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Azimuthal angular distributions in EDDE as a spin-parity analyser and glueball filter for the LHC.

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Abstract: Exclusive Double Diffractive Events (EDDE) are analysed as the source of the information about the central system. The experimental possibilities for the exotic particles searches are considered. From the reggeized tensor current picture some azimuthal angle dependences were obtained to fit the data from the WA102 experiment and to make predictions for the LHC collider.

Keywords: Phenomenological Models, Spin and Polarization Effects, Hadronic Colliders.

Contents

1. Introduction

Since the early time the process of the exclusive production of central systems of particles with quasi-diffractively scattered initial particles was considered as an important source of the information about the high-energy dynamics of strong interactions both in theory and experiment.

If one takes only one particle produced, this is the first "genuinely" inelastic process which not only retains a lot of features of the elastic scattering but also shows clearly how the initial energy is being transformed into the secondary particles. General properties of such amplitudes were considered in reference [1].

The theoretical consideration of these processes on the basis of Regge theory goes back to [2]. Some new interest was related to the possibly good signals of centrally produced Higgs bosons and heavy quarkonia [3].

As Pomerons are the driving force of the processes in question it is natural to expect that the glueball production will be favored, if one knows that Pomerons are mostly gluonic objects [4]. The central glueball production was suggested as the possible origin of the total cross section rise [5]. One of the early proposal for experimental investigations of centrally produced glueballs in EDDE was made in [6].

As to the most recent experimental results one has to mention the experiment WA102 [7]–[9].

It was proposed [10] that the process of single-particle diffractive production can be used as a filter to separate $q\bar{q}$-states from glueballs due to a special dependence on the azimuthal angle.

In this paper we study the process of single particle production in double diffractive events in the framework of Regge picture (based on Lorentz tensor reggeized exchanges) taking into account the absorption effects both in the initial and the final states. With parameters fixed from the fit to the WA102 data we give some predictions for the production of various J^{PC} states at the LHC. The main conclusion from WA102 that $q\bar{q}$-states and glueballs have distinct ϕ-dependence (taking into account their possible mixing) remains true, in our approach, for the LHC energy, though the functional form of the ϕ-distributions changes due to significant absorption effects.

2. EDDE kinematics and cross-sections

Here we consider the process $p+p \to p+X+p$, where X is a particle or a system of particles, spin and parity of which are fixed. In the kinematics which corresponds to the double Regge limit (see figure 1) the light-cone representation $(+,-;\perp)$ for momenta of colliding and scattered particles is the following:

$$\begin{aligned}
p_1 &= \left(\sqrt{\frac{s}{2}}, \frac{m^2}{\sqrt{2s}}, \mathbf{0}\right) \\
p_2 &= \left(\frac{m^2}{\sqrt{2s}}, \sqrt{\frac{s}{2}}, \mathbf{0}\right) \\
p_1' &= \left((1-\xi_1)\sqrt{\frac{s}{2}}, \frac{\mathbf{\Delta}_1^2+m^2}{(1-\xi_1)\sqrt{2s}}, -\mathbf{\Delta}_1\right) \\
p_2' &= \left(\frac{\mathbf{\Delta}_2^2+m^2}{(1-\xi_2)\sqrt{2s}}, (1-\xi_2)\sqrt{\frac{s}{2}}, -\mathbf{\Delta}_2\right). \qquad (2.1)
\end{aligned}$$

$\xi_{1,2}$ are the fractions of the momenta of protons carried by reggeons. Characters printed in bold are used for two-dimensional vectors. From the above notations we can obtain the relations:

$$\begin{aligned}
t_{1,2} &= \Delta_{1,2}^2 \simeq -\frac{\mathbf{\Delta}_{1,2}^2(1+\xi_{1,2})+\xi_{1,2}^2 m^2}{1-\xi_{1,2}} \simeq -\mathbf{\Delta}_{1,2}^2\,, \qquad \xi_{1,2}\to 0 \\
\cos\phi_0 &= \frac{\mathbf{\Delta}_1\mathbf{\Delta}_2}{|\mathbf{\Delta}_1||\mathbf{\Delta}_2|} \\
M_\perp^2 &= \xi_1\xi_2 s \simeq M_X^2+|t_1|+|t_2|+2\sqrt{t_1t_2}\cos\phi_0\,, \qquad 0\le\phi_0\le\pi \\
(p_1+\Delta_2)^2 &= s_1 \simeq \xi_2 s \\
(p_2+\Delta_1)^2 &= s_2 \simeq \xi_1 s \\
s_1s_2 &= sM_\perp^2 \qquad (2.2)
\end{aligned}$$

The physical region of double diffractive events is defined by the following kinematical cuts:

$$0.01\,\mathrm{GeV}^2 \le |t_{1,2}| \le 1\ \mathrm{GeV}^2\,, \qquad (2.3)$$

$$\xi_{\min} \simeq \frac{M_X^2}{s\xi_{\max}} \le\ \xi_{1,2}\ le\xi_{\max} = 0.1\,. \qquad (2.4)$$

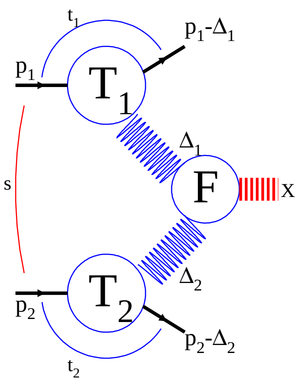

Figure 1: The process $p + p \to p + X + p$. The absorption in the initial and final pp-channels is not shown.

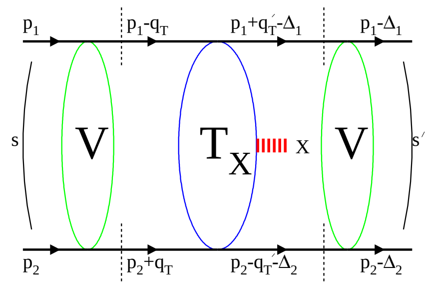

Figure 2: The full unitarization of the process $p + p \to p + X + p$.

The discussion on the choice of cuts (2.3), (2.4) for diffractive events and references to other authors were given in [11, 12]. We can write the relations in terms of $y_{1,2}$ and y_X (respectively rapidities of the hadrons and of the system X). For instance:

$$\begin{aligned} \xi_{1,2} &\simeq \frac{M_X}{\sqrt{s}} e^{\pm y_X} \,, \\ |y_X| \leq y_0 &= \ln\left(\frac{\sqrt{s}\xi_{\max}}{M_X}\right) , \\ |y_{1,2}| &= \frac{1}{2}\ln\frac{(1-\xi_{1,2})^2 s}{m^2 - t_{1,2}} \geq 9 \,. \end{aligned} \tag{2.5}$$

The cross-section of this process can be written as

$$\frac{d\sigma}{dt_1 dt_2 d\phi_0 dy_X} \simeq \frac{\pi |T^{\mathrm{Unit.}}_{pp\to pXp}|^2}{8s^2(2\pi)^5} \,, \tag{2.6}$$

where $T^{\mathrm{Unit.}}_{pp\to pXp}$ is the amplitude of the process, which can be calculated from the "bare" amplitude of figure 1 by the unitarization procedure depicted in figure 2, where

$$\begin{aligned} T_X &= T_{pp\to pXp} \,, \\ V(s, \mathbf{q}_T) &= 4s(2\pi)^2 \delta^2(\mathbf{q}_T) + 4s \int d^2\mathbf{b} e^{i\mathbf{q}_T \mathbf{b}} \left[e^{i\delta_{pp\to pp}} - 1\right] , \\ T_X^{\mathrm{Unit.}}(p_1, p_2, \Delta_1, \Delta_2) &= \frac{1}{16ss'} \int \frac{d^2\mathbf{q}_T}{(2\pi)^2} \frac{d^2\mathbf{q}'_T}{(2\pi)^2} V(s, \mathbf{q}_T) T_X(p_1 - q_T, p_2 + q_T, \Delta_{1T}, \Delta_{2T}) \cdot \\ &\quad \cdot V(s' \,,\, \mathbf{q}'_T) \,, \\ \Delta_{1T} &= \Delta_1 - q_T - q'_T \,, \\ \Delta_{2T} &= \Delta_2 + q_T + q'_T \,, \end{aligned} \tag{2.7}$$

and $\delta_{pp\to pp}$ can be found in [13]. V represents "soft" rescattering effects for the initial and final states, i.e. multi-Pomeron exchanges. These "outer" unitarity corrections can reduce the integrated cross-section. It depends on the kinematical cuts and the nature of the produced system X.

3. Calculation of the "bare" amplitude

In order to calculate peculiar momentum transfers and the azimuthal angle dependence we use the amplitudes with meson exchanges of arbitrary spins with subsequent reggeization.

Basic elements of such approach are the vertex functions

$$T^{\mu_1\cdots\mu_J}(p,\Delta) = \langle p-\Delta|I^{\mu_1\cdots\mu_J}|p\rangle \tag{3.1}$$

and

$$F^{\mu_1\ldots\mu_{J_1},\nu_1\ldots\nu_{J_2}}_{\alpha_1\ldots\alpha_J}(\Delta_1,\Delta_2,p_X) = \int d^4x\; d^4y \mathbf{e}^{-i\Delta_1 x - i\Delta_2 y}\cdot \\ \cdot\langle 0|T^* I^{\mu_1\ldots\mu_{J_1}}(x)\, I^{\nu_1\ldots\nu_{J_1}}(y)\, I_{\alpha_1\ldots\alpha_J}(0)\,|0\rangle\,, \tag{3.2}$$

where $I^{\mu_1\cdots\mu_J}$ is the current operator related to the hadronic spin-J field operator,

$$\left(\Box + m_J^2\right)\Phi^{\mu_1\cdots\mu_J}(x) = I^{\mu_1\cdots\mu_J}(x)\,. \tag{3.3}$$

The amplitude $T_{pp\to pXp}$ (fig. 1) is built with vertices $T^{\mu_1\cdots\mu_{J_1}}, T^{\nu_1\cdots\nu_{J_2}}, F^{\mu_1\cdots\mu_{J_1},\,\nu_1\cdots\nu_{J_2}}_{\alpha_1\cdots\alpha_J}$ and the propagators $d(J,t)/(m^2(J)-t)$ which have the poles at

$$m^2(J)-t=0\,,\qquad \text{i.e. } J=\alpha(t)\,, \tag{3.4}$$

after an appropriate analytic continuation of the signatured amplitudes in J. We assume that these poles, where α is the Pomeron trajectory, give the dominant contribution at high energies after having taken the corresponding residues. Regge-cuts are generated by the unitarization.

For vertex functions $T_{1,2}$ we can obtain the following tensor decomposition:

$$T^{\mu_1\cdots\mu_J}(p,\Delta) = T_0(\Delta^2)\sum_{n=0}^{\left[\frac{J}{2}\right]} Y_J^n \mathcal{T}^{\mu_1\ldots\mu_J}_{00J,\,n}\,, \tag{3.5}$$

$$Y_J^n = \frac{2^n(2(J-n))!J!D_p^{2n}}{(J-n)!(2J)!} \tag{3.6}$$

$$D_p^\mu = 2p^\mu - \Delta^\mu\,,\qquad D_p^2 = 4m^2-\Delta^2\,, \tag{3.7}$$

that satisfies Rarita-Schwinger conditions:

$$T^{\mu_1\ldots\mu_i\ldots\mu_j\ldots\mu_J} = T^{\mu_1\ldots\mu_j\ldots\mu_i\ldots\mu_J} \tag{3.8}$$

$$\Delta_{\mu_i}T^{\mu_1\ldots\mu_i\ldots\mu_J} = 0 \tag{3.9}$$

$$g_{\mu_i\mu_j}T^{\mu_1\ldots\mu_i\ldots\mu_j\ldots\mu_J} = 0 \tag{3.10}$$

Tensor structures $\mathcal{T}^{\mu_1\ldots\mu_J}_{00J,\,n}$ satisfy only two conditions (3.8),(3.9) and consist of the elements D_p^μ and $G^{\mu\nu}$:

$$G^{\mu\nu} = -g^{\mu\nu} + \frac{\Delta^\mu\Delta^\nu}{\Delta^2}\,. \tag{3.11}$$

$$\mathcal{T}^{\mu_1\ldots\mu_J}_{00J,\,n} = D_p^{(\mu_1}\cdots D_p^{\mu_{J-2n}}G^{\mu_{J-2n+1}\mu_{J-2n+2}}\cdots G^{\mu_{J-1}\mu_J)} \tag{3.12}$$

Let us assume the fusion of two particles with spins J_1 and J_2 into a particle with spin J. The general structure $F^{\mu_1\ldots\mu_{J_1},\,\nu_1\ldots\nu_{J_2}}_{\alpha_1\ldots\alpha_J}(\Delta_1,\Delta_2,p_X)$ (see figure 1) should satisfy the conditions (3.8–3.10) on each group of indices. Since the contraction of the vertex with structures $T^{\mu_1\ldots\mu_{J_1}}(p_1,\Delta_1)$, $T^{\nu_1\ldots\nu_{J_2}}(p_2,\Delta_2)$ and polarization tensor $e^{\alpha_1\ldots\alpha_J}(p_X)$ of the X particle leads to vanishing of some terms in F, the remainder can be constructed from

$$p_X^{\mu_i}, p_X^{\nu_j},\ \Delta_1^{\alpha_k}(\text{or }\Delta_2^{\alpha_k}),\ g^{\mu_i\nu_j}, g^{\mu_i\alpha_k},\ g^{\nu_j\alpha_k} \tag{3.13}$$

for tensors and additional terms

$$\Lambda_X^{\mu_i\nu_j\alpha_k} = p_X^\rho \epsilon^{\rho\mu_i\nu_j\alpha_k}\,, \tag{3.14}$$

$$\Lambda_n^{\mu_i\nu_j\alpha_k} = \Delta_n^\rho \epsilon^{\rho\mu_i\nu_j\alpha_k}\,, \tag{3.15}$$

$$Q_n^{\lambda\kappa} = \Delta_n^\rho p_X^\sigma \epsilon^{\rho\sigma\lambda\kappa} \to \Delta_1^\rho\Delta_2^\sigma\epsilon^{\rho\sigma\lambda\kappa}\,, \tag{3.16}$$

$$n=1 \text{ or } 2\,,\qquad (\lambda\kappa) = (\mu_i\nu_j)\,,(\mu_i\alpha_k)\,,(\nu_j\alpha_k)$$
$$i\le J_1\,,\qquad j\le J_2\,,\qquad k\le J$$

for pseudo-tensors [14]. Let $J_1 \le J_2$, and let us consider several cases.

- $J^P = 0^+$

$$F^{\mu_1\ldots\mu_{J_1},\nu_1\ldots\nu_{J_2}}(\Delta_1,\Delta_2,p_X) = \sum_{k=0}^{J_1} f_k\Big(p_X^{\mu_1}\cdot\ldots\cdot p_X^{\mu_{J_1-k}}\cdot g^{\mu_{J_1-k+1}\nu_{J_2-k+1}}\cdot\ldots\cdot g^{\mu_{J_1}\nu_{J_2}}\cdot \cdot p_X^{\nu_{J_2-k}}\cdot\ldots\cdot p_X^{\nu_1}\Big)\,. \tag{3.17}$$

After the tensor contraction we obtain the expansion of the type

$$\sum_{n=0}^{\left[\frac{J_i}{2}\right]} \frac{(-1)^n C^n_{J_i} C^{2n}_{J_i}}{C^{2n}_{2J_i}} \left(\sqrt{\frac{\xi_i^2(m^2-t_i/4)}{-t_i}}\,\frac{M_X^2-t_1-t_2}{M_\perp^2}\right)^{2n}. \tag{3.18}$$

In the kinematical region (2.3),(2.4) $\xi_i^2 m^2/|t_i| \ll 1$ and we can keep only the first term of the expansion (3.18). The tensor product is given by

$$T^{J_1}(p_1,\Delta_1)\otimes F^{J_1,\,J_2,\,0^+}(\Delta_1,\Delta_2)\otimes T^{J_2}(p_2,\Delta_2) \sim s_1^{J_1}s_2^{J_2}\sum_{k=0}^{J_1}\frac{f_k 2^k}{M_\perp^{2k}}\,. \tag{3.19}$$

- $J^P = 0^-$

$$F^{\mu_1\ldots\mu_{J_1},\,\nu_1\ldots\nu_{J_2}}(\Delta_1,\Delta_2,p_X) = Q_n^{\mu_1\nu_1}\sum_{k=0}^{J_1-1} f_k\Big(p_X^{\mu_2}\cdot\ldots\cdot p_X^{\mu_{J_1-k}}\cdot \cdot g^{\mu_{J_1-k+1}\nu_{J_2-k+1}}\cdot\ldots\cdot g^{\mu_{J_1}\nu_{J_2}}\cdot \cdot p_X^{\nu_{J_2-k}}\cdot\ldots\cdot p_X^{\nu_2}\Big), \tag{3.20}$$

$$\begin{aligned} T^{J_1}(p_1,\Delta_1)\otimes F^{J_1,\,J_2,\,0^-}(\Delta_1,\Delta_2)\otimes \\ \otimes T^{J_2}(p_2,\Delta_2) &\sim 4Q_n^{\mu_1\nu_1}p_1^{\mu_1}p_2^{\nu_1}s_1^{J_1-1}s_2^{J_2-1}\sum_{k=0}^{J_1-1}\frac{f_k2^k}{M_\perp^{2k}} \\ &\simeq [\boldsymbol{\Delta}_1\times\boldsymbol{\Delta}_2]\cdot s_1^{J_1}s_2^{J_2}\sum_{k=0}^{J_1-1}\frac{f_k 2^{k+1}}{M_\perp^{2k+2}} \end{aligned} \tag{3.21}$$

- $J^P = 1^-$

$$F^{\mu_1\dots\mu_{J_1},\,\nu_1\dots\nu_{J_2},\,\alpha}(\Delta_1,\Delta_2,p_X) = g^{\alpha\mu_1}\sum_{k=0}^{J_1-1} f_k\Big(p_X^{\mu_2}\cdot\ldots\cdot p_X^{\mu_{J_1-k}}\cdot \cdot g^{\mu_{J_1-k+1}\nu_{J_2-k+1}}\cdot\ldots\cdot g^{\mu_{J_1}\nu_{J_2}}\cdot \cdot p_X^{\nu_{J_2-k}}\cdot\ldots\cdot p_X^{\nu_1}\Big) + +g^{\alpha\nu_1}\sum_{k=0}^{J_1} f_{J_1+k}\Big(p_X^{\mu_1}\cdot\ldots\cdot p_X^{\mu_{J_1-k}}\cdot \cdot g^{\mu_{J_1-k+1}\nu_{J_2-k+1}}\cdot\ldots\cdot g^{\mu_{J_1}\nu_{J_2}}\cdot \cdot p_X^{\nu_{J_2-k}}\cdot\ldots\cdot p_X^{\nu_2}\Big) + +\Delta_n^\alpha\sum_{k=0}^{J_1} f_{2J_1+k+1}\Big(p_X^{\mu_1}\cdot\ldots\cdot p_X^{\mu_{J_1-k}}\cdot \cdot g^{\mu_{J_1-k+1}\nu_{J_2-k+1}}\cdot\ldots\cdot \cdot g^{\mu_{J_1}\nu_{J_2}}\cdot \cdot p_X^{\nu_{J_2-k}}\cdot\ldots\cdot p_X^{\nu_1}\Big), \tag{3.22}$$

$$T^{J_1}(p_1,\Delta_1)\otimes F^{J_1,\,J_2,\,1^-}(\Delta_1,\Delta_2)\otimes \otimes T^{J_2}(p_2,\Delta_2)\sim s_1^{J_1}s_2^{J_2}\left[\frac{2p_1^\alpha}{s_1}\sum_{k=0}^{J_1-1}\frac{f_k 2^k}{M_\perp^{2k}}+\frac{2p_2^\alpha}{s_2}\sum_{k=0}^{J_1}\frac{f_{J_1+k}2^k}{M_\perp^{2k}}+ + \Delta_n^\alpha\sum_{k=0}^{J_1}\frac{f_{2J_1+1+k}2^k}{M_\perp^{2k}}\right]. \tag{3.23}$$

It is easy to show from the general form of the tensor decompositions, that after the reggeization procedure one obtains the structure of the amplitude, which is similar to the special case $J_1 = J_2 = 1$. In this case it is convenient to use the following bose-symmetric form of the tensor

$$F^{\mu\nu,\,\alpha}(\Delta_1,\Delta_2,p_X) = f_0 g^{\alpha\mu}\Delta_1^\nu + \bar f_0 g^{\alpha\nu}\Delta_2^\mu + \left(f_1\Delta_1^\alpha + \bar f_1\Delta_2^\alpha\right)\Delta_2^\mu\Delta_1^\nu + + \left(f_2\Delta_1^\alpha + \bar f_2\Delta_2^\alpha\right) g^{\mu\nu}\,, \tag{3.24}$$

and the tensor product

$$T^{J_1}(p_1,\Delta_1)\otimes F^{J_1,\,J_2,\,1^-}(\Delta_1,\Delta_2)\otimes \otimes T^{J_2}(p_2,\Delta_2)\sim s_1^{J_1}s_2^{J_2}\left[\frac{2p_1^\alpha}{s_1}f_0 + \frac{2p_2^\alpha}{s_2}\bar f_0 + \left(f_1+\frac{f_2}{M_\perp^2}\right)\Delta_1^\alpha + + \left(\bar f_1 + \frac{\bar f_2}{M_\perp^2}\right)\Delta_2^\alpha\right], \tag{3.25}$$

where $\bar f(t_1,t_2) = f(t_2,t_1)$.

- $J^P = 1^+$

$$
\begin{aligned}
F^{\mu_1\dots\mu_{J_1},\nu_1\dots\nu_{J_2},\alpha}(\Delta_1,\Delta_2,p_X) = \; & \Lambda_X^{\mu_1\nu_1\alpha}\sum_{k=0}^{J_1-1} f_k \Big(p_X^{\mu_2}\cdot\ldots\cdot p_X^{\mu_{J_1-k}}\cdot \\
& \cdot g^{\mu_{J_1-k+1}\nu_{J_2-k+1}}\cdot\ldots\cdot g^{\mu_{J_1}\nu_{J_2}}\cdot \\
& \cdot p_X^{\nu_{J_2-k}}\cdot\ldots\cdot p_X^{\nu_2}\Big) + \\
& +\Lambda_n^{\mu_1\nu_1\alpha}\sum_{k=0}^{J_1-1} f_{J_1+k}\Big(p_X^{\mu_2}\cdot\ldots\cdot p_X^{\mu_{J_1-k}}\cdot \\
& \cdot g^{\mu_{J_1-k+1}\nu_{J_2-k+1}}\cdot\ldots\cdot \\
& \cdot g^{\mu_{J_1}\nu_{J_2}}\cdot \\
& \cdot p_X^{\nu_{J_2-k}}\cdot\ldots\cdot p_X^{\nu_2}\Big) + \\
& +Q_n^{\nu_1\alpha} f_{2J_1}\Big(g^{\mu_1\nu_2}\cdot\ldots\cdot g^{\mu_{J_1}\nu_{J_1+1}}\cdot \\
& \cdot p_X^{\nu_{J_1+2}}\cdot\ldots\cdot p_X^{\nu_{J_2}}\Big) + \\
& +Q_n^{\mu_1\alpha}\sum_{k=0}^{J_1-1} f_{2J_1+k+1}\Big(p_X^{\mu_2}\cdot\ldots\cdot p_X^{\mu_{J_1-k}}\cdot \\
& \cdot g^{\mu_{J_1-k+1}\nu_{J_2-k+1}}\cdot\ldots\cdot \\
& \cdot g^{\mu_{J_1}\nu_{J_2}}\cdot \\
& \cdot p_X^{\nu_{J_2-k}}\cdot\ldots\cdot p_X^{\nu_1}\Big), \quad (3.26)
\end{aligned}
$$

$$
\begin{aligned}
& T^{J_1}(p_1,\Delta_1)\otimes F^{J_1,\,J_2,\,1^+}(\Delta_1,\Delta_2)\otimes \\
& \otimes T^{J_2}(p_2,\Delta_2) \sim s_1^{J_1}s_2^{J_2}\Bigg[\frac{2p_1^{\mu_1}p_2^{\nu_1}p_X^{\rho}\epsilon^{\rho\alpha\mu_1\nu_1}}{s}\sum_{k=0}^{J_1-1}\frac{f_k 2^{k+1}}{M_\perp^{2k+2}} + \\
& + \frac{2p_1^{\mu_1}p_2^{\nu_1}\Delta_n^{\rho}\epsilon^{\rho\alpha\mu_1\nu_1}}{s}\sum_{k=0}^{J_1-1}\frac{f_{J_1+k}2^{k+1}}{M_\perp^{2k+2}} + \\
& + \frac{2p_2^{\nu_1}p_X^{\rho}\Delta_n^{\sigma}\epsilon^{\rho\sigma\alpha\nu_1}}{s_2}\frac{f_{2J_1}2^{J_1}}{M_\perp^{2J_1}} + \quad (3.27) \\
& + \frac{2p_1^{\mu_1}p_X^{\rho}\Delta_n^{\sigma}\epsilon^{\rho\sigma\alpha\mu_1}}{s_1}\sum_{k=0}^{J_1-1}\frac{f_{2J_1+1+k}2^{k}}{M_\perp^{2k}}\Bigg].
\end{aligned}
$$

For $J_1 = J_2 = 1$

$$
\begin{aligned}
& T^{J_1}(p_1,\Delta_1)\otimes F^{J_1,J_2,1^+}(\Delta_1,\Delta_2)\otimes \\
& \otimes T^{J_2}(p_2,\Delta_2) \sim s_1^{J_1}s_2^{J_2}\Bigg[\frac{4p_1^{\sigma}p_2^{\lambda}\left(f_0\Delta_1^{\rho}-\bar f_0\Delta_2^{\rho}\right)}{sM_\perp^2} + \\
& + 2\Delta_1^{\rho}\Delta_2^{\sigma}\left(\frac{f_1p_1^{\lambda}}{s_1}-\frac{\bar f_1 p_2^{\lambda}}{s_2}\right)\Bigg]\epsilon^{\rho\sigma\lambda\alpha} = \\
& = 2s_1^{J_1}s_2^{J_2}\left[p_1^{\lambda}P_2^{\rho}\Delta_1^{\sigma}+p_2^{\lambda}P_1^{\rho}\Delta_2^{\sigma}\right]\epsilon^{\rho\sigma\lambda\alpha}\,, \quad (3.28)
\end{aligned}
$$

$$P_1 = -\frac{\bar{f}_1}{s_2}\Delta_1 + \frac{\bar{f}_0}{sM_\perp^2}2p_1 \ ,$$
$$P_2 = -\frac{f_1}{s_1}\Delta_2 + \frac{f_0}{sM_\perp^2}2p_2 \ . \tag{3.29}$$

- $J^P = 2^+$. For simplicity we consider only the case $J_1 = J_2 = 1$.

$$\begin{aligned} F^{\mu\nu,\alpha_1\alpha_2}(\Delta_1,\Delta_2,p_X) &= f_0 g^{\alpha_1\mu} g^{\alpha_2\nu} + f_1 g^{\alpha_1\mu}\Delta_1^{\alpha_2}\Delta_1^{\nu} + \\ &+ \bar{f}_1 g^{\alpha_1\nu}\Delta_2^{\alpha_2}\Delta_2^{\mu} + \\ &+ \left(f_2\Delta_1^{\alpha_1}\Delta_1^{\alpha_2} + \bar{f}_2\Delta_2^{\alpha_1}\Delta_2^{\alpha_2}\right) g^{\mu\nu} + \\ &+ \left(f_3\Delta_1^{\alpha_1}\Delta_1^{\alpha_2} + \bar{f}_3\Delta_2^{\alpha_1}\Delta_2^{\alpha_2}\right)\Delta_2^{\mu}\Delta_1^{\nu} \ , \end{aligned} \tag{3.30}$$

$$\begin{aligned} T^{J_1}(p_1,\Delta_1) \otimes F^{J_1,J_2,2^+}(\Delta_1,\Delta_2) \otimes & \\ \otimes T^{J_2}(p_2,\Delta_2) \sim s_1^{J_1} s_2^{J_2} \Bigg[& f_0 \frac{4p_1^{\alpha_1}p_2^{\alpha_2}}{sM_\perp^2} + f_1\frac{2p_1^{\alpha_1}\Delta_1^{\alpha_2}}{s_1} + \\ & + \bar{f}_1\frac{2p_2^{\alpha_1}\Delta_2^{\alpha_2}}{s_2} + \\ & + \frac{2\left(f_2\Delta_1^{\alpha_1}\Delta_1^{\alpha_2} + \bar{f}_2\Delta_2^{\alpha_1}\Delta_2^{\alpha_2}\right)}{M_\perp^2} + \\ & + \left(f_3\Delta_1^{\alpha_1}\Delta_1^{\alpha_2} + \bar{f}_3\Delta_2^{\alpha_1}\Delta_2^{\alpha_2}\right)\Bigg] . \end{aligned} \tag{3.31}$$

Everywhere in the above expressions $f_k = f_k^{J^P}(t_1, t_2, M_X^2)$.

4. Azimuthal angle dependence

In this section we will obtain the general structure of the azimuthal angle dependence for different J^P states in the tensor current picture. If we assume that the dominant contribution is given by the Regge poles $\alpha_1(t_1)$ and $\alpha_2(t_2)$ then the amplitude of the process can be written in the following form

$$T_{pp\to pXp} \sim \eta(\alpha_1(t_1))\eta(\alpha_2(t_2))\left[T^{J_1} \otimes F^{J_1,\,J_2,\,J^P} \otimes T^{J_2}\right]_{\substack{J_1\to\alpha_1 \\ J_2\to\alpha_2}} \ , \tag{4.1}$$

where $J_i \to \alpha_i$ means the usual procedure of the analytical continuation to the complex J-plane and taking residues at Regge poles. $\eta(\alpha_i)$ are signature factors.

In the most important case, when the main contribution is given by one Pomeron trajectory, $\alpha_1 = \alpha_2 = \alpha_{\mathbb{P}}(0)$ and the "bare" amplitude squared is proportional to following expressions:

- $J^P = 0^+$

$$|T_{pp\to pXp}|^2 \sim (M_\perp^2)^{2(\alpha_{\mathbb{P}}(0)-1)}(f_0 M_\perp^2 + 2f_1)^2 \tag{4.2}$$

- $J^P = 0^-$

$$|T_{pp\to pXp}|^2 \sim (M_\perp^2)^{2(\alpha_{\mathbb{P}}(0)-1)} f_0 t_1 t_2 \sin^2\phi_0 \tag{4.3}$$

- $J^P = 1^-$. In this case we assume the existence of a C-odd vacuum trajectory, "Odderon", $\alpha_{\mathbb{O}}(t)$.

$$
\begin{aligned}
|T_{pp\to pXp}|^2 &\sim (M_\perp^2)^{\alpha_{\mathbb{P}}(0)+\alpha_{\mathbb{O}}(0)-2}(\mathcal{F}_0 M_\perp^4 + \mathcal{F}_1 M_\perp^2 + \mathcal{F}_2)\,, && (4.4)\\
\mathcal{F}_0 &= \frac{f_0^{S\,2}}{M_X^2} + f_0^A f_1^A + \frac{(t_1-t_2) f_1^A f_0^S}{M_X^2} + \frac{f_1^{A\,2}\lambda}{4M_X^2}\,,\\
\mathcal{F}_1 &= f_0^{A\,2} - f_0^{S\,2} - f_0^A f_2^A + \frac{(t_1-t_2) f_2^A f_0^S}{M_X^2} + \frac{f_1^A f_2^A \lambda}{2M_X^2}\,,\\
\mathcal{F}_2 &= \frac{f_2^{A\,2}\lambda}{4M_X^2}\,, && (4.5)\\
\lambda &= \lambda(M_X^2, t_1, t_2) = M_X^4 + t_1^2 + t_2^2 - 2M_X^2 t_1 - 2M_X^2 t_2 - 2t_1 t_2\,,\\
f_k^S &= f_k + \bar{f}_k\,, f_k^A = f_k - \bar{f}_k\,.
\end{aligned}
$$

- $J^P = 1^+$

$$
\begin{aligned}
|T_{pp\to pXp}|^2 &\sim (M_\perp^2)^{2(\alpha_{\mathbb{P}}(0)-1)}(\mathcal{F}_0 M_\perp^4 + \mathcal{F}_1 t_1 t_2 \sin^2\phi_0 + \mathcal{F}_2)\,, && (4.6)\\
\mathcal{F}_0 &= (f_1 \mathbf{\Delta}_1 - \bar{f}_1 \mathbf{\Delta}_2)^2\,, && (4.7)\\
\mathcal{F}_1 &= \frac{2(s_2 f_1 - s_1 \bar{f}_1)^2}{s} + 4M_\perp^2 f_1 \bar{f}_1 - \frac{4(f_0+\bar{f}_0)^2}{M_X^2}\,, && (4.8)\\
\mathcal{F}_2 &= 4(f_0 \mathbf{\Delta}_1 - \bar{f}_0 \mathbf{\Delta}_2)^2\,, && (4.9)
\end{aligned}
$$

- $J^P = 2^+$

$$
\begin{aligned}
|T_{pp\to pXp}|^2 &\sim (M_\perp^2)^{2(\alpha_{\mathbb{P}}(0)-1)}(\mathcal{F}_0 M_\perp^4 + \mathcal{F}_1 M_\perp^2 + \mathcal{F}_2)\,, && (4.10)\\
\mathcal{F}_0 &= \frac{f_1^{S\,2} - 12 f_0 f_3^S}{24} + \frac{f_1^S(-4f_0 + f_1^A(t_1-t_2) + 2\lambda f_3^S)}{12M_X^2} +\\
&+ \frac{1}{24M_X^4}\Big[16 f_0^2 + 4f_0\left(4f_1^A(t_1-t_2) + f_3^S(3(t_1-t_2)^2-\lambda)\right) +\\
&\qquad + 4f_3^S f_1^A(t_1-t_2)\lambda + f_3^{S\,2}\lambda +\\
&\qquad + f_1^{A\,2}\left((t_1-t_2)^2 + 3\lambda\right)\Big]\,, && (4.11)\\
\mathcal{F}_1 &= -\frac{1}{3} f_0\left(f_1^S + 3f_2^S\right) +\\
&+ \frac{f_2^S}{6M_X^4}\left[\lambda^2 f_3^S + 2\lambda\left(-f_0 + f_1^A(t_1-t_2)\right) + 6f_0(t_1-t_2)^2\right] +\\
&+ \frac{1}{6M_X^2}\Big[\lambda\left(2f_1^S f_2^S + 2f_0 f_3^S + 3(f_1^{S\,2} - f_1^{A\,2})/4\right) - 4f_0^2 -\\
&\qquad - 2f_0 f_1^A(t_1-t_2)\Big]\,, && (4.12)\\
\mathcal{F}_2 &= \frac{\left(2f_0 + \frac{f_2^S\lambda}{M_X^2}\right)^2}{6}\,. && (4.13)
\end{aligned}
$$

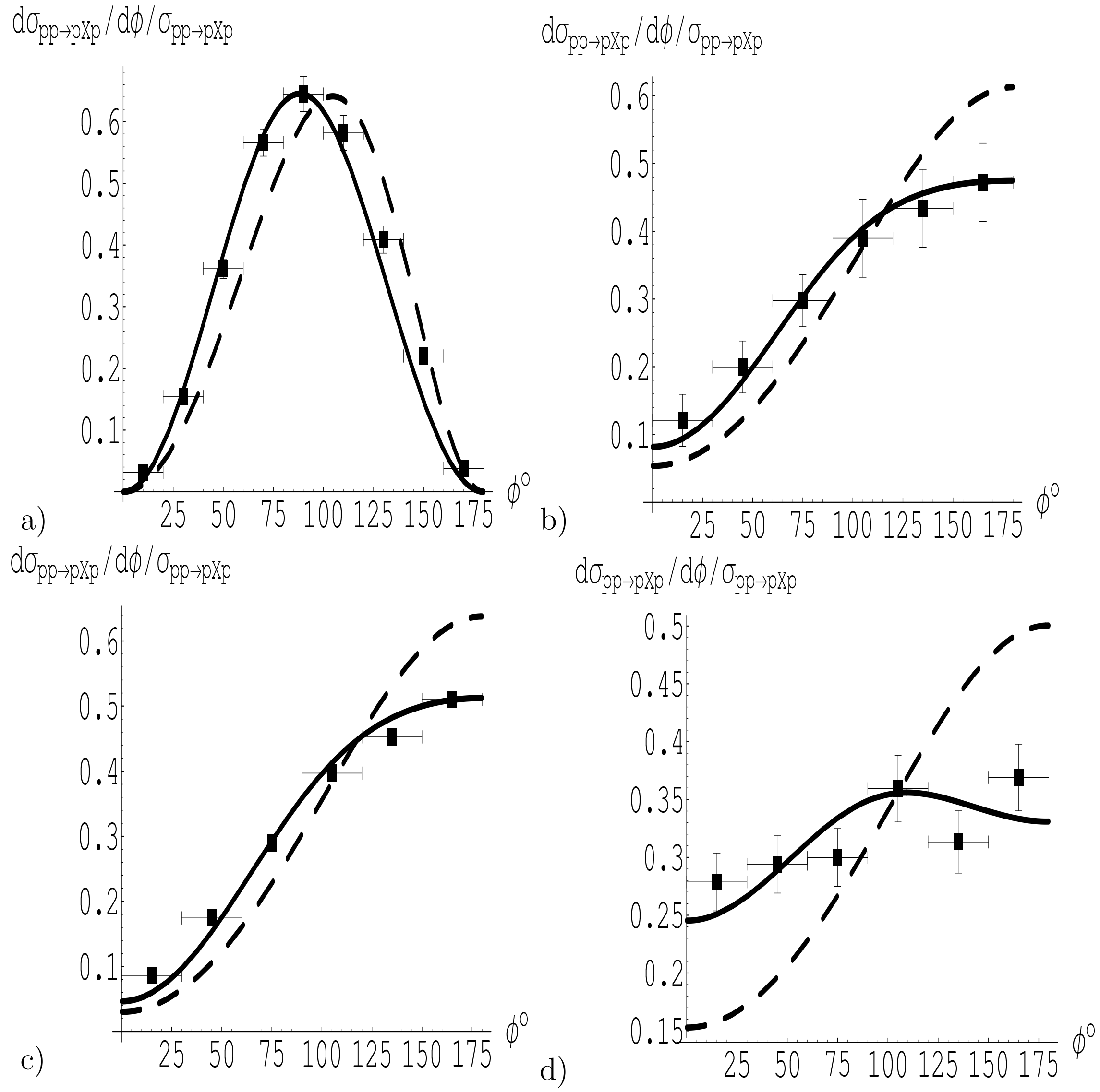

Figure 3: Experimental data from WA102. The dashed curve represents the "bare" cross-section and the solid one is the unitarized result. a) η', 0^{-+}; b) $f_1(1285)$, 1^{++}, all t_i; c) $f_1(1285)$, $|t_1 - t_2| < 0.2\,\mathrm{GeV}^2$; d) $f_1(1285)$, $|t_1 - t_2| > 0.4\,\mathrm{GeV}^2$.

Similar formulae for the differential cross-sections were obtained by other authors. In reference [15] results were obtained from the assumption that the Pomeron acts as a 1^+ conserved or nonconserved current. It was shown in [16] and (with more detailed investigations) in [17] that the same result follows from the simple Regge behaviour of helicity amplitudes. Experimental data are in good agreement with these predictions.

A good description of those data in the framework of our approach (taking into account the absorption) is obtained (figures 3,4).

In addition to the present description there are some models of the "glueball" production based on the "instanton" dynamics [18].

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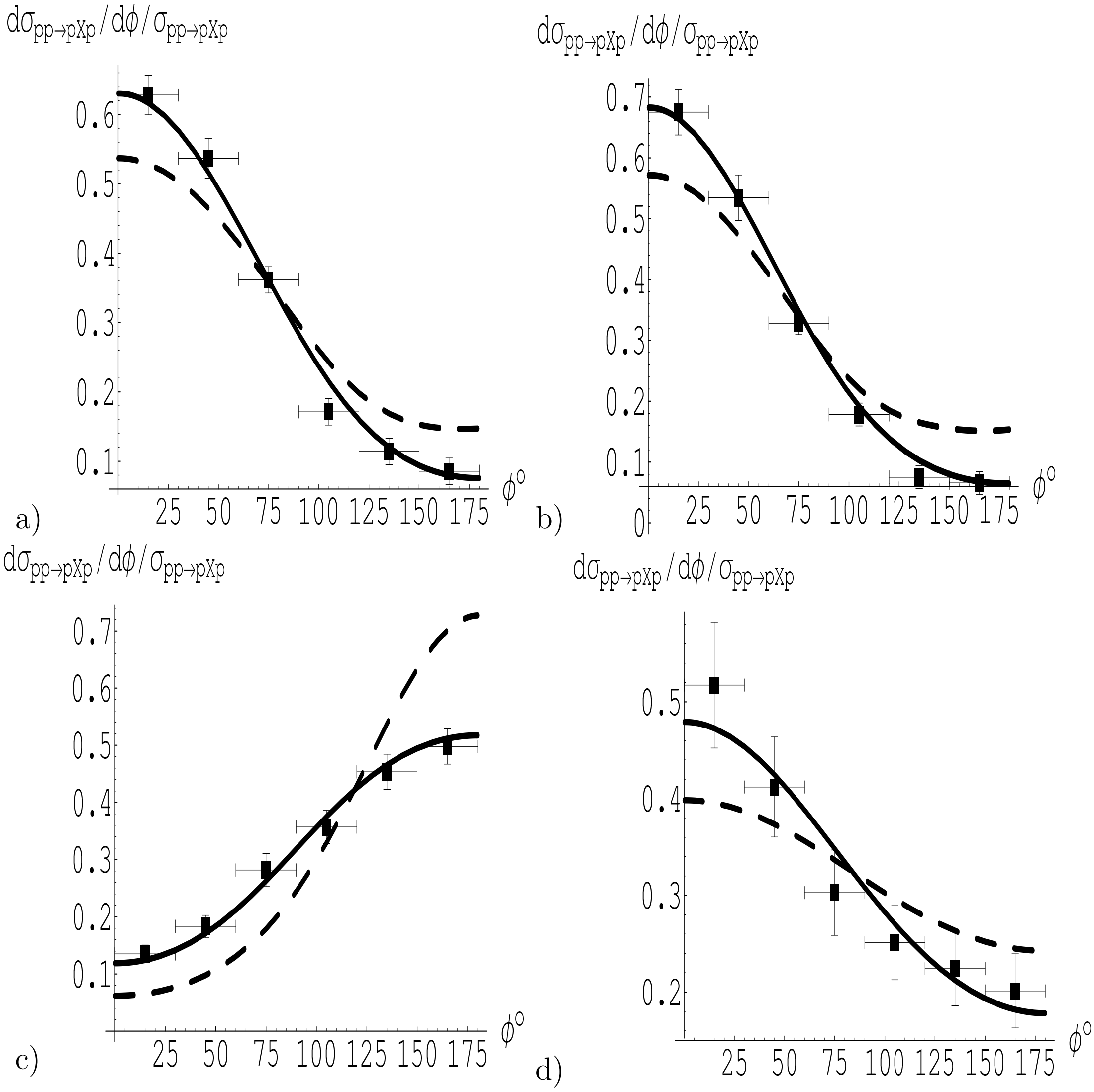

Figure 4: Experimental data from WA102 averaged over all measured t_i values. The dashed curve represents the "bare" cross-section and the solid one is the unitarized result. a) $f_0(980)$, 0^{++}; b) $f_0(1500)$, 0^{++}; c) $f_2(1270)$, 2^{++}; d) $f_2(1950)$, 2^{++}.

5. WA102 and predictions for the LHC

It is important to stress the fact, that at the WA102 energies absorptive effects are not so significant, and the azimuthal angle dependence looks like the "bare" one. We can use this fact to simplify the fitting procedure, that has been already done by the WA102 collaboration. Only at large values of $\mathrm{dP}_\perp = |\mathbf{\Delta}_1 - \mathbf{\Delta}_2|$ the process of "soft" rescattering can change the picture violently (see figure 3d).

In figures 3,4 we show the data from WA102 [7] and our curves for the "bare" and unitarized amplitudes. One can see that all the features of the average ϕ_0-dependence are consistent with the data. It makes possible to predict the azimuthal angle behaviour at higher energies and use these predictions as a spin-parity analyser.

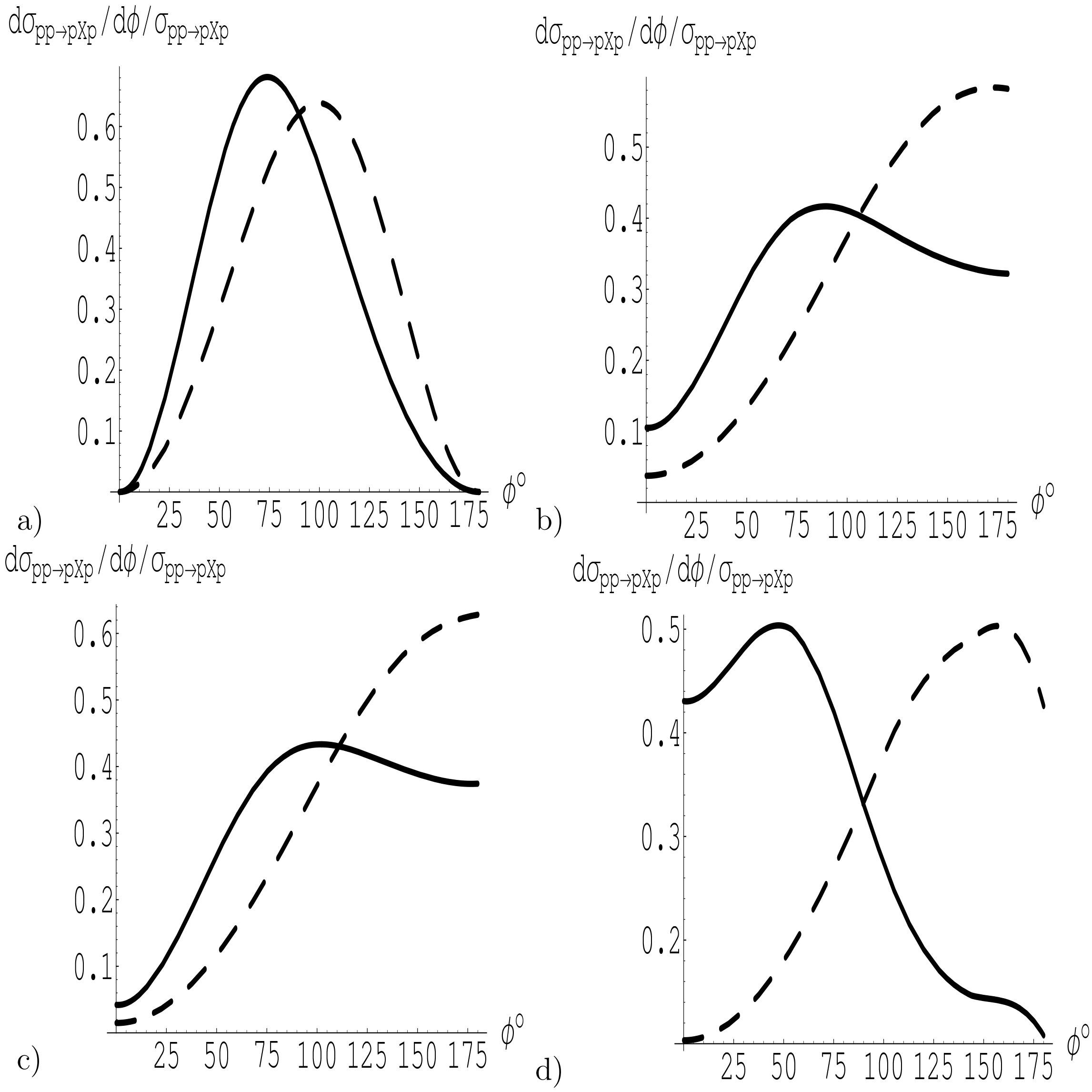

Figure 5: Results for the LHC energies. a) η', 0^{-+}; b) $f_1(1285)$, 1^{++}, all t_i; c) $f_1(1285)$, $|t_1 - t_2| < 0.1$ GeV^2; d) $f_1(1285)$, $|t_1 - t_2| > 0.2$ GeV^2;

The main property is that unitarization adds up to the shift or distortion of the "bare" cross-section towards small angles and to the reduction of its value. The difference is more significant at the LHC than at the WA102 energies, and we should take it into account necessarily. The "soft" survival probability is $0.25 \to 0.3$ for WA102 and $0.05 \to 0.1$ for the LHC. It depends on the mass M_X and on the kinematical cuts.

Features concerning each particle are the same as was mentioned in reference [15]:

- for the η' meson the kinematical distortion because of the different reference frame is totally compensated by the unitarization at the WA102 energy (figure 3a), and for the LHC the peak in the cross-section is shifted towards 65^o (figure 5a).

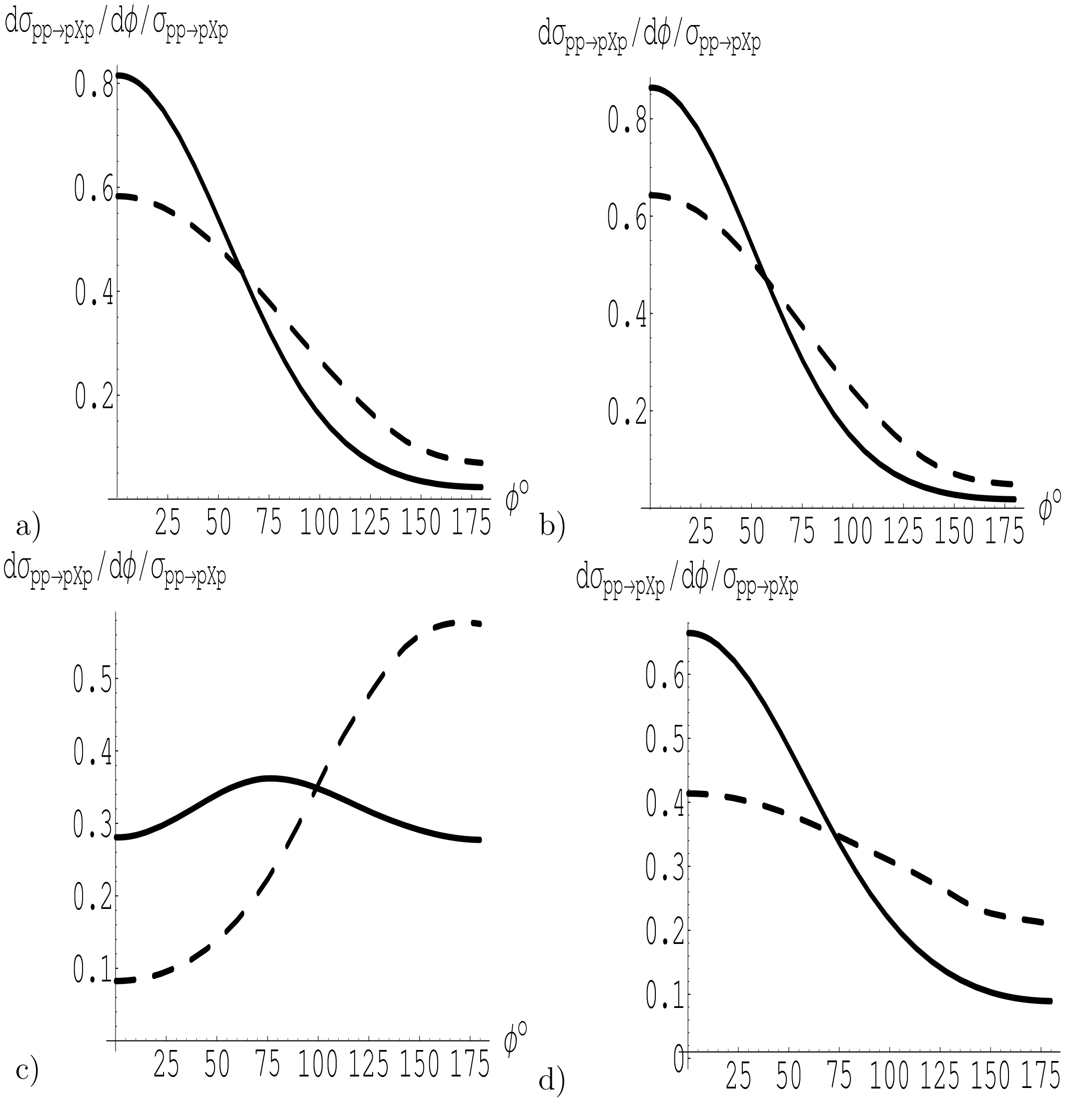

Figure 6: Results for the LHC energies. a)$f_0(980)$, 0^{++}; b)$f_0(1500)$, 0^{++}; c)$f_2(1270)$, 2^{++}; d) $f_2(1950)$, 2^{++};

- for the $\mathrm{f}_1(1285)$ we have almost a "flat" distribution at large values of $|t_1 - t_2|$ (figure 3d), since in the simplest case its cross-section is proportional to $\mathrm{dP}_\perp^2 = (\boldsymbol{\Delta}_1 - \boldsymbol{\Delta}_2)^2$. For the LHC we obtain a more stronger suppression for large angles (figure 5d).

- the difference between 0^{++} $q\bar{q}$ ($\mathrm{f}_2(1270)$) and non-$q\bar{q}$ ($\mathrm{f}_0(980)$, $\mathrm{f}_0(1500)$, $\mathrm{f}_2(1900)$) states at WA102 (figure 4) changes due to unitarization effects. For the $q\bar{q}$, the azimuthal angle dependence becomes almost "flat" (figure 6c), and for the non-$q\bar{q}$ mesons we see a shrinkage of the peak at $\phi_0 = 0^o$. The same is valid for 2^{++} mesons. To separate $q\bar{q}$ states from "glueball" candidates we can use the filter which was proposed in refs. [9, 10]. It was shown that the quantity $R = N(\mathrm{dP}_\perp < 0.2\ \mathrm{GeV})/N(\mathrm{dP}_\perp > 0.5\ \mathrm{GeV})$ is large (~ 1) for "glueball" candidates and becomes small for $q\bar{q}$ states due to the differencies in the dynamics of particles production.

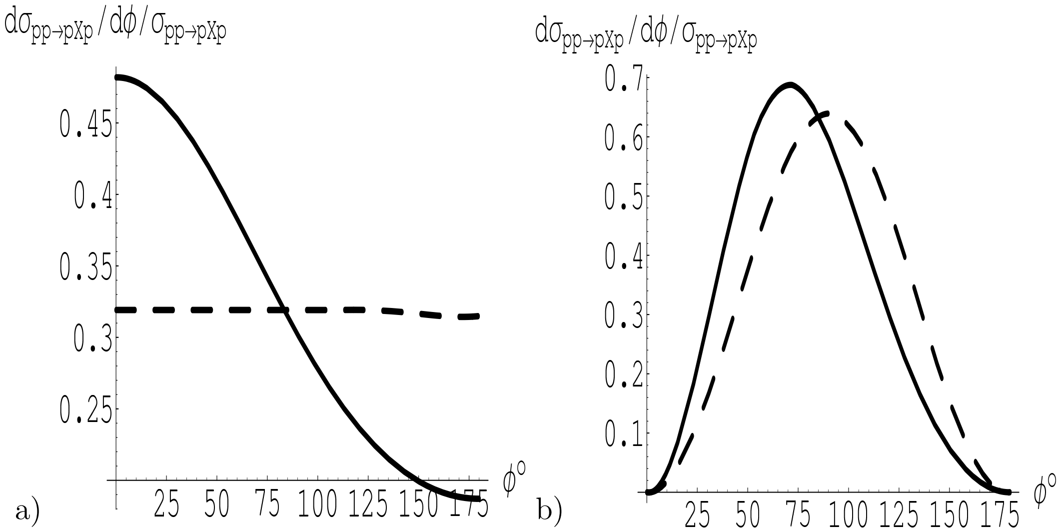

Figure 7: Examples of the azimuthal angle dependence for large mass states. a) 1^{--}, $M_X =$ 50 GeV; b) 0^{-+}, $M_X = 50$ GeV;

For the low mass particles, which can be produced in EDDE, total cross-sections are estimated to be of the order $1 \div 30$ μb at the LHC. Cross-sections for the "glueball" candidates $f_0(1500)$ and $f_2(1950)$ are about 30 μb (depends on the LHC kinematics and may be larger) and the effective slope is 10. The typical value of ξ is of the order $M_X/\sqrt{s} \sim 10^{-4}$. Experimental possibilities of azimuthal angle measurements for particles with masses less than 1 GeV (like η' and $f_0(980)$) seem to be doubtful. As to other light mesons, measuremets are possible at low luminocities only, when appropriate resolution on ξ and ϕ_0 can be achieved. This has to be studied in futher Monte-Carlo simulations.

It is possible to apply the method to large mass particles. It was argued in [19] that a heavy glueball ("knot") 1^{--} with the mass close to 50 GeV may exist. One can see from (4.2–4.10) that in this case, the ϕ_0 dependence is simply defined by the unitarization only, since $M_\perp^2 \simeq M_X^2 = const$. In this case, we have a good tool to check models for the "soft" rescattering. Examples are depicted in figure 7.

6. Conclusions

Detailed investigations of the azimuthal angle dependence in EDDE can help to solve several important problems:

- to check different models for "soft" processes and to study the real pattern of the interaction.
- to understand the difference in the dynamics of production of the $q\bar{q}$ and non-$q\bar{q}$ states and their possible filtering.
- to determine the quantum numbers of new produced particles.

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