Exclusive production of $\omega$ meson in proton-proton collisions at high energies

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Abstract

First we calculate cross section for the $\gamma p \rightarrow \omega p$ reaction from the threshold to very large energies. At low energies the pion exchange is the dominant mechanism. At large energies the experimental cross section can be well described within the $k_t$-factorization approach by adjusting light-quark constituent mass. Next we calculate differential distributions for the $pp \rightarrow pp\omega$ reaction at RHIC, Tevatron and LHC energies for the first time in the literature. We consider photon-pomeron (pomeron-photon), photon-pion (pion-photon) as well as diffractive hadronic bremsstrahlung mechanisms. The latter are included in the meson/reggeon exchange picture with parameters fixed from the known phenomenology. Interesting rapidity distributions are predicted. The hadronic bremsstrahlung contributions dominate at large (forward, backward) rapidities. At small energies the photon-pomeron contribution is negligible compared to the bremsstrahlung contributions. It could be, however, easily identified at large energies at midrapidities. Absorptions effects are included and discussed. Our predictions are ready for verification at RHIC and LHC.

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I. INTRODUCTION

The mechanism of exclusive production of mesons at high energies became recently a very active field of research (see [1] and references therein). The recent works concentrated on the production of $\chi_c$ mesons (see e.g. [2] and references therein) where the QCD mechanism is similar to the exclusive production of the Higgs boson. The latter process is an alternative to the inclusive production of the Higgs boson. In the case of heavy vector quarkonia ($J/\Psi$, $\Upsilon$) the dominant mechanism is photon-pomeron (pomeron-photon) fusion (see e.g. [3, 4]) which can be calculated in the QCD language.

The mechanism of exclusive light vector meson production was almost not studied in the literature, exception is $\phi$ meson [5]. In the present paper we consider exclusive production of the isoscalar $\omega$ meson. This process was studied before only close to its production threshold. Various theoretical models (see Refs. e.g. [6–10]) were developed to describe the experimental data (see Refs. [11]). Here the dominant mechanisms are meson exchange processes as well as the $\omega$-meson bremsstrahlung driven by meson exchanges.

How the situation changes at high-energy is interesting but was not studied so far. While at low energy the meson exchanges ($\pi, \rho, \omega, \sigma$) are the driving t-channel exchanges for the $\omega$ bremsstrahlung, at high energy their role is taken over by the pomeron exchange. The latter will be treated here purely phenomenologically. A similar hadronic bremsstrahlung-type mechanism is the Deck-mechanism for diffractive production of $\pi N$ final states in $pp$ collisions [12], for a review, see e.g. [13].

In the present paper we intend to make predictions for being in operation colliders RHIC, Tevatron and LHC. The hadronic bremsstrahlung mechanisms are expected to be enhanced for exclusive production of $\omega$ meson compared to other vector mesons as the $g_{\omega NN}$ coupling constant is known to be large from low-energy phenomenology [10, 14]. We will also show how important are the photoproduction mechanisms discussed previously in the context of exclusive heavy vector quarkonium production [3, 4]. The photoproduction mechanism constitutes a background for odderon-pomeron exchanges possible in the discussed reaction. So far odderon exchange was discussed for the exclusive $J/\Psi$ and $\Upsilon$ production [15]. The predicted by QCD odderon exchange was searched for in different reactions. No clear evidence was found so far. We shall comment on the issue in the Result Section.

II. PHOTOPRODUCTION MECHANISM FOR $\gamma p \rightarrow \omega p$

A. Pomeron exchange

Let us concentrate on the $\gamma p \rightarrow \omega p$ reaction which is a building block for the $pp \rightarrow pp\omega$ reaction. Photoproduction of the vector meson in photon-proton collisions is very interesting from both experimental and theoretical side. The corresponding cross sections have been measured by the ZEUS Collaboration at HERA at virtuality of photon $Q^2 \simeq 0$ GeV$^2$ for $\omega$ photoproduction [16] and at large values $Q^2$ for $\omega$ electroproduction $ep \rightarrow e\omega p$ [17]. The amplitude for this reaction is shown schematically in Fig. The Pomeron exchange is modelled by a pQCD gluon ladder. The details how to calculate the amplitude are explained
The imaginary part of the amplitude for the $\gamma p \rightarrow \omega p$ process is written as:

$$\Im m \mathcal{M}_{\lambda,\lambda_V}(W, t = -\Delta^2, Q^2) = W^2 \frac{C_V \alpha_{em}}{4\pi^2} \int \frac{d\kappa^2}{\kappa^4} \alpha_s(q^2) F(x_1, x_2, \kappa_1, \kappa_2) \left(2 \int_0^1 \frac{dz}{z(1-z)} I_{\lambda,\lambda_V}(z, k^2, \kappa_1, \kappa_2, Q^2)\right),$$

where the transverse momenta of gluons coupled to the $q \bar{q}$ pair can be written as $\kappa_1 = \kappa + \Delta^2/2$ and $\kappa_2 = -\kappa + \Delta^2/2$. $\Delta^2$ is the (transverse) momentum transfer squared and $k$ is the transverse momentum of the (anti-)quark. The quantity $F(x_1, x_2, \kappa_1, \kappa_2)$ is the off diagonal unintegrated gluon distribution. Explicit expressions for $I_{\lambda,\lambda_V}$ can be found in Ref. [18].

In the forward scattering limit, i.e. for $\Delta^2 = 0$, azimuthal integrations can be performed analytically. The following representation for the imaginary part of the amplitude for the transverse polarization for forward photoproduction $\gamma p \rightarrow \omega p$ process is used:

$$\Im m \mathcal{M}(W, \Delta^2 = 0, Q^2 = 0) = W^2 \frac{C_V \alpha_{em}}{4\pi^2} 2 \int_0^1 \frac{dz}{z(1-z)} \int_0^\infty \pi d\kappa^2 \psi_V(z, k^2)$$

$$\times \int_0^\infty \frac{\pi d\kappa^2}{\kappa^4} \alpha_s(q^2) F(x_{eff}, \kappa^2) \left(A_0(z, k^2) W_0(k^2, \kappa^2) + A_1(z, k^2) W_1(k^2, \kappa^2)\right),$$

where

$$A_0(z, k^2) = m_q^2 + \frac{k^2 m_q}{M + 2m_q},$$

$$A_1(z, k^2) = \left(z^2 + (1-z)^2 - (2z-1)^2 \frac{m_q}{M + 2m_q}\right) \frac{k^2}{k^2 + \epsilon^2},$$

$$W_0(k^2, \kappa^2) = \frac{1}{k^2 + \epsilon^2} - \frac{1}{\sqrt{(k^2 - \epsilon^2 - \kappa^2)^2 + 4\epsilon^2 k^2}},$$

$$W_1(k^2, \kappa^2) = 1 - \frac{k^2 + \epsilon^2}{2k^2} \left(1 + \frac{k^2 - \epsilon^2 - \kappa^2}{\sqrt{(k^2 - \epsilon^2 - \kappa^2)^2 + 4\epsilon^2 k^2}}\right),$$
and $M$ is the invariant mass of the constituent $q\bar{q}$ system

$$M = \frac{k^2 + m_q^2}{z(1-z)}, \quad (2.7)$$

where $z$ and $(1-z)$ are fractions of the longitudinal momentum of the $\omega$-meson carried by a quark and antiquark, respectively.

The diagonal unintegrated gluon distribution can be emulated by taking the ordinary gluon distribution at $F(x_{eff}, \kappa^2)$, where $x_{eff} = c_{skewed}(m^2_\omega/W^2)$ with $c_{skewed} = 0.41 \ [18]$. The forward unintegrated gluon distribution is taken from the work of Ivanov-Nikolaev \[19\], where it was found in the analysis of the deep-inelastic scattering data. The charge-isospin factor $c_V$ is

$$c_\omega = 1/\sqrt{2(e_u + e_d)} = 1/(3\sqrt{2}).$$

In our calculation we choose the scale of the QCD running coupling constant $\alpha_S$ at $q^2 = \max\{\kappa^2, k^2 + m_q^2\}$.

The full amplitude for the $\gamma p \to \omega p$ process is given as

$$M(W, \Delta^2, Q^2 = 0) = (i + \rho) \Im m M(W, \Delta^2 = 0, Q^2 = 0) \exp \left( -\frac{B(W)\Delta^2}{2} \right), \quad (2.8)$$

where $\rho$ is a ratio of real to imaginary part of the amplitude and $B(W)$ is the slope parameter dependent on the photon-proton center-of-mass energy and is parametrized as

$$B(W) = B_0 + 2\alpha'_{eff} \ln \left( \frac{W^2}{W_0^2} \right), \quad (2.9)$$

with: $W_0 = 95$ GeV, $B_0 = 11$ GeV$^{-2}$, $\alpha'_{eff} = 0.25$ GeV$^{-2}$ \[20\].

Our amplitude is normalized to the total cross section:

$$\sigma(\gamma p \to \omega p) = \frac{1 + \rho^2}{16\pi B(W)} \left| 3m \frac{M(W, \Delta^2 = 0, Q^2 = 0)}{W^2} \right|^2. \quad (2.10)$$

The radial light-cone wave function of the vector meson can be regarded as a function of three-momentum $p = (\vec{p}, p_z)$, where $\vec{p} = \vec{k}$, $p_z = (2z - 1)M/2$ then

$$\psi_V(z, \vec{k}^2) \to \psi_V(p^2), \quad \frac{dzd^2\vec{k}}{z(1-z)} \to \frac{4d^3p}{M}, \quad p^2 = \frac{M^2 - 4m_q^2}{4}. \quad (2.11)$$

In our calculation we use a Gaussian wave function, representing a standard harmonic-oscillator type quark model, which turned out to be superior over a Coulomb wave function (which a power-law tail in momentum space) for $J/\Psi, \ U$ and $\phi$ mesons exclusive photoproduction \[3\ [5\]

$$\psi_V(p^2) = N \exp \left( -\frac{p^2a_1^2}{2} \right). \quad (2.12)$$

The parameter $a_1$ is obtained by fitting to the decay electronic width

$$\Gamma(V \to e^+e^-) = \frac{4\pi\alpha^2 m_V^2 \cdot g_V}{3m^3_\omega}, \quad (2.13)$$

where $\Gamma(\omega \to e^+e^-) = 0.6$ keV \[21\] and imposing the normalization condition

$$1 = \frac{N A_\pi}{(2\pi)^3} \int_0^\infty p^2 dp 4M\psi^2_V(p^2). \quad (2.14)$$
In our calculation we use leading-order approximation, i.e. we neglect a possible NLO $K$-factor. The parameter $g_V$ can be expressed in terms of the $\omega$-meson wave function as $[18]$

$$g_V = \frac{8N_c}{3} \int \frac{d^3p}{(2\pi)^3} (M + m_q) \psi_V(p^2).$$  \hspace{1cm} (2.15)

FIG. 2: Total cross section for the photoproduction $\gamma p \rightarrow \rho^0 p$ (left panel) and $\gamma p \rightarrow \omega p$ (right panel) processes as a function of the photon-proton center-of-mass energy. In the calculation of the $\pi$-exchange mechanism the Gaussian wave function of the $\rho^0$ and $\omega$ mesons is used. At low energies $\pi$-exchange is the dominant mechanism. The curves are described in the text. Our results are compared with the HERA data [16, 22–25] (solid marks) and with a compilation of low energy data [26, 27] (open circles).

Having in view theoretical uncertainties in defining light quark mass it is treated here as a model parameter. In Fig.2 we show the total cross section for the exclusive $\gamma p \rightarrow \rho^0 p$ (left panel) and $\gamma p \rightarrow \omega p$ (right panel) processes as a function of the $\gamma p$ center-of-mass energy $W_{\gamma p}$ for the photon virtuality $Q^2 = 0$ GeV$^2$. Our results for exclusive $\rho^0$ and $\omega$ mesons production are compared with the corresponding experimental data. For the $\rho^0$ meson we present results for three different values of the $u$ and $d$ quark masses assumed here to be identical. The dashed line (bottom) is for $m_q = 0.33$ GeV, the dotted line (top) for $m_q = 0.22$ GeV and the thick solid line (fitted to experimental data) for $m_q = 0.3$ GeV. Because the results for $m_q = 0.3$ GeV give the best description of experimental data, this mass will be used in further calculations. In calculation the Gaussian wave function is used. We see that it gives quite good description of the high-energy $\omega$-meson data. At low energies the pion exchange mechanism dominates [28, 29]. This will be discussed in the following subsection.
FIG. 3: Diagram with the \( \pi \)-exchange for exclusive photoproduction \( \gamma p \rightarrow \omega p \).

### B. Pion exchange

The amplitude for the \( \pi \)-exchange can be written as:

\[
\mathcal{M}^{\pi_{\pi}^{\text{exch.}}} = g_{\omega\pi\gamma} F_{\omega\pi\gamma}(t) \varepsilon^\beta \mu \lambda k_\mu k'_\nu \varepsilon_\beta(k, \lambda, \gamma) \varepsilon^*(k', \lambda, \omega) \\
\times g_{\pi NN} F_{\pi NN}(t) \frac{1}{t - m_\pi^2} \bar{u}(p_{N'}, \lambda_{N'}) \gamma_5 u(p_N, \lambda_N).
\]

(2.16)

The \( g_{\omega\pi\gamma} \) coupling constant in the formula above is obtained from the \( \omega \) partial decay width through the relation:

\[
\Gamma(\omega \rightarrow \pi^0\gamma) = \mathcal{BR}(\omega \rightarrow \pi^0\gamma) \cdot \Gamma_{\text{tot}} = \frac{g_{\omega\pi\gamma}^2}{96\pi} \times m_\omega \left(1 - \frac{m_\pi^2}{m_\omega^2}\right)^3.
\]

(2.17)

Taking experimental partial decay width \( \Gamma(\omega \rightarrow \pi^0\gamma) \) from [21] we get \( g_{\omega\pi\gamma} \approx 0.7 \text{ GeV}^{-1} \) which is consistent with the values used in Refs. [29, 30]. The pion-nucleon coupling constant \( g_{\pi NN} \) is relatively well known [31]. In our calculations the coupling constant \( g_{\pi NN}^2/4\pi = 13.5 \).

We describe the low energy data shown in Fig. 2 (right panel) with \( \Lambda_{\text{mon}} \approx 0.7 \text{ GeV} \) for the monopole form factors by the dashed line

\[
F(t) = \frac{\Lambda_{\text{mon}}^2 - m_\pi^2}{\Lambda_{\text{mon}}^2 - t}
\]

(2.18)

or \( \Lambda_{\text{exp}} \approx 0.8 \text{ GeV} \) for the exponential form factors by the solid line

\[
F(t) = \exp \left( \frac{t - m_\pi^2}{\Lambda_{\text{exp}}^2} \right).
\]

(2.19)

The cut-off parameters obtained from the fit are significantly smaller than e.g. those used in the Bonn model [14]. Such soft form factors may be due to active coupling with the \( \pi N \) and \( \rho N \) channels not included explicitly both here nor in the literature. The pion exchange describes only angular distributions at forward angles. At larger angles there are other mechanisms as nucleon exchanges or \( s \)-channel nucleon resonances [29, 32]. A more refined analysis in the peak region would require description of new very precise CLAS Collaboration data [33] for full range angular distributions. Such an analysis would need to include also channel couplings discussed above.

The form factors found here will be used when discussing \( \gamma\pi^0 \) and \( \pi^0\gamma \) exchanges in the \( pp \rightarrow pp\omega \) reaction.

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1 Please note different normalization convention of the coupling constant in all the papers.
III. THE AMPLITUDES FOR THE \( pp \to pp\omega \) REACTION

A. \( \gamma P \) and \( P\gamma \) exchanges

\[ M(\vec{p}_1, \vec{p}_2) = \int \frac{d^2\vec{k}}{(2\pi)^2} S_{el}(\vec{k}) M^{(0)}(\vec{p}_1 - \vec{k}, \vec{p}_2 + \vec{k}) \]

\[ = M^{(0)}(\vec{p}_1, \vec{p}_2) - \delta M(\vec{p}_1, \vec{p}_2), \quad (3.1) \]

where

\[ S_{el}(\vec{k}) = (2\pi)^2 \delta^{(2)}(\vec{k}) - \frac{1}{2} T(\vec{k}), \quad T(\vec{k}) = \sigma_{tot}^{pp}(s) \exp\left(-\frac{1}{2} B_{el}\vec{k}^2\right). \quad (3.2) \]

Here \( \vec{p}_1 \) and \( \vec{p}_2 \) are the transverse momenta of outgoing protons (RHIC, LHC) or proton and antiproton (Tevatron). In practical evaluations we take \( B_{el} = 14 \text{ GeV}^{-2} \), \( \sigma_{tot}^{pp} = 52 \text{ mb} \) for the RHIC energy \( W = 200 \text{ GeV} \), \( B_{el} = 17 \text{ GeV}^{-2} \), \( \sigma_{tot}^{pp} = 76 \text{ mb} \) for the Tevatron energy \( W = 1.96 \text{ TeV} \) and \( B_{el} = 21 \text{ GeV}^{-2} \), \( \sigma_{tot}^{pp} = 100 \text{ mb} \) for the LHC energy \( W = 14 \text{ TeV} \).

The Born-amplitude (without absorptive correction) can be written in the form of a two-dimensional vector (corresponding to the two transverse (linear) polarizations of the final state vector meson) \[ \text{(3.3)} \] as

\[ M^{(0)}(\vec{p}_1, \vec{p}_2) = e_1 \frac{2}{z_1 t_1} \mathcal{F}_{\lambda_1\lambda_1}(\vec{p}_1, t_1) \mathcal{M}_{\gamma^*h_2\to Vh_2}(s_2, t_2, Q_1^2) \]

\[ + e_2 \frac{2}{z_2 t_2} \mathcal{F}_{\lambda_2\lambda_2}(\vec{p}_2, t_2) \mathcal{M}_{\gamma^*h_1\to Vh_1}(s_1, t_1, Q_2^2), \quad (3.3) \]

where \( \mathcal{M}_{\gamma^*h_2\to Vh_2}(s_2, t_2, Q_1^2) \) and \( \mathcal{M}_{\gamma^*h_1\to Vh_1}(s_1, t_1, Q_2^2) \) are the amplitudes for photoproduction discussed above (see \( (2.18) \)). Because of the presence of the Dirac electromagnetic form factor of the proton/antiproton only small \( Q_1^2 \) and \( Q_2^2 \) enter the amplitude for the hadronic process. This means that in practice one can put \( Q_1^2 = Q_2^2 = 0 \text{ GeV}^2 \) for the \( \gamma^*p \to Vp \) amplitudes. We used the assumption of s-channel helicity conservation in the \( \gamma \to \omega \) transition, \( \lambda_\gamma = \lambda_V \).
The absorptive correction for the amplitude have the form:

\[
\delta M(p_1, p_2) = \int \frac{d^2 \vec{k}}{(2\pi)^2} T(\vec{k}) M^{(0)}(\vec{p}_1 - \vec{k}, \vec{p}_2 + \vec{k}).
\]  

(3.4)

The differential cross section is expressed in terms of the amplitude \( M \) as

\[
d\sigma = \frac{1}{512\pi^4 s^2} |M|^2 dy dt_1 dt_2 d\phi.
\]  

(3.5)

where \( y \) is rapidity of the \( \omega \) meson, \( t_{1,2} \equiv -\vec{p}_{1,2}^2 \) and \( \phi \) is the azimuthal angle between transverse momenta \( \vec{p}_1 \) and \( \vec{p}_2. \)

B. \( \gamma\pi^0 \) and \( \pi^0\gamma \) exchanges

As shown in Fig.2 the QCD mechanism discussed in subsection II A does not describe the huge close-to-threshold enhancement of the cross section. This indicates a presence of another mechanisms of omega photoproduction. Neutral pion exchange is the best candidate which describes the low energy data as discussed in subsection II B. Therefore for the \( pp \rightarrow pp\omega \) reaction we should include also photon-pion and pion-photon exchanges. The underlying mechanisms are shown in Fig.5.

![Diagrams with the \( \gamma\pi^0 \) and \( \pi^0\gamma \) exchange amplitudes in the \( pp \rightarrow pp\omega \) reaction.](image)

The amplitudes for the two new processes can be easily written as:

\[
\mathcal{M}_{\gamma\pi^0 \text{-exch.}}^{\lambda_a \lambda_b \rightarrow \lambda_1 \lambda_2 \lambda_3} = e F_1(t_1) \bar{u}(p_1, \lambda_1)\gamma^\alpha u(p_a, \lambda_a) \\
\times \frac{-g_{\alpha\beta}}{t_1} g_{\omega\pi^0\gamma} F_{\gamma\pi \rightarrow \omega}(t_1, t_2) \varepsilon^{\beta\mu\nu\lambda} q_1 \mu p_3 \nu \varepsilon^\lambda_3(p_3, \lambda_3) \\
\times g_{\pi^0NN} F_{\pi NN}(t_2) \frac{1}{t_2 - m_\pi^2} \bar{u}(p_2, \lambda_2) i\gamma_5 u(p_b, \lambda_b),
\]  

(3.6)

\[
\mathcal{M}_{\pi^0\gamma \text{-exch.}}^{\lambda_a \lambda_b \rightarrow \lambda_1 \lambda_2 \lambda_3} = g_{\pi^0NN} F_{\pi NN}(t_1) \frac{1}{t_1 - m_\pi^2} \bar{u}(p_1, \lambda_1) i\gamma_5 u(p_a, \lambda_a) \\
\times \frac{-g_{\alpha\beta}}{t_2} g_{\omega\pi^0\gamma} F_{\gamma\pi \rightarrow \omega}(t_2, t_1) \varepsilon^{\beta\mu\nu\lambda} q_2 \mu p_3 \nu \varepsilon^\lambda_3(p_3, \lambda_3) \\
\times e F_1(t_2) \bar{u}(p_2, \lambda_2)\gamma^\alpha u(p_b, \lambda_b),
\]  

(3.7)

\(^2\) In the following for brevity we shall use notation \( t_{1,2} \) which means \( t_1 \) or \( t_2.\)
where $F_1(t_{1,2})$ are the Dirac electromagnetic form factors of participating protons. The $g_{\omega\pi^0\gamma}$ constant was obtained from the omega partial decay width as discussed in subsection II B.

The coupling of the pion to the nucleon is described by the pion-nucleon coupling constant $g_{\pi NN}$ and the corresponding form factor is taken in the exponential form:

$$F_{\pi NN}(t_{1,2}) = \exp\left(\frac{t_{1,2} - m_\pi^2}{\Lambda_{\pi NN}^2}\right).$$  \hspace{1cm} (3.8)

The central vertices involve off-shell particles. Here the $\gamma\pi^0\gamma$ and $\pi^0\gamma$ form factors are taken in the following factorized form:

$$F_{\gamma\pi^0\omega}(t_1, t_2) = \frac{m_\rho^2}{m_\rho^2 - t_1} \exp\left(\frac{t_2 - m_\pi^2}{\Lambda_{\omega\pi\gamma}^2}\right).$$  \hspace{1cm} (3.9)

The factor describing the virtual photon coupling is taken as in the vector dominance model.

In practical calculations we take: $\Lambda_{\pi NN} = 0.8$ GeV and $\Lambda_{\omega\pi\gamma} = 0.8$ GeV as found from the fit to the $\gamma p \rightarrow \omega p$ experimental data.

At high-energies often light-cone form factors are used instead of the $t_1$ or $t_2$ dependent ones discussed above (see Eq. (3.8)). In such an approach the pion is rather a constituent of the initial proton. Then the form factors are parametrized in terms of the squared invariant masses of the $\pi N$ system:

$$M_{2,\pi N}^2(z_2, p_{2t}) = \frac{m_N^2 + p_{2t}^2}{z_2} + \frac{m_\pi^2 + p_{2t}^2}{1 - z_2},$$

$$M_{1,\pi N}^2(z_1, p_{1t}) = \frac{m_N^2 + p_{1t}^2}{z_1} + \frac{m_\pi^2 + p_{1t}^2}{1 - z_1},$$  \hspace{1cm} (3.10)

where the longitudinal momentum fractions of outgoing protons with respect to the initial protons can be calculated from energies and $z$-components of momenta of participating protons

$$z_2 = (p_{20} - p_{2z})/(p_{00} - p_{0z}),$$

$$z_1 = (p_{10} + p_{1z})/(p_{00} + p_{0z}).$$  \hspace{1cm} (3.11)

The light-cone form factors are parametrized then as

$$F_{\pi NN}(M_{2,\pi N}^2) = \exp\left(-\frac{M_{2,\pi N}^2 - m_N^2}{2\Lambda_{LC}^2}\right),$$

$$F_{\pi NN}(M_{1,\pi N}^2) = \exp\left(-\frac{M_{1,\pi N}^2 - m_N^2}{2\Lambda_{LC}^2}\right).$$  \hspace{1cm} (3.12)

The parameter $\Lambda_{LC}$ in the light-cone parametrization was fitted in Ref. 35 to the data on forward nucleon production and the value $\Lambda_{LC} = 1.1$ GeV was found.

The amplitude for processes shown in Fig. 5 is calculated numerically for each point in the phase space. In calculating cross section we perform integration in $\log_{10}(p_{1t})$ (for $\gamma\pi$-exchange) and $\log_{10}(p_{2t})$ (for $\pi\gamma$-exchange) instead in $p_{1t}$ and $p_{2t}$.
IV. HADRONIC BREMSSTRAHLUNG MECHANISMS

A. The amplitude in the standard approach

The strong coupling of the $\omega$ meson to the nucleon causes that the hadronic bremsstrahlung mechanisms become important. The bremsstrahlung mechanisms for exclusive production of $\omega$ discussed here are shown schematically in Fig. 6. In the case of $\omega$ production the diagrams with intermediate nucleon resonances are negligible (see [21]). Because at high energy the pomeron is the driving mechanism of bremsstrahlung it is logical to call the mechanisms diffractive bremsstrahlung to distinguish from the low-energy bremsstrahlung driven by meson exchanges.

It is straightforward to evaluate the contribution of diagrams shown in Fig. 6. The Born amplitudes read:

\[
\mathcal{M}_{\lambda_1\lambda_2\lambda_3 \rightarrow \lambda_1\lambda_2\lambda_3}^{(a)} = u(p_1, \lambda_1)\varepsilon_{\mu}^{*}(p_3, \lambda_3)\gamma^{\mu}S_{N}(p_1^{*})u(p_2, \lambda_2) g_{\omega NN} F_{\omega NN^*}(p_1^{*2}) F_{\omega NN^*}(p_2^{*2}) \\
\times is_{ab}C_{F}^{NN}(s_{ab}/s_0)\alpha_{F}(t_2)^{-1} \exp\left(\frac{B_{F}^{NN}t_2}{2}\right) \delta_{\lambda_2\lambda_3},
\]

\[
\mathcal{M}_{\lambda_1\lambda_2\lambda_3 \rightarrow \lambda_1\lambda_2\lambda_3}^{(b)} = \bar{u}(p_2, \lambda_2)\varepsilon_{\mu}^{*}(p_3, \lambda_3)\gamma^{\mu}S_{N}(p_2^{*})u(p_1, \lambda_1) g_{\omega NN} F_{\omega NN^*}(p_2^{*2}) F_{\omega NN^*}(p_1^{*2}) \\
\times is_{ab}C_{F}^{NN}(s_{ab}/s_0)\alpha_{F}(t_1)^{-1} \exp\left(\frac{B_{F}^{NN}t_1}{2}\right) \delta_{\lambda_1\lambda_3},
\]

\[
\mathcal{M}_{\lambda_1\lambda_2\lambda_3 \rightarrow \lambda_1\lambda_2\lambda_3}^{(c)} = \bar{u}(p_1, \lambda_1)S_{N}(p_1^{*2})\varepsilon_{\mu}^{*}(p_3, \lambda_3)\gamma^{\mu}u(p_2, \lambda_2) g_{\omega NN} F_{\omega NN^*}(p_1^{*2}) F_{\omega NN^*}(p_2^{*2}) \\
\times is_{12}C_{F}^{NN}(s_{12}/s_0)\alpha_{F}(t_2)^{-1} \left(\begin{array}{c}s_{13} \\ s_{1h}\end{array}\right)\alpha_{N}(p_1^{*2})^{-\frac{1}{2}} \exp\left(\frac{B_{F}^{NN}t_2}{2}\right) \delta_{\lambda_2\lambda_3},
\]

\[
\mathcal{M}_{\lambda_1\lambda_2\lambda_3 \rightarrow \lambda_1\lambda_2\lambda_3}^{(d)} = \bar{u}(p_2, \lambda_2)S_{N}(p_2^{*2})\varepsilon_{\mu}^{*}(p_3, \lambda_3)\gamma^{\mu}u(p_1, \lambda_1) g_{\omega NN} F_{\omega NN^*}(p_2^{*2}) F_{\omega NN^*}(p_1^{*2}) \\
\times is_{12}C_{F}^{NN}(s_{12}/s_0)\alpha_{F}(t_1)^{-1} \left(\begin{array}{c}s_{23} \\ s_{2h}\end{array}\right)\alpha_{N}(p_2^{*2})^{-\frac{1}{2}} \exp\left(\frac{B_{F}^{NN}t_1}{2}\right) \delta_{\lambda_1\lambda_3}.
\]
The diagrams for the interaction with emitted $\omega$:

$$\mathcal{M}^{(e)}_{\lambda_a\lambda_b \rightarrow i_1, i_2, i_3} = \bar{u}(p_1, \lambda_1)\gamma^\mu u(p_0, \lambda_0)S_{\mu\nu}(t_1)\varepsilon^{\nu*}(p_3, \lambda_3) g_{\omega NN} F_{\omega*NN}(t_1) F_{\omega*\omega}(t_1)$$

$$\times is_{23} C_{\omega N}^N \left( \frac{s_{23}}{s_0} \right) \frac{\alpha_F(t_2) - 1}{\alpha_F(t_1) - 1} \exp \left( \frac{B_{\omega N}^N t_2}{2} \right) \delta_{i_2 i_3}, \quad (4.5)$$

$$\mathcal{M}^{(f)}_{\lambda_a\lambda_b \rightarrow i_1, i_2, i_3} = \bar{u}(p_2, \lambda_2)\gamma^\mu u(p_0, \lambda_0)S_{\mu\nu}(t_2)\varepsilon^{\nu*}(p_3, \lambda_3) g_{\omega NN} F_{\omega*NN}(t_2) F_{\omega*\omega}(t_2)$$

$$\times is_{13} C_{\omega N}^N \left( \frac{s_{13}}{s_0} \right) \frac{\alpha_F(t_2) - 1}{\alpha_F(t_1) - 1} \exp \left( \frac{B_{\omega N}^N t_1}{2} \right) \delta_{i_1 i_3}, \quad (4.6)$$

where $s_0 = 1 \text{ GeV}^2$ and $s_{th} = (m_N + m_\omega)^2$.

In the above equations $u(p_i, \lambda_i)$, $\bar{u}(p_f, \lambda_f)$ are the Dirac spinors (normalized as $\bar{u}(p)u(p) = 2m_N$) of the initial and outgoing protons with the four-momentum $p$ and the helicities $\lambda_i$. The propagators of nucleons and $\omega$ meson can be written as

$$S_N(p_{i,2f}^2) = \frac{i(p_{1f,2f}^* \gamma^\mu + m_N)}{p_{1f,2f}^* - m_N^2},$$

$$S_N(p_{i,2i}^2) = \frac{i(p_{1i,2i}^* \gamma^\mu + m_N)}{p_{1i,2i}^* - m_N^2},$$

$$S_{\mu\nu}(t) = \frac{-g_{\mu\nu} + g_{\mu q} g_{\nu q} m_q^2}{t - m_q^2}, \quad (4.7)$$

where $t_{1,2} = (p_{a,b} - p_{1,2})^2 = q_{1,2}^2$, $p_{i,2f}^2 = (p_{a,b} - p_3)^2$, $p_{1f,2f}^2 = (p_{1,2} + p_3)^2$ are the four-momenta squared of objects in the middle of diagrams and $s_{ij} = (p_i + p_j)^2$ are squared invariant masses of the $(i, j)$ system.

The factor $g_{\omega NN}$ is the omega nucleons coupling constant. Different values have been used in the literature $[14]$. In our calculations the coupling constant is taken as $g_{\omega NN}^2/4\pi = 10$. Similar value was used in Refs. $[3, 10]$.

Using the known strength parameters for the $NN$ and $\pi N$ scattering fitted to the corresponding total cross sections (the Donnachie-Landshoff model $[36]$) we obtain $C_{NN}^{NN} = 21.7$ mb and $C_{\omega N}^{NN} = C_{\pi N}^{NN} = 13.63$ mb. The pomeron/reggeon trajectory determined from elastic and total cross sections is taken in the linear approximation in $t$ ($\alpha(t) = \alpha(0) + \alpha' t$)

$$\alpha_P(t) = 1.0808 + 0.25 t, \quad \alpha_\omega(t) = 0.5 + 0.9 t, \quad (4.8)$$

where the values of the intercept $\alpha(0)$ and the slope of the trajectory $\alpha'$ are also taken from the Donnachie-Landshoff model $[36]$ for consistency. The slope parameter can be written as

$$B(s) = B_0 + 2\alpha' \ln \left( \frac{s}{s_0} \right). \quad (4.9)$$

In our calculation we use $B_0$: $B_{\omega N}^{NN} = 5.5 \text{ GeV}^{-2}$ and $B_{NN}^{NN} = 9 \text{ GeV}^{-2}$.

The extra factors $F_{\omega NN}$ and $F_{NN}$ (or $F_{\omega N}$) allow for modification when one of the nucleons or the $\omega$-meson is off its mass shell. We parametrize all the form factors in the
following exponential form:

\[
F_{\omega NN}(p_{1f,2f}^2) = \exp \left( -\frac{(p_{1f,2f}^2 - m_N^2)}{\Lambda^2} \right), \quad F_{PNN}(p_{1f,2f}^2) = \exp \left( -\frac{(p_{1f,2f}^2 - m_N^2)}{\Lambda_{PNN}^2} \right),
\]

\[
F_{\omega NN}(p_{i1,2i}^2) = \exp \left( \frac{p_{i1,2i}^2 - m_N^2}{\Lambda^2} \right), \quad F_{PNN}(p_{i1,2i}^2) = \exp \left( \frac{p_{i1,2i}^2 - m_N^2}{\Lambda_{PNN}^2} \right),
\]

\[
F_{\omega NN}(t_{1,2}) = \exp \left( \frac{(t_{1,2} - m_\omega^2)}{\Lambda^2} \right), \quad F_{P\omega\omega}(t_{1,2}) = \exp \left( \frac{(t_{1,2} - m_\omega^2)}{\Lambda_{P\omega\omega}^2} \right). \tag{4.10}
\]

In general, the cut-off parameters are not known but could be fitted to the (normalized) experimental data. From our general experience in hadronic physics we expect \( \Lambda \approx \Lambda_{PNN} \approx \Lambda_{P\omega\omega} = 1 \text{ GeV} \). We shall discuss how the uncertainties of the form factors influence our final results.

Since the amplitudes given by formulas \([4.5, 4.6]\) are as if for \( \omega \) meson exchanges they are corrected by the factors \( \alpha_{3}(t_{1,2})^{-1} \) to reproduce the high-energy Regge dependence. We improve also the parametrization of the amplitudes \([4.3, 4.4]\) by the factors \( \alpha_{3}(p_{i1,2i})^{-1/2} \), where the degenerate nucleon trajectory is \( \alpha_{N}(p_{i1,2i}^2) = -0.3 + \alpha'_N p_{i1,2i}^2 \), with \( \alpha'_N = 0.9 \text{ GeV}^{-2} \).

We have chosen a representation for the polarization vectors of the \( \omega \)-meson in the helicity states \( \lambda_3 = 0, \pm 1 \). The polarization vectors are parametrized, in a frame where \( p = (E_3, p_3 \cos \phi \sin \theta, p_3 \sin \phi \sin \theta, p_3 \cos \theta) \), as

\[
\varepsilon(p_3, 0) = \frac{E_3}{m_\omega} \left( \frac{p_3}{E_3}, \cos \phi \sin \theta, \sin \phi \sin \theta, \cos \theta \right),
\]

\[
\varepsilon(p_3, \pm 1) = \frac{1}{\sqrt{2}} \left( 0, i \sin \phi \mp \cos \theta \cos \phi, -i \cos \phi \mp \cos \theta \sin \phi, \pm \sin \theta \right). \tag{4.11}
\]

It is easy to check that they fulfill the relation \( p^\mu \varepsilon_\mu(p, \lambda) = 0 \).

**B. \( \omega \)-production as a diffractive excitation of the \( \omega p \)-Fock state**

The exclusive production of \( \omega \)-mesons in the fragmentation region of either proton can also be understood as a diffractive excitation of a two-body \( \omega p \)-Fock state of the physical proton. This is best formalized by a Fock-state decomposition of the protons light-cone wave function in terms of meson-baryon Fock states. A comprehensive treatment of meson-cloud effects with applications to deep-inelastic scattering and baryon form factors within this framework has been developed in \([35, 38]\), for a review and references see \([39]\). For the problem at hand, we can write schematically

\[
|p\rangle_{\text{phys}} = \sqrt{Z} \left( |p\rangle_{\text{bare}} + \int dz d^2k_\perp \Psi_{\omega p}(z, k_\perp) |p(1 - z, -k_\perp); \omega(z, k_\perp) + \ldots \right). \tag{4.12}
\]

Here, the bare proton state represents, for example, a three-quark core of the physical proton, \( \Psi_{\omega p} \) is the light-cone wave function of the \( \omega p \)-Fock-state. The \( \omega \)-meson in the two-body Fock-state carries a fraction \( z \) of light-cone plus-momentum of the physical proton and
transverse momentum $\vec{k}_\perp$; for simplicity helicity labels are suppressed. The invariant mass of the virtual $\omega p$ system is then given as

$$M^2_{\omega p} = \frac{\vec{k}_\perp^2 + m_\omega^2}{z} + \frac{\vec{k}_\perp^2 + m_N^2}{(1 - z)},$$

(4.13)

and enters the radial part of the wave function in terms of the $\omega NN$-form factor

$$F_{\omega NN}(M^2_{\omega p}) = \exp\left(-\frac{M^2_{\omega p} - m_N^2}{2\Lambda_{LC}^2}\right).$$

(4.14)

The parameter $\Lambda_{LC}^2$ which controls the momentum distribution of $\omega$-mesons in the Fock-state is taken as $\Lambda_{LC} = 1.1$ GeV [35].

In accordance with the classic Good-Walker formalism [40], diffractive excitation of the $\omega p$-state now occurs because interactions of the bare proton and the two-body $\omega N$-state differ. We can write the $\omega p$ scattering state as:

$$|\omega p\rangle_{\text{scatt}} = \left(\hat{S}_{\omega p} - \hat{S}_{\omega} \hat{S}_p\right)|\omega p\rangle,$$

(4.15)

where $\hat{S}_{\omega p}$ and $\hat{S}_p$ are the elastic scattering matrices for the $\omega p$ and $p$ interactions with the target. Assuming, that the $S$-matrix of the two-body state factorizes, $\hat{S}_{\omega p} = \hat{S}_\omega \hat{S}_p$, one can show that Eq.(4.15) generates precisely the diagrams a), c), e) of Fig.6. Diagrams b), d), f) can be obtained by an obvious symmetrization. In the practical evaluation, these diagrams will give similar expressions in momentum space as the ones obtained in the reggeized field theory model (the “standard approach” discussed above), modulo the absence of Regge-factors and the careful replacement of all $\omega NN$-form factors by their light-cone counterparts given in Eq.(4.14).

Notice that this description of diffractive dissociation, which treats the $\omega$-meson as a non-perturbative parton of the proton has a good physical motivation only in the fragmentation region of the proton(s). When the $\omega$-meson is produced in the central rapidity domain, the reggeization of the crossed channel exchanges must be taken into account. For Reggeon exchanges however the light-cone wave function formalism described above is ill defined [41]. Therefore, for a description of midrapidity $\omega$ production, one would have to add the reggeized $\omega$ exchange. We do not do this here, as the final result would not differ much from the reggeized field theory diagrams (the “standard approach”). At rapidities close to the proton fragmentation region the difference between the “standard approach” and the light-cone wave function treatment can serve as an indicator for the model dependence of our predictions for this particular soft process.

Finally let us note, that at the high energies of interest the deviation from factorization

$$\delta \hat{S} = \hat{S}_{\omega p} - \hat{S}_\omega \hat{S}_p,$$

(4.16)

is quantified by the shadowing or absorption correction to which we now turn.

C. Absorption effects

The absorption effect for the hadronic bremsstrahlung contributions requires a short comment. Since in practice for the pomeron exchanges in diagrams a) - d) we use phenomenological interactions which effectively describe the total and elastic data an additional use of
absorption would be a double counting. This is not the case for diagrams e) and f) where the interaction is between $\omega$-meson and proton. Consequently in the latter case we include absorption effect in full analogy to that described in section about photoproduction. This is illustrated in Fig. 7.

V. RESULTS

In the present section we present differential distributions for three different energies: $W = 200$ GeV (RHIC), $W = 1960$ GeV (Tevatron) and $W = 14$ TeV (LHC). This includes rapidity and transverse momentum of $\omega$ meson distributions as well as azimuthal correlations between outgoing protons.

In Fig. 8 we present differential cross sections $d\sigma/dW_{13}$ for the $pp \rightarrow pp\omega$ reaction at $W = 14$ TeV for the hadronic bremsstrahlung mechanisms. The left panel is for results with Mandelstam variable dependents $\omega NN$ form factors and with reggeization included while the light-cone approach correspond to the right panel. The thick solid line presents the result for the coherent sum of all amplitudes shown in Fig. 6.

In Fig. 8 we present differential cross sections $d\sigma/dW_{13}$ for the $pp \rightarrow pp\omega$ reaction at $W = 14$ TeV. We show results with Mandelstam variable dependent form factors (left panel), which we will call standard in the following, and with light-cone form factors (right panel). In the left panel we show results for the standard spin-$1/2$ propagators in diagrams a) and c) as well as with reggeization $[42]$. The long dashed, dashed and dotted lines correspond to contributions from diagrams a), c) and e), respectively. The thick solid line presents the coherent sum of all amplitudes. The light-cone form factors lead to much steeper dependence.
of the cross section on $W_{13}$ ($W_{23}$) than the standard form factors. The reggezation leads to an extra damping of the large $W_{13}$ ($W_{23}$) cross section.

In Fig. 9 we present the role of the form factors and reggezation for differential distributions in the $\omega$ meson rapidity and transverse momentum as well as for azimuthal angle correlation between outgoing protons. The distribution in rapidity is closely related to that for $W_{13}$ ($W_{23}$). As seen from the middle panels the reggezation makes the distribution steeper in the $\omega$ meson transverse momentum.

In Fig. 10 we present rapidity distribution of the $\omega$ meson in the two approaches for different energies. In the first approach we use the standard $\omega NN$ form factors (upper panels) and in the second approach we use the light-cone form factors (bottom panels) for the omega-nucleon-nucleon coupling. The distributions for the standard form factors extend more towards midrapidities. We show the $\gamma IP$ ($IP\gamma$), $\gamma\pi^0$ ($\pi^0\gamma$) as well as diffractive bremsstrahlung mechanisms. At “low” energy (RHIC) the discussed hadronic bremsstrahlung mechanisms dominate over the $\gamma IP$ and $IP\gamma$ ones. The cross section for the hadronic bremsstrahlung contribution is two-orders of magnitude bigger than that for the ($\gamma IP$, $IP\gamma$) contribution. The latter mechanism is known to be the dominant one for $J/\Psi$ and $\Upsilon$ meson production [3, 4]. A recent analysis at the Tevatron seems to confirm this claim [37]. Increasing the center-of-mass energy the hadronic bremsstrahlung components move to large rapidities. The $\gamma\pi^0$ (left peak) and the $\pi^0\gamma$ (right peak) components are separated. The separation
rapidity means also lack of interference effects which is very different compared to the $\gamma I P$ ($I P \gamma$) mechanism.\footnote{The interference between the two mechanisms $\gamma I P$ and $I P \gamma$ is proportional to $e_1 e_2 (\vec{p}_1 \cdot \vec{p}_2)$ and introduces a charge asymmetry as well as an angular correlations between the outgoing protons.}

At LHC energy at midrapidities the photoproduction mechanisms with $I P$ exchange dominate over the hadronic bremsstrahlung ones. We predict a narrow plateau around $y \approx 0$ and a significant increase when going to large $|y|$. Experimental observation of the increase would confirm the bremsstrahlung mechanisms discussed here. Only at the highest LHC energy the region of very small rapidities is free of the hadronic bremsstrahlung contributions.

The difference between the results with standard and light-cone form factors illustrates theoretical uncertainties. While the hadronic bremsstrahlung contributions are subjected to rather large theoretical uncertainties (see discussion above), the $\gamma I P$ ($I P \gamma$) contributions are fairly precisely estimated. Deviations from the pQCD contribution at midrapidities may be caused by either the difficult to predict hadronic bremsstrahlung contributions or by the
very interesting pomeron-odderon contributions. The rise of the cross section with increasing $|y|$ would be a clear signal of the hadronic bremsstrahlung contributions, while a sizeable deviation of the cross section normalization a potential signal of the odderon exchange.

![Graphs showing differential cross sections](image)

FIG. 11: Differential cross section $d\sigma/dp_t$ for the $pp(\bar{p}) \rightarrow pp(\bar{p})\omega$ reaction at $W = 200, 1960, 14000$ GeV in the full rapidity range. Here the reggeized propagators of omega and nucleons are used. The dashed lines present the contribution without absorption, while the thick solid lines include the absorption.

In Fig.11 we show the distribution in the $\omega$ meson transverse momentum. In this case the integration is done over full range of meson rapidities. The thin lines are for the Born level calculations while the thick lines include effect of absorption. The hadronic bremsstrahlung contributions calculated with the light-cone form factors are steeper than those for the standard form factors. The distribution of the photon-pomeron contribution for the $pp$ scattering is somewhat different than that for the $pp$ scattering. This is caused by different signs of the interference terms (different combination of electric charges). The distribution of the $\gamma\pi^0 (\pi^0\gamma)$ contribution (green dash-dotted line) is very similar to that of the $\gamma IP (IP\gamma)$ contribution (blue lines).

![Graphs showing differential cross sections](image)

FIG. 12: Differential cross section $d\sigma/d\phi_{12}$ for the $pp(\bar{p}) \rightarrow pp(\bar{p})\omega$ reaction at $W = 200, 1960, 14000$ GeV in the full rapidity range. Here the reggeized propagators of omega and nucleons are used. The dashed lines present the contribution without absorption, while the thick solid lines include the absorption.

Whether the $\gamma\pi^0$ mechanism can be identified requires further studies. What are other specific features of this mechanism?
In Fig.12 we show distribution in relative azimuthal angle between outgoing protons. For the $\gamma\pi^0$ mechanism the maximum occurs at $\phi_{12} \approx \pi/2$ which is dictated by a specific tensorial coupling $\gamma\pi^0 \rightarrow \omega$. The azimuthal distribution for the $\gamma\pi^0$ mechanism is very different than for the hadronic bremsstrahlung contributions which peak at $\phi_{12} = \pi$, especially for the light-cone form factors. In principle, the azimuthal angle correlations could be used therefore to separate the different mechanisms. One can clearly see that the absorption effects lead to extra decorrelation in azimuth compared to the Born-level result. In Fig.12 we show rapidity-integrated results. In general the azimuthal angle correlations are rapidity dependent. Quite different distributions for the $\gamma\Pi^0 (\Pi\gamma)$ contribution have been predicted for the Tevatron and RHIC or LHC. The correlation function is for this mechanism caused totally be the interference of the $\gamma\Pi^0$ and $\Pi\gamma$ contributions (see [3]).

FIG. 13: Differential cross section $d\sigma/dp_t$ for the $pp(\bar{p}) \rightarrow pp(\bar{p})\omega$ reaction at $W = 200, 1960, 14000$ GeV for the limited rapidity range $-1 < y_\omega < 1$. Here the reggeized propagators of omega and standard form factors are used. The dashed lines present the contribution without absorption, while the thick solid lines include the absorption.

FIG. 14: Differential cross section $d\sigma/d\phi_{12}$ for the $pp(\bar{p}) \rightarrow pp(\bar{p})\omega$ reaction at $W = 200, 1960, 14000$ GeV for the limited rapidity range $-1 < y_\omega < 1$. Here the reggeized propagators of omega and standard form factors are used. The dashed lines present the contribution without absorption, while the thick solid lines include the absorption.

The distributions in the full (pseudo)rapidity range are rather theoretical and may be difficult to measure. One may expect that in practice only limited range of (pseudo)rapidity
around \( y_\omega = 0 \) will be available experimentally. Therefore, as an example, we have made an extra calculation for a limited rapidity range. In Fig.13 we show transverse momentum distributions for \(-1 < y_\omega < 1\). Here, as can be seen from Fig.10, it is enough to include only the hadronic bremsstrahlung diagrams e) and f). In this case standard form factors are used only. Please note (see Fig.10) that in the case of light-cone form factors the hadronic bremsstrahlung mechanism does not contribute to the restricted rapidity region. For comparison we show the contributions of photoproduction mechanisms which are calculated fairly precisely as discussed before. This is very useful in the context of the searches for odderon.

Finally in Fig.14 we show angular correlations between outgoing protons for \(-1 < y_\omega < 1\). In the case of light-cone form factors only the photoproduction mechanism contributes. Testing such distributions together with rapidity distributions could provide therefore new information on the mysterious odderon exchange.

VI. CONCLUSIONS

In this paper we have calculated the cross section for \( \gamma p \rightarrow \omega p \) reaction at high-energy within a QCD-inspired model. A good description of the HERA experimental data has been achieved, comparable as for the \( J/\Psi \) and \( \phi \) mesons in our previous works. In the present paper the Gaussian wave function was used with parameters adjusted to reproduce the electronic decay width of \( \omega \) meson.

This model is used then to predict the cross sections for \( pp \rightarrow pp\omega \) and \( p\bar{p} \rightarrow p\bar{p}\omega \) reactions at high-energies for the first time in the literature. In contrast to the \( J/\Psi \) and \( \phi \) exclusive production, in the case of the \( \omega \) meson different hadronic bremsstrahlung processes are possible due to large nonperturbative coupling of the \( \omega \) meson to the nucleon. At high energy there is a class of diffractive bremsstrahlung processes never considered in the literature.

At low energies the hadronic bremsstrahlung contributions dominate over the photoproduction ones if the standard Mandelstam-dependent form factors are used. With increasing energy the hadronic bremsstrahlung contributions move in rapidity to the fragmentation regions. At high energies the photoproduction mechanisms dominate at midrapidities. We predict a short plateau at midrapidities due to the photoproduction mechanism and a significant increase towards fragmentation regions (large \( |y_\omega| \)) due to the \( \omega \) bremsstrahlung. The identification of the increase would be a confirmation of the hadronic bremsstrahlung effects discussed here. However, this may be not simple experimentally. The precisely evaluated photoproduction mechanism constitutes a background for the odderon exchange searches.

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