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DIFFRACTIVE PRODUCTION OF C-EVEN MESONS IN γp AND ep COLLISIONS* **

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We consider the processes of diffractive charge-parity $C=+1$ 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. The possibility to study this reactions at ep collider HERA is discussed briefly.

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1. Odd charge-parity exchange

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. Studies of these reactions are of great interest since the opposite charge-parities C of photon ($C=-1$) and M^+ meson cause definite C-parity $C=-1$ of the t- channel exchange. Consequently 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.

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Reggeons having $C = -1$ are the ω trajectory and the odderon. According to fits to data on total hadron–hadron cross sections the intercept of ω trajectory is close to 0.5 [1] and its contribution decreases with energy. The odderon is the C odd partner of pomeron with an intercept ≥ 1 which has been introduced in phenomenology long time ago [2]. It is assumed that, like the pomeron, the odderon is related to gluonic degrees of freedom. The odderon exchange results in a difference between, for example, the total pp and $p\bar{p}$ cross sections which does not decrease with increasing 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 the odderon is described by exchange of three interacting reggeized gluons (hard odderon), see [3]. In Born approximation, *i.e.* considering the pure three–gluon exchange in the t-channel, the intercept of the odderon is equal to 1. The effect of gluon interactions is expected to increase of odderon intercept [4]. Therefore the hard odderon should manifest itself in hard diffractive processes at high energies.

Here we shall consider the diffractive production of $C = +1$ mesons in Born approximation, *i.e.* involving the exchange of three noninteracting gluons in the t-channel.

2. Impact factors and meson wave functions

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 a heavy quark (for example, in η_C production) or the virtuality Q^2 of the photon in the production of light mesons. 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 the transfered momentum $|t| \geq m_\rho^2$ the process with photon disintegration shown in Fig. 2 becomes more important than the elastic one (Fig. 1). In the region of large t ($|t| \gg \Lambda_{\rm QCD}^2$) the cross section for this process can be expressed in terms of the photon-quark cross section and the quark densities inside the proton.

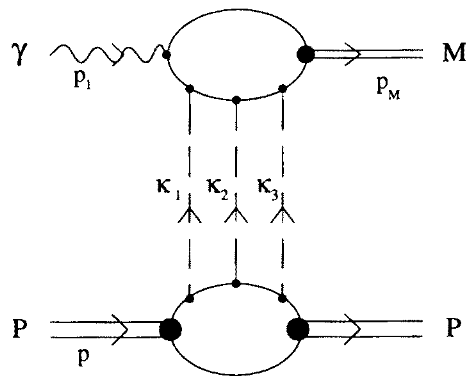

Fig. 1. The kinematics of elastic diffractive process $\gamma^* p \to Mp$.

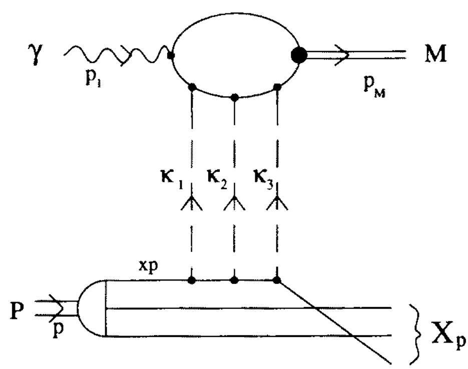

Fig. 2. The hard diffractive process $\gamma^* p \to MX_p$. X_p are the hadrons originating from proton disintegration.

In this case there is no soft scale and we deal with a purely hard diffractive process where the perturbative odderon should contribute. The amplitude determined by three-gluon exchange can be written in impact representation as an integral over the gluon transverse momenta $\vec{\kappa}_i$ [5]

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

The kinematical notations are given in Fig. 1. Here $s = (q+p)^2$ and $t = (p_1 - p_M)^2 = -\vec{q}^{\,2}$ 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 $J_{\gamma^* M}$ for the $\gamma^* \to M$ transition can be derived within perturbative QCD. It can be expressed in terms of the amplitude $J_{\gamma^* q\bar{q}}$ for $q\bar{q}$ production which has the form [5]

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

In the above expression $e^\mu = (0, \vec{e}, 0)$ is the polarization vector of the photon. Furthermore we use the notations $\alpha = e^2/4\pi = 1/137$ and $\alpha_s = g^2/4\pi$. $Q_q e$ is the quark charge and d^{abc} are the totally symmetric structure constants for the 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 the quark longitudinal momentum. The virtuality of photon is $p_1^2 = -Q^2$, m is the quark mass. The quark spinors are denoted by $u_1 = u(q_1)$, $u_2 = u(q_2)$.

We will neglect the transverse motion of quarks inside the meson (collinear approximation). Then the $q\bar{q} \to M$ transition is described by the Eq. (2) in which the quark spinors are substituted by the meson wave function $\varphi_M(\xi)$, $\xi = 2x - 1$ (see *e.g.* [7,8]). One obtains the substitutions:

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

$$
Q_q \bar{u}_1 \ldots u_2 \to \frac{Q_T}{4N} \int_{-1}^{+1} d\xi f_T \varphi_T(\xi)\, \mathrm{Tr}\,(\ldots \hat{p}_3)\ ; \qquad (3)
$$

(b) for the pseudoscalar mesons P

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

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

According to Ref. [7] we adopt for numerical calculations the following parametrizations for the wave functions and the coupling constants:

(a) for light tensor mesons T

$$\varphi_T(\xi) = \tfrac{15}{4}\xi\,(1-\xi^2)\ ; \quad f_T = 85\,\mathrm{MeV}\ ; \qquad (5)$$
$$Q_{a_2} = \tfrac{1}{\sqrt{2}}\ , \quad Q_{f_2} = \tfrac{1}{3\sqrt{2}}\ , \quad Q_{f_2'} = \tfrac{1}{3}\ ,$$

where an ideal mixing is assumed, *i.e.* f' consists of $s\bar{s}$-states only.

(b) for light pseudoscalar mesons P

$$\varphi_P(\xi) = \begin{cases} \tfrac{3}{4}(1-\xi^2) & \text{asymptotical form}\,, \\ \tfrac{15}{4}\xi^2(1-\xi^2) & \text{Chernyak–Zhitnitsky form}\,; \end{cases} \qquad (6)$$

$$f_\pi = 133\,\mathrm{MeV}\ , \quad f_\eta = 150\,\mathrm{MeV}\ , \quad f_{\eta'} = 110\,\mathrm{MeV}\ ; \qquad (7)$$
$$Q_\pi = \tfrac{1}{\sqrt{2}}\ , \quad Q_\eta = 0.38\ , \quad Q_{\eta'} = 0.14\ ,$$

where the standard $\eta\eta'$- mixing is taken into account.

(c) for the η_C meson

$$\varphi_{\eta_C} = \delta(\xi)\ , \quad f_{\eta_C} = 400\ \mathrm{MeV}\ . \qquad (8)$$

The photon width of η_C is expressed through f_{η_C} as $\Gamma_{\eta_C\to\gamma\gamma} = e_c^4 f_{\eta_C}^2/(4\pi m_{\eta_C})$ $\approx$ 7 keV [9]. The usage of a delta function for the wave function of heavy mesons 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.

Introducing the dimensionless vectors $\vec{r}_i$, $\vec{\kappa}_i = \frac{1}{2}\,(\vec{r}_i + \vec{n})\,|\,\vec{q}\,|$, $\vec{n} = \vec{q}/|\,\vec{q}\,|$ and taking into account the conditions of the collinearity of quark momenta the impact-factors read

$$J_{\gamma^* M}\,(\vec{\kappa}_{1,2}, \vec{q}) = eg^3 \left(\frac{d^{\,\mathrm{abc}}}{4N}\right) Q_M \frac{f_M}{|\,q\,|} \begin{cases} \left(\vec{e}\cdot\vec{F}_T\right) & \text{for} \quad M = T \\ \left[\vec{e}\times\vec{F}_P\right] & \text{for} \quad M = P\,, \end{cases} \qquad (9)$$

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

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 J_{pp} which describes the lower part shown in Fig. 1 cannot be calculated in the region of small momentum transfers ($|t| \leq \Lambda^2_{\rm QCD}$) 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.

For J_{pp} we use the expression derived by Fukugita and Kwieciński within the eikonal model [10]. In our conventions it reads

$$J_{pp}(\vec{\kappa}_1,\vec{\kappa}_2,\vec{\kappa}_3) = \bar{g}^3\left(\frac{d^{\rm abc}}{4N}\right)\times 3\left[F(\vec{q},0,0) - \sum_{i=1}^{3} F(\vec{\kappa}_i,\vec{q}-\vec{\kappa}_i,0) + 2F(\vec{\kappa}_1,\vec{\kappa}_2,\vec{\kappa}_3)\right], \quad (10)$$

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

$$F(\vec{\kappa}_1,\vec{\kappa}_2,\vec{\kappa}_3) = \frac{A^2}{A^2 + \frac{1}{2}[(\vec{\kappa}_1-\vec{\kappa}_2)^2 + (\vec{\kappa}_2-\vec{\kappa}_3)^2 + (\vec{\kappa}_3-\vec{\kappa}_1)^2]} \quad (11)$$

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 [10, 11]. The normalization of Eq. (10) can be understood in the following way. The three-gluon impact factor of point like fermions does not depend on the gluon momenta. Hence this coupling is given by [5]

$$J^{3G}_{qq} = g^3\left(\frac{d^{\rm abc}}{4N}\right). \quad (12)$$

The first term in brackets in Eq. (10) results from diagrams where all the 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 $3\,g^3\left(\frac{d^{\rm abc}}{4N}\right)$ 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 $1 = 2/3 + 2/3 - 1/3$ and the gauge group factor is replaced by 1. The other terms in brackets in Eq. (10) describe the contribution of the other diagrams where gluons couple to different quark lines. The whole expression (10) is gauge invariant and obeys Bose symmetry. In terms of dimensionless variables the proton impact factor reads

$$J_{pp} = 3\bar{g}^3 \left(\frac{d^{\rm abc}}{4N} \right) F_p \,, \tag{13}$$

$$F_p = \left[\frac{1}{1+4z} - \sum_{i=1}^{3} \frac{1}{1+z(3\vec{r}_i^{\,2}+1)} + \frac{2}{1+\frac{z}{2}[(\vec{r}_1-\vec{r}_2)^2+(\vec{r}_2-\vec{r}_3)^2+(\vec{r}_3-\vec{r}_1)^2]} \right] \tag{14}$$

with $z = \frac{\vec{q}^{\,2}}{m_p^2}$, $\vec{r}_3 = -\vec{r}_1 - \vec{r}_2 - \vec{n}$.

Substituting the expressions for impact factors into Eq. (1) the following formula for the amplitude is obtained

$$\mathcal{M}_{\gamma^* p \to Mp} = e g^3(\mu^2) \bar{g}^3 Q_M \frac{5}{18\pi^2} \frac{s f_M}{|t|^{3/2}} I \cdot \begin{cases} [\vec{e} \times \vec{n}] & \text{for} \quad M = P \\ (\vec{e} \times \vec{n}) & \text{for} \quad M = T \,, \end{cases} \tag{15}$$

where

$$I = \frac{1}{12\pi^2} \int \frac{d^2 r_1 d^2 r_2}{(\vec{r}_1+\vec{n})^2 (\vec{r}_2+\vec{n})^2 (\vec{r}_1+\vec{r}_2)^2} F_M F_p \,. \tag{16}$$

In the case of η_C production Eq. (15) can be compared with the results reported in [6]. Our amplitude (15) is 3 times smaller.

As already mentioned above, 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 in terms of the cross section of photon-quark scattering and the quark densities inside proton (Fig. 2)

$$\frac{d\sigma}{dt\, dx} = \sum_{q_i, \bar{q}_i} \frac{d\hat{\sigma}_{\gamma q \to Mq}}{dt} \left[q_i(x,t) + \bar{q}_i(x,t) \right] . \tag{17}$$

It is important to notice that the three-gluon C-odd odderon does not couple to gluons in the proton [12].

3. Results on diffractive meson production

The results which are presented below have been obtained by numerical evaluation of the integral I in Eq. (15) and of a similar integral for hard diffractive scattering on quarks. The hard cross section $d\hat{\sigma}_{\gamma q \to Mq}/dt$ was calculated from the amplitude (20) with the replacement $g^3 \bar{g}^3 \to g^6(\mu^2)$, where μ is the scale for the evaluation of the running strong coupling constant.

Let us start with the discussion of the results for elastic η_C production. In the calculations we use the scale $\mu^2 = m_c^2$, $m_c = 1.4$ GeV and $\Lambda^{(4)}_{\rm QCD} = 0.2$ GeV (for 4 flavors). 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 and decreases 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. [6] which is approximately 4 times larger than ours. The reasons for this discrepancy are the following: *(i)* as discussed above, the amplitude (20) turns out to be smaller in normalization by a factor of 3 compared to the result quoted in [6], and *(ii)* we use $\alpha_s(m_c^2) \approx 0.387$ instead of 0.3 as done in [6]. Our estimate for the total cross section is about 47 pb, it decreases with increasing Q^2.

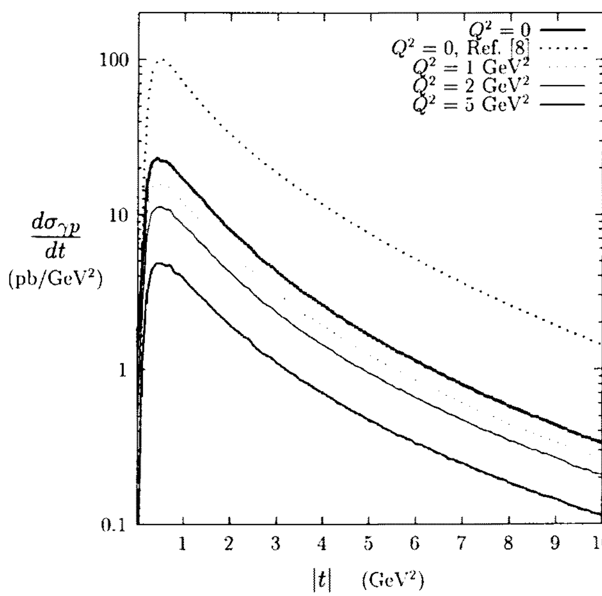

Fig. 3. Differential cross section for $\gamma^* p \to \eta_C p$.

Now let us discuss the results for large momentum transfer t when the proton disintegrates. In this case the coupling constant has been evaluated at the scale $\mu^2 = |t|$. The cross section of photon-quark scattering was convoluted with the GRV parametrization [13] of the quark densities in the proton, see Eqs. (17). The results are given in Fig. 4. Comparing these figure with the corresponding cross sections for elastic η_C production (Fig. 3) one can see that the 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 $x \geq 0.1$.

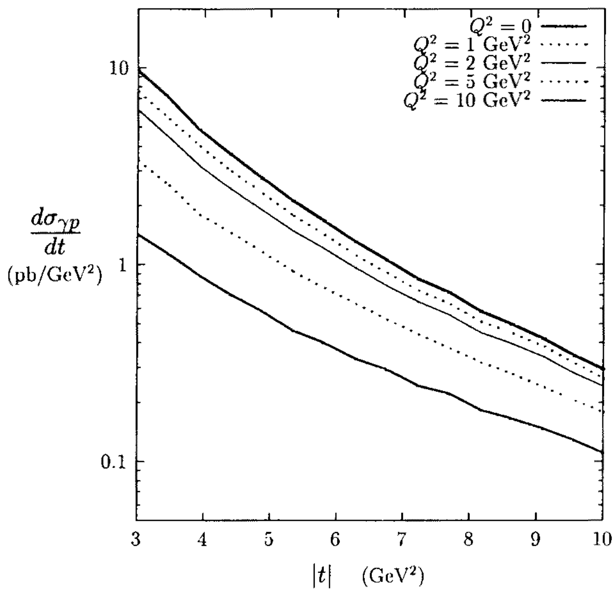

Fig. 4. Differential cross section for inelastic process $\gamma^* p \to \eta_C X_p$.

Calculating light meson production we have an additional nontrivial integration over the variable ξ, describing the longitudinal motion of light quarks. We cite here in Table I our results for the total cross sections of π^0 and a_2 meson production. The cross sections for the production of other mesons can be readily obtained by changing corresponding coupling constants given in Eqs. (5), (7). The strong coupling constant has been evaluated at the scale $\mu^2 = Q^2 + |t|$. In the case of π^0 production the asymptotic form of the wave function (6) has been used. The cross sections of π^0 and 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 with increasing Q^2.

TABLE I

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	4.2
π^0	-	-	100	8.4	1.1
a_2	-	-	31	3.6	0.68

Performing the calculations for the diffractive production of light mesons we have to take care careful about the lower limit of the photon virtuality

Q^2 which to ensure the self-consistency of our perturbative approach. To get an idea about this limit the numerical calculations of the cross sections have been done keeping in our formulae a non zero quark mass m which can be considered as the measure of non-perturbative effects. We have chosen two typical values of the constituent quark masses $m = 0.2$ GeV and $m = 0.3$ GeV. In Table II the results for the ratio of the total cross sections $\sigma(m^2)/\sigma(m^2 = 0)$ for tensor meson production are given for some values of photon virtualities Q^2. It can be concluded that perturbative calculations should be reliable for $Q^2 \gtrsim 2\text{GeV}^2$.

TABLE II

Ratio of the total cross sections $\sigma(m^2)/\sigma(m^2 = 0)$ for elastic tensor meson production for different photon virtualities.

quark mass / Q^2 [GeV^2]	1	2	5	10
$m = 200$ MeV	0.69	0.90	0.98	0.99
$m = 300$ MeV	0.50	0.74	0.92	0.95

4. Discussion

We would like to stress that the numerical results presented here are based on amplitudes the essential part of which have been derived from perturbative QCD. In particular these are the impact representation (1) and the virtual photon impact factor (2). As further inputs we need the meson wave functions and couplings which are well known from phenomenology.

The proton impact factor in the quasi-elastic diffraction which we have adopted following [10] can be motivated by perturbative calculations. However there is no hard scale in that part of the process and the form of the proton impact factor enters as an additional assumption. The form of J_{pp} (proposed by Fukugita and Kwieciński [10]) exhibits the following features: it is gauge invariant, Bose symmetric and incorporates the parameter $\bar{g}$ (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 the differential cross section in the region of small momentum transfers and also not the value of total cross section since it is saturated mostly by the region of small t. But it might be that at large momentum transfers $|t| \gg \Lambda^2_{\rm QCD}$ this impact factor has to be changed since it does not obey the dimensional counting rules.

In inelastic diffraction being characterized by proton disintegration the quark distributions in the proton take the role played by the proton impact factor before. This non-perturbative input is, of course, much better 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 $(s/\Lambda^2)^{2(\lambda_{\rm odd}-1)}$. Using a value for the odderon intercept of $\lambda_{\rm odd} = 1.1$ and $\Lambda^2 \sim m_{\eta_c}^2, Q^2$ this factor can be as large as 5 at HERA energies.

The value of the total η_C photoproduction cross section in Born approximation is 47 pb. Similar cross sections are obtained for the production of light mesons by virtual photons, if the photon virtuality Q^2 is of the order m_c^2. It might be difficult to observe reactions with such cross sections at HERA. On the other hand the situation is much better if the odderon intercept is greater than 1 resulting in cross sections being a few times larger. Therefore we expect that experimental searches for diffractive production of $C = +1$ mesons at HERA can bring new interesting results relevant for the longstanding and controversial questions about the odderon.

Note added: Recently we have learned about the revised version of Ref. [6], where the normalization of the amplitude is now in agreement with our result (15).

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