Direct calculation of the triple-pomeron coupling for diffractive deep inelastic scattering and real photoproduction

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We present a unified direct evaluation of the triple-pomeron coupling $A_{3\mathrm{P}}(Q^2)$ for diffractive real photoproduction ($Q^2=0$) and deep inelastic scattering at large Q^2 in the framework of the dipole approach to the generalized BFKL pomeron. We demonstrate how the reaction mechanism changes from diffraction excitation of the constituent quarks of the photon at $Q^2=0$ to excitation of the octet-octet color dipole state of the virtual photon at large Q^2. The small phenomenological value of $A_{3\mathrm{P}}(0)$, which was a mystery, is related to the small nonperturbative correlation radius $R_c \approx 0.3$ fm for the perturbative gluons. We confirm the early expectations of weak Q^2 dependence of the dimensional coupling $A_{3\mathrm{P}}(Q^2)$ and predict that it rises by the factor ~ 1.6 from real photoproduction to deep inelastic scattering. © *1995 American Institute of Physics.*

1. INTRODUCTION

A salient feature of diffraction dissociation $a+p \to X+p'$ of ($a=h$) hadrons and ($a=\gamma$) real photons ($Q^2=0$) is the so called triple-pomeron regime

$$\frac{M^2}{\sigma_{\mathrm{tot}}(ap)} \left. \frac{d\sigma_D(a\to X)}{dt dM^2} \right|_{t=0} \approx A_{3\mathrm{P}} \qquad (1)$$

with the approximately energy- and projectile-independent dimensional triple-pomeron coupling $A_{3\mathrm{P}}$, which holds at moderately high energies such that the photoabsorption and hadronic cross-sections are approximately constant (Ref. 1; for a review see Ref. 2). Here t is the (p,p') momentum transfer squared and the mass M of the excited state satisfies $m_p^2 \ll M^2 \ll W^2$, where W is the total c.m.s. energy. The FNAL data on the diffractive real photoproduction give[1)] $A_{3\mathrm{P}}(Q^2=0) \approx 0.16\ \mathrm{GeV}^{-2}$.[3] In the language of Regge theory, $A_{3\mathrm{P}}$ measures the projectile-pomeron cross-section[1]

$$\sigma_{\mathrm{tot}}(a\mathrm{P}; M^2 \gg m_p^2) = \frac{16\pi M^2}{\sigma_{\mathrm{tot}}(pp)} \left. \frac{d\sigma_D(a\to X)}{dt dM^2} \right|_{t=0}$$
$$\approx 16\pi A_{3\mathrm{P}} \frac{\sigma_{\mathrm{tot}}(ap)}{\sigma_{\mathrm{tot}}(pp)} \approx 0.08\sigma_{\mathrm{tot}}(ap). \qquad (2)$$

Why $A_{3\mathrm{P}}$ is small and $\sigma_{\mathrm{tot}}(a\mathrm{P})$ is more than an order in magnitude smaller than $\sigma_{\mathrm{tot}}(aN)$, is one of outstanding mysteries of the pomeron.

An entirely new process of diffraction dissociation of the virtual photon, $\gamma^*+p \to X+p'$, in deep inelastic scattering (DIS) for $x=Q^2/(Q^2+W^2) \ll 1$, where Q^2 is the virtuality of the photon, is being studied at HERA,[4] and soon new data will be available on the triple-pomeron region $\beta = Q^2/(Q^2+M^2) \ll 1$. In diffractive DIS, too, one can operationally define the triple-pomeron coupling

$$\frac{M^2+Q^2}{\sigma_{\mathrm{tot}}(\gamma^* p)} \left. \frac{d\sigma_D(\gamma^* \to X)}{dt dM^2} \right|_{t=0} \approx A_{3\mathrm{P}}(Q^2). \qquad (3)$$

The dimensional quantity $A_{3\mathrm{P}}(Q^2)$ is nonperturbative. The subject of the present paper is the evaluation and study of the Q^2 dependence of the Born approximation for $A_{3\mathrm{P}}(Q^2)$ in terms of the nonperturbative parameter of the dipole cross-section approach to the QCD pomeron,[5–7] the correlation radius R_c for the perturbative gluons. We find that in both real photoproduction and deep inelastic scattering at moderate energy ν and/or moderately small x, the $A_{3\mathrm{P}}(Q^2)$ is approximately the same, thus confirming earlier expectation.[5–9]

In this paper we operationally define, and calculate, $A_{3\mathrm{P}}(Q^2)$ as the normalization of the term $\propto 1/(Q^2+M^2)$ in the diffraction mass spectrum at moderately large energy and moderately large M^2. The convenient variable is the fraction x_{P} of the proton momentum taken away by the pomeron, $x_{\mathrm{P}} = (M^2+Q^2)/(W^2+Q^2) \ll 1$. The final-state proton p' carries a fraction $1-x_{\mathrm{P}}$ of the beam proton momentum and is separated from the hadronic debris X of the photon by a large (pseudo)rapidity gap $\Delta\eta \approx \log(1/x_{\mathrm{P}}) \gg 1$. In real photoproduction and hadronic interactions, the pomeron exchange has been shown to dominate if $x_{\mathrm{P}} \lesssim x_{\mathrm{P}}^c = (0.05-0.1)$ holds and/or the rapidity-gap cut satisfies $\Delta\eta \gtrsim \Delta\eta_c = (2.5-3)$. Here the triple-pomeron regime corresponds to high c.m.s. energy of the $a\mathrm{P}$ interaction, $M^2 \gg m_p^2$.[1,2] In DIS it requires $\beta \ll 1$ and/or $M^2 \gg Q^2$. Because of the important kinematical relationship $x_{\mathrm{P}}\beta = x$, even at HERA neither x_{P} nor β can be made asymptotically large, and instead one probes very sub-asymptotic properties of the dipole pomeron.[10] For this reason, in the present analysis we focus on the transition from real photoproduction at FNAL energies[3] to DIS at moderately small x, in the Born approximation for $A_{3\mathrm{P}}(Q^2)$; recent

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treatments[11,12] consider the asymptotic regime of $1/x_{\rm P}$, $1/\beta \gg 1$ in a somewhat related approach.

The presentation is organized as follows. In Sec. 2 we briefly review how diffractive DIS is described in terms of the diffraction excitation of multiparton Fock states of the photon, which interact with the target photon by dipole BFKL pomeron exchange. In Sec. 3 we show that in diffractive DIS at large $Q^2 \gg 1/R_c^2$, the virtual photon acts as an octet-octet color dipole of size $\sim R_c$. We demonstrate how small sizes $R_c \approx 0.3$ fm provide a natural small scale for $A_{3P}(Q^2)$ as suggested by lattice QCD studies (for a recent review see Ref. 13). The case of real photoproduction is studied in Sec. 4. Here the underlying mechanism of diffraction dissociation into large masses is the excitation of qg ($\bar{q}g$) "clusters" ("constituent" quarks) of size R_c in the real photon, and we find $A_{3P}(0)$, which agrees well with the experimental determination. This is the first direct evaluation of A_{3P} and the first instance in which DIS and real photoproduction processes are shown to share the nonperturbative dimensional coupling, $[A_{3P}] = [\mathrm{GeV}]^{-2}$, which does not scale with $1/Q^2$. Furthermore, in the scenario[10,14] for the dipole cross section, we predict a slight rise (by a factor ~ 1.6) in $A_{3P}(Q^2)$ from real photoproduction to DIS. In Sec. 5 we summarize our main results.

2. DIPOLE POMERON DESCRIPTION OF DIFFRACTIVE DIS

We rely upon the microscopic dipole pomeron description of diffractive DIS.[5–8] Diffraction excitation of the lowest $q\bar{q}$ Fock state of the photon has the cross-section (hereafter we focus on the dominant diffraction dissociation of transverse photons)

$$\left.\frac{d\sigma_D(\gamma^* \to X)}{dt}\right|_{t=0} = \int dM^2 \left.\frac{d\sigma_D(\gamma^* \to X)}{dt dM^2}\right|_{t=0} = \frac{1}{16\pi}\int_0^1 dz \int d^2\mathbf{r} |\Psi_{\gamma^*}(Q^2,z,r)|^2 \sigma^2(x,r). \tag{4}$$

Here $\mathbf{r}$ is the transverse separation of the quark and antiquark in the photon, z and $1-z$ are partitions of the lightcone momentum of the photon between the quark and antiquark, $\sigma(x,r)$ is the dipole cross-section for interaction of the $q\bar{q}$ dipole with the proton target (hereafter we use $\sigma(x,r)$ of Refs. 10, 14), and the dipole distribution in the transverse polarized photon $|\Psi_{\gamma^*}(Q^2,z,r)|^2$ derived in Ref. 8 equals

$$|\Psi_{\gamma^*}(Q^2,z,r)|^2 = \frac{6\alpha_{em}}{(2\pi)^2}\sum_i^{N_f} e_i^2\{[z^2+(1-z)^2] \times \varepsilon^2 K_1^2(\varepsilon r) + m_i^2 K_0^2(\varepsilon r)\}, \tag{5}$$

where α_{em} is the fine structure constant, e_i is the quark charge in units of the electron charge, m_i is the quark mass,

$$\varepsilon^2 = z(1-z)Q^2 + m_i^2, \tag{6}$$

and $K_\nu(x)$ is a modified Bessel function of the second kind. The mass spectrum for the $q\bar{q}$ excitation was calculated in Ref. 5; it steeply decreases with M^2:

$$\left.\frac{d\sigma_D}{dM^2 dt}\right|_{t=0} \sim \frac{M^2}{(Q^2+M^2)^3}. \tag{7}$$

The component $\propto 1/(M^2+Q^2)$ of the mass spectrum comes from the diffraction excitation of the $q\bar{q}g$ Fock state of the photon containing the soft gluon which carries the fraction $z_g \ll 1$ of the lightcone momentum of the photon and gives rise to $M^2 + Q^2 \propto Q^2/z_g \gg Q^2$. Let $\mathbf{r}, \boldsymbol{\rho}_1$ and $\boldsymbol{\rho}_2 = \boldsymbol{\rho}_1 - \mathbf{r}$ be the $\bar{q}-q$, $g-q$ and $g-\bar{q}$ separations in the impact parameter (transverse size) plane. Then, in the triple-pomeron regime of $x_{\rm P}$, $\beta \ll 1$,

$$(Q^2+M^2)\left.\frac{d\sigma_D}{dt dM^2}\right|_{t=0} = \int dz d^2\mathbf{r} d^2\boldsymbol{\rho}_1 \{z_g |\Phi(\mathbf{r},\boldsymbol{\rho}_1,\boldsymbol{\rho}_2,z,z_g)|^2\}_{z_g=0} \times \frac{\sigma_3^2(x_{\rm P},r,\rho_1,\rho_2) - \sigma^2(x_{\rm P},r)}{16\pi} \tag{8}$$

in which the square of the 3-parton wave function $|\Phi|^2$ equals[6,7]

$$|\Phi(\mathbf{r},\boldsymbol{\rho}_1,\boldsymbol{\rho}_2,z,z_g)|^2 = \frac{1}{z_g}\frac{1}{3\pi^3}|\Psi_{\gamma^*}(Q^2,z,r)|^2 \mu_G^2 \left| g_S(R_1) K_1(\mu_G \rho_1) \times \frac{\boldsymbol{\rho}_1}{\rho_1} - g_S(R_2) K_1(\mu_G \rho_2)\frac{\boldsymbol{\rho}_2}{\rho_2}\right|^2. \tag{9}$$

Here $g_S(r)$ is the running color charge, $\alpha_S(r) = g_S(r)^2/4\pi$, and the arguments of the color charges are $R_i = \min\{r, \rho_i\}$. In the wave function (9), $\mu_G K_1(\mu_G r)\boldsymbol{\rho}/\rho$ emerges as $\nabla_\rho K_0(\mu_G \rho)$, where $K_0(\mu_G \rho)$ is precisely the two-dimensional Coloumb–Yukawa screened potential. This makes self-evident the interpretation[6,7,10,15] of $R_c = 1/\mu_G$ as the correlation (propagation) radius for perturbative gluons. The 3-body interaction cross-section equals[6,7]

$$\sigma_3(r,\rho_1,\rho_2) = \tfrac{9}{8}[\sigma(\rho_1)+\sigma(\rho_2)] - \tfrac{1}{8}\sigma(r). \tag{10}$$

Hereafter, for of brevity we write $\sigma(r)$ for $\sigma(x_{\rm P}^0, r)$. The ordering of sizes in the three-parton $q\bar{q}g$ Fock state, $r < \rho_{1,2}$ and/or $r > \rho_{1,2}$, changes with the photon's virtuality Q^2, leading to the Q^2-dependent underlying mechanism of diffraction dissociation. We start our analysis with the case of DIS.

3. A_{3P} (Q^2) IN DIFFRACTIVE DIS

Because $K_\nu(z) \sim \exp(-z)$ holds at large z, by virtue of (5), (6), the typical size of the $q\bar{q}$ dipole satisfies $r^2 \lesssim R_{q\bar{q}}^2 = 1/\varepsilon^2 \propto 1/Q^2$ and in the standard leading-log Q^2 approximation, the dominant contribution to the diffraction cross-section (8) comes from $r^2 \ll \rho_1^2 \approx \rho_2^2 \sim R_c^2$. In this region we have

$$\sigma_3(r,\rho_1,\rho_2) \approx \tfrac{9}{4}\sigma(\rho) \gg \sigma(r), \quad \rho = \tfrac{1}{2}(\rho_1+\rho_2),$$

and the virtual photon interacts as an effective octet-octet color dipole of size $\rho \sim R_c$, with the $q\bar{q}$ pair acting as an octet color charge. Also, in this region

$$\mu_G^2\left|K_1(\mu_G\rho_1)\frac{\boldsymbol{\rho}_1}{\rho_1}-K_1(\mu_G\rho_2)\frac{\boldsymbol{\rho}_2}{\rho_2}\right|^2\simeq\frac{r^2}{\rho^4}\cdot\mathcal{F}(\mu_G\rho), \tag{11}$$

the 3-parton wave function factors and (8) takes on the factored form

$$(Q^2+M^2)\left.\frac{d\sigma_D}{dtdM^2}\right|_{t=0}=\int dzd^2\mathbf{r}|\Psi_{\gamma*}(Q^2,z,r)|^2\frac{16\pi^2}{27}\alpha_S(r)r^2\frac{1}{2\pi^4}\left(\frac{9}{8}\right)^3\times\int d\rho^2\left[\frac{\sigma(\rho)}{\rho^2}\right]^2\mathcal{F}(\mu_G\rho). \tag{12}$$

Here the form factor

$$\mathcal{F}(z)=z^2[K_1^2(z)+zK_1(z)K_0(z)+\tfrac{1}{2}z^2K_0^2(z)]$$

satisfies $\mathcal{F}(0)=1$ and $\mathcal{F}(z)\propto\exp(-2z)$ at $z>1$.

The factorization (12) is a crucial starting point for the description of diffractive DIS in terms of the well defined triple-pomeron coupling, which emerges in quite a nontrivial fashion. Indeed, the first factor in the r.h.s of (12) is nearly identical to $\sigma_{\rm tot}(\gamma^*p)$ in the Born approximation, because in the latter (for 3 active flavors)[6,7]

$$\sigma(x_{\rm P},r)\approx\frac{16\pi^2}{27}r^2\alpha_S(r)\log\frac{1}{\alpha_S(r)}\approx\frac{16\pi^2}{27}r^2\alpha_S(r) \tag{13}$$

and, to within a logarithmic factor $\sim\log[1/\alpha_S(Q^2)]\sim1$, we have

$$\int dzd^2\mathbf{r}|\Psi_{\gamma*}(Q^2,z,r)|^2\frac{16\pi^2}{27}\alpha_S(r)r^2\approx\sigma_{\rm tot}(\gamma^*p)=\int dzd^2\mathbf{r}|\Psi_{\gamma*}(Q^2,z,r)|^2\sigma(r). \tag{14}$$

Consequently, (13) becomes equivalent to (3) with

$$A_{3\rm P}(Q^2)\sim A_{3\rm P}^*=\frac{1}{2\pi^4}\left(\frac{9}{8}\right)^3\int d\rho^2\left[\frac{\sigma(\rho)}{\rho^2}\right]^2\mathcal{F}(\mu_G\rho). \tag{15}$$

(We shall return to a detailed calculation of $A_{3\rm P}(Q^2)$ in Sec. 5.) The factorization (12) holds simultaneously for all the $q_i\bar{q}_ig$ states and we predict that $A_{3\rm P}(Q^2)$ is independent of the flavor i. $A_{3\rm P}^*$ is the nonperturbative quantity dominated by $\rho\sim R_c$. Making use of (13), for $R_c=0.3$ we obtain the order-of-magnitude estimate $A_{3\rm P}^*\sim 1/16R_c^2\sim0.1\ \mathrm{GeV}^{-2}$.

Equation (13) describes the Born term of the perturbative contribution $\sigma^{(pt)}(x,r)$ to the dipole cross-section $\sigma(x,r)$ from the exchange by perturbative gluons. This perturbative dipole cross-section $\sigma^{(pt)}(x,r)$ is a solution of the generalized BFKL equation[6,7,10,15] and dominates $\sigma(x,r)$ at small size $r\ll R_c$. It rapidly rises towards large $1/x$, dominating the observed growth and giving a good quantitative description of the proton structure function at HERA.[10] The phenomenological description of $\sigma(x,r)$ at large dipole sizes $r\gtrsim R_c$ and moderate $1/x$ requires introduction of the nonperturbative component $\sigma^{(npt)}(r)$ of the dipole cross-section,[10,14,15] which is expected to have a weak energy dependence and must be inferred from experimental data [the scenario[10,14] introduces an energy-independent $\sigma^{(npt)}(r)$]. Here we wish to recall that real photoproduction of the J/Ψ and exclusive leptoproduction of the ρ^0 at $Q^2\sim10\ \mathrm{GeV}^2$ probe the (predominantly nonperturbative) dipole cross-section at $r\sim3/m_c\sim0.5\ \mathrm{fm}\lesssim2R_c$.[14,16,17] Real and weakly virtual ($Q^2\lesssim10\ \mathrm{GeV}^2$) photoproduction of the open charm probes the (predominantly perturbative) dipole cross-section at $r\sim1/m_c\sim R_c/2$.[10,15] The proton structure function $F_2^p(x,Q^2)$ probes the dipole cross-section over a broad range of radii from $r\sim1$ fm down to $r\sim0.02$ fm. The successful quantitative description of the corresponding experimental data in Refs. 10, 14, 16, 17 implies that we have a reasonably good (conservatively, $\lesssim$15–20%) understanding of the dipole cross-section at $r\sim R_c$ of the interest for evaluation of $A_{3\rm P}^*$. Quantitatively, for $r\sim R_c$ the dipole cross-section $\sigma(x_{\rm P}^0,r)$ receives contributions from the exchange by perturbative gluons (13), and from the nonperturbative component $\sigma^{(npt)}(r)$ in approximately a ratio of 1:2. The numerical calculation with the dipole cross-section of Ref. 14 gives $A_{3\rm P}^*=0.56\ \mathrm{GeV}^{-2}$.

4. REAL PHOTOPRODUCTION: DIFFRACTION EXCITATION OF CONSTITUENT QUARKS

For real photons ($Q^2=0$) in the dipole distribution (5), (9) the typical $q\bar{q}$ size is large,

$$r^2\sim R_{q\bar{q}}^2\approx\frac{1}{m_q^2}\gg R_c^2. \tag{16}$$

Here m_q is still another nonperturbative parameter to which the diffractive DIS is sensitive, and here too we rely upon experiment. With the same dipole cross-section $\sigma(x,r)$, the choice of $m_q=0.15$ GeV leads to a good quantitative description of the real photoabsorption cross-section,[14,5] of color transparency effects and the total cross-section of exclusive leptoproduction of vector mesons at moderate Q^2,[16,17] and of nuclear shadowing in DIS on nuclei.[18] By virtue of the inequality (16), the dipole distribution $|\Phi|^2$ will be dominated by configurations with $\rho_1^2\lesssim R_c^2\ll\rho_2^2\sim r^2$ and $\rho_2^2\lesssim R_c^2\ll\rho_1^2\sim r^2$ and takes on the factored form first considered in:[5]

$$|\Phi(\mathbf{r},\boldsymbol{\rho}_1,\boldsymbol{\rho}_2,z,z_g)|^2=\frac{1}{z_g}\frac{4}{3\pi^2}|\Psi_{\gamma*}(Q^2,z,r)|^2\mu_G^2[\alpha_S(\rho_1)K_1^2(\mu_G\rho_1)+\alpha_S(\rho_2)K_1^2(\mu_G\rho_2)]. \tag{17}$$

We recover a kind of constituent quark model with qg and/or $\bar{q}g$ clusters of size $\rho\lesssim R_c$, where the square of the qg wave function of the constituent quark is $\propto(1/z_g)\alpha_S(\rho)K_1^2(\mu_G\rho)$. Diffraction dissociation of the photon $\gamma^*\to q^+\bar{q}+g$ proceeds via diffraction excitation of the constituent (anti)quark of the photon $q(\bar{q})\to q(\bar{q})+g$. For the size ordering (16), Eq. (10) leads to $\sigma_3(r,\rho_1,\rho_2)\approx\sigma(r)+(9/8)\sigma(\rho)$, the factorization $\sigma_3^2(r,\rho_1,\rho_2)-\sigma^2(r)\approx(9/4)\sigma(r)\sigma(\rho)$, where $\rho=\min\{\rho_i\}$, and the factored form of the diffraction cross-section (8):

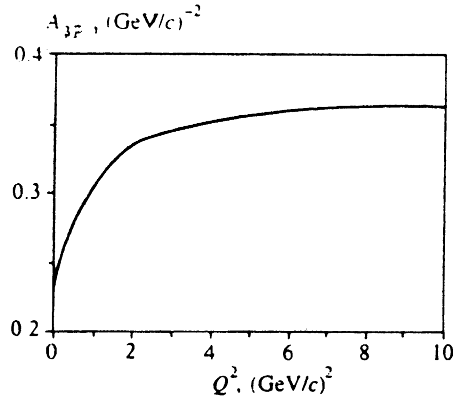

FIG. 1. Our prediction for the Q^2-dependence of $A_{3P}(Q^2)$.

$$M^2 \frac{d\sigma_D}{dtdM^2}\bigg|_{t=0} \approx \int dz d^2\mathbf{r} |\Psi_{\gamma^*}(Q^2=0,z,r)|^2 \sigma(r) \frac{3}{8\pi^2} \times \int d\rho^2 \frac{\sigma(\rho)}{\rho^2} f^2(\mu_G \rho) = \sigma_{\rm tot}(\gamma p) A_{3P}(0), \tag{18}$$

in which

$$\sigma_{\rm tot}(\gamma p) = \int dz d^2\mathbf{r} |\Psi_{\gamma^*}(Q^2,z,r)|^2 \sigma(r), \tag{19}$$

$$A_{3P}(0) \approx \frac{3}{8\pi^2} \int d\rho^2 \alpha_S(\rho) \frac{\sigma(\rho)}{\rho^2} f^2(\mu_G \rho) \sim \frac{1}{18} R_c^2, \tag{20}$$

and $f(z) = zK_1(z)$. We see that, to within numerical and logarithmic factors ~ 1, we obtain $A_{3P}(Q^2) \approx A_{3P}(0)$, which was anticipated some time ago.[5–9] As a matter of fact, the exact large-r behaviour of $\sigma(r)$ is not the main point here. We only should assume that (19) reproduces the observed total photoabsorption cross-section. Here we wish to mention that with the dipole cross-section of Ref. 14 we find $\sigma_{\rm tot}(\gamma p) = 108 \mu b$, in good agreement with the Fermilab data.[19] A direct calculation from (8)–(10) gives $A_{3P}(0) = 0.23$ GeV^{-2}, in reasonable agreement with the experimental determination $A_{3P}(0) \approx 0.16$ GeV^{-2}.[3]

5. Q^2-DEPENDENCE OF $A_{3P}(Q^2)$; DISCUSSION OF RESULTS

To obtain more insight into the Q^2 dependence of the triple-pomeron coupling, we present here the results of a direct evaluation of $A_{3P}(Q^2)$ from Eqs. (3), (8), (14) and (19) (for a detailed comparison of the cross-section calculations with the experiment see Refs. 10, 14, 17. We consider $x_P = 0.03$ and in the case of DIS, we take $x = 0.004$. In this range of x and moderate $Q^2 \lesssim 10$ GeV^2, the proton structure function is approximately flat versus $1/x$. The equivalent c.m.s. energy in the real photoproduction can be estimated as $W^2 \sim m_p^2/x \sim 250$ GeV^2, which corresponds to the energy range of the FNAL experiment.[3]

Our results for the Q^2-dependence of the triple-pomeron coupling $A_{3P}(Q^2)$ are presented in Fig 1. The main feature of $A_{3P}(Q^2)$ is its weak Q^2-dependence; still, we predict a slight growth (by a factor ~ 1.6) of $A_{3P}(Q^2)$ from $A_{3P}(Q^2 = 0) = 0.23$ GeV^{-2} to the DIS value $A_{3P}(Q^2 \gtrsim Q^{*2}) = 0.36$ GeV^{-2}. This rise of $A_{3P}(Q^2)$ take place for $Q^2 \lesssim Q^{*2} \approx (2-3)$ GeV^2 and describes the transition from the regime of diffraction dissociation of the constituent quark of Sec. 4 to the regime of diffraction dissociation of the octet-octet dipole state of the photon of Sec. 3. The transition takes place when the $q\bar{q}$ size $R_{q\bar{q}}$ and the size $\sim R_c$ of the constituent quark become comparable, $R_{q\bar{q}} = 1/\varepsilon \approx 2/\sqrt{Q^2} \sim R_c$, i.e., at $Q^{*2} \sim 4/R_c^2 = 3$ GeV^2, in agreement with the results shown in figure. The predicted Q^2-dependence of $A_{3P}(Q^2)$ can be tested experimentally at HERA and should shed light on our understanding of dipole cross-section at $r \sim R_c$. Notice that $A_{3P}(0)$ is a linear functional of $\sigma(R_c)$, whereas $A_{3P}(Q^2 \gtrsim Q^{*2})$ is quadratic. Hence the conservative ($\sim$15–20%) uncertainty in our present knowledge of $\sigma(R_c)$ implies a conservative theoretical uncertainty of $\sim$15–20% and $\sim$30–40% in the predicted values of $A_{3P}(0)$ and $A_{3P}(Q^2 \gtrsim Q^{*2})$, respectively. This theoretical uncertainty can further be reduced with future higher-accuracy experimental determinations of the dipole cross-section.

The indirect experimental evidence for the weak Q^2-dependence of $A_{3P}(Q^2)$ comes from nuclear shadowing in DIS on nuclei. Diffraction excitation of large masses contributes to nuclear shadowing at $x \ll 10^{-2}$, and in Ref. 18 it was shown that calculations using the photoproduction value of $A_{3P}(0)$ are in good agreement with the experiment. Crude evaluations[5,20] of the total rate of diffractive DIS, using the photoproduction value of $A_{3P}(0)$, are also consistent with the HERA data.[4] More detailed calculations of the rate of diffractive DIS is presented in Ref. 21.

Here we only consider the Born term of the $1/(Q^2+M^2)$ mass spectrum. The factorization of the Q^2 and x_P dependence in Eq. (12) allows one to introduce the description of diffraction dissociation in terms of the structure function of the pomeron. The mass spectrum $\propto 1/(M^2+Q^2)$ corresponds to the $q\bar{q}$ sea structure function of the pomeron,[5–7] the effects of the Q^2 evolution for $Q^2 \gg Q^{*2}$ are considered in Ref. 21. Still another effect not considered here is the so-called absorption corrections which will slightly reduce $A_{3P}(Q^2)$ (for the relevant formalism see Refs. 6,7). The absorption correction is typically of order $\sigma_3/8\pi B$, where B is the diffraction slope. In real photoproduction we have $\sigma_3 \sim \sigma(r \sim 1/m_q)$ and $\sigma_3/8\pi B \approx (\sigma_{\rm el}/\sigma_{\rm tot})_{\pi N} \ll 1$. In DIS the absorption correction will be still smaller for small values of $\sigma_3 \sim (9/4)\sigma(r \sim R_c)$. Therefore, the increase of $A_{3P}(Q^2)$ [cf. $A_{3P}(0)$] is stable against absorption corrections and is an interesting prediction from the dipole cross-section model.

To summarize, we have presented the first direct evaluation of the Born approximation for the nonperturbative triple-pomeron coupling A_{3P}, which is relevant to the subasymptotic regime of diffraction production of moderately large masses at moderately small x of the fixed-target FNAL/CERN experiments and those at HERA. With the nonperturbative parameter of the model, the correlation (propagation) radius for the perturbative gluons $R_c \approx 0.3$ fm as used in other successful phenomenological application of the model, we found good agreement with the experimental determination of $A_{3P}(0)$. We have related the small numerical value of A_{3P}, which was a mystery, to the small value of R_c. We predict a slight rise of $A_{3P}(Q^2)$ with Q^2, by a factor ~ 1.6, from real photoproduction to DIS at $Q^2 \gtrsim 3$ GeV^2.

In a somewhat related approach to the BFKL pomeron at large N_c, in the scaling approximation α_S=const and $R_c=\infty$,[22] the triple-pomeron regime was considered also by Mueller and Patel.[11] These authors consider the case of large t at asymptotically large rapidity gap and energy and, because of the approximation $R_c=\infty$, find $A_{3P} \propto 1/\sqrt{-t}$. How to remove this divergence at $t=0$ by regularization of the notorious infrared sensitivity of the scaling BFKL pomeron,[22,23] remains an open issue; our Born approximation for the subasymptotic A_{3P} gives an explicit dependence on the nonperturbative infrared parameter of our specific model. In our analysis of the same approximation α_S=const, Bartels and Wüsthoff found that, for lack of the LLOA ordering of sizes in the asymptotic BFKL regime, the cut pomeron becomes the mixture of the two-gluon and four-gluon states, unlike the exchanged two-gluon pomerons.[12] Numerical analysis of the infrared sensitivity of the triple-pomeron vertex in the framework of this approach and phenomenological applications are not yet available.

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[1)]Apart from ≈5% statistical error and 16% normalization uncertainty,[3] this number contains ≲10% uncertainty from our extrapolation from $|t|=0.05$ GeV^2 to $t=0$ using the slope of the t-dependence as measured in Ref. 3.

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