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DIFFRACTIVE REACTIONS AND SHOWER FACTOR DEPENDENCE ON ENERGY AND MOMENTUM TRANSFER

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The analysis of all the existing experimental data on diffractive reactions in the region of small masses and momentum transfers is carried out in the model of complex angular moments taking into account the cuts. A phenomenological parametrization is obtained which describes correctly all the features of the diffraction cross sections and allows to separate the pomeron proton scattering cross section. Taking into account the contribution of cuts one obtains results strongly differing from the results obtained earlier in other works. The derived shower factors are function of energy and momentum transfer due to the strong difference between the slops of differential cross sections of elastic scattering and particle diffraction production. This conclusion, which is of interest by itself, allows to obtain also a correct description of the growth of the cross sections maintaining the condition $\alpha_p(0)=1$.

In the first section the shower factors (S. F) are considered in detail and it is shown the necessity of their revision taking into account the great number new experimental data on diffraction dissociation processes. In the second section we determine the quasieikonal model for elastic scattering and diffraction dissociation (D. D). In the third and fourth sections it is given the phenomenological parametrization of D.D. process amplitudes and their comparison with the experimental data. It is also shown that the new results on S.F influence the analysis of the elastic scattering.

I. SHOWER FACTORS

The shower factors have been first introduced by Kaydalov [1] in order to take into account the showers in the intermediate state of two reggeon cuts. The two reggeon cut may be presented as a sum of four diagrams (see Fig.1). Thus, it becomes evidient that the correct estimate of cut contribution into binary processes requires detail study of D. D processes. For this purpose three types of processes are essential (see Fig.2b,c). First let us consider only the case when the fixed invariant beam mass $S_x = M_x^2 = (q_1 + q_2 + \cdots + q_k)^2$ and momentum transfer $t_2 = (P_2' - P_2)^2 = (P_1 - \sum_{i=1}^{k} q_i)^2$ are small, i.e. $S_x \ll S$ and $t_2 \ll S_x$.

Due to the factorization of vacuum pole residues one may express the cross sections of two jet production by means of the cross sections of one jet and elastic scattering processes:

$$\frac{d^3\sigma}{dS_1\, dS_2\, dt} = \frac{d^2\sigma_1}{dS_1\, dt} \cdot \frac{d^2\sigma_2}{dS_2\, dt} \Big/ \frac{d\sigma^{el}}{dt} . \qquad (1)$$

Using the Gribov's diagram technics [2] one may determine the magnitude of two reggeon PP-cut in the elastic scattering:

$$M^{2P}(s,t) = \frac{i}{2\pi} \eta_P^2 \varepsilon_P^2 \int N_1 N_2 \left(\frac{s-u}{2s_0}\right)^{-\alpha_P'(\varkappa_a^2 + \varkappa_b^2)} \times e^{i\frac{\pi}{2}\alpha_P'(\varkappa_a^2 + \varkappa_b^2)} d^2\varkappa_b , \qquad (2)$$

where η_P is the signature factor at $\varkappa^2 = 0$, $\varkappa^2 = -t$.

$$\varepsilon_P = \exp\left(\alpha_P^0 \ln \frac{s-u}{2s_0}\right) \Big/ \sqrt{s(s-4m_P^2)}$$

and

$$N_m(\varkappa^2; S_m) = g_m^2(\varkappa^2) + \frac{1}{\pi}\int_{S_{0m}}^{S_1} \mathrm{Im}\, \tilde{N}_m(\varkappa^2; S_m)\, dS_m \; ; m = 1,2 \qquad (3)$$

$$S_1 = (P_1 + \varkappa)^2 , \quad S_2 = (P_2 + \varkappa)^2 ,$$

and $\widetilde{N}_m(\varkappa^2, S_m)$ is the pomeron elastic scattering amplitude on the particle m (All the notations and normalization of the experimentally measurable magnitudes are the same as in the work [3]). From the two reggeon unitarity condition [4] it follows:

$$\operatorname{Im} \widetilde{N}_m(\varkappa^2, S_m) = \pi \int |\mathcal{V}_m(\varkappa^2, S_m, \tau_m)|^2 d\tau_m \,, \tag{4}$$

where τ_m are variables characterizing the particles of the beam emitted by the particle m , while $\mathcal{V}_m(\varkappa^2, S_m, \tau_m)$ describes the transition of the particle m and P -pole into an intermediate state with mass $M_X = \sqrt{S_m}$. In the inclusive type experiments one measures only the momentum of the observed particle and carries out summation over all the particles of the beam (i.e. integration over the phase volume and summation over the number of particles). The magnitude

$$\sigma^{tot}_{\alpha a}(S_X, t) = \frac{1}{2S_X} \sum_n \frac{1}{\pi} \int |\mathcal{V}_n(S_X, t, \tau)|^2 d\tau$$

is usually called as total cross section of the interaction of reggeon α with the particle a .

From expressions (2-4) it is seen that in order to obtain the correct magnitude of two pomeron cut it is necessary to know the cross sections of inelastic processes going via the vacuum pole.

Let us introduce the amplitudes corresponding to the diagrams shown in Fig.2:

$$M_0^{1P}(S, \varkappa^2) = \eta_P \varepsilon_P g^P_{11}(\varkappa^2) g^P_{22}(\varkappa^2) \left(\frac{S-u}{2S_0}\right)^{-\alpha'_P \varkappa^2} e^{i\frac{\pi}{2}\alpha'_P \varkappa^2} \tag{5}$$

is the ρ -pole contribution into the elastic scattering amplitude (Fig.2a),

$$T_m^{1P}(s,\varkappa^2,s_m,\tau_m)=\eta_\rho \varepsilon_\rho^m g_{mm}^\rho(\varkappa^2)\mathcal{V}_m^\rho(s_m,\varkappa^2,\tau_m)\left(\frac{s-u}{2s_m}\right)^{-\alpha'_\rho\varkappa^2}e^{i\frac{\pi}{2}\alpha'_\rho\varkappa^2} \tag{6}$$

is the amplitude of beam production by the particle m on the vacuum pole (Fig.2b),

$$T_{12}^{1P}(s,\varkappa^2,s_1,s_2,\tau_1,\tau_2)=$$
$$\eta_\rho\varepsilon_\rho^{12}\mathcal{V}_1^\rho(s_1,\varkappa^2,\tau_1)\mathcal{V}_2^\rho(s_2,\varkappa^2,\tau_2)\left(\frac{(s-u)s_0}{2s_1s_2}\right)^{-\alpha'_\rho\varkappa^2}e^{i\frac{\pi}{2}\alpha'_\rho\varkappa^2} \tag{7}$$

is the production amplitude of two beams (Fig.2c) where

$$\varepsilon_\rho^m=\frac{\exp\left(\alpha_\rho^0\ln\frac{s-u}{2s_m}\right)}{\sqrt{s(s-4m_\rho^2)}},\ \varepsilon_\rho^{12}=\frac{\exp\left(\alpha_\rho^0\ln\frac{(s-u)s_0}{2s_1s_2}\right)}{\sqrt{s(s-4m_\rho^2)}},$$

According to Fig.1 the amplitude of two pomeron cut will have the form

$$M^{2P}(s,\varkappa^2)=M_0^{2P}(s,\varkappa^2)+M_1^{2P}(s,\varkappa^2)+M_2^{2P}(s,\varkappa^2)+M_{12}^{2P}(s,\varkappa^2), \tag{8}$$

where

$$M_i^{2P}(s,\varkappa^2)=\int_{\tau_i}\int_{s_0}^{s_{10}}T_i^{2P}(s,\varkappa^2,s_i,\tau_i)\,ds_i\,d\tau_i\ ,\ i=1,2 \tag{9}$$

$$M_{12}^{2P}(s,\varkappa^2)=\int_{\tau_1}\int_{\tau_2}\int_{s_0}^{s_{10}}\int_{s_0}^{s_{20}}T_{12}^{2P}(s,\varkappa^2,s_1,s_2,\tau_1,\tau_2)\,ds_1\,ds_2\,d\tau_1\,d\tau_2\ , \tag{10}$$

The amplitudes M_0^{2P}, T_i^{2P}, T_{12}^{2P} are calculated by formula (2) whith corresponding values of N_i taken from the expression (3)-(7). The expression (8) may be rewritten in the following form:

$$M^{2P}(s,t)=M_0^{2P}(s,t)\left[1+\frac{M_1^{2P}(s,t)+M_2^{2P}(s,t)+M_{12}^{2P}(s,t)}{M_0^{2P}(s,t)}\right]. \tag{11}$$

Let us introduce the notation

$$C_0^{P}(s,t) = 1 + \frac{M_1^{2P}(s,t) + M_2^{2P}(s,t) + M_{12}^{2P}(s,t)}{M_0^{2P}(s,t)},$$

Then taking into account the factorization of residues and expression (1) the formula for $C_0^{P}(s,t)$ takes the following form:

$$C_0^{P}(s,t) = \left(1 + \frac{M_1^{2P}(s,t)}{M_0^{2P}(s,t)}\right)\left(1 + \frac{M_2^{2P}(s,t)}{M_0^{2P}(s,t)}\right). \tag{12}$$

The magnitude $C_0^{P}(s,t)$ determining the corrections to the eikonal amplitude $M_0^{2P}(s,t)$ due to particle production in the intermediate state is called as shower factor. In the work [1] S. F have been obtained from the analysis of the experimental data on D. D available at that time. The calculated values of C_0^{P} were independent of energy and momentum transfer and appeared to be equal to

$$C_{pp}^{P} = 1.35 \pm 0.1; \quad C_{\pi p}^{P} = 1.6 \pm 0.2; \quad C_{Kp}^{P} = 1.7 \pm 0.3. \tag{13}$$

The set of experimental data on D. D processes available up to 1971, being incomplete, allowed to make some simplifying assumptions which require certain revision in view of new experimental data:

1. At that time it was assumed that the radii of the vacuum amplitudes in the differential cross sections of elastic scattering and D. D are equal each to other . In fact, as the new experimental data show the slope of pomeron contribution into D. D amplitude depends very strongly on the mass M_x of the produced particles and it is some times greater than the elastic

slope in the region $M_X \sim 1.4$GeV.

2. It was assumed that the contribution of vacuum cuts into D. D processes and elastic scattering are about the same and, therefore, in the calculation of C^P according to the formula (12) the cut contribution is simply canceled. However, thoroughly considering the D. D differential cross sections at fixed masses in the region $M_X \leq$ 1.7GeV one can see that there is a deep at $t \approx -0.3(\mathrm{GeV/c})^2$ which is conditioned by the destructive contribution of vacuum cuts. A similar deep is observed in the elastic cross sections at $t \simeq -1.5(\mathrm{GeV/c})^2$. This phenomenon is in favour of the fact that cut contribution into D. D is much larger than into elastic scattering amplitude and taking them into account one obtains values essentially differing from that of (13).

3. It was usually assumed that the contribution of the interferece terms having a $1/\sqrt{S}$ behaviour is negligible. In fact, trying to connect the data at energies up to 30GeV with the new experimental points at energies of about 400GeV these terms appear to be rather essential. Moreover, it appears necessary to take into account also the pion pole contribution having a S^{-2} behaviour.

4. The D. D and elastic scattering amplitudes were considered as pure imaginary magnitudes. This assumption is not in good agreement with the new experimental situation in the energy region up to 400GeV.

Taking into account all the above mentioned one obtains S. F numerical values much greater than the ones given by (13) and they become complex and depend strongly on energy and momentum transfer which seems to be more essential.

2. QUASIEIKONAL MODEL FOR D.D PROCESSES

According to (11) the amplitude of two pomeron cut for elastic processes takes the following form:

$$M^{2P}(s,t) = C_0^P(s,t) M_0^{2P}(s,t). \tag{14}$$

In the same way one may write for the three pomeron cut:

$$M^{3P}(s,t) = C_0^P(s,t)\, C_1^P(s,t) M_0^{3P}(s,t), \tag{15}$$

where M_0^{2P} and M_0^{3P} are the corresponding eikonal amplitudes (in this section the index "0" always corresponds to eikonal approximation), while $C_1^P(s,t)$ corresponds to the amplitude of transition of beams of particles with masses M_k and M_i into beams with masses M_j and M_ℓ (Fig.2d) integrated over all the masses.

Similarly for n -pomeron cut one obtains the following expression:

$$M^{nP}(s,t) = C_0^P(s,t)\left(C_1^P(s,t)\right)^{n-2} M_0^{nP}(s,t). \tag{16}$$

The following relation between the eikonal and pole amplitudes takes place:

$$\chi_0^P(s,b) = 2i\int_0^\infty M_0^{1P}(s,\varkappa^2)\, J_0(\varkappa b)\, \varkappa\, d\varkappa, \tag{17}$$

where b is the impact parameter and $\chi_0^P(s,b)$ is the pole eikonal.

Taking into account all the cuts in the eikonal approximation the total amplitude has the following form:

$$M_0^P(s,t)=\int_0^\infty f_0^P(s,b)\,J_0(\varkappa b)\,b\,db, \tag{18}$$

where $f_0^P(s,b)$ is the partial amplitude in the impact parameter plane. In the eikonal approximation it has the following form:

$$f_0^P(s,b)=\frac{1}{2i}\left(e^{\chi_0^P(s,b)}-1\right). \tag{19}$$

If one takes into account the showers in the intermediate state the reggeon cuts are connected with the pole eikonal (17) by the following relations:

$$M^{1P}(s,t)=\frac{1}{2i}\int_0^\infty \chi_0^P(s,b)\,J_0(\varkappa b)\,b\,db,$$

$$M^{2P}(s,t)=\frac{1}{2i}\,C_0^P(s,t)\int_0^\infty \frac{(\chi_0^P(s,b))^2}{2!}\,J_0(\varkappa b)\,b\,db,$$

$$M^{3P}(s,t)=\frac{1}{2i}\,C_0^P(s,t)\,C_1^P(s,t)\int_0^\infty \frac{(\chi_0^P(s,b))^2}{3!}\,J_0(\varkappa b)\,b\,db, \tag{20}$$

. . . .

$$M^{nP}(s,t)=\frac{1}{2i}\,C_0^P(s,t)\left(C_1^P(s,t)\right)^{n-2}\int_0^\infty \frac{(\chi_0^P(s,b))^n}{n!}\,J_0(\varkappa b)\,b\,db.$$

For the convinience of numerical computations one may write the sum of the terms (20) in a compact form allowing to take into account all the reggeon cuts. Then the total amplitude taking into account all cuts will have the form:

$$M^P(s,t)=\int_0^\infty F^P(s,y^2,b)\,J_0(yb)\,b\,db\,\Big|_{y^2=t}, \tag{21}$$

where

$$F^{P}(s,y^2,b)=$$
$$\frac{1}{2i}\left\{\frac{C_0^P(s,y^2)}{(C_1^P(s,y^2))^2}\left[e^{C_1^P(s,y^2)\chi_0^P(s,b)}-1-C_1^P(s,y^2)\chi_0^P(s,b)\right]+\chi_0^P(s,b)\right\}. \quad (22)$$

The partial amplitude will take the form:

$$f(s,b)=\int_0^\infty\left(\int_0^\infty F^P(s,y^2,b)\,J_0(yb)\,b\,db\right)_{y^2=t} J_0(\varkappa b)\,\varkappa\,d\varkappa . \quad (23)$$

For D. D processes the pole amplitude has the form (6) and the pole eikonal corresponding to T^{1P} is equal to:

$$\chi_1^P(s,M_x,b)=2i\int_0^\infty T^{1P}(s,t,M_x)\,J_0(\varkappa b)\,\varkappa\,d\varkappa . \quad (24)$$

Then taking into account S. F the amplitudes of pole and corresponding cuts will have the following form:

$$T^{1P}(s,t,M_x)=\frac{1}{2i}\int_0^\infty \chi_1^P(s,b,M_x)\,J_0(\varkappa b)\,b\,db,$$

$$T^{2P}(s,t,M_x)=\frac{1}{2i}\,C^P(M_x,s,t)\sqrt{C_0^P(s,t)}\ \times$$
$$\int_0^\infty \chi_1^P(s,b,M_x)\,\chi_0^P(s,b)\,J_0(\varkappa b)\,b\,db,$$

$$T^{3P}(s,t,M_x)=\frac{1}{2i}\,C^P(M_x,s,t)\sqrt{C_0^P(s,t)}\,C_1^P(s,t)\ \times$$
$$\int_0^\infty \chi_1^P(s,b,M_x)\,\frac{(\chi_0^P(s,b))^2}{2!}\,J_0(\varkappa b)\,b\,db,$$

. . . .

$$T^{nP}(s,t,M_x)=\frac{1}{2i}\,C^P(M_x,s,t)\sqrt{C_0^P(s,t)}\left(C_1^P(s,t)\right)^{n-2}\ \times$$
$$\int_0^\infty \chi_1^P(s,b,M_x)\,\frac{(\chi_0^P(s,b))^n}{n!}\,J_0(\varkappa b)\,b\,db .$$

As in the case of expression (21) one may write this sum in a more compact form convinient for computation:

$$T^{P}(s,t,M_x)=\int_0^\infty F^{P}_{DD}(s,y^2,M_x,b)\,J_0(yb)\,b\,db\,\Big|_{y^2=t}\,, \tag{25}$$

where

$$F^{P}_{DD}(s,y^2,M_x,b)=\frac{1}{2i}\chi_1^{P}(s,b)\left\{\frac{C^{P}(M_x,s,y^2)\sqrt{C_0^{P}(s,y^2)}}{C_1^{P}(s,y^2)}\left(e^{C_1^{P}(s,y^2)\chi_0^{P}(s,b)}-1\right)+1\right\}. \tag{26}$$

The partial amplitude for D. D processes takes the form:

$$f_{DD}(s,b)=\int_0^\infty\left(\int_0^\infty F^{P}_{DD}(s,y^2,M_x,b)\,J_0(yb)\,b\,db\right)_{y^2=t} J_0(\varkappa b)\,\varkappa\, d\varkappa\,. \tag{27}$$

The coefficients $C^{P}(M_x,s,t)$ and $C_1^{P}(s,t)$ are connected by the relation:

$$C_1^{P}(s,t)=\int_{S_{01}}^{S_1} C^{P}(M_x,s,t)\,dM_x\,, \tag{28}$$

where $S_{01}=(m_p+m_\pi)^2$ and S_1 determine the lower boundary of the three reggeon region.

3. PHENOMENOLOGICAL DESCRIPTION OF D.D PROCESSES

In the last years the results of many experimental works devoted to the study of D. D processes in reaction $p+p\to p+X$ in a wide energy interval, from 6 up to 400GeV, and in the region of small momentum transfers and produced masses have been published [6-9]. Taking into account the cuts the amplitude of D. D processes will have the form shown in Fig.3. The corresponding differential cross sections are given in Fig.4.

Since reliable experimental data exist only for the process $pp\to pX$ we shall further consider only this reaction.

Besides, we shall consider only the spin non flip amplitude. The vacuum exchange amplitude for elastic pp-scattering has the following form:

$$M^{1P}(s,t)=\eta_P\varepsilon_P\left(g^P_{pp}(t)\right)^2\left(\frac{s-u}{2s_0}\right)^{\alpha'_P t}e^{-i\frac{\pi}{2}\alpha'_P t}. \tag{29}$$

Averaging the dependence on τ over which one always carries out integration in the final expressions for the vacuum amplitude of D. D processes we recieve the following :

$$T^{1P}(s,t)=\eta_P\varepsilon^x_P g^P_{pp}(t)\,\widetilde{\mathcal{V}}_P(t,M_x)\left(\frac{s-u}{2s_x}\right)^{\alpha'_P t}e^{-i\frac{\pi}{2}\alpha'_P t} \tag{30}$$

The magnitudes entering into (30) are defined in section I.

The vertex $\widetilde{\mathcal{V}}_P(t,M_x)$ of the scattering of pomeron with mass t on proton has been parametrized in the following way:

$$\widetilde{\mathcal{V}}_P(t,M_x)=A(M_x)\,e^{-R^2(M_x)\varkappa^2}. \tag{31}$$

According to the work [10] the pp-scattering vacuum pole residue must have the form:

$$g^P_{pp}(t)=a\left(e^{-R_1^2\varkappa^2}+\varkappa^2 B\,e^{-R_2^2\varkappa^2}\right). \tag{32}$$

Then the elastic pp-scattering amplitude will be:

$$M^{1P}(s,t)=\eta_P\varepsilon_P a^2\left(e^{-\lambda_1\varkappa^2}+2\varkappa^2 B\,e^{-\lambda_{12}\varkappa^2}+\varkappa^4B^2e^{-\lambda_2\varkappa^2}\right), \tag{33}$$

while the D. D amplitude becomes:

$$T^{1P}(s,t,M_x)=\eta_P\varepsilon^x_P a\,A(M_x)\left(e^{-\lambda_{x1}\varkappa^2}+\varkappa^2B\,e^{-\lambda_{x2}\varkappa^2}\right), \tag{34}$$

where

$$\lambda_i=2R_i^2+\alpha'_P\left(\ln\frac{s-u}{2s_0}-i\frac{\pi}{2}\right);\quad \lambda_{xi}=R_i^2+R^2(M_x)+\alpha'_P\left(\ln\frac{s-u}{2s_x}-i\frac{\pi}{2}\right);$$

$$\lambda_{12} = R_1^2 + R_2^2 + \alpha_P' \left(\ln \frac{s-u}{2s_0} - i\frac{\pi}{2} \right) ; \quad i = 1,2$$

The parameters of the elastic pp-scattering secondary trajectories P' and ω have been taken from the work [10]. The vacuum pole parameters for elastic scattering and D. D processes have been derived from the simultaneous comparison with the experimental data on these two types of reactions. The values of pomeron parameters are given in Table I.

The total amplitudes taking into account all the cuts have been calculated in quasieikonal approximation according to formulae (21) and (25). Finally, the D. D differential cross section may be presented in the form:

$$\frac{d^2\sigma}{dt\,dM_x} = 4\pi(\hbar c)^2 \left(|T^P(s,t,M_x)|^2 + \frac{B(M_x)}{\sqrt{s}} e^{\lambda_{x3} t} + \frac{D(M_x)}{s} e^{\lambda_{x4} t} \right) + \left(\frac{d^2\sigma}{dt\,dM_x} \right)_{\pi\pi P}, \tag{35}$$

where

$$\lambda_{x3,4} = R_x^2(M_x) + 2\ln\frac{s-u}{2s_x},$$

$B(M_x)$ corresponds to the interference of vacuum pole with secondary trajectories, while $D(M_x)$ corresponds to the contribution of secondary trajectories with $\alpha^0 = 0.5$. The last term corresponds to the one pion exchange contribution and is equal to [11] :

$$\left(\frac{d^2\sigma}{dt\,dM_x} \right)_{\pi\pi p} = \frac{g^2 \sigma_{tot}^{\pi N}(M_x^2)}{(4\pi)^2} \frac{2M_x^3 \varkappa^2}{s^2 (m_\pi^2 + \varkappa^2)^2} \left(\frac{s}{M_x^2} \right)^{-\alpha_\pi' \varkappa^2} e^{-R^2 \varkappa^2}, \tag{36}$$

where $g^2/4 = 15$, $R^2 = 3.3\ (\mathrm{GeV}/c)^{-2}$, $\alpha_\pi' = 1\ (\mathrm{GeV})^{-2}$;

Using the notations of formula (34) the zero mass pomeron scattering total cross section on proton may be presented in the form:

$$\sigma_{pp}^{tot}(M_x,0) = 4\pi(\hbar c)^2 \frac{1}{M_x^3} A^2(M_x). \qquad (37)$$

(As a matter of fact one must calculate all the integrals over in formula (17) in the interval from up to and not from 0. This problem has been studied in detail and it has been shown that at energies of order of 10 GeV the difference between the results of integrations from 0 and from t_{min} is $\sim 1\%$, while at energies higher than 20 GeV and $M_x <$ 5 GeV this difference is negligible).

Thus, the problem is to determine the magnitudes $A(M_x)$, $R^2(M_x)$, $C(M_x)$, $B(M_x)$, $D(M_x)$ and $R_x^2(M_x)$ for each given value of M_x.

4. DISCUSSION OF THE RESULTS

The obtained results are given in Figs. 5, 8–10. The parameter $R_x^2(M_x)$ is determined with large errors, therefore, as a value of this parameter we have taken 2.5 $(GeV/c)^{-2}$, the average of its values for all M_x. From the existing experimental data one can not determine the energy and momentum transfer dependence of the parameter $C(M_x,s,t)$, therefore, it has been assumed $C(M_x,s,t) = C(M_x)$. Since in the elastic scattering this parameter is essential only for nP-cuts at $n \geqslant 3$ this assumption has very little effect on the final result. The obtained parameters describe well all the set of the experimental data on D. D processes and furthermore they provide good results for elastic pp-scattering, in particular, they describe well the growth of the total cross sections as well as the differential cross sections up to $t \lesssim -2(\mathrm{GeV/c})^2$.

In Fig. 5a it is shown the pomeron proton total cross section $\sigma_{Pp}^{tot}(M_x, 0)$ in the region of small masses, calculated by formula (37) a) without taking and b) taking into account the cuts. As it is seen from figure quite different results are obtained in the cases a) and b), the results being different not only by magnitude (in the case b) σ_{Pp}^{tot} is much larger) but also by their structure. In the case a) there are two peaks at 1.4 and 1.7 GeV, while in the case b) there is only a wide maximum at $M_x \approx 1.7$ GeV. On the boundary of the three reggeon region the cross section $\sigma_{Pp}^{tot}(M_x > 2.5)$ is about 1mb and 3mb in the first and second cases, respectively, i.e. when one takes into account the cuts the effective three pomeron vertex G_{PPP} must be three times larger than in the pole case. This fact is in good agreement with the result of the work [15]. In the normalization of the works [12-14] such a σ_{Pp}^{tot} corresponds to a three pomeron vertex equal to $G_{PPP} = (7.2 \pm 0.2)$ mb/GeV^2 instead of 3.24 in [13].

In the three reggeon region σ_{Pp}^{tot} consists of the sum of the following terms (see Fig. 6). We call the sum of the diagrams 6a), 6b), 6c) and 6d) as effective three pomeron vertex. Therefore, the pure three pomeron vertex 6a) indeed will be much larger than 7.2 mb/GeV^2. Unfortunately, one can not obtain a correct estimate for the diagrams 6b), 6c), 6d) using the existing experimental data. Therefore, we further shall consider only the effective three pomeron vertex. However, using the available experimental results one can obtain the magnitude of the diagram 6e) which corresponds to the three reggeon vertex G_{PPR} . The last magnitude appeared to be small and G_{PPR} = =(1±0.5) mb/GeV^2 which is in good agreement with the previous analyses [12-14] . At large masses, $M_x \sim 10$ GeV, non-diffractive processes resulting in a rapid growth of $d^2\sigma/dtdM_x$ begin to admix to the diffraction processes. In the work [16] it has been shown that one can obtain a differential cross section growth in agreement with the experimental data by adding a background term of the form:

$$\frac{d^2\sigma}{dtdM_x}(\text{фон}) = \frac{(\hbar c)^2}{M_x S} \cdot 550 \left(\frac{M_x}{S}\right)^{0.1-1.5|t|} \qquad (38)$$

to the effective three pomeron vertex.

In Fig.5b) it is shown the square of the radius $R^2(M_x)$ of the vertex of diffraction particle production in the region of small masses. It depends strongly on the mass of produced particles, but it goes to 1 when $M_x > 2$ GeV. Fig. 5c) shows the shower factor $C(M_x)$ for diffraction production. It depends strongly on the mass M_x . However, at large values of M_x it goes to 1 as $1/M_x^{1/2}$ .This fact shows that in the region of large masses the contributions of the strengthened diagrams must be negligible compared with that of eikonals. Thus, there is no necessity to take into account the diagrams 7a), 7b) and 7c) and it is quite enough to consider only the diagrams of the type 7d).

Fig. 8 shows the interference term $B(M_x)$ which is the sum of the interference terms of ω and P' -trajectories with pomeron. In the region of small masses $B(M_x)$ is negative and depends strongly on energy. This fact shows that ω -pole is somewhat dominant compared with P'-pole. In the normalization of the work [13] the term $B(M_x)$ corresponds to the sum of three reggeon terms $2Re(G_{RPP} + 1/M_x\, G_{PRR})$. Since $B(M_x)=0$ at $M_x \approx 2.5$ GeV one can separate the G_{RPP} and G_{PRR} contributions each from other. They appeared to be equal to $2Re\, G_{RPP} = (7.2\pm2)$ mb/GeV^2 and $2Re\, G_{PRR} = (-17.3\pm3)$ mb/GeV^2, respectively. Only P' -pole can give contribution to the vertex G_{RPP} , while ω and P' give contrition to G_{PRR} . The ω -pole contribution is significantly larger than the P'-contribution as a result of which the constant G_{PRR} is negative.

The dependence of the term $D(M_x)$ on the mass of produced particles is given in Fig. 9. In the region of small masses $D(M_x)$ repeats on the whole the behaviour of πN-scattering cross section in the resonance region. It is seen the sharply expressed peaks in the mass region $M_x \sim$ 1236, 1500, 1700, 1800 and 2000 MeV. This fact is in favour of the assumption (following from the quark model) that f, ω, ρ and A_2 mesons are like π -meson by their main features and the fact that they are off-mass-shell gives small effect on their properties. In the region of large masses the sum $G_{RRR} + M_x\, G_{RRP}$ of three region constants corresponds to the term $D(M_x)$. One can easily calculate them using the available data and in the normalization [13] they are equal to $G_{RRR} = (62\pm8)$ mb/GeV^2 and $G_{RRP} = (8.5\pm2)$ mb/GeV^2. All the obtained values of the three reggeon constants in the normalization [13] are given in Table2.

Fig. 10 presents the differential cross sections $d^2\sigma/dtdM_X$ of D. D processes in the region of small masses and momentum transfers at energies 15.1; 270 and 400 GeV. As it is seen from figure there is a sufficiently good description of experimental data which witnesses the reliability of the obtained values of the parameters. At energies which are available today the peak in the differential cross sections is observed at $M_X \gtrsim$ 1.4 GeV, however, it must be shifted toward $M_X \approx 1.7$ GeV with the further increase of energy. This effect is conditioned by the fact that the pomeron proton cross section has a peak in the region $M_X \gtrsim$ 1.7 GeV.

In Fig. 11 it is shown the dependence of the slope of the differential cross section $d^2\sigma/dtdM_X$ on the mass of the produced particles at low momentum transfers. Though $R^2(M_X)$ is relatively smooth function of M_X the contribution of the cuts results in the fact that $b(M_X)$ becomes a function with sharp maxima and minima in good agreement with the existing experimental data.

5. ELASTIC SCATTERING

In the previous section a good description of D. D processes has been obtained with a sufficiently reliable parametrization of the vacuum amplitude of these processes. Now, let us consider briefly how the obtained results influence the elastic NN-scattering amplitude.

Since now we have the pomeron proton scattering vertex it becomes possible to calculate correctly the diagrams of Fig.1b), 1c) and 1d). Therefore, one can now calculate the S. F according to the formula (12). The three-dimension picture of Fig. 12 shows the dependence of the real part of the obtained S. F on energy and momentum transfer. As it is seen from figure for a fixed momentum transfer in the beginning $Re\, C_0^P$ increases rapidly with the increase of energy, achieves a maximum at $E \approx 50$ GeV, then begins to decrease slowly. Our value of $Re\, C_0^P$ at t =0 is much greater than that of the work [1] because taking into account the cuts the pomeron proton scattering vertex obtained by us appeared to be three times greater than that in pole model. Fixing the energy and varying t $Re\, C_0^P$ begins to decrease with the increase of t and it becomes even less than 1 when $t > 1$. This effect is conditioned by the fact that $Jm\, M_1^{2P}$ changes its sign at $t \approx 1$ earlier than $Jm\, M_0^{2P}$ and the amplitude ratio M_1^{2P}/M_0^{2P} in formula (12) becomes negative.

The obtained energy dependence of $Re\, C_0^P$ influences the total cross section. It results in the fact that the two pomeron cut having a negative sign is amplified differently at various energies. The amplification achieves a maximum in the region of the total cross section minimum and begins to decrease when the energy increases. This decrease together with the logarithmic decrease of the cut contribution results in a sharp growth of the total cross section in good agreement with existing experimental data (see Fig.13). The total cross sections insrease rapidly up to $E \approx 5000$ GeV, then the rate of the increase becomes smaller and the cross sections begin to go logarithmically to their asymptotical value.

Thus, it is possible to make conclusion that taking into account correctly the energy and momentum transfer dependence of S. P one obtains a rapid growth of cross sections even with $\alpha_P^0 = 1$ without making various additional assumptions, e.g. $\alpha_P^0 > 1$.

In conclusion the author expresses his gratitude to A.Ts. Amatuni for guidance and constant interest to the work and S.G. Matinian for useful discussions.

TABLE I

The Values of the Elastic pp-Scattering Vacuum Pole Parameters

$\alpha^2 = (7.131 \pm 0.05)\ \mathrm{GeV}^{-2}$

$B = (0.55 \pm 0.05)\ \mathrm{GeV}^{-4}$

$R^2_{0P} = (1.7 \pm 0.1)\ \mathrm{GeV}^{-2}$

$R^2_{1P} = (0.7 \pm 0.2)\ \mathrm{GeV}^{-2}$

$\alpha_P(t) = 1 + 0.5t$

TABLE II

The Three Reggeon Vertices in the Normalization [13]

Parameter	G_{ijk} mb/GeV^2	R^2_{ijk} (GeV)$^{-2}$
G_{PPP}	7.2 ± 0.2	5.4 ± 0.5
G_{RRP}	8.5 ± 2	5 ± 2
$2\mathrm{Re}\, G_{RPP}$	7.2 ± 2	5 ± 2
G_{PPR}	11.1 ± 0.5	5 ± 1
G_{RRR}	62 ± 8	5 ± 1
$2\mathrm{Re}\, G_{PRR}$	-17.3 ± 3	5 ± 2

FIGURE CAPTIONS

Fig.1. Two reggeon cut.

Fig.2. Main pole diagrams giving contribution to reggeon cuts.

Fig.3. D. D process amplitudes.

Fig.4. D.D differential cross section.

Fig.5. Dependece of the main parameters of vacuum amplituae on the mass of produced particles;

a) Pomeron proton scattering total cross section.

o------- pole model,

•_______ model taking into account the cuts,

b) Dependence of the pomeron amplitude radius on M_X.

c) Dependence of the shower factor $C(M_X)$ on M_X.

Fig.6. Effective three pomeron vertex.

Fig.7. Corrections to the effective three pomeron vertex.

Fig.8. Dependence of the interference term $B(M_X)$ on M_X.

Fig.9. Dependence of the term $D(M_X)$ on the mass of produced particles in the resonance region.

Fig.10. Differential cross section of the process $pp \to pX$ in the region of small produced masses and momentum transfers at energies 15.1; 270 and 400GeV.

Fig.11. Dependence of the slope of the differential cross section for the process $pp \to pX$ on the mass of produced particles.

Fig.12. Dependence of $Re\, C_0^{P}(s,t)$ on energy and momentum transfer.

Fig.13. Total cross sections of pp and $\bar{p}p$ -scattering.

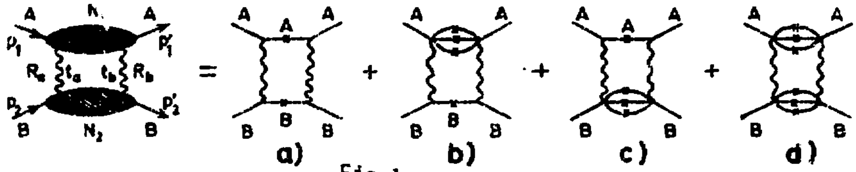


Fig.1

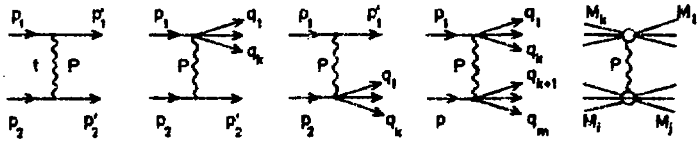


Fig.2

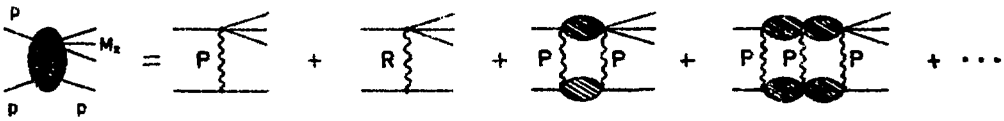
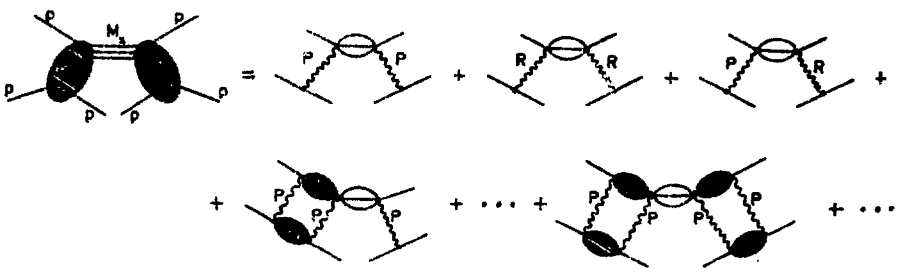


Fig.3



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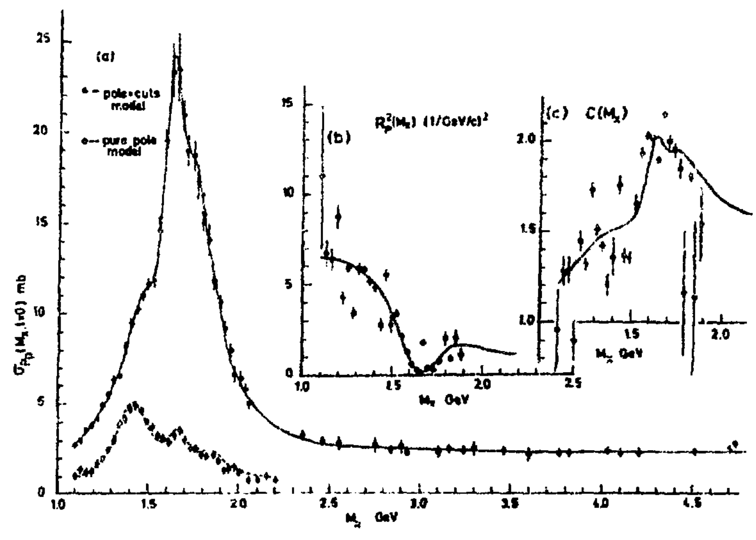

Fig.5

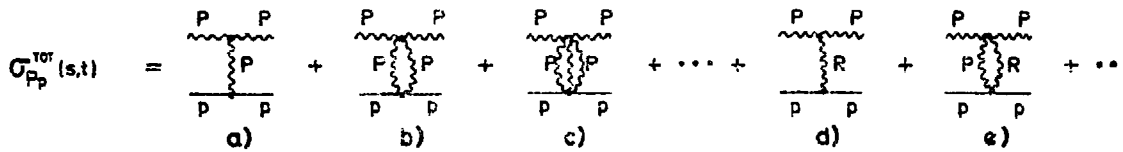

Fig.6

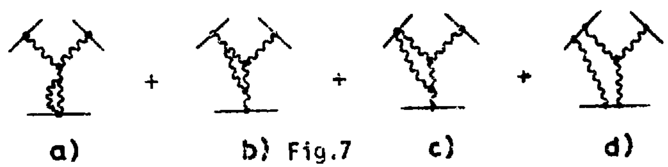

Fig.7

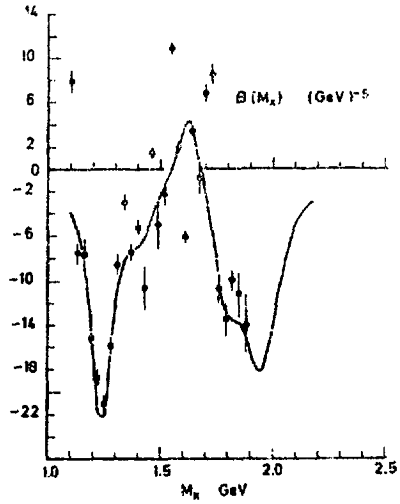
14
10
6
2
0
-2
-6
-10
-14
-18
-22
B(M_x) (GeV)^-5
1.0
1.5
2.0
2.5
M_x GeV

Fig.8

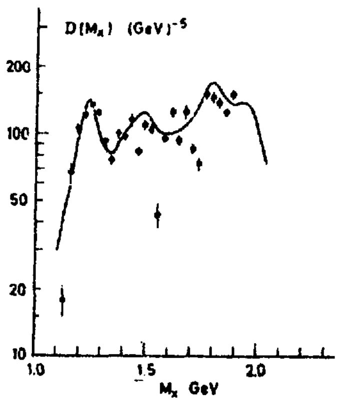
D(M_x) (GeV)^-5
200
100
50
20
10
1.0
1.5
2.0
M_x GeV

Fig.9

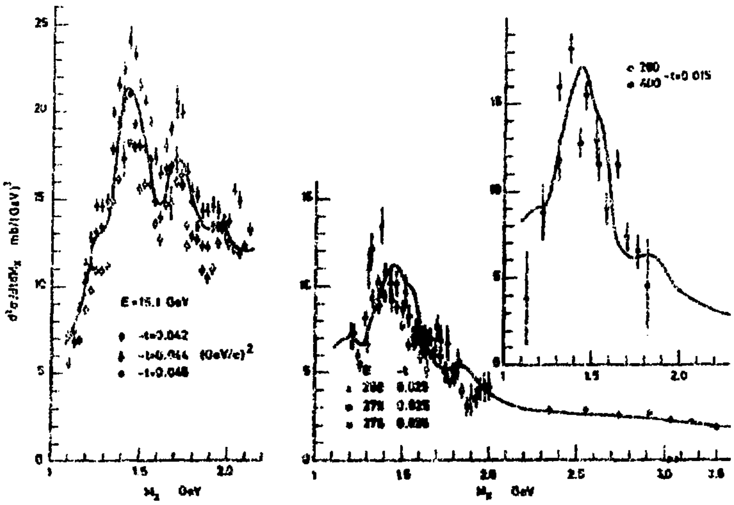

Fig. 10

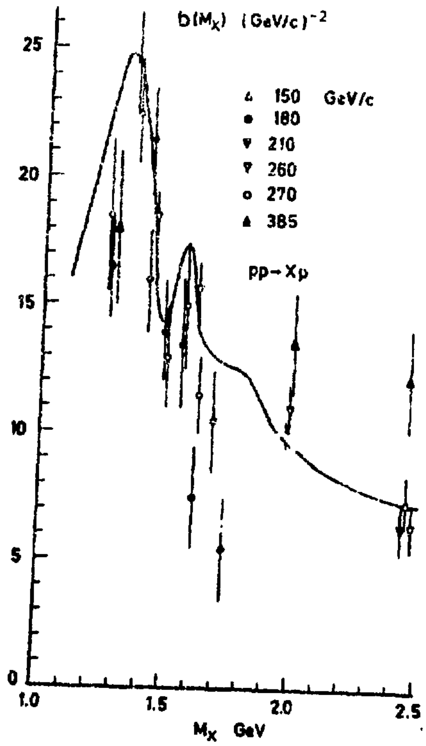

Fig.11

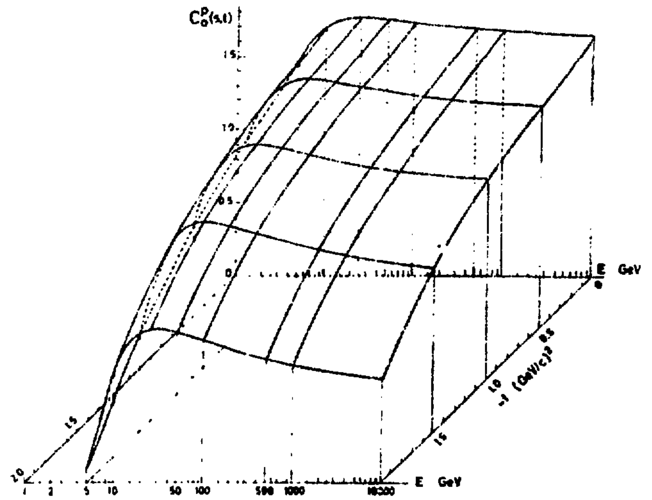

Fig.12

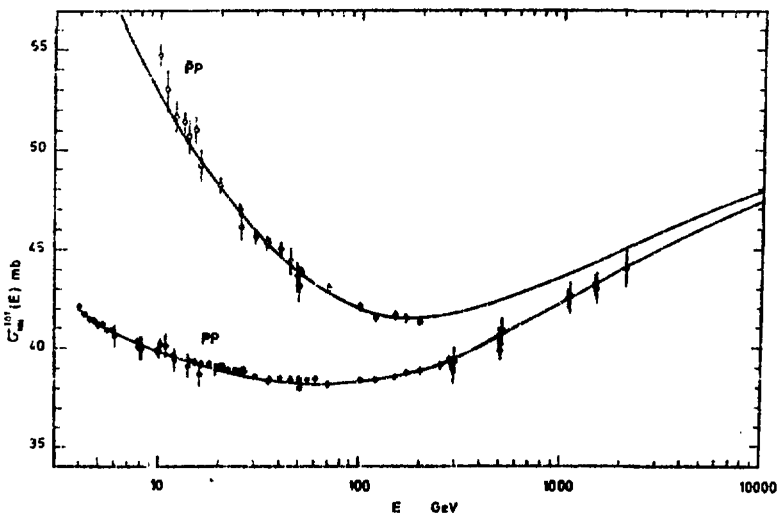

Fig.13

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ДИФРАКЦИОННОЕ РАССЕЯНИЕ И ЗАВИСИМОСТЬ КОЭФФИЦИЕНТОВ ЛИВНЕВОГО УСИЛЕНИЯ ОТ ЭНЕРГИИ И ПЕРЕДАННОГО ИМПУЛЬСА

(на английском языке)

Ереванский физический институт

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