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Direct calculation of the triple-pomeron coupling for diffractive DIS and real photoproduction

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Abstract

We present a unified direct evaluation of the triple-pomeron coupling $A_{3\mathbf{IP}}(Q^2)$ 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 colour dipole state of the virtual photon at large Q^2. The small phenomenological value of $A_{3\mathbf{IP}}(0)$, which was a mystery, is related to the small non perturbative correlation radius $R_c \approx 0.3\,\mathrm{fm}$ for the perturbative gluons. We confirm the early expectations of weak Q^2 dependence of the dimensionful coupling $A_{3\mathbf{IP}}(Q^2)$ and predict that it rises by a factor ~ 1.6 from real photoproduction to deep inelastic scattering.

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1 Introduction

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

$$\frac{M^2}{\sigma_{tot}(ap)} \cdot \frac{d\sigma_D(a \rightarrow X)}{dt dM^2}\bigg|_{t=0} \approx A_{3\mathbf{IP}} \qquad (1)$$

with the approximately energy- and projectile-independent dimensionful triple-pomeron coupling $A_{3\mathbf{IP}}$, which holds at moderately high energies such that the photoabsorption and hadronic cross sections are approximately constant ([1], for a review see [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 $A_{3\mathbf{IP}}(Q^2 = 0) \approx 0.16\,\mathrm{GeV}^{-2}$ [3]. (Apart from $\approx 5\%$ statistical error and 16% normalization uncertainty [3], this number contains $\lesssim 10\%$ uncertainty from our extrapolation from $|t| = 0.05\,\mathrm{GeV}^2$ to $t = 0$ using the slope of the t-dependence as measured in [3].) In the Regge theory language, $A_{3\mathbf{IP}}$ measures the projectile-pomeron cross section [1]

$$\sigma_{tot}(a\mathbf{IP}; M^2 \gg m_p^2) = \frac{16\pi M^2}{\sigma_{tot}(pp)} \cdot \frac{d\sigma_D(a \rightarrow X)}{dt dM^2}\bigg|_{t=0} \approx 16\pi A_{3\mathbf{IP}} \frac{\sigma_{tot}(ap)}{\sigma_{tot}(pp)} \approx 0.08\sigma_{tot}(ap)\,. \qquad (2)$$

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

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

$$\frac{M^2 + Q^2}{\sigma_{tot}(\gamma^* p)} \cdot \frac{d\sigma_D(\gamma^* \rightarrow X)}{dt\, dM^2}\bigg|_{t=0} \approx A_{3\mathbf{IP}}(Q^2). \qquad (3)$$

The dimensionful $A_{3\mathbf{IP}}(Q^2)$ is a non perturbative quantity. The subject of the present paper[1] is an evaluation and the study of the Q^2 dependence of the Born approximation for $A_{3\mathbf{IP}}(Q^2)$

[1] This paper updates earlier preprints KFA-IKP(Th)-1994-36, DFTT 41/94; we here present a more detailed discussion of the non perturbative aspects of the evaluation of the triple-pomeron coupling and add more references to the recent literature on the subject.

in terms of the non perturbative 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 both in the real photoproduction and deep inelastic scattering at moderate energy ν and/or moderately small x, the $A_{3\mathbf{IP}}(Q^2)$ is approximately the same, thus confirming the anticipations [5-9].

In this paper we operationally define, and calculate, $A_{3\mathbf{IP}}(Q^2)$ 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. A convenient variable is the fraction $x_{\mathbf{IP}}$ of proton momentum taken away by the pomeron, $x_{\mathbf{IP}} = (M^2+Q^2)/(W^2+Q^2) \ll 1$. The final-state proton p' carries the fraction $(1-x_{\mathbf{IP}})$ of the beam proton momentum and is separated from the hadronic debris X of the photon by a large (pseudo)rapidity gap $\Delta\eta \approx \log\frac{1}{x_{\mathbf{IP}}} \gg 1$. In real photoproduction and hadronic interactions, the pomeron exchange has been shown to dominate at $x_{\mathbf{IP}} \lesssim x_{\mathbf{IP}}^c$ =(0.05- -0.1) and/or the rapidity-gap cut $\Delta\eta \gtrsim \Delta\eta_c$ =(2.5- -3). Here the triple-pomeron regime corresponds to the high c.m.s. energy of the $a\mathbf{IP}$ 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 $x_{\mathbf{IP}}\beta = x$, even at HERA neither $x_{\mathbf{IP}}$ nor β can be made asymptotically large and one rather probes very subasymptotic 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 $A_{3\mathbf{IP}}(Q^2)$; recent works [11,12] consider the asymptotic regime of $1/x_{\mathbf{IP}},\ 1/\beta \gg 1$ in a somewhat related approach.

The presentation is organized as follows. In section 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 proton by the dipole BFKL pomeron exchange. In section 3 we show that in diffractive DIS at large $Q^2 \gg 1/R_c^2$, the virtual photon acts as an octet-octet colour dipole of size $\sim R_c$. We demonstrate how the small $R_c \approx 0.3\,\mathrm{fm}$, as suggested by lattice QCD studies (for a recent review see [13]), gives a natural small scale for $A_{3\mathbf{IP}}(Q^2)$. The case of real photoproduction is studied in section 4. Here the underlying mechanism of diffraction dissociation into large masses is an excitation of qg ($\bar{q}g$) 'clusters' ('constituent quarks') of size R_c in the real photon, and we find $A_{3\mathbf{IP}}(0)$, which agrees well with the experimental determination. This is the first direct evaluation of $A_{3\mathbf{IP}}$ and the first instance when DIS and real photoproduction processes are shown to share the non perturbative dimensionful coupling, $[A_{3\mathbf{IP}}] = [\mathrm{GeV}]^{-2}$, which does not scale with $1/Q^2$. Furthermore, in the scenario [10,14] for the dipole cross section, we predict a slight (by a factor ~ 1.6) rise of $A_{3\mathbf{IP}}(Q^2)$ from real photoproduction to DIS. In section 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 following cross section (hereafter we focus on the dominant diffraction dissociation of transverse photons):

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

Here $\vec{r}$ is the transverse separation of the quark and antiquark in the photon, z and $(1-z)$ are partitions of the photon light cone momentum between the quark and antiquark, $\sigma(x,r)$ is the dipole cross section for the 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 [8], equals

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

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 the modified Bessel function. The mass spectrum for the $q\bar{q}$ excitation was calculated in [5] and steeply decreases with M^2:

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

The $\propto 1/(M^2+Q^2)$ component 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 photon light cone momentum and gives rise to $(M^2+Q^2)\propto Q^2/z_g \gg Q^2$. Let $\vec{r},\vec{\rho}_1$ and $\vec{\rho}_2=\vec{\rho}_1-\vec{r}$ 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 $x_{\bf I\!P},\beta \ll 1$,

$$
(Q^2+M^2)\left.\frac{d\sigma_D}{dtdM^2}\right|_{t=0}
$$
$$
=\int dz\, d^2\vec{r}\, d^2\vec{\rho}_1\ \left\{z_g|\Phi(\vec{r},\vec{\rho}_1,\vec{\rho}_2,z,z_g)|^2\right\}_{z_g=0}\cdot\frac{\sigma_3^2(x_{\bf I\!P},r,\rho_1,\rho_2)-\sigma^2(x_{\bf I\!P},r)}{16\pi} \qquad (8)
$$

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

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

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

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

Hereafter, for the sake of brevity, $\sigma(r)$ stands for $\sigma(x^0_{\bf I\!P},r)$. 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 from the case of DIS. -

3 $A_{3{\bf I\!P}}(Q^2)$ in diffractive DIS

Because of $K_\nu(z)\sim\exp(-z)$ at large z, and by virtue of (5,6), the typical size of the $q\bar{q}$ dipole $r^2\lesssim R^2_{q\bar{q}}=1/\varepsilon^2\propto 1/Q^2$ and, to the standard leading-logQ^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 $\sigma_3(r,\rho_1,\rho_2)\approx\frac{9}{4}\sigma(\rho)\gg\sigma(r)$, where $\rho=\frac{1}{2}(\rho_1+\rho_2)$, and the virtual photon interacts as an effective octet-octet colour dipole of size $\rho\sim R_c$, with the $q\bar{q}$ pair acting as an octet colour charge. Also, in this region

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

the three-parton wave function factorizes and (8) takes on the factorized form

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

Here the form factor $\mathcal{F}(z) = z^2[K_1^2(z) + zK_1(z)K_0(z) + \frac{1}{2}z^2K_0^2(z)]$ satisfies $\mathcal{F}(0) = 1$ and $\mathcal{F}(z) \propto \exp(-2z)$ at $z > 1$.

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

$$\sigma(x_{\mathbf{IP}}, r) \approx \frac{16\pi^2}{27} r^2 \alpha_S(r) \log\left[\frac{1}{\alpha_S(r)}\right] \approx \frac{16\pi^2}{27} r^2 \alpha_S(r) \qquad (13)$$

and, modulo the logarithmic factor $\sim \log[1/\alpha_S(Q^2)] \sim 1$, we have

$$\int dz\, d^2\vec{r}\, |\Psi_{\gamma^*}(Q^2,z,r)|^2 \cdot \frac{16\pi^2}{27} \cdot \alpha_S(r) r^2 \approx$$
$$\sigma_{tot}(\gamma^* p) = \int dz\, d^2\vec{r}\, |\Psi_{\gamma^*}(Q^2,z,r)|^2 \sigma(r)\,. \qquad (14)$$

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

$$A_{3\mathbf{IP}}(Q^2) \sim A^*_{3\mathbf{IP}} = \frac{1}{2\pi^4}\cdot\left(\frac{9}{8}\right)^3 \cdot \int d\rho^2 \left[\frac{\sigma(\rho)}{\rho^2}\right]^2 F(\mu_G\rho)\,. \qquad (15)$$

[We will go back to a detailed calculation of $A_{3\mathbf{IP}}(Q^2)$ in section 5.] The factorization (12) holds simultaneously for all the $q_i\bar{q}_i g$ states and we predict independence of $A_{3\mathbf{IP}}(Q^2)$ of the flavour i. $A^*_{3\mathbf{IP}}$ is the non perturbative quantity dominated by $\rho \sim R_c$. Making use of (13), for $R_c = 0.3$ we obtain an order of magnitude estimate $A^*_{3\mathbf{IP}} \sim \frac{1}{16}R_c^2 \sim 0.1$ GeV^{-2}.

Equation (13) describes the Born term of the perturbative contribution $\sigma^{(pt)}(x,r)$ from the exchange by perturbative gluons to the dipole cross section $\sigma(x,r)$. This perturbative dipole cross section $\sigma^{(pt)}(x,r)$ is a solution of the generalized BFKL equation [6,7,10,15] and dominates the $\sigma(x,r)$ at small size $r \ll R_c$. It rapidly rises towards large $\frac{1}{x}$, 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 $\frac{1}{x}$ requires the introduction of the non perturbative 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 $Q^2 \sim 10\,\mathrm{GeV}^{-2}$, probe the (predominantly non perturbative) dipole cross section at $r \sim 3/m_c \sim 0.5\,\mathrm{fm} \lesssim 2R_c$ [14,16,17]. Real, and weakly virtual $Q^2 \lesssim 10\,\mathrm{GeV}^2$, photoproduction of the open charm probes the (predominantly perturbative) dipole cross section at $r \sim \frac{1}{m_c} \sim \frac{1}{2}R_c$ [10,15]. The proton structure function $F_2^p(x,Q^2)$ probes the dipole cross section in a broad range of radii from $r \sim 1\,\mathrm{fm}$ down to $r \sim 0.02\,\mathrm{fm}$. A successful quantitative description of the corresponding experimental data in [10,14,16,17] implies that we have a reasonably good (to a conservative accuracy $\lesssim$(15-20)%) understanding of the dipole cross section at $r \sim R_c$, which is of interest for the evaluation of $A^*_{3\mathbf{IP}}$. Quantitatively, at $r \sim R_c$ the dipole cross section $\sigma(x^0_{\mathbf{IP}}, r)$ receives approximately $1:2$ contributions from the exchange by perturbative gluons (13), and from the non perturbative component $\sigma^{(npt)}(r)$. The numerical calculation with the dipole cross section of Ref. [14] gives $A^*_{3\mathbf{IP}} = 0.56\,\mathrm{GeV}^{-2}$. – -

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,

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

Here m_q is still another non perturbative parameter to which the diffractive DIS is sensitive, and here too we rely upon the experiment. The choice of $m_q = 0.15\,\mathrm{GeV}$ leads, with the same dipole cross section $\sigma(x,r)$, to a good quantitative description of the real photoabsorption cross section [14,5], of colour transparency effects and total cross section of exclusive leptoproduction of vector mesons at moderate Q^2 [16,17], and of nuclear shadowing in DIS on nuclei [18]. Because of 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 factorized form first considered in [5]:

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

We recover a sort of constituent quark model with the qg and/or $\bar{q}g$ clusters of size $\rho \lesssim R_c$, with the square of the qg wave function of the 'constituent' quark $\propto \frac{1}{z_g}\alpha_S(\rho)K_1^2(\mu_G\rho)$. 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 $\sigma_3(r,\rho_1,\rho_2) \approx \sigma(r)+\frac{9}{8}\sigma(\rho)$, the factorization $\sigma_3^2(r,\rho_1,\rho_2)-\sigma^2(r) \approx \frac{9}{4}\sigma(r)\sigma(\rho)$, where $\rho=\min\{\rho_i\}$, and the factorized form of the diffraction cross section (8):

$$M^2\left.\frac{d\sigma_D}{dtdM^2}\right|_{t=0} \approx \int dz\, d^2\vec{r}\, |\Psi_{\gamma^*}(Q^2=0,z,r)|^2\sigma(r)$$
$$\times\frac{3}{8\pi^2}\cdot\int d\rho^2\left[\frac{\sigma(\rho)}{\rho^2}\right]f^2(\mu_G\rho) = \sigma_{tot}(\gamma p)A_{3\mathbf{I\!P}}(0)\,, \quad (18)$$

in which

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

$$A_{3\mathbf{I\!P}}(0) \approx \frac{3}{8\pi^2}\cdot\int d\rho^2\alpha_S(\rho)\left[\frac{\sigma(\rho)}{\rho^2}\right]f^2(\mu_G\rho) \sim \frac{1}{18}R_c^2\,, \quad (20)$$

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

5 Q^2 dependence of $A_{3\mathbf{I\!P}}(Q^2)$. Discussion of results.

To have more insight into the Q^2 dependence of the triple-pomeron coupling, we present here the results of a direct evaluation of $A_{3\mathbf{I\!P}}(Q^2)$ from Eq.s (3),(8),(14),(19) (for a detailed comparison of the cross section calculations with the experimental results see [10,14,17]). We consider $x_{\mathbf{I\!P}}=0.03$ and in the case of DIS, we take $x=0.004$. In this range of x and moderate

$Q^2 \lesssim 10\,\mathrm{GeV}^2$, the proton structure function is approximately flat as a function of $\frac{1}{x}$. The equivalent c.m.s. energy in the real photoproduction can be estimated as $W^2 \sim m_p^2/x \sim 250\,\mathrm{GeV}^2$, which corresponds to the energy range of the FNAL experiment [3].

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

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

Here we only considered the Born term of the $1/(Q^2 + M^2)$ mass spectrum. The factorization of the Q^2 and $x_{\mathbf{IP}}$ dependence in Eq. (12) allows one to introduce the description of diffraction dissociation in terms of the structure function of the pomeron. The $\propto 1/(M^2+Q^2)$

mass spectrum corresponds to the $q\bar{q}$ sea structure function of the pomeron [5-7], the effects of the Q^2 evolution at $Q^2 \gg Q^{*2}$ are considered elsewhere [21]. Still another effect not considered here is the so-called absorption corrections, which will slightly reduce $A_{3\mathbf{I\!P}}(Q^2)$ (for the relevant formalism see [6,7]). The absorption correction is typically of the order of $\sigma_3/8\pi B$, where B is the diffraction slope. In real photoproduction, $\sigma_3 \sim \sigma(r \sim \frac{1}{m_q})$ from which we can obtain the $\sigma_3/8\pi B \approx (\sigma_{el}/\sigma_{tot})_{\pi N} \ll 1$; in DIS the absorption correction will be smaller still for the small value of $\sigma_3 \sim \frac{9}{4}\sigma(r \sim R_c)$. Therefore, the increase of $A_{3\mathbf{I\!P}}(Q^2)$ cf. $A_{3\mathbf{I\!P}}(0)$ is stable against absorption corrections and is an interesting prediction from the dipole cross section model.

To summarize, we presented the first direct evaluation of the Born approximation for the non perturbative triple-pomeron coupling $A_{3\mathbf{I\!P}}$, 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 at HERA. With the non perturbative parameter of the model, the correlation (propagation) radius for the perturbative gluons $R_c \approx 0.3\,\mathrm{fm}$ as used in other successful phenomenological applications of the model, we found good agreement with the experimental determination of $A_{3\mathbf{I\!P}}(0)$. We related the small numerical value of $A_{3\mathbf{I\!P}}$, which was a mystery, to the small value of R_c. We predict a slight rise of $A_{3\mathbf{I\!P}}(Q^2)$ with Q^2, by a factor ~ 1.6, from real photoproduction to DIS at $Q^2 \gtrsim 3\,\mathrm{GeV}^2$.

In a somewhat related approach to the BFKL pomeron at large N_c, in the scaling approximation $\alpha_S = \mathrm{const}$ 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 $A_{3\mathbf{I\!P}} \propto 1/\sqrt{-t}$. How this divergence at $t = 0$ of the asymptotic theory is removed by a regularization of the notorious infrared sensitivity of the scaling BFKL pomeron [22,23], remains an open issue; our Born approximation for the subasymptotic $A_{3\mathbf{I\!P}}$ gives an explicit dependence on the non perturbative infrared parameter of our specific model. In their analysis of the same approximation $\alpha_S = \mathrm{const}$, Bartels and Wüsthoff found that, because of the lack of the LLQA 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]. A numerical analysis of the infrared sensitivity of the triple-pomeron vertex in the framework of this approach, and in its phenomenological applications is not yet available.

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Figure caption

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