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Diffractive $C = +1$ Meson Production in $\gamma^* p$ Collisions as a Signal of Odderon Exchange[1]

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Abstract: We calculate the cross section of diffractive $C = +1$ neutral meson production in virtual photon proton collision at high energies. Due to the opposite C parities of photon and meson M^+ ($M^+ = \eta_C, \pi^0, a_2$) this process probes the t-channel $C = -1$ odderon exchange, which is described here as noninteracting three–gluon exchange. Estimates for the cross section of inelastic diffractive $\gamma^* p \to \eta_C X_p$ process are presented . The total photoproduction cross section of η_C meson photoproduction is found to be 47 pb. The cross sections for the diffractive production of light mesons (π^0, a_2) in $\gamma^* p$ collisions are of the same order if $Q^2 \sim m_C^2$.

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1 Introduction

The increase of luminosity at HERA offers the possibility of experimental investigations of diffractive processes of photo- and electroproduction of neutral charge parity even mesons $M^+ = \pi^0, a_2, \eta_C$ at high energies. (see, for example, [1]). Studies of these reactions are of great interest since the opposite C parities of photon ($C = -1$) and M^+ meson cause definite C parity $C = -1$ of the t- channel exchange. Therefore the pomeron exchange being most important at high energies does not contribute there. Therefore, these processes allow for clean investigations of other Regge trajectories which are characterized by negative $C = -1$ parity.

The reggeons having $C = -1$ are ω trajectory and odderon. According to the fit to data for the total hadron–hadron cross sections the intercept of ω trajectory is small (near 1/2) [2] and its contribution decreases with energy. The odderon is the C odd partner of pomeron with an intercept ≥ 1 and was introduced long time ago [3]. It is assumed that the odderon is related to gluonic degrees of freedom. The odderon exchange results in a difference between, for example, total pp and $p\vec{p}$ cross sections which does not decrease with the energy. Up to now there is no indication of that difference in present high energy pp and $p\bar{p}$ data. Also odderon effects have not been observed in other soft hadronic collisions.

The reasons for the absence of the odderon exchange in soft hadronic reactions remains to be understood since in perturbative QCD the pomeron and the odderon appear on the equal footing. In Leading Logarithmic Approximation (LLA) of perturbative QCD the pomeron corresponds to the exchange of two interacting reggeized gluons (hard pomeron), whereas odderon is described by exchange of three interacting reggeized gluons (hard odderon), see [4]. In Born approximation, i.e. considering the pure three–gluon exchange in the t- channel, the intercept of odderon is equal to 1. The effect of gluon interactions is expected to increase of odderon intercept [5]. Therefore the hard odderon should manifest itself in hard diffractive processes at high energies.

The idea to study odderon exchanges in the diffractive meson production was considered by Schäfer, Mankiewicz and Nachtmann in Ref. [6] where the odderon exchange was treated purely phenomenologically, in analogy to the one-photon exchange. The possibility to probe the hard odderon in processes of light meson production in photon-photon collisions at large t was considered in [7]. Recently exclusive η_C photoproduction on proton was studied by Czyżewski *et al.* in [8] describing the odderon as three-gluon exchange.

In this paper we shall calculate the diffractive production of $C = +1$ mesons in Born approximation, i.e. involving the exchange of three noninteracting gluons in the t-channel. We use assumptions similar to the ones made in [8], though our method of calculation is different. For the calculations the impact factors obtained in [7] are used. The results for heavy mesons are compared with [8] and discrepancies are discussed. Furthermore diffractive production of light mesons

$M^+ = \pi^0, a_2, \ldots$ in virtual photon-proton interactions and the process of diffractive production of mesons at large t, where the proton disintegrates are studied.

2 Method of calculation

In the following we consider diffractive meson production processes which can be characterized by at least one hard scale. This scale is either the mass of the heavy quark in the case of η_C production or the virtuality of photon Q^2 for light mesons production.

The quasi-elastic production of a M^+ meson is shown in Fig.1. The mean transverse distance of the quarks in the upper block of Fig.1 is $\sim 1/m_c, 1/Q$. The other essential (soft) scale in our problem is the mean transverse separation between the quarks in proton $\sim 1/m_\rho$ (lower part of Fig.1). With the growth of transfered momentum $|t| \geq m_\rho^2$ the process with photon disintegration shown at Fig.2 becomes more probable than the elastic one (Fig.1). In the region of large t ($|t| \gg \Lambda^2_{\rm QCD}$) the cross section for this process can be expressed through the photon-quark cross section and the quark densities inside the proton. In this case there is no soft scale and we deal with a purely hard diffractive process where the perturbative odderon should contribute.

We do not discuss here the production of light mesons by real photons. Those are purely soft processes and their cross sections should be larger than the cross sections of hard processes $\gamma^* p \to \pi^0(a_2)p$, $\gamma^* p \to \eta_C p$.

The amplitude determined by three–gluon exchange can be written in impact representation as an integral over the gluon transverse momenta $\vec{\kappa}_i$ [7]

$$\mathcal{M} = \frac{s}{3(2\pi)^4} \int \frac{d^2\kappa_1\, d^2\kappa_2}{\vec{\kappa}_1^{\,2}\, \vec{\kappa}_2^{\,2}\, \vec{\kappa}_3^{\,2}} J_{\gamma^* M}(\vec{\kappa}_1, \vec{\kappa}_2; \vec{q}) \cdot J_{pp}(\vec{\kappa}_1, \vec{\kappa}_2; \vec{q}) \ . \tag{1}$$

The kinematical notations are given in Fig.1. Here $s = (q+p)^2$, $t = (p_1 - p_M)^2 = -\vec{q}^{\,2}$ is the momentum transfer with $\vec{q} = \vec{\kappa}_1 + \vec{\kappa}_2 + \vec{\kappa}_3$. The impact factors $J_{\gamma M}$ and J_{pp} correspond to the upper and lower part of Fig. 1, correspondingly.

The impact factor of $\gamma^* \to M$ transition, $J_{\gamma^* M}$, can be derived from QCD. It can be expressed through the amplitude for the $q\bar{q}$ production $J_{\gamma^* q\bar{q}}$ which has the form [7]

$$\begin{aligned}
J_{\gamma^* q\bar{q}} &= e g^3 Q_q \left(\frac{d^{abc}}{4N}\right) \bar{u}_1 \,[\,\cdots\,]\, \frac{\hat{p}_2}{s} u_2 \\
[\,\cdots\,] &= mR\hat{e} - 2x(\vec{Q}\vec{e}) - \hat{Q}\hat{e} \\
R &= \left(\frac{1}{m_0^2 + \vec{q}_1^{\,2}} - \sum_{i=1}^{3} \frac{1}{m_0^2 + (\vec{\kappa}_1 - \vec{q}_1)^2}\right) - (\vec{q}_1 \leftrightarrow \vec{q}_2) \\
\vec{Q} &= \left(\frac{\vec{q}_1}{m_0^2 + \vec{q}_1^{\,2}} + \sum_{i=1}^{3} \frac{\vec{\kappa}_i - \vec{q}_1}{m_0^2 + (\vec{\kappa}_1 - \vec{q}_1)^2}\right) + (\vec{q}_1 \leftrightarrow \vec{q}_2)
\end{aligned}$$

$$m_0^2 \;=\; m^2 + Q^2 x(1-x) \tag{2}$$

In the above equation $e^\mu = (0, \vec{e}, 0)$ is the polarization vector of the photon. We denote also as usual $\alpha = e^2/4\pi = 1/137$, $\alpha_s = g^2/4\pi$, $Q_q e$ is the quark charge and d^{abc} are totally symmetric structure constant for $SU(N)$ color group. $\vec{q}_1$ and $\vec{q}_2$ are the transverse momenta of quark and antiquark, $0 < x < 1$ is the fraction of quark longitudinal momentum. The virtuality of photon is $p_1^2 = -Q^2$, m is the quark mass. $u_1 = u(q_1)$, $u_2 = u(q_2)$ are the quark spinors. It is convenient to introduce the variables

$$\xi = 2x - 1 \;\;, \quad \nu = \frac{4m^2 + Q^2(1-\xi^2)}{\vec{q}^{\,2}} \;. \tag{3}$$

We will neglect the transverse motion of quarks inside meson (the collinear approximation). Then the $q\bar{q} \to M$ transition is described by the Eq. (2), in which quark spinors are substituted by the meson wave function (see e.g. [9, 10]) One gets

(a) for the tensor meson T with helicity $\lambda = 0$

$$Q_q \bar{u}_1 \ldots u_2 \to \frac{Q_T}{4N} \int\limits_{-1}^{+1} d\xi\, f_T \varphi_T(\xi) Tr\,(\ldots \hat{p}_3) \tag{4}$$

(b) for the pseudoscalar mesons P

$$Q \bar{u}_1 \ldots u_2 \to \frac{Q_P}{4N} \int\limits_{-1}^{+1} d\xi\, f_P \varphi_P(\xi) Tr\left(\ldots \gamma^5 \hat{p}_3\right) . \tag{5}$$

The trace is performed over vector and color indices. The quantity Q_M ($M = P, T$) is the average quark charge in the meson M. For instance $|\pi^0> = (|u\bar{u}> - |d\bar{d}>)/\sqrt{2}$, therefore

$$Q_\pi = \frac{1}{\sqrt{2}}(+\frac{2}{3} - (-\frac{1}{3})) = \frac{1}{\sqrt{2}} \;.$$

It should be noticed that the production of tensor mesons with the helicity $\lambda = \pm 1, 2$ is suppressed in LLA at least by a factor $\sim \left(\Lambda^2_{\rm QCD}/Q^2\right)$ [7].

According to Ref.[9] we adopt for numerical calculations the following parametrizations for the wave functions and the coupling constants:

(a) for the all light tensor mesons T

$$\begin{array}{c} \varphi_T(\xi) = \frac{15}{4}\xi\,(1-\xi^2)\;; \quad f_T = 85\,\mathrm{MeV}\;\; ; \\ Q_{a_2} = \frac{1}{\sqrt{2}}\;, \quad Q_{f_2} = \frac{1}{3\sqrt{2}}\;, \quad Q_{f_2'} = \frac{1}{3} \end{array} \tag{6}$$

(assumed here that the mixing to be an ideal one, i.e. supposing f' to be consisted from $s\bar{s}$- quark only).
(b) for the light pseudoscalar mesons P

$$\varphi_P(\xi) = \begin{cases} \frac{3}{4}(1-\xi^2) & \text{asymptotical form} \\ \\ \frac{15}{4}\xi^2(1-\xi^2) & \text{Chernyak-Zhitnitsky form} \end{cases} \quad ; \tag{7}$$

$$f_\pi = 133\,\mathrm{MeV}\ , \quad f_\eta = 150\,\mathrm{MeV}\ , \quad f_{\eta'} = 110\,\mathrm{MeV}\ ; \tag{8}$$
$$Q_\pi = \frac{1}{\sqrt{2}}\ , \quad Q_\eta = 0.38\ , \quad Q_{\eta'} = 0.14$$

(the standard $\eta\eta'$- mixing is taking into account).
(c) for η_C meson

$$\varphi_{\eta_C} = \delta(\xi)\ , f_{\eta_C} = 400\ \mathrm{MeV}\ . \tag{9}$$

The photon width of η_C is expressed through f_{η_C} as $\Gamma_{\eta_C\to\gamma\gamma} = \frac{e_c^4 f_{\eta_C}^2}{4\pi m_{\eta_C}} = 7$ KeV [14].

We use delta function for the wave function of heavy mesons. This is equivalent to the non-relativistic approximation. In the case of light meson production we take into account the nontrivial longitudinal motion of the light quarks.

The conditions of the collinearity of quark momenta are

$$\vec{q}_1 = \frac{1}{2}(1+\xi)\vec{q}, \quad \vec{q}_2 = \frac{1}{2}(1-\xi)\vec{q}, \tag{10}$$

Taking into account these relations we insert Eqs. (4,5) into Eq. (2) to obtain the impact-factors $J_{\gamma T}$, $J_{\gamma P}$

$$J_{\gamma^* T}\left(\vec{\kappa}_{1,2}, \vec{q}\right) = eg^3 \left(\frac{d^{abc}}{4N}\right) \frac{Q_T}{2} \int\limits_{-1}^{+1} d\xi f_T \varphi_T(\xi)\xi\vec{Q}\vec{e} \tag{11}$$

$$J_{\gamma^* P}\left(\vec{\kappa}_{1,2}, \vec{q}\right) = eg^3 \left(\frac{d^{abc}}{4N}\right) \frac{Q_P}{2} \int\limits_{-1}^{+1} d\xi f_P \varphi_P(\xi) \left[\vec{e}\times\vec{Q}\right] \tag{12}$$

Let us introduce dimensionless vectors $\vec{r}_i$ instead of $\vec{\kappa}_i$:

$$\vec{\kappa}_i = \frac{1}{2}\left(\vec{r}_i + \vec{n}\right) |\ \vec{q}\ |, \quad \vec{n} = \vec{q}/|\ \vec{q}\ | \tag{13}$$

With the use of these variables the impact-factors (11,12) can be written as

$$J_{\gamma^* M}\left(\vec{\kappa}_{1,2}, \vec{q}\right) = eg^3 \left(\frac{d^{abc}}{4N}\right) Q_M \frac{f_M}{|\ q\ |} \begin{cases} \left(\vec{e}\vec{F}_T\right) & \text{for} \quad M=T \\ \left[\vec{e}\times\vec{F}_P\right] & \text{for} \quad M=P \end{cases} \tag{14}$$

$$\vec{F}_M = \int\limits_{-1}^{+1} d\xi \left\{ \left[\frac{\vec{n}(1+\xi)}{\nu + (1+\xi)^2} + \sum_{i=1}^{3} \frac{\vec{r}_i - \vec{n}\xi}{\nu + (\vec{r}_i - \vec{n}\xi)^2} \right] + [\xi \leftrightarrow -\xi] \right\}$$
$$\cdot \left\{ \begin{array}{lll} \xi\varphi_T(\xi) & \text{for} & M = T \\ \varphi_P(\xi) & \text{for} & M = P \end{array} \right.$$

It should be noted that $\vec{F}_M(\vec{r}_i, \vec{n}) \to 0$ at $\vec{r}_i \to -\vec{n}$.

The impact factor which describes the lower part of Fig.1 (J_{pp}) cannot be calculated in the region of small momentum transfers ($|t| \leq \Lambda^2_{\text{QCD}}$) within perturbative QCD since the mean transverse distance between the quarks inside the proton is large. Therefore, a phenomenological ansatz is used for this quantity. There are several restrictions on this impact factor imposed by gauge invariance: J_{pp} must vanish when $\vec{\kappa}_i$ $(i = 1, 2, 3)$ tends to zero. This property ensures the convergence of the integral in Eq. (1). Bose symmetry demands that J_{pp} should be symmetric under interchange of t-channel gluons. The impact factors $J_{\gamma q\bar{q}}$, $J_{\gamma M}$ discussed above satisfy both those requirements.

For J_{pp} we use the expression derived by Fukugita and Kwieciński within the eikonal model [11]. In our conventions it reads

$$J_{pp}(\vec{\kappa}_1, \vec{\kappa}_2, \vec{\kappa}_3) =$$
$$\bar{g}^3 \left(\frac{d^{abc}}{4N} \right) \cdot 3 \cdot \left[F(\vec{q}, 0, 0) - \sum_{i=1}^{3} F(\vec{\kappa}_i, \vec{q} - \vec{\kappa}_i, 0) + 2F(\vec{\kappa}_1, \vec{\kappa}_2, \vec{\kappa}_3) \right] \quad (15)$$

where $\vec{q} = \vec{\kappa}_1 + \vec{\kappa}_2 + \vec{\kappa}_3$ and

$$F(\vec{\kappa}_1, \vec{\kappa}_2, \vec{\kappa}_3) = \frac{A^2}{A^2 + \frac{1}{2}[(\vec{\kappa}_1 - \vec{\kappa}_2)^2 + (\vec{\kappa}_2 - \vec{\kappa}_3)^2 + (\vec{\kappa}_3 - \vec{\kappa}_1)^2]} \quad (16)$$

with $\bar{g}^2/(4\pi) \approx 1$ and $A = m_\rho/2$. This eikonal model describes correctly the values of total hadronic cross sections [12, 11]. The normalization in Eq. (15) can be understood in the following way. The three-gluon impact factor of point like fermions does not depend on gluon momenta. Hence this coupling is simply the number [7]

$$J^{3G}_{qq} = g^3 \left(\frac{d^{abc}}{4N} \right) . \quad (17)$$

The first term in square brackets in Eq. (15) results from diagrams where all three t-channel gluons couple to the same quark line. When q is equal to zero the interaction with the t-channel gluons does not disturb quark motion and, due to the normalization of proton wave function, this part of the proton impact factor is equal to $3 \cdot g^3 \left(\frac{d^{abc}}{4N} \right)$ where the factor 3 counts the number of constituent quarks inside the proton. In the case of the proton electromagnetic form factor this coefficient 3 will be changed to 1 due to the fractional electric charges of quarks

$1 = 2/3 + 2/3 - 1/3$. The other terms in square brackets of Eq. (15) describe the contribution of the other diagrams (where gluons couple to different quark lines). The whole expression (15) is gauge invariant and obeys Bose symmetry. In terms of dimensionless variables the proton impact factor reads

$$J_{pp} = 3\bar{g}^3 \left(\frac{d^{abc}}{4N} \right) F_p \tag{18}$$

$$\begin{aligned} F_p \;=\; & \left[\frac{1}{1+4z} - \sum_{i=1}^{3} \frac{1}{1+z(3\vec{r}_i^{\,2}+1)} + \right. \\ & \left. \frac{2}{1+\frac{z}{2}[(\vec{r}_1-\vec{r}_2)^2+(\vec{r}_2-\vec{r}_3)^2+(\vec{r}_3-\vec{r}_1)^2]} \right] \end{aligned} \tag{19}$$

with $z = \frac{\bar{g}^2}{m_\rho^2}$, $\vec{r}_3 = -\vec{r}_1 - \vec{r}_2 - \vec{n}$.

Substituting the impact factors (14, 18) into Eq. (1) the following expression for the amplitude is obtained

$$\mathcal{M}_{\gamma^* p \to Mp} = e g^3(\mu^2) \bar{g}^3 Q_M \frac{5}{18\pi^2} \frac{s f_M}{|t|^{3/2}} I \cdot \left\{ \begin{array}{lll} [\vec{e} \times \vec{n}] & for & M = P \\ (\vec{e} \cdot \vec{n}) & for & M = T \end{array} \right. \tag{20}$$

where

$$I = \frac{1}{12\pi^2} \int \frac{d^2 r_1 d^2 r_2}{(\vec{r}_1+\vec{n})^2(\vec{r}_2+\vec{n})^2(\vec{r}_1+\vec{r}_2)^2} F_M F_p \tag{21}$$

In the case of η_C production Eq.(20) can be compared with the results reported in [8]. Our amplitude (20) is 3 times smaller.

As already mentioned in the Introduction, at large t the diffractive process with proton dissociation $\gamma p \to MX$ dominates over the elastic one. In this case the cross section can be expressed through the cross section of photon-quark scattering and the quark densities inside proton (Fig.2)

$$\frac{d\sigma}{dt\, dx} = \sum_{q_i, \bar{q}_i} \frac{d\hat{\sigma}_{\gamma q \to Mq}}{dt} \left[q_i(x,t) + \bar{q}_i(x,t) \right] \tag{22}$$

It is important to note that the three-gluon C-odd odderon does not couple to gluons in the proton [13].

3 Discussion of the results

The results which are presented below were obtained by numerical evaluation of the integral I in Eq. (20), and a similar integral for hard diffractive scattering on quark, with the replacement $g^3\bar{g}^3 \to g^6(\mu^2)$. Let us start with the discussion of the results for elastic η_C production.

In the calculations we use the scale $\mu^2 = m_c^2$, $m_c = 1.5$ GeV and $\Lambda_{\rm QCD}^{(4)} = 0.2$ GeV . Our results for the unpolarized differential cross section for the process $\gamma^* p \to \eta_C p$ are given in Fig.3. This cross section vanishes at $t = 0$, it reaches the maximum at $|t| \approx 0.5$ GeV falling down for larger $|t|$. The differential cross section decreases rapidly with the growth of photon virtuality. In Fig. 3 we show also the result for $Q^2 = 0$ reported in Ref. [8] which is approximately 4 times larger than ours. We would expect here a factor 9 difference since the amplitude (20) is 3 times smaller than the one quoted in [8].

In Fig. 4 we present the cross section integrated over the region $0 < |t| < t_0$ as a function of t_0. We see that the main part of total cross section originates from the region of small $|t|$. For real photoproduction the total cross section is 47 pb , it decreases with growing Q^2.

Let us discuss now the results for large momentum transfer t when the proton disintegrates. In this case coupling constant was evaluated at the scale $\mu^2 = |t|$. The cross section of photon-quark scattering was convoluted with the GRV parametrization [15] of the quark densities, see Eqs. (22). The results are given in Figs. 5 and 6. Comparing these figures with the corresponding cross sections for elastic η_C production (Figs. 3, 4) one can see that process $\gamma^* p \to \eta_C X_p$ is not considerably suppressed. For example its cross section is 11 pb for $|t| > 3$ GeV^2 and $x \geq 0.1$.

Calculating light meson production we have an additional nontrivial integration over variable ξ, describing the longitudinal motion of light quarks. The results for the cross sections of π^0, a_2 meson production are shown in Figs. 7 and 8. The cross sections for other meson production can be readily obtained by changing corresponding coupling constants given in Eqs. (6,8). The strong coupling constant was evaluated at the scale $\mu^2 = Q^2 + |t|$. In the case of π^0 production the asymptotic form of the wave function (7) was used. Comparing Eqs. (14) and (5,6) it can be seen that the formula for the amplitude of tensor meson production is similar to one for the production of pseudoscalar meson described by the Chernyak-Zhitnitsky wave function. They differ only in the meson coupling constants and polarization dependence. The unpolarized cross sections in that case differ by the coupling constants only, e.g. $d\sigma_\pi = d\sigma_{a2}(\frac{f_\pi}{f_{a_2}})^2$. Comparing the plots for π^0 and a_2 it is seen that the shape of differential cross sections depends strongly on the form of meson wave function. The position of the dip is very sensitive to the form of wave function of the light meson, it shifts to larger t with increasing photon virtuality Q^2.

In Tab. 1 the cross sections of diffractive meson production are compared. The cross sections of π^0, a_2 meson production are comparable to the cross section of η_C production for $Q^2 \approx m_c^2$. They fall off very rapidly when Q^2 increases.

In the calculations of the elastic process $\gamma^* p \to M^+ p$ there is an important non-perturbative input – the three–gluon proton impact factor. The form of J_{pp} (proposed by Fukugita and Kwiecinski [11]) which we use looks very natural: it

Table 1: Table of cross sections for diffractive meson production (for light mesons $|t| > t_0 = 3$ GeV^2). The cross sections are given in pb.

meson / Q^2 [GeV^2]	0	1	2	5	10
η_C	47	33	24	11	-
π^0	-	-	100	8.4	1.1
a_2	-	-	24	2.8	0.52

is gauge invariant, Bose symmetric and incorporates the parameters $\bar{g}$ (frozen strong coupling) and a parameter related to the size of the proton $A \approx m_\rho/2$. We suppose that the uncertainty related to its form does not affect seriously our value of differential cross section in the region of small momentum transfers, and the value of total cross section since it saturates mostly by the region of small t (see Figs. 3, 4). But it might be that at large transfers $|t| >> \Lambda^2_{\mathrm{QCD}}$ this impact factor has to be changed since it does not obey the dimensional counting rules. Note that there is no such uncertainty in the case of inelastic process $\gamma p \to M X_p$. The non-perturbative input, the quark densities in the proton, is known. Due to that this process has less theoretical uncertainties, but its cross section in the region $t_0 < |t|, t_0 = 2 \div 3$ GeV is smaller than the total cross section for the elastic process.

Our calculations are done in the Born approximation. One can try to estimate roughly the effects related to the interaction of t-channel gluons by multiplying the formulae for cross sections by the Regge factor $(\frac{s}{\Lambda^2})^{2(\lambda_{odd}-1)}$. Using a value for the odderon intercept of $\lambda_{odd} = 1.1$ and $\Lambda^2 \sim m^2_{\eta_c}, Q^2$ this factor can be as large as 5 at HERA energies.

The values of total cross sections which we calculated in the Born approximation are 47 pb for η_C photoproduction and similar values are obtained for production of light mesons by virtual photon, if its virtuality $Q^2 \sim m_c^2$. These cross sections are small and it might be difficult to observe reactions with such cross sections at HERA. The situation is much better if the odderon intercept is greater than 1 and hence the cross sections would be a few times larger. Therefore it seems that the experimental searches for $C = +1$ mesons diffractive production at HERA can bring new interesting results.

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Figures

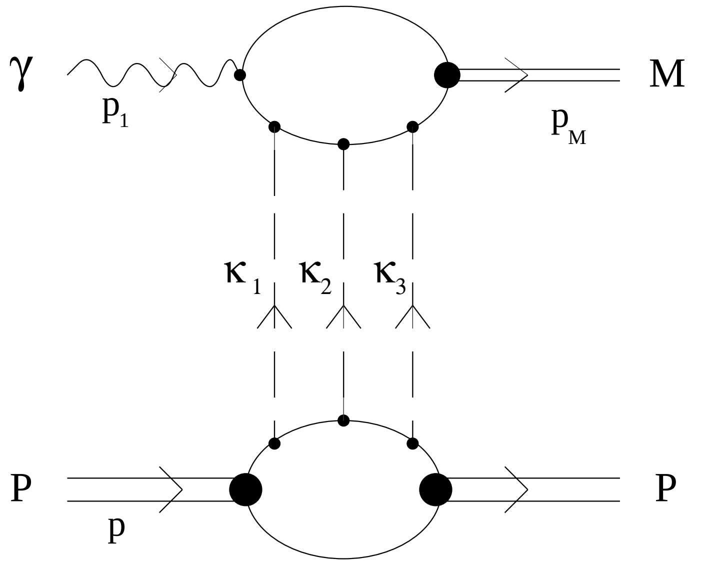

Figure 1: The kinematics of elastic diffractive process $\gamma^{*}p \to Mp$.

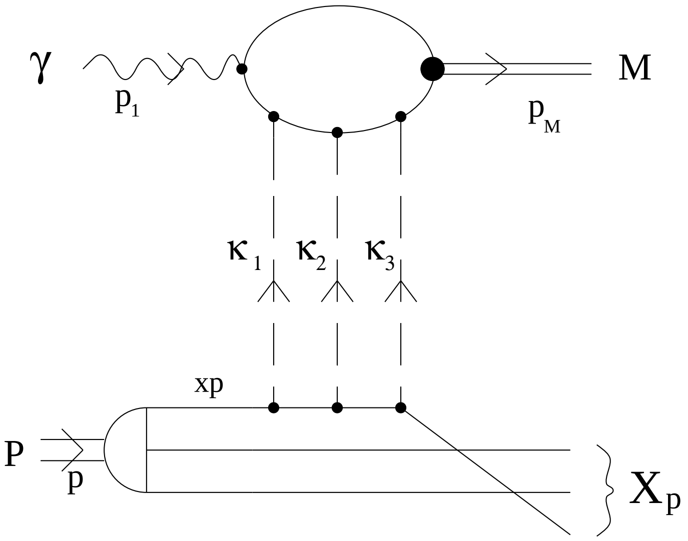

Figure 2: The hard diffractive process $\gamma^* p \to \eta_C X_p$. X_p are the hadrons on that proton disintegrates.

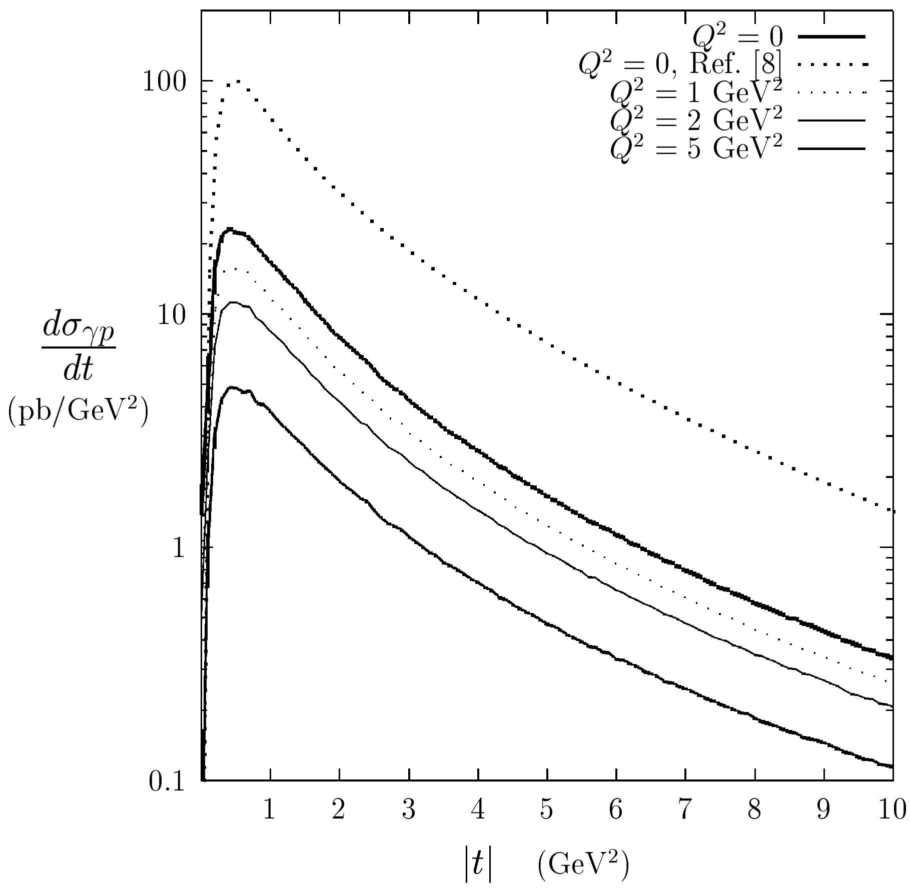

Figure 3: Differential cross section for $\gamma^* p \to \eta_C p$. The strong coupling constant was evaluated with the scale $\mu^2 = m_c^2$ and $\Lambda^{(4)}_{\rm QCD} = 0.2$ GeV.

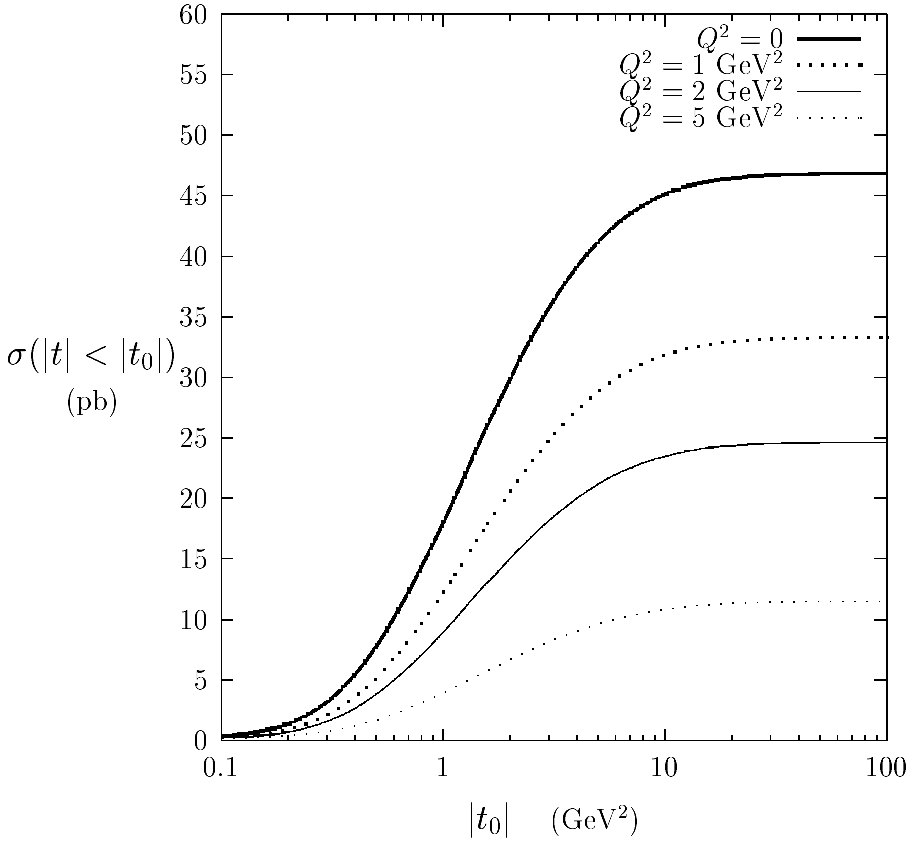

Figure 4: Total cross section for the process $\gamma^* p \to \eta_C p$ integrated over the region $0 < |t| < |t_0|$ as a function of $|t_0|$.The strong coupling constant was evaluated with the scale $\mu^2 = m_c^2$ and $\Lambda^{(4)}_{\mathrm{QCD}} = 0.2$ GeV.

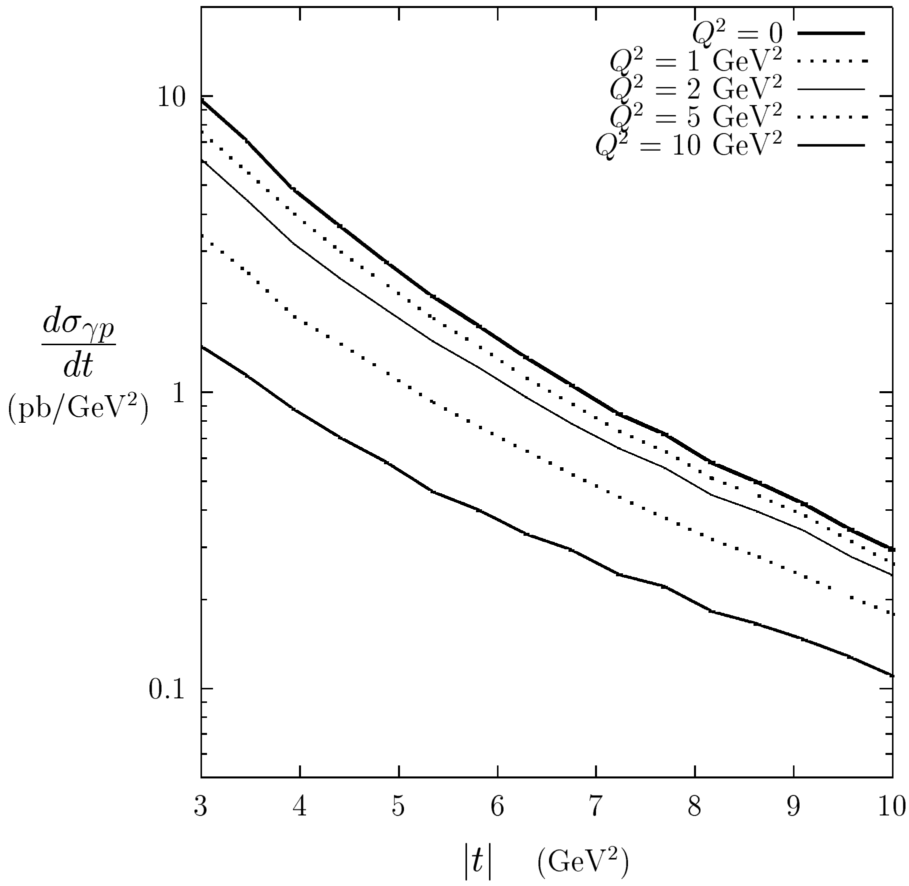

Figure 5: Differential cross section for inelastic process $\gamma^* p \to \eta_C X_p$. The strong coupling constant was evaluated with the scale $\mu^2 = |t|$ and $\Lambda^{(4)}_{\rm QCD} = 0.2$ GeV. The integral over the quark densities in the proton was limited to $x \geq 0.1$.

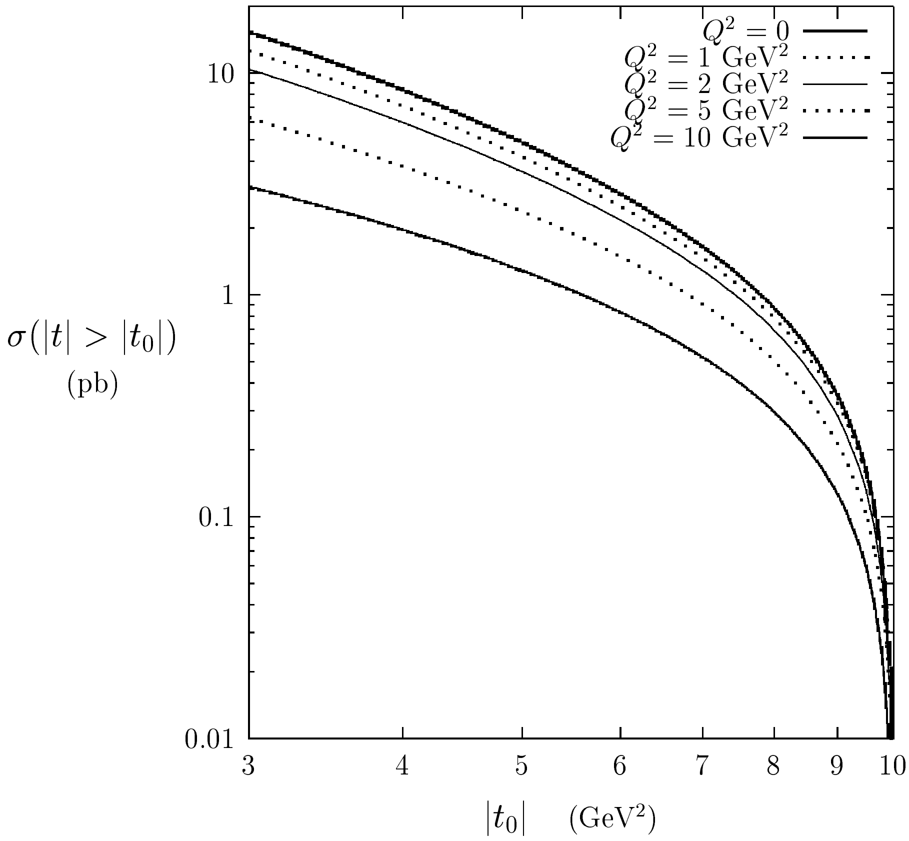

Figure 6: Total cross section for the process $\gamma^* p \to \eta_C X_p$ integrated over the region $|t_0| < |t|$. The integral over the quark densities in the proton was limited to $x \geq 0.1$. The strong coupling constant was evaluated with the scale $\mu^2 = |t|$ and $\Lambda_{\rm QCD}^{(4)} = 0.2$ GeV.

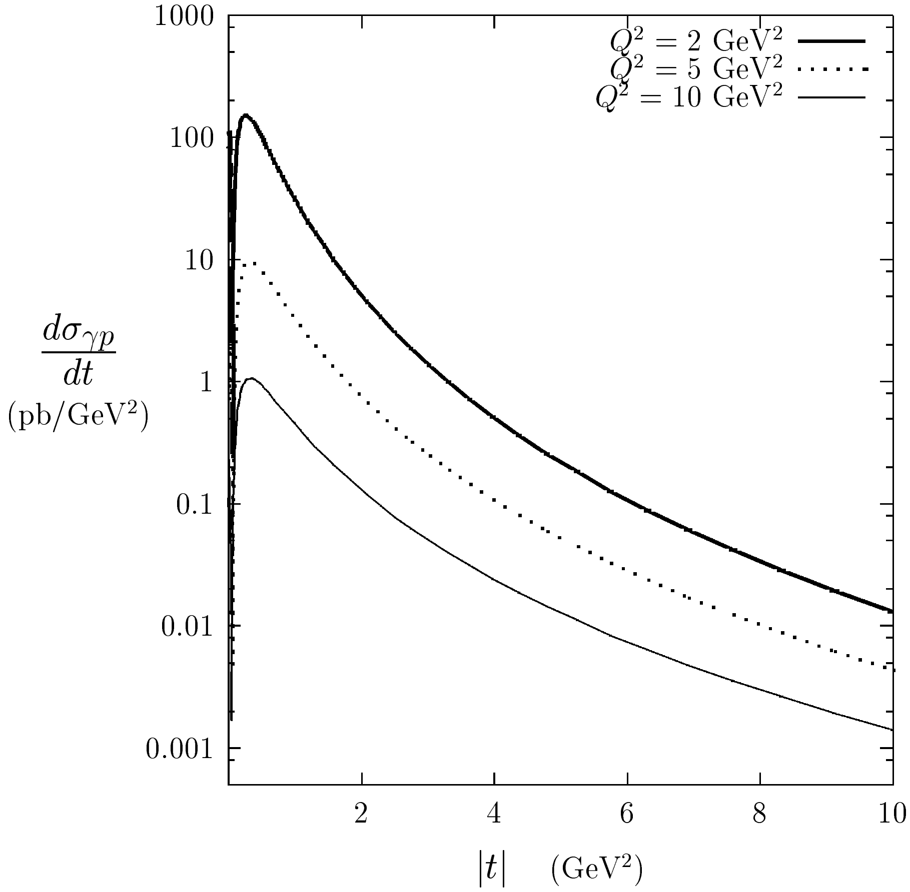

Figure 7: The differential cross section for the process $\gamma^* p \to \pi^0 p$. The pion wave function was approximated by asymptotical form. The strong coupling constant was evaluated with the scale $\mu^2 = Q^2 + |t|$ and $\Lambda^{(4)}_{\mathrm{QCD}} = 0.2$ GeV.

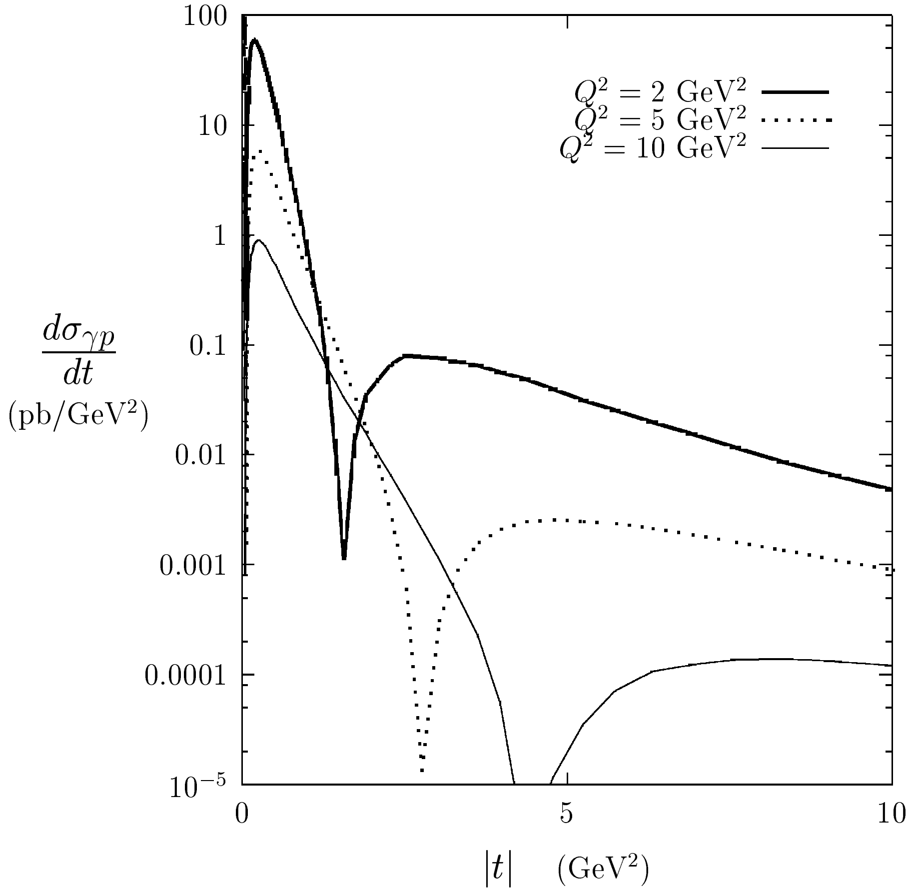

Figure 8: The differential cross section for the process $\gamma^* p \to a_2 p$. The strong coupling constant was evaluated with the scale $\mu^2 = Q^2 + |t|$ and $\Lambda^{(4)}_{\rm QCD} = 0.2$ GeV.