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A.A.GRIGORYAN, A.B.KAIDALOV, G.N.KHACHATRYAN

DISPERSION SUM RULES AND RESIDUES OF ISOSCALAR REGGEONS

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А.А.ГРИГОРЯН, А.Б.КАЙДАЛОВ*), Г.Н.ХАЧАТРЯН

ДИСПЕРСИОННЫЕ ПРАВИЛА СУММ И ВЫЧЕТЫ ИЗОСКАЛЯРНЫХ РЕДЖЕОНОВ

Сверхсходящиеся правила сумм для амплитуд рассеяния реджеонов на частицах применяются к исследованию спиральных вершин связи бозонных реджеонов α_S с изоспином I =0 ($S=\mathbb{P}$, f и ω) с барионами N и Δ_{33}. Предсказания для вершин $N\alpha_S N$ согласуются с экспериментальными данными. Результаты обобщены на случай связи α_S с произвольными частицами a и b.

Ереванский физический институт

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*) Институт теоретической и экспериментальной физики.

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A.A.GRIGORYAN, A.B.KAIDALOV*), G.N.KHACHATRYAN

DISPERSION SUM RULES AND RESIDUES OF ISOSCALAR REGGEONS

The helicity structure of the vertices $N\alpha_i^s N$ and $\Delta_{33}\,\alpha_i^s\,\Delta_{33}$ (α_i^s are the boson reggeons with isospin I=0; $S_i = \mathbb{P}, f, \omega$) is investigated using the superconvergent sum rules for the reggeon-particle scattering amplitudes. The predictions for the vertex $N\alpha_i^s N$ are in agreement with experiments. The results are generalized to the case of couplings of α_i^s with arbitrary particles.

Yerevan Physics Institute

Yerevan 1980

*) Institute of Theoretical and Experimental Physics, Moscow

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YEREVAN PHYSICS INSTITUTE

A.A.GRIGORYAN, A.B.KAIDALOV, G.N.KHACHATRYAN

DISPERSION SUM RULES AND RESIDUES OF ISOSCALAR REGGEONS

Yerevan 1980

1. Introduction

The dispersion sum rules for the reggeon-particle scattering amplitudes (αa -amplitudes) play an important role in the theoretical investigation of the hadron interactions at high energies. The superconvergent sum rules (SSR) which allow to get a number of relations between different Regge residues are of special interest.

The general method of extracting SSR for the scattering of reggeons on particles with arbitrary spins, based on the helicity amplitudes in the infinite momentum frame (IMF), has been proposed in the previous papers [1,2]. It has been shown that two types of the SSR exist:

a) The SSR which are valid when $\alpha_e - \alpha_i - \alpha_k < -1$ where α_i and α_k are the scattered trajectories; α_e is the rightmost singularity in the j-plane with the given t-channel quantum numbers. These SSR do not depend on spins of particles.

b) The number of the superconvergent amplitudes increase substantially when the spins of external particles are taken into account, and the additional SSR which are valid if $\alpha_e - \alpha_i - \alpha_k - m < -1$ $(m = 1, 2 \ldots)$ arise. In the papers [2,3] the formalism of SSR has been used to investigate the helicity couplings of I=1 (ρ, A_2 π) boson reggeons with baryons N (938) and Δ_{33} (1232). It has been shown that this approach allows one to obtain the relations which are in good agreement with experimental data.

The predictive power of SSR has been demonstrated in Refs. [4,5], where the existence of the whole series of exotic baryon resonances with isospins I = 5/2 and spins $S = I$ was predicted.

The properties of the first resonance from this series with I = 5/2 (called E_{55}) have been analysed in detail. It has been shown that SSR allow one to determine unambiguously the spin (S = 5/2), P-parity ($\eta_E = +$) and helicity vertices $G^{a\alpha_i E}_{\lambda a \lambda_E}$ ($a = \Delta_{33}, E_{55}$; $i = \rho, A_2 \pi$) of the resonance E_{55}.

In this work we apply the method of SSR to the investigation of I=0 boson reggeons ($\mathbb{P}$, f, ω) couplings with the nucleon N(938) and Δ_{33}(1232)-isobar.

In Sec.2 we give the expressions for the contribution of resonances to the sum rules and analyse the asymptotic behaviour of the sum rules in the case when at least one of the scattered trajectories has isospin I=0.

In Sec.3 we consider the SSR for the processes $\alpha^s a \to \alpha^{s,v} b$, where α^s are isoscalar ($s = \mathbb{P}, f, \omega$) trajectories and α^v are isovector ($V = \rho, A_2$) ones;

Saturation of these SSR by the contributions of N and Δ_{33} $S(u)$ channel states leads to a set of equations for helicity couplings $G^{a d^s a}_{\lambda_a \lambda'_a}$ $(a = N, \Delta_{33})$ which has a non-trivial solution. The results are generalized to the case of arbitrary particles a and b . The relation between the spin-flip residue of $\mathbb{P}$ - $G^{N d \mathbb{P} N}_{1/2\, -1/2}$ and vertices of inelastic diffraction is established. The restriction on the quantity $G^{N d \mathbb{P} N}_{1/2\, -1/2}$ is found.

In Sec.4 we discuss the obtained results.

2. Superconvergent Sum Rules and Scattering of Trajectories with I = 0.

The superconvergent dispersion relations for the amplitude $T^{d_i d_\kappa}_{ab}(q_i^2, q_\kappa^2, q^2, \nu)$ of the $d_i a \to d_\kappa b$ -scattering have the form

$$\int_0^\infty d\nu\, \mathrm{Im}\, T^{d_i d_\kappa (-)}_{ab}(\nu, q_i^2, q_\kappa^2, q^2) = 0 \tag{2.1}$$

where $\nu = (S - U)/4M_N$, $T^{d_i d_\kappa (-)}_{ab}$ is the crossing-odd amplitude of da -scattering,

$$T^{d_i d_\kappa (-)}_{ab} = T^{d_i d_\kappa (s)}_{ab} - T^{d_i d_\kappa (u)}_{ab}$$

For the particles (a, b) with spin it is convenient to construct the SSR using the helicity amplitudes in the IMF. The reggeon vertices and contribution of resonances to the sum rules for helicity amplitudes have a simple form in the IMF.

The particle-reggeon-particle helicity vertex in the IMF has the form

$$\Gamma^{a\alpha b}_{\lambda_a \lambda_b}(q) = q^{\perp|\lambda_a - \lambda_b|} e^{i\varphi(\lambda_a - \lambda_b)} G^{a\alpha b}_{\lambda_a \lambda_b}(q^{\perp 2}) \tag{2.2}$$

where $\vec{q}^{\perp}$ is the transverse component of the momentum transfer $q = p_b - p_a$; φ is the angle between $\vec{q}^{\perp}$ and X-axis.

The residues $G^{a\alpha b}_{\lambda_a \lambda_b}(q^2)$ have no kinematic singularities in $q^{\perp 2}$. We shall consider in what follows the region of small ($q^{\perp 2}$) and discuss the quantities $G^{a\alpha b}_{\lambda_a \lambda_b}(0)$. P- and $G(T)$-invariances impose the constraints on the couplings

$$G^{a\alpha b}_{\lambda_a \lambda_b} = \sigma P \eta_a \eta_b (-1)^{S_a - S_b - \lambda_a + \lambda_b} G^{a\alpha b}_{-\lambda_a - \lambda_b} \tag{2.3}$$

$$G^{a\alpha b}_{\lambda_a \lambda_b} = G(-1)^{I} P(-1)^{\lambda_a - \lambda_b} \tilde{G}^{b\alpha a}_{\lambda_b \lambda_a} \tag{2.4}$$

where σ, P, G and I are the signature, P-parity, G-parity and isospin of α (nonstrange boson reggeons are considered); η_i and S_i is the P-parity and spin of particle i.

The coupling $\tilde{G}^{b\alpha a}_{\lambda_b \lambda_a}$ describes the inverse transition $b \to a$.

If $\alpha_e - \alpha_i - \alpha_k < -1$, then the following crossing-odd SSR which connect the residues $G^{c\alpha f}_{\lambda_c \lambda_f}$ take place

$$\sum_{I_s} x_{ts} \sum_{d_s} G^{a d_i d_s}_{\lambda a\ \lambda a+n}{}^{(0)} \quad G^{d_s d_\kappa b}_{\lambda a+n\ \lambda b} -$$

$$-\sum_{I_u} \chi_{tu} \sum_{d_u} G^{a d_\kappa d_u}_{\lambda a\ \lambda b-n}{}^{(0)} \quad G^{d_u d_i b}_{\lambda b-n\ \lambda b}{}^{(0)} = 0 . \tag{2.5}$$

SSR (2.5) correspond to definite t-channel isospin state I_t ; $X_{ts} \equiv X(I_t, I_s)$ ($X_{tu} \equiv X(I_t, I_u)$ is the isotopic crossing matrix which projects the state with the definite s(u)-channel isospin $I_s(I_u)$ onto the state with the given t-channel isospin I_t .

In (2.5) the quantities $G^{c d_j f}_{\lambda_c \lambda_f}$ are the reduced residues They are connected with the physical residues in the following way

$$\left(G^{c d_j f}_{\lambda_c \lambda_f}\right)^{phys} = \varepsilon_c\, \varepsilon_j\, \varepsilon_f\, (-1)^{I_j - I_c + m_f} \begin{pmatrix} I_j & I_c & I_f \\ m_j & m_c & -m_f \end{pmatrix} G^{c d_j f}_{\lambda_c \lambda_f} \tag{2.6}$$

where $I_p\,(m_p)$ is the isospin (3-rd projection of isospin) of particle p ; ε_p is the phase connecting the particle state with the basic isospin state; $\begin{pmatrix} I_1 & I_2 & I_3 \\ m_1 & m_2 & m_3 \end{pmatrix}$ is the $3jm$-symbol.

In (2.5) the quantity n is integer and satisfies the condition

$$\min(|\lambda_a|, |\lambda_b|) \le |\lambda_a + n| \ (|\lambda_b - n|) \le \max(|\lambda_a|, |\lambda_b|) \ . \tag{2.7}$$

SSR (2.5) are valid for the amplitudes $T^{\alpha_i \alpha_k (\sigma_e P_e)}_{\lambda a \, \lambda b}$ with fixed naturality $\sigma_e P_e$ at $\tau \equiv (\sigma_i P_i)(\sigma_K P_K)(\sigma_e P_e) = +1$. If $\tau = -1$, the SSR connecting $G^{c\alpha j f}_{\lambda c \, \lambda f}$ take place only when $\lambda_a = \lambda_b$:

$$\sum_{I_s} \chi_{ts} \sum_{d_s} \left\{ G^{a d_i d_s}_{\lambda a \, \lambda a-1}(0) \; G^{d_s d_K b}_{\lambda a-1 \, \lambda a}(0) - G^{a d_i d_s}_{\lambda a \, \lambda a+1}(0) G^{d_s d_K b}_{\lambda a+1 \, \lambda a}(0) \right\} + \tag{2.8}$$

$$+ \sum_{I_u} \chi_{tu} \sum_{d_u} \left\{ G^{a d_K d_u}_{\lambda a \, \lambda a-1}(0) \; G^{d_u d_i b}_{\lambda a-1 \, \lambda a}(0) - G^{a d_K d_u}_{\lambda a \, \lambda a+1}(0) \; G^{d_u d_i b}_{\lambda a+1 \, \lambda a}(0) \right\} = 0$$

Apart from the SSR (2.5) and (2.8), which are valid if $\alpha_e - \alpha_i - \alpha_K < -1$, for the particles with spins, the additional sum rules (the so called "obligatory" superconvergent sum rules (OSSR)) arise. They are valid under weaker restrictions on the position of α_e in the j-plane, i.e. $\alpha_e - \alpha_i - \alpha_K < 0$ [1,2]. Saturating the OSSR by the resonances one can get the following equations

$$\left\{ \sum_{I_s} \chi_{ts} \sum_{d_s} G^{a d_i d_s}_{\lambda a \, \lambda a+n}(0) \quad G^{d_s d_K b}_{\lambda a+n \, \lambda b}(0) \right. -$$

$$\left. \sum_{I_u} \chi_{tu} \sum_{d_u} G^{a d_K d_u}_{\lambda a \, \lambda b-n}(0) \; G^{d_u d_i b}_{\lambda b-n \, \lambda b}(0) \right\} \Big/ \binom{|n|}{|\lambda a - \lambda b|} = \tag{2.9}$$

$$= \left\{ \sum_{I_s} \chi_{ts} \sum_{d_s} G^{a d_i d_s}_{\lambda a \, \lambda a+n'}(0) \quad G^{d_s d_K b}_{\lambda a+n' \, \lambda b}(0) \right. -$$

$$\left. \sum_{I_u} \chi_{tu} \sum_{d_u} G^{a d_K d_u}_{\lambda a \, \lambda b-n'}(0) \quad G^{d_u d_i b}_{\lambda b-n' \, \lambda b}(0) \right\} \Big/ \binom{|n'|}{|\lambda a - \lambda b|}$$

In Eq.(2.9) n and n' are integers which satisfy the conditions (2.7) and $n \neq n'$; $\binom{|m|}{|\lambda_a - \lambda_b|}$ is the binomial coefficient.

Let us consider now the SSR for the processes

$$\alpha_i^S \, a \to \alpha_K^S \, b \tag{2.10}$$

and

$$\alpha_i^S \, a \to \alpha_K^V \, b \tag{2.11}$$

$(S_j = P, f, \omega; \; V_K = \rho, A_2; \; \alpha_j^{S,V}(0) \geq 0,5$

In order to solve the problem on superconvergency one must determine the position of the rightmost j-plane singularity α_e which can contribute asymptotically to the corresponding crossing-odd amplitude

For this purpose it is convenient to introduce the quantity

$$G^R = (-1)^{\alpha} G_6 \tag{2.12}$$

This quantity is positive for all the well-known Regge poles (except for mysterious A_1 -pole with $\alpha_{A_1}(0) < 0$) and it is natural therefore to assume that the poles with $G^R = -1$ (if they exist) have $\alpha(0) < 0$. Let us find G^R for the crossing-odd amplitudes. At high energies ($\nu \gg M_N$) these amplitudes are determined by the contribution of the singularities with $G_e = -G_i G_K$ and have the usual Regge behaviour $\nu^{\alpha_e - \alpha_i - \alpha_K}$ and

$$G_e^R = -(-1)^{I_e + I_i + I_K} G_i^R G_K^R \qquad (2.13)$$

It follows from (2.13) (for reggeons with $G_i^R G_K^R = +1$) that

$$G_e^R = -1$$

if, at least one of the reggeons α_i, α_K has isospin I=0 and the asymptotic behaviour of the sum rules (2.5) and (2.8) for the processes (2.10) and (2.11) is determined by the singularities with $\alpha_e(0) < 0$ provided only Regge poles are taken into account.

Let us find now the position of the moving cuts which correspond to the exchange of any number of Regge poles and have $G^R = -1$. Note that the intercept of $R\mathbb{P}\ldots\mathbb{P}$ cuts (R-reggeon),which have the same values of I, G and G - parity as the pole R, is equal to the intercept of pole R.

Consider the $R_1 \ldots R_n$ cuts. For these cuts

$$G_e^R = (-1)^{I_e + I_1 + I_2 + \ldots + I_n} G_1^R \ldots G_n^R$$

A simple analysis in the case of $\alpha^s a \to \alpha^s b$ scattering shows that (for nonstrange R_i) the minimum number of reggeons needed to construct the state with I=0 and $G_e^R = -1$ is equal to 3 and hence the maximum $\alpha_e(0) \lesssim -0{,}5$. The analogous analysis in the case of scattering $\alpha^s a \to \alpha^v b$ shows that the rightmost singularity corresponds to the exchange of two reggeons with I=1 (ρ, A_2) and consequently has $\alpha_e(0) \lesssim 0$.

Thus for the scattering of highlying trajectories with

$\alpha_e(0) \gtrsim 0.5$ the SSR (2.5) and (2.8) take place if at least one of the scattered trajectories is isoscalar.

Turn now to the OSSR (2.9). The asymptotic behaviour of these sum rules is determined by the "noncanonical" $\nu^{\alpha_e - \alpha_i - \alpha_K - m}$ $(m = 1, 2 \ldots)$ contributions of reggeons. The symmetry properties of these contributions depend on . The determination of m for each of the sum rules demands the special investigation and we shall always choose the minimum value of $m = 1$ assuming that OSSR exist if $\alpha_e - \alpha_i - \alpha_K < 0$ (independently of the sign of σ_e and consequently of G_e^R).

Taking into account the above analysis we can see that SSR (2.9) take place for all the reactions (2.10) and (2.11) except for $\alpha_i a \to \alpha_i b$ $(i = \omega, f)$, since in this case the rightmost singularity with I=0 is α_P and then $\alpha_P(0) - 2\alpha_i(0) \gtrsim 0$

3. Coupling of the Isoscalar Reggeons with Particles.

In Ref. [3] the method of SSR has been applied to the investigation of the I=1 reggeons (ρ, A_2, π) couplings with baryons N and Δ_{33} . The reactions $\alpha_i^{v} N \to \alpha_K^{v} a$ ($a = N, \Delta_{33}$; $i, K = \rho, A_2, \pi$) have been considered and the SSR for these processes have been saturated by the contributions of N and Δ_{33} . It has been shown that such approach gives the relations between Regge residues $G_{\lambda_a \lambda_b}^{a \alpha_j b}$ ($a, b = N, \Delta_{33}$; $j = \rho, A_2\ \pi$) which agree well with experimental data.

In accordance with this scheme let us consider now the processes

$$\alpha_i^{s,v} a \rightarrow \alpha_k^{s,v} a \qquad (a = N,\ \Delta_{33}), \tag{3.1}$$

$$\alpha_i^{s} N \rightarrow \alpha_k^{v} \Delta_{33} \tag{3.2}$$

and saturate the SSR for these processes by the contributions of nucleon and Δ_{33} -isobar. Taking into account (2.5), (2.8), (2.9) and the analysis of SSR in Sec.2 one can obtain a set of equations for the helicity residues $G^{a\alpha_j^s a}_{\lambda a \lambda' a}$ which is given in Table 1 *). The solution of this set has the following features.

a) The residues of isoscalar poles with single flip are equal to zero

$$G^{a\alpha_j^s a}_{\lambda a\, \lambda a \pm 1} = 0 \tag{3.3}$$

In order to find the relations for the other helicity residues we shall use the information on the residues $G^{a\alpha^v b}_{\lambda a \lambda b}$ which has been obtained in Ref. [3], namely taking into account that residues with single flip $G^{a\alpha^v b}_{\lambda a\, \lambda a \pm 1}$ $(a, b = N,\ \Delta_{33})$ are nonzero we get that

b) In the case of isoscalar reggeons the nonflip residues are are not equal to zero

*) In Table 1 and in what follows $G^{a\alpha_j b}_{\lambda a \lambda b} \equiv G^{a\alpha_j b}_{\lambda a \lambda b}(0)$

$$G^{a\alpha^s_j a}_{\lambda_a \lambda_a} \neq 0 \tag{3.4}$$

and satisfy the following relations

$$G^{\Delta\alpha^s_j \Delta}_{1/2\ 1/2} = G^{\Delta\alpha^s_j \Delta}_{3/2\ 3/2}$$

(3.5)

$$\sqrt{2}\ G^{N\alpha^s_j N}_{1/2\ 1/2} = G^{\Delta\alpha^s_j \Delta}_{1/2\ 1/2}$$

For the physical vertices (3.5) acquire the form

$$\left(G^{N\alpha^s_j N}_{1/2\ 1/2}\right)^{Phys} = \left(G^{\Delta\alpha^s_j \Delta}_{1/2\ 1/2}\right)^{Phys} = \left(G^{\Delta\alpha^s_j \Delta}_{3/2\ 3/2}\right)^{Phys} \tag{3.6}$$

These relations can be generalized to the case of arbitrary particles if one assumes that for the poles with I=0 the nondiagonal transitions are small as compared with the diagonal ones. This assumption formulated for the first time in [6] agrees with experimental data on the inelastic diffraction processes: at high energies the cross sections of the production of resonances via the I=0 t-channel exchange are at least one order smaller than those of the corresponding elastic reactions. The matrices of the transitions $a\alpha^s_i b$ are diagonal ($a = b$) in this approximation and in SSR one must take into account only one intermediate state in s- and u-channels (see Fig.1).

Analysing the SSR for the arbitrary particle scattering (as it was done in the case of N and Δ_{33}) it is not difficult to get that the general solution has the form

$$G^{a\alpha_i^s b}_{\lambda_a \lambda_b} = A_i\, \delta_{ab}\, \delta_{\lambda_a \lambda_b} \tag{3.7}$$

i.e. the vertices of poles with I=0 are independent of the particle b type and of the spin *). This result was for the first time obtained in Ref. [6] , and then discussed in Ref. [7] on the basis of t-channel approach to the sum rules for αa -scattering.

For the $\mathbb{P}$ -pole one can deduce the relation between the spin-flip residue $G^{N\alpha_{\mathbb{P}}N}_{1/2\,-1/2}$ and inelastic transitions on the vacuum pole. For this purpose it is convenient to consider the reaction $\alpha_{\mathbb{P}} N \to \alpha_{\mathbb{P}} N$. Since initial and final reggeons are identical, the helicity amplitudes are completely crossing-symmetric. So kinematic singularities free amplitude $\tilde{T}^{(\tau=-1)}_{1/2\,1/2}$ for which the SSR (2.1) hold is completely antisymmetric. Therefore for this amplitude the following SSR exists

$$\left(G^{N\alpha_{\mathbb{P}}N}_{1/2\,-1/2}\right)^2 = -\int_{\nu_h}^{\infty} \mathrm{Im}\, \tilde{T}^{\alpha_{\mathbb{P}}\alpha_{\mathbb{P}}\,(\tau=-1)}_{1/2\;1/2}(\nu, q_1^2, q_2^2, q^2=0) \tag{3.8}$$

where the sign $\sim$ means the extraction of the kinematic

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*) Th coefficient A_i in (3.7) can, in principle, depend on the baryon number and strangeness of particle a . The consideration of the sum rules for the reactions involving the baryon and strange trajectories scattering can clear up this dependence.Such consideration is not subject of this work.

factor $\vec{\Pi}_0[\vec{q}_1\ \vec{q}_2]$ from the corresponding amplitude (see [1]). According to the unitarity condition the imaginary part of $\tilde{T}^{\alpha_P \alpha_P}_{1/2\ 1/2}(\tau=-1)$ is determined by the inelastic vertices $G^{N\alpha_P N^*}_{\lambda_N \lambda_N^*}$ with helicity flip. Note that the contribution of resonances with spin S=1/2 to the integral is positive, therefore the relation (3.8) can be satisfied only in the presence of the states with $S \geqslant 3/2$. The corresponding helicity vertices cannot be found from the data on the cross sections of the unpolarized nucleons diffraction dissociation. However using the information on these cross sections one can obtain the upper limit on $G^{N\alpha_P N}_{1/2\ 1/2}$

$$\left(G^{N\alpha_P N}_{1/2\ -1/2}(q^2)\cdot M_N \big/ G^{N\alpha_P N}_{1/2\ 1/2}(q^2)\right)^2 < 0.1 \qquad (3.9)$$

at $|q^2| < 0.5\ (\mathrm{GeV/c})^2$.

4. Discussion

The main purpose of our work was the analysis of the superconvergent relations for the scattering of boson reggeons α_i^S with I=0 ($S_i = \mathbb{P},\ f,\ \omega$) and the application of these relations to the investigation of the structure of helicity residues $G^{a\alpha_i^S b}_{\lambda_a \lambda_b}$. It is shown that the residues have the form (3.7) if the transitions $a\alpha_i^S b$ with $a \neq b$ are small.

Let us review now the experimental consequences of the considered approach:

a) The total cross sections of any of particle pairs which carry the same baryon number and strangeness should be equal at high energy and not depend on spin states

$$\sigma^{tot}_{\pi N} = \sigma^{tot}_{\rho N} = \sigma^{tot}_{A} = \sigma^{tot}_{\pi \Delta_{33}} = \sigma^{tot}_{\rho \Delta_{33}} = \dots ;$$

$$\sigma^{tot}_{KN} = \sigma^{tot}_{K^* N} = \sigma^{tot}_{K^{**} N} = \dots ; \qquad (4.1)$$

$$\sigma^{tot}_{NN} = \sigma^{tot}_{N^* N} = \sigma^{tot}_{N \Delta_{33}} = \sigma^{tot}_{\Delta_{33} \Delta_{33}} = \dots .$$

b) The polarization of particles should not change in the elastic scattering processes at high energies and small $|t|$ [6]

c) The residues of f- and ω -poles are proportional to the residue of $\mathbf{P}$ -pole.

It is worth emphasizing that the conclusion on the total cross sections equality was obtained also in the theory with asymptotically constant cross section [8].

The accuracy of the results a) - c) depends on the smallness of nonpole intermediate states contribution i the SSR. The quoted consequences a) - c) accord with existing experimental data. For example, the analysis of the spin dependence of the elastic processes shows the smallness of the spin-flip residues $G^{N \alpha_i^s N}_{1/2 \; 1/2}$. The proportionality of the $\mathbf{P}$ and f -pole residues allows one to prove the f-dominant model of vacuum pole which is confirmed by the analysis of the binary reactions in Regge poles framework.

The indirect arguments in favour of the relations (3.6) can be given also. Let us consider the reactions $P^-_p \to P^0_n$ and $P^+_p \to P^0 \Delta^{++}_{33}$, where $P^\pm$ are pseudoscalars $\pi^\pm$, $K^\pm$;

$P^\circ = \pi^\circ, \eta, K^\circ, \bar{K}^\circ$.

At high energy the cross sections of these reactions are determined by the contribution of Regge poles ρ , A_2 and moving cuts. The following relation between these processes has been predicted [3] in the framework of pole model

$$\frac{d\sigma}{dt}(P^+_p \to P^\circ \Delta^{++}_{33}) = \frac{3}{2}\,\frac{d\sigma}{dt}(P^-_p \to P^\circ n) \tag{4.2}$$

Though the prediction has been made for the region of small $|t|$, it turns out that the (4.2) works well up to $|t| \approx 1$ (GeV/c)^2. This fact is not trivial since at $|t| \gtrsim 0.1$ the relations based on the pole model can be destroyed by the contribution of moving cuts (see Fig.2). It is not difficult, however, to be convinced (by the simple counting of the helicity states in Fig.2) that the cuts do not destroy the (4.2) if the relations (3.3) and (3.6) take place in the sufficiently large region of t.

In conclusion let us emphasize that the results of this and previous [1-5] works testify the fruitfulness of the approach based on the superconvergent sum rules as a method for the theoretical investigation of the Regge pole residues.

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Amplitude	τ	Equations
		$\alpha_i^s N \to \alpha_k^s N$
1/2 1/2	-	$G^{N\alpha_i^s N}_{1/2\,-1/2} \cdot G^{N\alpha_k^s N}_{1/2\,-1/2} = 0$
		$\alpha_i^s N \to \alpha^v N$
1/2 1/2	-	$G^{N\alpha_i^s N}_{1/2\,-1/2} \cdot G^{N\alpha^v N}_{1/2\,-1/2} = 0$
		$\alpha_i^s N \to \alpha^v \Delta_{33}$
1/2 1/2		$G^{N\alpha^v\Delta}_{1/2\,1/2}\left[G^{N\alpha_i^s N}_{1/2\,1/2} - \frac{1}{\sqrt{2}} G^{\Delta\alpha_i^s\Delta}_{1/2\,1/2}\right] = 0$
1/2 -1/2	-	$G^{N\alpha_i^s N}_{1/2\,-1/2}\, G^{N\alpha^v\Delta}_{1/2\,-1/2} - \frac{1}{\sqrt{2}}\left(G^{N\alpha^v\Delta}_{1/2\,-1/2}\, G^{\Delta\alpha_i^s\Delta}_{1/2\,-1/2} - G^{N\alpha^v\Delta}_{1/2\,3/2}\, G^{\Delta\alpha_i^s\Delta}_{1/2\,3/2}\right)$
1/2 -1/		$G^{N\alpha^v\Delta}_{1/2\,-1/2}\left[G^{N\alpha_i^s N}_{1/2\,1/2} - \frac{1}{\sqrt{2}} G^{\Delta\alpha_i^s\Delta}_{1/2\,1/2}\right] = 0$ $G^{N\alpha^v\Delta}_{1/2\,1/2}\left[G^{N\alpha_i^s N}_{1/2\,-1/2} - \frac{1}{\sqrt{2}} G^{\Delta\alpha_i^s\Delta}_{1/2\,-1/2}\right] = 0$
1/2 3/2		$G^{N\alpha^v\Delta}_{1/2\,3/2}\left[G^{N\alpha_i^s N}_{1/2\,1/2} - \frac{1}{\sqrt{2}} G^{\Delta\alpha_i^s\Delta}_{3/2\,3/2}\right] = 0$ $G^{N\alpha^v\Delta}_{1/2\,1/2} \cdot G^{\Delta\alpha_i^s\Delta}_{1/2\,3/2} = 0$
1/2 -3/2		$G^{N\alpha^v\Delta}_{1/2\,-3/2}\left[G^{N\alpha_i^s N}_{1/2\,1/2} - \frac{1}{\sqrt{2}} G^{\Delta\alpha_i^s\Delta}_{3/2\,3/2}\right] = 0$ $G^{N\alpha_i^s N}_{1/2\,-1/2}\, G^{N\alpha^v\Delta}_{1/2\,3/2} + \frac{1}{\sqrt{2}} G^{N\alpha^v\Delta}_{1/2\,-1/2}\, G^{\Delta\alpha_i^s\Delta}_{1/2\,3/2} = 0$ $G^{N\alpha^v\Delta}_{1/2\,1/2}\quad G^{\Delta\alpha_i^s\Delta}_{1/2\,-3/2}$

Table 1 (to be continued)

		$\alpha_i^s \Delta_{33} \to \alpha_K^s \Delta_{33}$
1/2 1/2	-	$G^{\Delta\alpha_i^s\Delta}_{1/2\,-1/2}\, G^{\Delta\alpha_K^s\Delta}_{1/2\,-1/2} - G^{\Delta\alpha_i^s\Delta}_{1/2\,3/2}\, G^{\Delta\alpha_K^s\Delta}_{1/2\,3/2} = 0$
1/2 3/2		$G^{\Delta\alpha_K^s\Delta}_{1/2\,3/2}\left[G^{\Delta\alpha_i^s\Delta}_{1/2\,1/2} - G^{\Delta\alpha_i^s\Delta}_{3/2\,3/2}\right] = 0$
1/2 -3/2		$G^{\Delta\alpha_K^s\Delta}_{1/2\,-3/2}\left[G^{\Delta\alpha_i^s\Delta}_{1/2\,1/2} - G^{\Delta\alpha_i^s\Delta}_{3/2\,3/2}\right] = 0$ $G^{\Delta\alpha_i^s\Delta}_{1/2\,-1/2}\, G^{\Delta\alpha_K^s\Delta}_{1/2\,3/2} - G^{\Delta\alpha_K^s\Delta}_{1/2\,-1/2}\, G^{\Delta\alpha_i^s\Delta}_{1/2\,3/2} = 0$
3/2 3/2	-	$G^{\Delta\alpha_i^s\Delta}_{1/2\,3/2}\, G^{\Delta\alpha_K^s\Delta}_{1/2\,3/2} = 0$
		$\alpha_i^s \Delta_{33} \to \alpha^v \Delta_{33}$
1/2 1/2	-	$G^{\Delta\alpha_i^s\Delta}_{1/2\,-1/2}\, G^{\Delta\alpha^v\Delta}_{1/2\,-1/2} - G^{\Delta\alpha_i^s\Delta}_{1/2\,3/2}\, G^{\Delta\alpha^v\Delta}_{1/2\,3/2} = 0$
1/2 3/2		$G^{\Delta\alpha^v\Delta}_{1/2\,3/2}\left[G^{\Delta\alpha_i^s\Delta}_{1/2\,1/2} - G^{\Delta\alpha_i^s\Delta}_{3/2\,3/2}\right] = 0$ $G^{\Delta\alpha_i^s\Delta}_{1/2\,3/2}\left[G^{\Delta\alpha^v\Delta}_{1/2\,1/2} - G^{\Delta\alpha^v\Delta}_{3/2\,3/2}\right] = 0$
1/2 -3/2		$G^{\Delta\alpha^v\Delta}_{1/2\,3/2}\left[G^{\Delta\alpha_i^s\Delta}_{1/2\,1/2} - G^{\Delta\alpha_i^s\Delta}_{3/2\,3/2}\right] = 0$ $G^{\Delta\alpha_i^s\Delta}_{1/2\,-3/2}\left[G^{\Delta\alpha^v\Delta}_{1/2\,1/2} - G^{\Delta\alpha^v\Delta}_{3/2\,3/2}\right] = 0$ $G^{\Delta\alpha_i^s\Delta}_{1/2\,-1/2}\, G^{\Delta\alpha^v\Delta}_{1/2\,3/2} - G^{\Delta\alpha^v\Delta}_{1/2\,1/2}\, G^{\Delta\alpha_i^s\Delta}_{1/2\,3/2} = 0$
3/2 3/2	-	$G^{\Delta\alpha_i^s\Delta}_{1/2\,3/2}\, G^{\Delta\alpha^v\Delta}_{1/2\,3/2} = 0$

Table 1 (continuation)

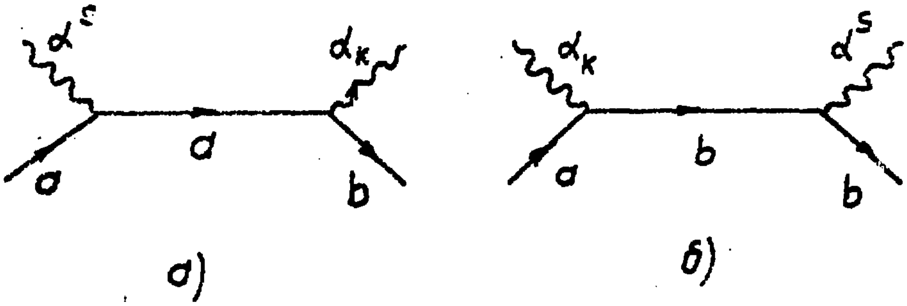

Fig. 1

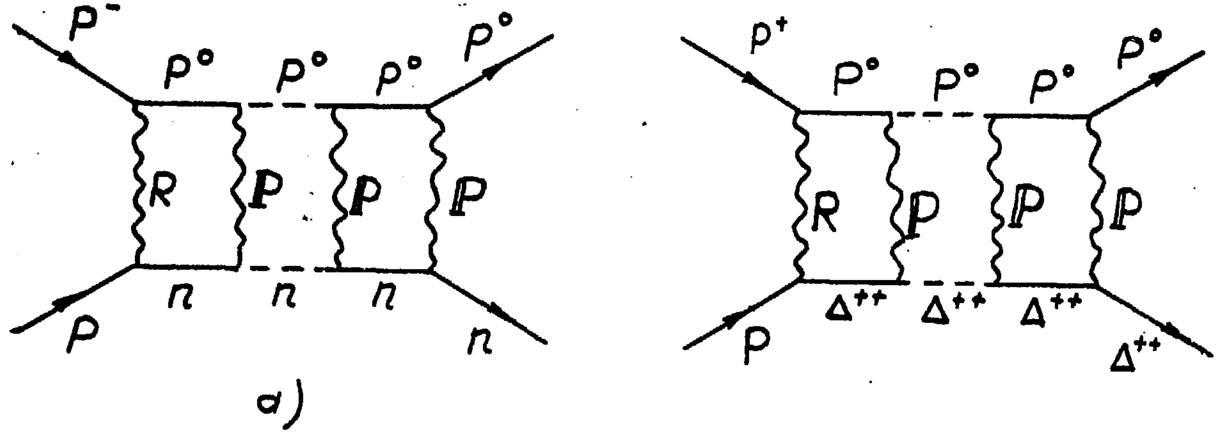

Fig. 2

FIGURE CAPTIONS

Fig.1 Contribution of resonances to the sum rules for the processes $\alpha^{s} a \to \alpha_{K} b$; a) in the s-channel, b) in the u-channel.

Fig.2 Contribution of moving cuts $R\,\mathbb{P} \ldots \mathbb{P}$ to the reactions;
a) $\rho^{-} p \to \rho^{0} n$ and b) $\rho^{+} p \to \rho^{0} \Delta_{33}^{++}$

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А.А.ГРИГОРЯН, А.Б.КАЙДАЛОВ,
Г.Н.ХАЧАТРЯН

ДИСПЕРСИОННЫЕ ПРАВИЛА СУММ И ВЫЧЕТЫ ИЗОСКАЛЯРНЫХ РЕДЖЕОНОВ

(на английском языке)

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Тех.редактор А.С.Абрамян

Заказ694 ВФ-05265 Тираж 299

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