the initial proton has momentum $(P, P)$. Correspondingly $Q^2 = 4x^2 P^2$. The process of diffraction can be viewed as the virtual photon scattering off a color singlet $q\bar{q}$ pair with the interacting parton carrying a light-cone fraction $\alpha$ and the spectator parton carrying a light-cone fraction $x_1$. We assume here the local correspondence in rapidity space between partons and hadrons. The mass of the produced system, $M$ is given by

$$M^2 = (p_{\gamma} + p_{x_1} + p_{\alpha})^2 =$$

$$Q^2 + 4P^2(\alpha + x_1) + \frac{x}{x_1} \quad (39)$$

In the approximation that gluon radiation is neglected (parton model) $\alpha = x$ and the mass of the diffractionally produced system $M$ is

$$M^2 = Q^2 x_1 / x. \quad (40)$$

The differential cross section for production of mass $M$ follows from equation (34).

$$\Gamma \int dx_1 \delta(1 - x_1 - xM^2/Q^2) \frac{1}{(Q^2 + M^2)^2} =$$

$$\Gamma \int dx_1 \delta(x - \alpha) \delta(x_1 - xM^2/Q^2) \frac{1}{(1 + x_1/x)^2}. \quad (41)$$

Here $\Gamma$ is the factor which includes the density of correlated color singlet pairs and the cross section for interaction of the photon with the parton. The total cross section for diffractive dissociation comes out to be proportional to $\frac{1}{Q^2}$.

$$\int \frac{d\sigma^{AJM}}{dM^2} dM^2 = \frac{\Gamma}{Q^2} \int \frac{dx_1}{x} \frac{1}{(1 + x_1/x)^2} = \frac{\Gamma}{Q^2}. \quad (42)$$

We do not restrict the integration over $x_1$ in equation (42) since the major contribution comes from the region of $x_1 \sim x$. Thus we can formulate diffraction in the infinite momentum frame as a manifestation of short rapidity range color correlation between partons in the nonperturbative parton wave function of the nucleon. To calculate the $Q^2$ evolution in QCD we have to take into account that the parton with momentum fraction $\alpha$ has its own structure at higher $Q^2$ resolution and that the $\gamma^* \rightarrow$ scatter off constituents of the "parent" parton. This is the usual evolution with $Q^2$ which can be accounted for in the same way as in the QCD evolution equations by the substitution

$$\Gamma \delta(x - \alpha) \rightarrow P \sum_j c_j^2 d_{j\alpha}^{pert} \left( \frac{x}{\alpha}, Q^2, Q_0^2 \right)$$

where $d_{j\alpha}^{pert}(x, Q^2, Q_0^2)$ are the structure functions of the parent parton. This effect leads to the change of the relationship between $x_1$ and $x$ resulting from parton bremsstrahlung. After performing the integral over $x_1$ we obtain

$$\frac{d\sigma^{soft + QCD}}{dM^2} =$$

$$\frac{P}{Q^4} \int_x^1 \frac{dx}{\alpha} \sum_j c_j^2 d_{j\alpha}^{pert} \left( \frac{x}{\alpha}, Q^2, Q_0^2 \right) d_{j\alpha}^{nonpert} \left( \alpha, Q_0^2 \right)$$

$$\frac{1}{(2 - \alpha/x + M^2/Q^2)^2} \times$$

$$\theta \left( \ln(1 - x) - \frac{xM^2}{Q^2} \right) \theta \left( 1 - \frac{\alpha}{x} - \frac{M^2}{Q^2} \right). \quad (43)$$
Here $P$ denotes the probability of diffractive scattering in a soft interaction and $d_{ij}^{\text{pert}}(\alpha, Q_0^2)$ is the parton distribution in the soft component producing diffraction (compare discussion in the previous section). The $\theta$ function term reflects the condition that diffraction in the nonperturbative domain is possible only for

$$0 \leq \frac{M^2}{s} = \frac{\alpha + x_1 - x}{(1 - x)} \equiv \lambda \sim 0.05 - 0.1 \ .$$

(44)

9.1. Qualitative pattern of $x$ and $Q^2$ dependence of diffraction.

It is easy to see that discussed equations lead to the leading twist diffraction. To see the pattern of the $x, Q^2$ dependence we can assume that $d_{ij}^{\text{pert}}(x, Q^2) = \frac{d}{x}$ and $d_{ij}^{\text{nonpert}}(x, Q_0^2) = \frac{d}{2}$. It follows from equation (43) that for $x \ll \lambda$ the ratio $\frac{d_{\text{nonpert}}}{d_{\text{pert}}}$ does not depend on $x$. One can also see that the characteristic gap interval is

$$\Delta y = \ln \frac{s}{M m_p} = \ln \frac{1}{x} + \ln \left( \frac{Q^2}{M m_p} \right) .$$

(45)

The second term $\ln \frac{Q^2}{M m_p}$ increases with $Q^2$ in the parton model, while the scaling violation tends to reduce this increase since the mean value of $M^2/Q^2$ at fixed $x$ increases with $Q^2$.

There are several qualitative differences between the QCD improved soft diffraction and the parton model (AJM).

(i) Due to QCD evolution the number of diffractively produced hard jets and the average transverse momentum of diffractively produced hadrons should increase with $Q^2$.

(ii) The distribution of $M^2/Q^2$ becomes broader in QCD with increasing $Q^2$.

(iii) While in the parton model the cross section for the interaction of the longitudinally polarized virtual photon is a higher twist effect, in QCD diffraction is a leading twist for any polarization of the virtual photon. The final state in the case of longitudinally polarized photons should contain at least 3 jets, two of them should have large transverse momenta comparable with $Q$.

(iv) The $x$ dependence of the soft component is likely to be faster than for soft Pomeron as seen in $pp$ scattering both due to smaller screening corrections and due to contribution of configurations of sizes somewhat smaller than normal hadron sizes. In the laboratory frame of the target these configurations correspond to $q\bar{q}$ pairs with $p_t \sim 0.5 \div 1\text{GeV}/c$.

9.2. Connection with the Ingelman-Schlein Model

Ingelman and Schlein have suggested to treat hard diffractive processes using the concept of parton distribution in the Pomeron \cite{22}. In this approach one calculates the light-cone fraction of the target carried by the Pomeron, $x_P$, and light-cone fractions of the Pomeron momentum carried by quarks and gluons, $\beta$. It is assumed that parton distributions in the Pomeron, $\beta q_{P}(\beta, Q^2)$, $\beta g_{P}(\beta, Q^2)$ are independent of $x_P$ and the transverse momentum of the recoil nucleon. For the process of inclusive deep inelastic diffraction $\beta$ is simply related to the observables,

$$\beta = \frac{Q^2}{Q^2 + M_X^2} .$$

(46)

The $Q^2$ evolution of the total cross section of diffraction as considered in the previous subsections is consistent with the expectation of the Ingelman-Schlein model (though the final states are not necessarily the same). The aligned jet model in this case serves as a boundary condition defining parton distributions in the Pomeron at intermediate $Q_0^2$ above which QCD evolution takes place. The aligned jet model corresponds to the quark distribution in the Pomeron

$$\beta q_{P}(\beta, Q_0^2) \propto \beta .$$

(47)

It follows from the discussion in the end of section 8 that taking into account the non-diagonal transitions in the aligned jet model and the triple Pomeron contribution would make the distribution flatter. A similar, rather flat, distribution is expected for gluons for these $Q^2$. This expectation of the aligned jet model is different from the counting rule anzatz of \cite{27}: $\beta q_{P}(\beta, Q_0^2) \propto (1 - \beta)$. Note also that since the density of gluons at small $x$ is larger than the $q\bar{q}$ density the effective gluon density in the Pomeron should be larger that the charged parton density already at $Q^2 \sim Q_0^2$.

10. Non-universality of the pomeron in QCD.

Theoretical considerations of soft diffractive processes have demonstrated that ordinary hadrons contain components of very different interaction strength \cite{31}. This includes configurations which interact with cross sections much larger than the average one and configurations which interact with very small cross sections, described by equation (3) for a meson projectile.

The presence in hadrons of various configurations of partons having different interaction cross sections with a target is in evident contradiction with the idea of a universal vacuum pole where universal factorization is expected. At the same time it is well known that the Pomeron pole approximation is not self-consistent. The vacuum pole should be accompanied by a set of Pomeron cuts \cite{22}. For the sum of the Pomeron pole and the Pomeron cuts no factorization is expected. Thus the S
matrix description and the QCD description are not in variance. We shall enumerate now where and how to search for the non universality of the effective Pomeron understood as the sum of the pomeron pole and the Pomeron cuts.

It is natural to distinguish two basic manifestations of the non universality of the effective Pomeron trajectory, $\alpha_P(t) \approx \alpha_0 + \alpha'_t$, a different energy dependence of the interaction cross section, which is characterized by a different value of $\alpha_0$, and a different rate of the Gribov diffusion, which would manifest itself in different values of $\alpha'$.

10.1. Non-universality of the energy dependence.

To study the non universality of $\alpha_0$ it is necessary to study the energy dependence of the electroproduction of vector mesons as a function of $Q^2$. Up to now only two results are known, $\alpha_0 \sim 1.08$ from the $\rho$ meson photoproduction [3] and $\alpha_0 \sim 1.30$ as estimated from NMC and HERA data at $Q^2 \sim 10 \text{ GeV}^2$ [11, 12]. The key question is at what $Q^2$ a significant rise of $\alpha_0$ starts – this will give a direct information on the transition region from soft to hard physics. A fast increase of $F_{2p}(x, Q^2)$ at small $x$ observed at $Q^2$ as low as 1.5 GeV$^2$ indicates that the rise may occur already at $Q^2 \sim 3 \text{ GeV}^2$. The same question applies for production of heavier $\phi$ and $J/\Psi$ mesons. Since the $J/\Psi$ meson is a small object one may speculate that in this case the rise could start already for photoproduction (the experimental data indicate that the slope of the $J/\Psi$ exclusive photoproduction cross section is close to the value given by the two-gluon form factor of the nucleon). The practical problem for a quantitative analysis is that no accurate data on exclusive $J/\Psi$ photoproduction at fixed target energies are available at the moment. Inclusive fixed target data where the $J/\Psi$ meson carries practically the whole momentum of the projectile photon which are used to extract the exclusive channel seem to be significantly contaminated by the contribution of the reaction $\gamma + p \rightarrow J/\Psi + X$ which is peaked at $x_F \equiv p_{J/\Psi}/p_\gamma$ close to 1.

10.2. Non-universality of the $t$-dependence.

The slope of the effective Pomeron trajectory $\alpha'$ should decrease with increasing $Q^2$. This is because the Gribov diffusion in the impact parameter space, which leads to finite $\alpha'$ [13], becomes inessential in the hard regime. This is a consequence of the increase with energy of the typical transverse momenta of partons. Thus for the reactions $\gamma^* + N \rightarrow V + N$ the effective $\alpha'$ should decrease with increasing $Q^2$ while a universal Pomeron exchange approximation predicts for the energy dependence of the slope

$$B(s) = B(s_0) + 2\alpha' \ln \left( \frac{s}{s_0} \right)$$

with $\alpha' \sim 0.25 \text{ GeV}^{-2}$. It is possible to look for this effect by comparing the HERA and the NMC data on the $\rho$ meson production. The universal Pomeron model predicts that the slope should change by $\Delta B \sim 2 \text{ GeV}^{-2}$ between NMC energies where $B \sim 4 \div 5 \text{ GeV}^{-2}$ [10] and HERA energies while in the perturbative domain a much weaker change of the slope is expected.

The slope of the effective Pomeron trajectory $\alpha'$ may depend on the flavor. It should decrease with the mass of flavor. Thus it would be very important to measure the effective $\alpha'$ for the diffractive photoproduction of $\rho$, $\phi$ and $J/\Psi$. If PQCD is important for $J/\Psi$ photoproduction one would expect a smaller increase of the slope with energy in this case.

10.3. Non-universality of the gap survival probability.

The presence of configurations of different size in hadrons (photons) should also manifest itself in the non universality of the gap survival probability in the two jet events. Since the probability of gap survival is determined by the intensity of the soft interaction of the projectile with the target, the survival probability should increase with increase of $Q^2$, and at fixed $Q^2$ it should be larger for the heavy $q\bar{q}$ components of the photon. Also, the gap survival probability in the photon case should be substantially larger than that observed in $p\bar{p}$ collisions at FNAL collider [54]. This reflects the difference between $\sigma_{\text{tot}}(p\bar{p}) \approx 80 \text{ mb}$ and the effective cross section for the interaction of the hadronic components of $\gamma(\gamma^*)$ with nucleon of $\leq 30 \text{ mb}$.

Observation of non-universality discussed here will shed light on the structure of the effective Pomeron operating in strong interactions and will help to address the question about the major source of the increase of the total cross section of $p\bar{p}$ interaction — soft physics or hard physics of small size configurations.

10.4. Non-universality of diffraction dissociation

Since the object which couples to the nucleon in the hard coherent processes is different from soft Pomeron one may expect a difference between the value of the ratio $d\sigma^{*+p-\rho+X}/d\rho d\epsilon$ at $t=0$ and similar ratio for soft processes. Qualitatively, one may expect that since the coupling of effective Pomeron in hard processes is more local the ratio of diffraction dissociation and elastic cross sections should be substantially smaller for hard processes, at least for small excitation masses.
11. Summary.

We have demonstrated that color coherent phenomena should play in QCD a rather direct role both in the properties of hadrons and in the high energy collisions. It seems now that recent experimental data confirm some of the rather nontrivial predictions of QCD and help to elucidate such old problems as the origin of the Pomeron pole and the Pomeron cuts in the Reggeon Calculus. Thus we expect that the investigation of coherent hard and soft diffractive processes may be the key in obtaining a three dimensional image of hadrons, in helping to search for new forms of hadron matter at accelerators and in understanding the problem of inter-nucleon forces in nuclei. Forthcoming high luminosity studies of diffraction at HERA which will include among other things the detection of the diffracting nucleon and the $\sigma_L - \sigma_T$ separation would greatly help in these studies.

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Figure 1. The total longitudinal cross section, $\sigma_{\gamma^* N \rightarrow p N}$, calculated from Eq. (11) for several recent parameterizations of the gluon density in comparison with experimental data from ZEUS [11] (full circles) and NMC [10] (squares). Typical parameters for the $\rho$-meson wave functions as discussed above are taken ($\langle k_t^2 \rangle^{1/2} = 0.45 \text{GeV}/c$). We set $\eta_V = 3$ and parameterize the dependence of the differential cross section on the momentum transfer in exponential form with $B \approx 5 \text{ GeV}^{-2}$. Note that a change of $T^2(Q^2)$ in the range corresponding to $\langle k_t^2 \rangle^{1/2}$ between 0.3 GeV/c and 0.6 GeV/c introduces an extra scale uncertainty of $0.7 \div 1.4$. 

$$Q^2 = 8.8 \text{ GeV}^2$$

$$Q^2 = 16.9 \text{ GeV}^2$$