Possible probe of the QCD odderon singularity through the quasidiffractive η_c production in $\gamma\gamma$ collisions

L. Motyka
Institute of Physics, Jagellonian University, Cracow, Poland

J. Kwieciński
Department of Theoretical Physics, H. Niewodniczański Institute of Nuclear Physics, Cracow, Poland

(Received 27 April 1998; revised manuscript received 24 June 1998; published 28 October 1998)

The reactions $\gamma\gamma\to\eta_c\eta_c$ and $\gamma\gamma\to\eta_c+X$ are discussed within the three-gluon-exchange model. We give predictions for the differential cross sections and discuss the feasibility of measuring these processes at CERN LEP2 and Cornell TESLA. The total cross sections were estimated to be approximately equal to 40 fb and 120 fb for $\gamma\gamma\to\eta_c\eta_c$ and $\gamma\gamma\to\eta_c+X$, respectively, assuming exchange of elementary gluons that corresponds to the odderon with intercept equal to unity. These values can be enhanced by a factor equal to 2.4 and 2.9 for LEP2 and TESLA energies if the odderon intercept is equal to 1.07. The estimate of cross sections $\sigma(e^+e^-\to e^+e^-\eta_c\eta_c)$ and $\sigma(e^+e^-\to e^+e^-\eta_cX)$ for antitagged e^+ and e^- is also given. [S0556-2821(98)03421-3]

PACS number(s): 13.60.Le, 12.38.−t, 13.65.+i

The dominant contribution in the high-energy limit of perturbative QCD is given by the exchange of interacting gluons [1,2]. The fact that the gluons have spin equal to unity automatically implies that the cross sections corresponding to the exchange of elementary gluons are independent of the incident energies while the interaction between gluons leads to an increasing cross section. In addition to the Pomeron [1–5] one also expects the presence of the so-called ''odderon'' singularity [1–3,6,7]. In the leading logarithmic approximation the Pomeron is described by the Balitskii-Fadin-Kuraev-Lipatov (BFKL) equation which corresponds to the sum of ladder diagrams with Reggeized gluon exchange along the ladder. The odderon is described by three-gluon exchange. Unlike the Pomeron which corresponds to vacuum quantum numbers and so to positive charge conjugation the odderon is characterized by $C=-1$ (and $I=0$); i.e., it carries the same quantum numbers as the ω Regge pole. The (phenomenologically) determined intercept λ_ω of the ω Regge pole is approximately equal to 1/2 [8]. The novel feature of the odderon singularity corresponding to the gluonic degrees of freedom is the potentially very high value of its intercept $\lambda_{odd}\gg\lambda_\omega$. The exchange of the three (noninteracting) gluons alone generates singularity with the intercept equal to unity while the interaction between gluons described in the leading logarithmic approximation by the Bartels-Kwieciński-Praszałowicz (BKP) equation [6,7,9] can move the odderon intercept above [10,11] or below [12,13] unity.[1] The energy dependence of the amplitudes corresponding to $C=-1$ exchange becomes similar to the diffractive ones which are controlled by Pomeron exchange.

Possible tests of both the QCD perturbative Pomeron as well as of the odderon have to rely on (semi)hard processes where the presence of a hard scale can justify the use of perturbative QCD. A very useful measurement in this respect is the very-high-energy exclusive photoproduction (or electroproduction) of heavy charmonia (i.e., J/ψ [14,15] or η_c [16–18], etc., for probing the QCD Pomeron or odderon, respectively). The estimate of the odderon contribution to the photoproduction (or electroproduction) of even charge conjugation mesons does, however, require model assumptions about the coupling of the three-gluon system to a proton. It would therefore be useful to consider a process which could in principle be calculated entirely within perturbative QCD. The relevant measurement which fulfills those criteria is the exclusive quasidiffractive production of even C charmonia in $\gamma\gamma$ collisions or, to be precise, the processes $\gamma\gamma\to\eta_c\eta_c$ or $\gamma\gamma\to\eta_c+X$ [19,20]. The main purpose of our paper is to present a theoretical and phenomenological description of double η_c production in high-energy $\gamma\gamma$ collisions and of the process $\gamma\gamma\to\eta_c+X$ with a large rapidity gap between η_c and the hadronic system X assuming the three-gluon-exchange mechanism. In our paper we shall follow the formalism developed in Ref. [20] where the production of pseudoscalar mesons in $\gamma\gamma$ collisions within the three-gluon-exchange mechanism is discussed with great detail.

The kinematics of the three-gluon-exchange diagram to the processes $\gamma^*(q_1)+\gamma^*(q_2)\to\eta_c\eta_c$ and $\gamma^*(q_1)+\gamma^*(q_2)\to\eta_c+X$ is illustrated in Fig. 1(a) and Fig. 1(b).

The amplitude M^{ij} for the process $\gamma^*(q_1)+\gamma^*(q_2)\to\eta_c\eta_c$ which corresponds to the transverse polarization of

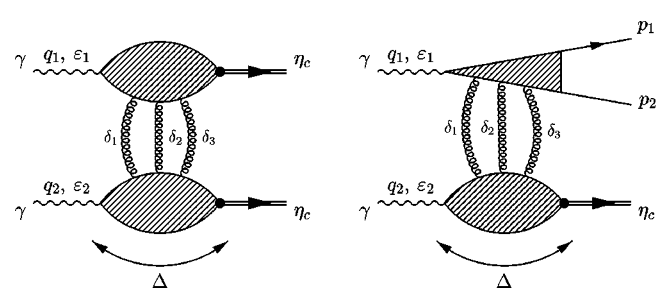

FIG. 1. Kinematics of the three-gluon-exchange mechanism of the process $\gamma\gamma\to\eta_c\eta_c$ (a) and $\gamma\gamma\to\eta_c+X$ (b).

[1]The results of Ref. [11] are, however, in conflict with a recent analysis [12].

both photons can be written in the following way:

$$M^{ij}=W^2C_{qq}\int \frac{d^2\boldsymbol{\delta}_1 d^2\boldsymbol{\delta}_2}{\boldsymbol{\delta}_1^2\boldsymbol{\delta}_2^2\boldsymbol{\delta}_3^2}\Phi^i_{\gamma\eta_c}(Q_1^2,\boldsymbol{\delta}_1,\boldsymbol{\delta}_2,\boldsymbol{\Delta})\times\Phi^j_{\gamma\eta_c}(Q_2^2,-\boldsymbol{\delta}_1,-\boldsymbol{\delta}_2,-\boldsymbol{\Delta}), \tag{1}$$

where i and j are the polarization indices,

$$C_{qq}=\frac{10}{9\pi N_c^2}\alpha_s^3(m_c^2), \tag{2}$$

and $\boldsymbol{\Delta}$ and $\boldsymbol{\delta}_i$ denote the transverse components of Δ and δ_i while $Q_i^2=-q_i^2$.

The relevant diagrams describing the impact factor $\Phi^{i,j}_{\gamma\eta_c}$ are given in Refs. [17, 20]. In the nonrelativistic approximation one gets the following expression for $\Phi^i_{\gamma\eta_c}$:

$$\Phi^i_{\gamma\eta_c}=F_{\eta_c}\sum_{k=1}^{2}\epsilon_{ik}\left[\frac{\boldsymbol{\Delta}^k}{\bar{M}_1^2+\boldsymbol{\Delta}^2}+\sum_{s=1}^{3}\frac{2\boldsymbol{\delta}_s^k-\boldsymbol{\Delta}^k}{\bar{M}_1^2+(2\boldsymbol{\delta}_s-\boldsymbol{\Delta})^2}\right], \tag{3}$$

where

$$F_{\eta_c}=\sqrt{\frac{4M_{\eta_c}\Gamma_{\eta_c\to\gamma\gamma}}{\alpha q_c^2}}, \tag{4}$$

$$\bar{M}_1^2=4m_c^2+Q_1^2. \tag{5}$$

In Eqs. (3), (4), (5) m_c, M_{η_c}, $\Gamma_{\eta_c\to\gamma\gamma}$, and q_c denote the charmed quark mass, the mass of the η_c, the η_c radiative width, and the charge of the charm quark, respectively. The formula for the impact factor $\Phi^j_{\gamma\eta_c}$ corresponding to the lower vertex is given by Eq. (3) after changing the polarization index i into j, after reversing the sign of $\boldsymbol{\delta}_s$ and of $\boldsymbol{\Delta}$ in this equation, and after changing $\bar{M}_1^2$ into $\bar{M}_2^2$ given by Eq. (5) with Q_2^2 instead of Q_1^2.

For $Q_1^2=Q_2^2\simeq 0$ it is convenient to represent the amplitude M^{ij} in the following way:

$$M^{ij}=W^2\left(M_1g^{ij}+M_2\frac{\boldsymbol{\Delta}^i\boldsymbol{\Delta}^j}{\boldsymbol{\Delta}^2}\right). \tag{6}$$

The corresponding formula for the differential cross section averaged over the transverse photons polarizations for the process $\gamma\gamma\to\eta_c\eta_c$ reads

$$\frac{d\sigma}{dt}=\frac{1}{64\pi}[(M_1+M_2)^2+M_1^2]. \tag{7}$$

For real $\gamma-s$ we have, of course, set $Q_{1,2}^2=0$.

The process $\gamma\gamma\to\eta_c+X$ is given by the diagram of Fig. 1(b) and the amplitude which corresponds to this diagram can be written as

$$M_{\gamma\gamma\to\eta_c+X}=W^2\tilde{C}_{qq}\int\frac{d^2\boldsymbol{\delta}_1 d^2\boldsymbol{\delta}_2}{\boldsymbol{\delta}_1^2\boldsymbol{\delta}_2^2\boldsymbol{\delta}_3^2}\times\Phi_{\gamma\eta_c}(Q_2^2,-\boldsymbol{\delta}_1,-\boldsymbol{\delta}_2,-\boldsymbol{\Delta})\times\Phi_{\gamma X}(Q_1^2,\boldsymbol{\delta}_1,\boldsymbol{\delta}_2,\boldsymbol{\Delta},p_1,p_2), \tag{8}$$

where

$$\Phi_{\gamma\eta_c}=\sum_{i,j=1}^{2}\epsilon_2^i\epsilon_{ij}\Phi^j_{\gamma\eta_c}, \tag{9}$$

$$\tilde{C}_{qq}=\frac{10}{9\pi N_c^2}[\alpha_s(m_c^2)\tilde{\alpha}_s]^{3/2}. \tag{10}$$

The coupling constant $\tilde{\alpha}_s$ is defined by the hard scale characteristic for the upper vertex in Fig. 1(b). The diagrams defining impact factor $\Phi_{\gamma X}$ are given in Ref. [20]. They give the following expression for $\Phi_{\gamma X}$:

$$\Phi_{\gamma X}(Q_1^2,\boldsymbol{\delta}_1,\boldsymbol{\delta}_2,\boldsymbol{\Delta},p_1,p_2)=-ieq_f\bar{u}(p_1)[m_fR\not{\epsilon}_1+2z\bar{Q}\epsilon_1+\bar{\not{Q}}\not{\epsilon}_1]\not{q}_2v(p_2), \tag{11}$$

where

$$\bar{\boldsymbol{Q}}=\left[\frac{\boldsymbol{p}_1}{m_f^2+\boldsymbol{p}_1^2}+\sum_{s=1}^{3}\frac{\boldsymbol{\delta}_s-\boldsymbol{p}_1}{m_f^2+(\boldsymbol{\delta}_s-\boldsymbol{p}_1)^2}\right]+(\boldsymbol{p}_1\to\boldsymbol{p}_2), \tag{12}$$

$$R=\left[\frac{-1}{m_f^2+\boldsymbol{p}_1^2}+\sum_{s=1}^{3}\frac{1}{m_f^2+(\boldsymbol{\delta}_s-\boldsymbol{p}_1)^2}\right]+(\boldsymbol{p}_1\to\boldsymbol{p}_2). \tag{13}$$

In Eqs. (11), (12), (13) $\bar{Q}$ is a four-vector with transverse components $\bar{\boldsymbol{Q}}$ and vanishing longitudinal components, m_f denotes the mass of the (light) quark produced as the pair of $q\bar{q}$ jets, and q_f is its charge while z is the component of the photon four-momentum q_1 carried by a quark jet. The four-momenta p_1 and p_2 denote the four-momenta of the quark (antiquark) in the final state [see Fig. 1(b)] and $\boldsymbol{p}_1$, $\boldsymbol{p}_2$ are their transverse parts, respectively. When calculating the cross sections we shall make the ''equivalent quark approximation'' [20,21] which corresponds to the approximation of setting m_f equal to zero and to retaining the dominant term in $\bar{\boldsymbol{Q}}$: i.e.,

$$\bar{\boldsymbol{Q}}=\boldsymbol{p}_1/\boldsymbol{p}_1^2+\boldsymbol{p}_2/\boldsymbol{p}_2^2. \tag{14}$$

The remaining terms in Eq. (12) can also be large for $\boldsymbol{\delta}_s\simeq\boldsymbol{p}_i$ but their dependence on $\boldsymbol{\delta}_s$ is such that after the integration over $\boldsymbol{\delta}_s$ performed in Eq. (8) they are suppressed in comparison to the leading terms.

The differential cross section for the process $\gamma\gamma\to\eta_c+X(q\bar{q})$ averaged over photon polarizations is given by the following formula:

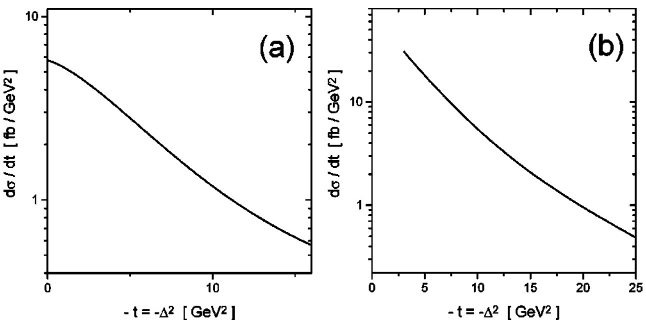

FIG. 2. Differential cross sections of the processes $\gamma\gamma\to\eta_c\eta_c$ (a) and $\gamma\gamma\to\eta_c+X$ (b) for $Q_1^2\simeq Q_2^2\simeq 0\ \text{GeV}^2$.

$$d\sigma=\frac{(2\pi)^4}{2W^2}\sum_{\{\cdots\}} M^{+}_{\gamma\gamma\to\eta_c+X} M_{\gamma\gamma\to\eta_c+X} \times dPS_3(\gamma\gamma\to\eta_c+X(q\overline{q})), \qquad (15)$$

where $\{\cdots\}$ stands for incident photons helicities, outgoing light quark colors, flavors, and polarizations and $dPS_3(\gamma\gamma\to\eta_c+X(q\overline{q}))$ is the standard parametrization of the Lorentz-invariant three-body phase space. The decomposition property

$$dPS_3(\gamma\gamma\to\eta_c+X(q\overline{q})) =(2\pi)^3\ dM_X^2\ dPS_2(\gamma\gamma\to\eta_c+X)dPS_2(X\to q\overline{q}) \qquad (16)$$

is employed and integration over the invariant mass squared M_X^2 of the $q\overline{q}$ system is performed. The remaining integration over the two-body phase space $dPS_2(X\to q\overline{q})$ in the applied approximation cannot be extended to the regions $\boldsymbol{p}_i^2\simeq 0$ since the integral would then be divergent. We do therefore introduce, following Refs. [20, 21], a physical cutoff μ which is defined by the light quark constituent masses and we set for its magnitude $\mu=0.3$ GeV. The integral is bounded from above by the condition $\boldsymbol{p}_i^2<\boldsymbol{\Delta}^2$ which assures the diffractive nature of the process. Thus we obtain a logarithmic expression $\ln|t|/\mu^2$ arising from the integral $\int_{\mu^2}^{|t|} d\boldsymbol{p}_1^2/\boldsymbol{p}_1^2$. Finally the differential cross section reads

$$\frac{d\sigma}{dt}=\frac{1}{64\pi}\left|\tilde{C}_{qq}\int\frac{d^2\boldsymbol{\delta}_1 d^2\boldsymbol{\delta}_2}{\boldsymbol{\delta}_1^2\boldsymbol{\delta}_2^2\boldsymbol{\delta}_3^2} \times\Phi_{\gamma\eta_c}(Q_2^2,-\boldsymbol{\delta}_1,-\boldsymbol{\delta}_2,-\boldsymbol{\Delta},p_1,p_2)\right|^2\frac{2\alpha\overline{e^2}}{3\pi}\ln\frac{|t|}{\mu^2}, \qquad (17)$$

where $\overline{e^2}=N_c(e_u^2+e_d^2+e_s^2)=2$ and e_i are the charges of the light quarks u, d, and s, respectively.

In Fig. 2(a) we show the differential cross section of the process $\gamma\gamma\to\eta_c\eta_c$. Unlike the photoproduction of η_c on a nucleon target this cross section does not vanish at $t=0$. For the integrated cross section we get $\sigma^{tot}_{\eta_c\eta_c}=43$ fb. The differential cross section for the process $\gamma\gamma\to\eta_c+X$ is presented in Fig. 2(b). The integrated cross section is now $\sigma^{tot}_{\eta_c X}(|t|>3\ \text{GeV}^2)=120$ fb. In our calculation we set $\alpha_s(m_c^2)=0.38$, $\tilde{\alpha}_s=0.3$, $m_c=1.4$ GeV, $M_{\eta_c}=2.98$ GeV, and $\Gamma_{\eta_c\to\gamma\gamma}=7$ keV.

The calculated cross sections are energy independent since they correspond to the exchange of three elementary and noninteracting gluons. In this approximation the odderon singularity has its intercept λ_{odd} equal to unity. Interaction between gluons can boost this intercept above unity and one can take approximately this effect into account by multiplying the cross sections by the enhancement factor

$$A_{enh}(W^2)=(W^2/M_{\eta_c}^2)^{2(\lambda_{odd}-1)}. \qquad (18)$$

The $\gamma\gamma$ system in e^+e^- collisions has a continuous spectrum. The c.m. energy squared W^2 of the $\gamma\gamma$ system is

$$W^2=z_1z_2s, \qquad (19)$$

where z_i are the energy fractions of electrons (positrons) carried by the exchanged photons and s denotes the c.m. energy squared of the e^+e^- system. The distribution of those energy fractions is given by the standard flux factor $f_{\gamma/e}(z,Q_{min},Q_{max})$ of the virtual photons. In the equivalent photon approximation it is given by the following formula [22]:

$$f_{\gamma/e}(z,Q_{min},Q_{max})=\frac{\alpha}{2\pi}\left[\frac{1+(1-z)^2}{z}\ln\frac{Q_{max}^2}{Q_{min}^2} -2m_e^2 z\left(\frac{1}{Q_{min}^2}-\frac{1}{Q_{max}^2}\right)\right]. \qquad (20)$$

Q_{min}^2 and Q_{max}^2 in Eq. (20) denote the minimal and maximal values of the photon virtuality and m_e is electron mass. For antitagged experiments the former is given by the kinematical limit $Q_{min}^2=m_e^2z^2/(1-z)$ and the latter by the antitagging condition $\theta_{e^\pm}<\theta_{max}$ which gives $Q_{max}^2=(1-z)E_{beam}^2\theta_{max}^2$ where $\theta_{e^\pm}$ denotes the scattering angle of the scattered $e^\pm$. Following Ref. [22] we set $\theta_{max}=30$ mrad. The cross section for the process $e^+e^-\to e^+e^-+Y$, which for antitagged $e^\pm$ corresponds to the production of the hadronic state Y in the collision of almost real (virtual) photons, is given by the following convolution integral:

$$\sigma_{e^+e^-\to e^+e^-+Y}=\int_0^1 dz_1\int_0^1 dz_2\Theta(W^2-W_{Y0}^2)\sigma_{\gamma\gamma\to Y}(W^2) \times f_{\gamma/e}(z_1,Q_{min},Q_{max}) \times f_{\gamma/e}(z_2,Q_{min},Q_{max}). \qquad (21)$$

In our calculation we set $W_{Y0}^2=M_{\eta_c}^2/z_{max}$ with $z_{max}=0.05$ so that the $\gamma\gamma$ system is in the high-energy (i.e., Regge) region. In order to estimate the effective enhancement factor due to $\lambda_{odd}>1$ we have compared the convolution integrals

$$I(s,\lambda_{odd})=\int_0^1 dz_1\int_0^1 dz_2\Theta(z_{max}-M^2_{\eta_c}/(z_1z_2s))\times\left(\frac{z_1z_2s}{M^2_{\eta_c}}\right)^{2(\lambda_{odd}-1)}\times f_{\gamma/e}(z_1,Q_{min},Q_{max})f_{\gamma/e}(z_2,Q_{min},Q_{max}) \tag{22}$$

for $\lambda_{odd}=1$ with that calculated for $\lambda_{odd}=1.07$.

The ratio $A=I(s,\lambda_{odd}=1.07)/I(s,\lambda_{odd}=1)$ should give the expected enhancement factor for the given value of s. We get $A=2.4$ and $A=2.9$ for the CERN e^+e^- collider LEP2 and Cornell TeV Energy Superconducting Linear Collider (TESLA) energies respectively. The relatively small change in A with increasing s is caused by the fact that the convolution integral is dominated by small values of z_i.

The magnitude of the estimated cross sections $\sigma(e^+e^-\to e^+e^-\eta_c\eta_c)$ and $\sigma(e^+e^-\to e^+e^-\eta_cX)$ for antitagged e^+ and e^- in the the final state are summarized in Table I. We give values of those cross sections for two different c.m. energies of incident leptons corresponding to LEP2 and TESLA energies and for two different values of the odderon intercept, $\lambda_{odd}=1$ and $\lambda_{odd}=1.07$. It may be seen from this table that the cross sections are very small and so it may in particular be difficult to measure them with presently available luminosity at LEP2 [22].

TABLE I. The estimated cross sections $\sigma(e^+e^-\to e^+e^-\eta_c\eta_c)$ and $\sigma(e^+e^-\to e^+e^-\eta_cX)$.

$\sqrt{s}$ [GeV]	$\lambda_{odd}-1$	$\sigma(e^+e^-\to e^+e^-\eta_c\eta_c)$ [fb]	$\sigma(e^+e^-\to e^+e^-\eta_cX)$ [fb]
180	0	0.7	1.9
180	0.07	1.6	4.5
500	0	2.1	5.8
500	0.07	6.1	17.0

To sum up we have applied the formalism developed in Ref. [20] to the quantitative analysis of the quasidiffractive processes $\gamma\gamma\to\eta_c\eta_c$ and $\gamma\gamma\to\eta_cX(q\overline{q})$ within the three-gluon-exchange mechanism. The main merit of those processes is that the corresponding cross sections can be, in principle, calculated within perturbative QCD. We have estimated the corresponding cross sections for the processes $e^+e^-\to e^+e^-\eta_c\eta_c$ and $e^+e^-\to e^+e^-\eta_cX$ with antitagged e^+e^- which were found to be within the range 0.7–17 fb depending upon the incident c.m. energy $\sqrt{s}$ and on the magnitude of the odderon intercept.

L.M. is grateful to J. Czyżewski for inspiring discussions, to the Foundation for Polish Science for financial support, and to the DESY Theory Division for hospitality. This research has been supported in part by the Polish Committee for Scientific Research grants 2 P03B 89 13 and 2 P03B 044 14.

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