ON INELASTIC PROCESSES DUE TO VARIOUS TYPES OF MOVING POLES

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Inelastic peripheral high energy interactions due to exchanges of various types of "poles" (pions, kaons and vacuum reggeon) are considered from the unified viewpoint of amplitude poles in the complex angular momentum plane. It is demonstrated that in the region in which effects related to the motion of poles describing these quanta are important, the overwhelming part of the contribution is from the vacuum reggeon. However, the contribution from this whole region decreases logarithmically with energy and the production multiplicity is small. In the region which yields the main contribution to the cross section these pole motion effects do not play any role. Here the main contribution is from one-pion exchange as computed in the one-meson approximation. This circumstance is related to the difference in kinematics of elastic and inelastic processes.

Some criteria are considered which can be employed for distinguishing between inelastic processes due to vacuum reggeon exchange and processes due to meson exchange.

1. The method of moving poles (MPM) in the plane of complex orbital momenta [1] makes it possible to determine the asymptotic behavior of the elastic-scattering amplitude. 1) It is customary to state here that the process is due to the exchange of one "vacuum pole" or "reggeon." This yields, through the optical theorem, the asymptotic behavior of the total cross section. Until recently, the constancy of the total cross section of inelastic processes at high energy was the only information that could be extracted within the framework of this method concerning the inelastic processes. Yet the question of inelastic processes is most important in the theory of particle interaction.

It is known that elastic scattering, particularly diffraction scattering, can be accompanied by inelastic processes. An example is the so-called diffraction generation of particles [3]. From the point of view of the MPM this process must be regarded as an inelastic interaction, due to the exchange of one vacuum reggeon [4,5].

On the other hand, the theory of peripheral interaction [6,7] deals with inelastic processes due to the exchange of a single quantum: pion, kaon, etc. (This method of analysis is called the one-meson approximation—OMA). In the method of complex orbital momenta, these quanta must also be regarded as "moving poles" with quantum numbers (parity, strangeness, etc.) which differ from the quantum numbers of the vacuum 2).

This raises the possibility of regarding all peripheral inelastic collisions (both "diffraction" and one-meson) from a unified point of view. In the present paper we consider the interaction of high-energy nucleons and discuss the following questions:

1. What is the relative role of different Regge trajectories in inelastic processes?

2. What is the character of the inelastic processes of various types occurring at high energies?

2. In examining inelastic processes we shall make use of the diagram method in the MPM, which has already been used by Ter-Martirosyan [4] and Frautschi [5] (who have also considered the validity of this method). The gist of this method consists in the following. Secondary particles are divided into two groups (or two "jets"), with several particles in each (in particular, one of these groups may contain one particle—we shall call such a process a single-jet process).

It is assumed that the process can be regarded as an ordinary "four-point diagram," i.e., the production of two heavier particles with masses $\mathfrak{M}_1 = \sqrt{s_1}$ and $\mathfrak{M}_2 = \sqrt{s_2}$ from two nucleons of mass m (s_1 and s_2 are the squares of the total

1) We leave aside the question of the universality of the MPM, which has become particularly acute following recent experiments on scattering of pions by protons [2].

2) We note that effects connected with the motion of the pole describing the quantum that effects the exchange were not taken into account in [6,7]. The role of these effects will be discussed below.

energy of the particles of the first and second groups, respectively, in their c.m.s.). We can then use the concept of crossed channel (t-channel), in which two particles with masses m and $\mathfrak{M}_1$ go over into two particles with masses m and $\mathfrak{M}_2$.

The cosine of the angle in the t channel can in this case be calculated from the general rules, and has the form ($x \equiv k^2 = -t$)

$$\cos\theta_t \equiv z = \frac{(s_1 - m^2 + x)(s_2 - m^2 + x) - 2x(s - 2m^2)}{[(s_1 - m^2 + x)^2 + 4m^2x]^{1/2}\,[(s_2 - m^2 + x)^2 + 4m^2x]^{1/2}}. \quad (1)$$

It is assumed that the interaction is carried by a single quantum, which is set in correspondence to the moving pole in the MPM. The diagrams of this process are shown in Fig. 1 (single-jet process) and Fig. 2 (two-jet process). The amplitude is calculated by the Feynman rule, but with the following differences [4,5,8].

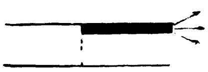

FIG. 1.

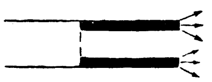

FIG. 2.

The propagation function D(x) is replaced by the quantity $-\pi\alpha I_i(t)/2$, where $I_i(t) = \sin^{-1}(\pi l_i(t)/2)$, $l_i(t)$ is the i-th Regge trajectory. It is customary to assume that when $|t| \ll m^2$, all the trajectories are straight lines with equal slopes, i.e.,

$$l_i(t) = b_i + \alpha t, \qquad \alpha = \text{const} \sim m^{-2}.$$

The quantities b_i are then determined from the condition $l_i(m_i^2) = J_i$, where m_i and J_i are the mass and spin of the quantum. For particles with zero spin (for example pions and kaons), this yields $b_\pi = -\mu^2/m^2$ and $b_k = -m_k^2/m^2$; μ and m_k are the masses of the pion and kaon. It is easy to see that when $x \to -m_i^2$ the quantity I_i has a pole $I_i \sim (x + m_i^2)^{-1}$, i.e., it behaves like $D_i(x)$.

In addition, it is necessary to introduce in the MPM an additional factor $P_{l_i(t)}(z)$—Legendre polynomials of order l_i of the cosine of the angle in the t-channel.

Starting from the foregoing rules, we can write down the differential cross sections for single-jet and two-jet nucleon-nucleon interaction processes at high energies in the form

$$\partial^2\sigma_i^{(1)}/\partial s_1\partial x = \tfrac{1}{4}\pi f_i^2 m^2\alpha^2 g_i^2(x)\, s^{-2}\,[(s_1 - m^2 + x)^2 + 4m^2x]^{1/2} \times \sigma_i(s_1, x)\,|I_i(x)|^2 P^2_{l_i(x)}(z), \quad (2)$$

$$\partial^3\sigma_i^{(2)}/\partial s_1\,\partial s_2\,\partial x = (3\alpha^2/32\pi s^2)\,[(s_1 - m^2 + x)^2 + 4m^2x]^{1/2} \times [(s_2 - m^2 + x)^2 + 4m^2x]^{1/2}\sigma_i(s_1, x)\,\sigma(s_2, x) \times |I(x)|^2 P^2_{l_i(x)}(z). \quad (3)$$

Here $\sigma_i^{(1,2)}$—cross section of single-and two-jet processes (Figs. 1, 2) due to the exchange of the i-th "reggeon"; f_i—is the reggeon—nucleon coupling constant in the unexcited vertex of the diagram of Fig. 1, and $g_i(x)$ is the "form factor" in this vertex. The quantities $\sigma_i(s,x)$ have the same meaning as in the OMA, and stand for the cross sections of the interactions between the virtual reggeon and the nucleon. When $I_i(x) = 2D_i(x)/\pi\alpha$ and $P_{l_i}(z) \to 1$ these expressions coincide with those of the OMA.

Expressions (2) and (3) can be recast in different form, which is more usual from the point of view of the MPM. Indeed, putting

$$\pi f^2 m^2\alpha^2 g^2(x)\,\sigma(s_1, x) = C_1(s_1, x),$$

$$(3\alpha^2/32\pi)\,\sigma(s_1, x)\,\sigma(s_2, x) = C_2(s_1, s_2, x),$$

we can rewrite (1) and (2) in the form

$$\partial^2\sigma_i^{(1)}/\partial s_1\partial x = C_1 s^{-2}\,[(s_1 - m^2 + x)^2 + 4m^2x]^{1/2}\,|I_i(x)|^2 P^2_{l_i(x)}(z), \quad (4)$$

$$\partial^3\sigma_i^{(2)}/\partial s_1\partial s_2\partial x = C_2 s^{-2}\,[(s_1 - m^2 + x)^2 + 4m^2x]^{1/2} \times [(s_2 - m^2 + x)^2 + 4m^2x]^{1/2}\,|I_i(x)|^2 P^2_{l_i(x)}(z). \quad (5)$$

The quantities C_1 and C_2 are proportional to the squares of the residues of the partial amplitude $f_i(l, t)$ which is encountered in elastic scattering. The arguments s_1 and s_2 reflect the dependence of the residues on the "masses" of the excited nucleons.

We note that (2)—(5) contain as unknown functions $\sigma_i(s, x)$, $C_1(s, x)$, and $C_2(s_1, s_2, x)$. The $\sigma_i(s, x)$ are known only for $x = -m_i^2$, where they represent the cross sections for the interaction between real particles; the quantities $C_{1,2}$ can be regarded as known only for $s_1 = s_2 = m^2$, where they are related to the elastic scattering amplitude.

Thus, when we use expressions (2) and (3) or (4) and (5), the question arises of extrapolating $\sigma(s, x)$ or $C(s, x)$ to regions where they are not directly observable. In this connection there is a large quantitative difference between the different Regge trajectories. Indeed, in the physical region of the inelastic processes $x > 0$; the main contribution is made by the region $|x| \ll m^2$ (see below). Therefore, if the pole of the function $I_i(x)$ is located near the point $x = 0$ (this occurs, for example, in the pion trajectory, where $x_p = -\mu^2$), expressions (2) and (3) are more "physical," since it is more convenient to use the quantities $\sigma(s_i, -\mu^2)$ and continue them into the region $x > 0$. On the other hand, if the pole of $I_i(x)$ is far from the point $x = 0$ (as, for example, in the case of a vacuum reggeon, when it is located near $x \sim -1\ \mathrm{BeV}^2$), expressions (5) and (4) are preferable, since the quantity $C_1(s = m^2, x)$ is known for the region $x > 0$ from elastic scattering. We shall not concern ourselves, however, with the continuation of the functions $\sigma(s,x)$ and $C(s,x)$, and focus attention on the factors $I_i(x)$ and $P_{l_i(x)}(z)$. In this connection, we shall assume henceforth that $C_{1,2}$ are constants of the order of unity.

The main condition for the effectiveness of MPM is that z should tend to infinity with increasing energy s. It is only in this region of values $z \to \infty$ that effects connected with the motion of the poles, corresponding to elementary particles, can appear. However, in inelastic processes, the kinematically allowed region of x, s, s_1, and s_2 is such that both large and small values of z are possible in it (details will follow). Therefore, under the condition $z \gg 1$ we can describe only some of the possible inelastic processes, involving part of the kinematically allowed region.

More specifically, we confine ourselves to the region

$$z \gg (s/2m^2)^\nu \gg 1. \tag{6}$$

Here ν—arbitrary parameter in the range $0 < \nu < 1$. We can always put here $P_{l_i(x)}(z) = z^{l_i(x)}$. The asymptotic energy region in which the effects connected with the motion of the pole appear, is determined from the condition

$$\alpha \ln z = \nu\alpha \ln (s/2m^2) \gg 1. \tag{7}$$

It follows therefore that the parameter ν cannot be close to zero, or else the asymptotic region of s will move too far away. If ν is equal to unity or is close to it, then condition (6) singles out too small a fraction of the inelastic processes (see below). We shall therefore put $\nu \sim (1-\nu) \sim 1$.

3. Before we proceed to direct calculations by Eqs. (4) and (5), let us consider the kinematics of the inelastic processes in nucleon-nucleon collisions. As is well known (see, for example [7]), the region of the variation of s_1, s_2 and x, which is allowed by the conservation laws for two-jet processes, lies between the planes $s_1 - m^2 = 2m\mu + \mu^2$ and $s_2 - m^2 = 2m\mu + m\mu^2$ (here μ is the mass of the pion, the lightest of the possible generated particles) and the surface given by the equation

$$(s_1 - m^2 + x)(s_2 - m^2 + x) - x(s - 4m^2) + m^2(s_1 - s_2)^2/s = 0 \tag{8}$$

for single-jet processes we should put $s_2 = m^2$ everywhere).

By substitution of (8) in (1) we can verify that the identity $|z| \equiv 1$ holds over the entire surface (8). In this lies the main difference between inelastic and elastic processes. In the latter $s_1 = s_2 = m^2$ and the value of z is

$$z = 1 - 2s/(x + 4m^2) \tag{9}$$

and tends to infinity asymptotically for forward scattering[3)].

Let us examine expression (1) for single-jet processes in two limiting cases.

1. Region where $s_1 - m^2 \gg x$. Here

$$z = \sqrt{\frac{x}{x+4m^2}}\left(1 - \frac{2s - 4m^2}{s_1 - m^2}\right). \tag{10}$$

This region can be arbitrarily called the region of large excitations.

2. Region where $s_1 - m^2 < x$. From (1) we have in this region

$$z = 1 - 2s/(x + 4m^2), \tag{11}$$

i.e., it has the same form as in the elastic process. We shall henceforth call it the region of small excitations.

Accordingly, for the two-jet processes we have:

1. Region of large excitations, where $s_{1,2} - m^2 \gg x$. Here

$$z \approx 1 - 2sx/s_1s_2. \tag{12}$$

2. Region of small excitations, where either both $s_{1,2}$ are small, $s_{1,2} - m^2 < x$, and then z is given by (11), or else one of the s is small and the other large: $s_1 - m^2 < x$, $s_2 - m^2 > x$. Then

$$z = \sqrt{\frac{x}{x+4m^2}}\left(1 + \frac{s_1 - m^2}{x} - \frac{2s - 4m^2}{s_2 - m^2}\right). \tag{13}$$

4. Let us consider the contributions from different regions. In single-jet processes, as can be readily verified, the region of small excitations plays an insignificant role in the asymptotic region of s. Indeed, the total cross section (after integration in formula (4) over the region under consideration) is

$$\sigma^{(1)}_{i,M} = \frac{C_1}{s^2} \int\limits_{x_{min}=2m\mu+\mu^2}^{\infty} dx \times \int\limits_{(m+\mu)^2}^{m^2+x} ds_1 |I_i(x)|^2 P^2_{l_i(x)}(z) [(s_1 - m^2 + x)^2 + 4m^2x]^{1/2}$$

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3)This difference is due to the fact that at the point $s_1 - m^2 = s_2 - m^2 = x = 0$ the limiting values of z depend on the path followed on approaching this point. For example, for $|s_1 - m^2|$, $|s_2 - m^2| < x \to 0$, Eq. (9) applies at the point $x = 0$. As $x < |s_{1,2} - m^2| \to 0$ we have $z = 1$ at $x = 0$. This circumstance is connected with the fact that at this point z is not an analytic function in its variables.

$$= \frac{C_1}{s^2} \int\limits_{2m\mu}^{\infty} dx \int\limits_{(m+\mu)^2}^{m^2+x} ds_1$$

$$\times \sqrt{4m^2 x} \sin^{-2}\left[\frac{\pi}{2}(b_i - \alpha x)\right]\left(\frac{s}{2m^2}\right)^{2(b_i - \alpha x)}$$

$$= C_1 \sqrt{\frac{\mu}{32m}}\left(\frac{s}{2m^2}\right)^{2[b_i - 1 - 2\alpha m\mu]}$$

$$\times \left(\ln \frac{s}{2m^2}\right)^{-2}\left[\sin \frac{\pi}{2}(b_i - 2\alpha m\mu)\right]^{-2}. \quad (14)$$

We see therefore that the contribution from this region decreases in power-law fashion in all cases, since $b_i \leq 1$ for all trajectories.

In the region of large excitations, condition (6) assumes a major role, and can be written in the form

$$z = \sqrt{\frac{x}{x+4m^2}}\left(1 - \frac{2s}{s_1}\right) \approx 4\sqrt{m^2 x}\left(\frac{s}{2m^2}\right)\frac{1}{s_1} \gg \left(\frac{s}{2m^2}\right)^{\nu}. \quad (15)$$

This determines the upper limit of s_1:

$$s_{1max} = 4\sqrt{m^2 x}\,(s/2m^2)^{1-\nu}. \quad (16)$$

We see therefore that when $\nu = 1$, $s_{1\,max}$ does not depend on s and is bounded[4)].

Integrating (4) with respect to s_1 and with respect to x, we get

$$\sigma_i^{(1)} = 2C_1\left[\sin \frac{\pi}{2} b_i\right]^{-2}\left(\frac{s}{2m^2}\right)^{-2\nu(1-b_i)}$$

$$\times \int\limits_{x_{min}}^{\infty} \frac{x\,dx}{(1 - b_i + \alpha x)}\left(\frac{s}{2m^2}\right)^{-2\nu\alpha x},$$

$$x_{min} = 4m^2 s^{-2}\,(s/2m^2)^{2\nu}. \quad (17)$$

x_{min} tends to zero with increasing s.

It is seen from (17) that the contributions from all the non-vacuum trajectories, for which $b_i < 1$, will decrease in power-law fashion with increasing s, although not as rapidly as the contributions from the small-excitation region:

$$\sigma_i^{(1)} = \frac{C_1}{4\nu^2\alpha^2(1-b_i)}\left(\ln \frac{s}{2m^2}\right)^{-2}\left(\frac{s}{2m^2}\right)^{-2\nu(1-b_i)}\left(\sin \frac{\pi b_i}{2}\right)^{-2}. \quad (18)$$

The contribution from the pion trajectory will be larger than that from the kaon trajectory, since $|b_\pi|$ is appreciably smaller than $|b_K|$. Namely, $\sigma_\pi \sim s^{-2\nu(1+\mu^2/m^2)}$, whereas $\sigma_K \sim s^{-2\nu(1+m_K^2/m^2)}$. In addition, the pre-exponential factor of the type $\sin^2(\pi b_i/2)$ will in the case of pion exchange be $(m_K^2/\mu^2)^2$ times larger than in the case of kaon exchange.

The vacuum trajectory for which $b_V = 1$ makes a logarithmically decreasing contribution to the cross section:

$$\sigma_V = C_1/\nu\alpha^2 \ln(s/2m^2).$$

The ratio of the contribution of the pion trajectory to that of the vacuum is

$$\frac{\sigma_\pi}{\sigma_V} = \left(\frac{s}{2m^2}\right)^{-2\nu(1+\mu^2/m^2)}\left[\left(1 + \frac{\mu^2}{m^2}\right)\frac{\pi^2}{4}\frac{\mu^4}{m^4} 2\nu \ln \frac{s}{2m^2}\right]^{-1}$$

$$\approx \frac{2m^4}{\nu\pi^2\mu^4 \ln(s/2m^2)}\left(\frac{s}{2m^2}\right)^{-2\nu}. \quad (19)$$

This ratio decreases with increasing energy in power-law fashion, but contains a large coefficient. Therefore there exists an energy interval in which the pion exchange is more significant than vacuum-reggeon exchange. At $\nu \approx 0.7$—0.8, this occurs at energies on the order of several times 10 BeV. It is possible that this is the reason why diffraction inelastic processes appear in pronounced fashion only at energy $E_{l.s} \sim 20$ BeV [9].

The situation in two-jet processes is analogous. The region of small excitations makes a small contribution; we shall not present an expression for this contribution.

In the region of large excitations, the condition

$$z = 2sx/s_1 s_2 \gg (s/2m^2)^{\nu} \quad (20)$$

determines the upper limits of integration with respect to s_1 and s_2. Taking this into account, the contribution to the cross section assumes the form

$$\sigma_i^{(2)} = C_2\left(\sin \frac{\pi}{2} b_i\right)^{-2}\left(\frac{s}{2m^2}\right)^{-2\nu(1-b_i)} \int\limits_{x_{min}}^{\infty}\left(\frac{s}{2m^2}\right)^{-2\nu\alpha x}\frac{x^2\,dx}{(1-b_i+\alpha x)}$$

$$\times \left[\ln \frac{x}{x_{min}} - \frac{1}{2(1 - b_i + \alpha x)}\right], \quad (21)$$

where $x_{min} = 2m^2 s^{-1}(s/2m^2)^{\nu}$.

Processes consisting of a large number of "jets," can be represented as two-jet processes by combining the jets in two groups; their contribution is given by (21). It follows therefore that, as in the preceding (single-jet) case, all but the vacuum trajectory make a contribution which decreases in power-law fashion. The vacuum trajectory makes a contribution that decreases logarithmically with the energy.

Thus, in the considered region of inelastic processes ($z \gg 1$) the main contribution is made by the process due to the exchange of the vacuum reggeon. However, it cannot ensure constancy of the total cross section of the inelastic processes. This is to be expected. Indeed, the cross section of the elastic scattering which gives rise to inelastic diffraction interactions due to exchange of

4)The situation is analogous to the case of small excitations. In this case there is no region of large excitations. We have therefore put $1 - \nu \sim 1$.

the vacuum reggeon decreases itself logarithmically with increasing energy[5)]. We can therefore conclude that the main contribution to the total cross section is made by the region of small $z \sim 1$. The moving pole method is not effective in this region, since the quantity $P_l(z)$ which is basic to it is small: $P_l(z) \sim 1$ for all trajectories. We need therefore estimate only the influence of the factors $|I_i(x)|^2 = \sin^{-2}[\pi(b_i - \alpha x)/2]$. At small x, the quantities I_i^2 are inversely proportional to the fourth powers of the masses of the quanta that carry the interaction. Therefore the main contribution will be made by the process due to the exchange of the lightest (among the strongly-interacting) particle, i.e., the pion. The same result is obtained in ordinary field theory if the propagation function as $x \to 0$ is represented in the form $D_i(x) = (x + m_i^2)^{-1}$. Thus, in the main region of the inelastic processes, at $z \sim 1$, the effects connected with the motion of the pole, describing the exchange quantum do not play any role. The previous results obtained in the OMA without account of the effects, remain in force here.

5. Let us ascertain the characteristic features of the inelastic processes in the region $z \geq (s/2m^2)^{\nu} \gg 1$, and let us compare them with the characteristics of the process in the region $z \sim 1$.

A. Values of the excitations s_1 and s_2 (on which the multiplicity, etc. depend) are limited from above by conditions (16) and (20). The effective values of s_1 and s_2 are close to the largest values, while the effective values of x are of the order of $[2\nu\alpha \ln(s/2m^2)]^{-1}$. Therefore $s_{1,\text{eff}} \sim m^2(s/2m^2)^{1-\nu}$ in single-jet processes and $(s_1 s_2)_{\text{eff}} \sim m^2(s/2m^2)^{1-\nu}$ in the case of excitation of both nucleons (we are presenting estimates with logarithmic accuracy).

The corresponding values for single-meson processes with $z \sim 1$ have been considered previously[6]. Their order of magnitude is

$$s_{1\,\text{eff, OMA}} \sim s, \qquad (s_1 s_2)_{\text{eff, OMA}} \sim s.$$

Thus, even "large" excitations are much smaller in the exchange of vacuum reggeons (by a factor $(s/2m^2)^{\nu}$) than excitations in single-meson processes. In this connection, the diffraction mechanism of the inelastic processes (exchange of vacuum reggeon) can be appreciable only for relatively "lean" jets.

B. The condition (6) imposes a limitation on the components of the momentum transfer k. Indeed, let us write x in the form $x = k^2 = k_\perp^2 + \kappa^2$, $\kappa^2 = k_\parallel^2 - k_0^2$ (where $k_\perp$, $k_\parallel$, and k_0 are the transverse, longitudinal, and time components of the four-momentum k in the c.m.s.). We take account of the fact that (see, for example,[6]) $\kappa^2 = s_1^2 m^2/s^2$ in single-jet processes, $\kappa^2 = s_1 s_2/s$ in two-jet processes, and $k_\perp^2 = s\theta^2$, where θ is the angle between the momenta of the primary particle and one of the jets. We can then represent z in the form $z = \sqrt{x/\kappa^2} = \sqrt{1 + k_\perp^2/\kappa^2}$ in single-jet processes and $z = x/\kappa^2 = 1 + k_\perp^2/\kappa^2$ in two-jet processes. The condition $z \gtrsim (s/2m^2)^{\nu} \gg 1$ means that in both cases $k_\perp^2 \gg \kappa^2$, and accordingly $\theta^2 \gg s_1 s_2/s$.

The existence of interactions in which this inequality is satisfied can be checked experimentally. By the same token, we can verify whether the interaction with $z \gg 1$ is actually realized.

We note also that in single-meson processes (in the OMA) $k_\perp^2 \sim \kappa^2$ and $\theta^2 \sim s_1 s_2/s$. This criterion can be used in the interpretation of the experimental data.

C. It must be noted that in the case of exchange of a vacuum reggeon an azimuthal asymmetry can arise in the distribution of the secondary particles. It is connected with the possibility of the transfer of the angular momentum and consequently, with the fact that the Treiman and Yang conditions[10] are not satisfied in this case. The occurrence of azimuthal asymmetry follows from the formulas obtained by Ter-Martirosyan[4], who, however, did not discuss this asymmetry. This circumstance can also be used to ascertain the mechanism of a given specific inelastic process.

In conclusion the authors are grateful to E. L. Feinberg for very fruitful discussions.

[5)]It must also be noted that if inelastic processes of this type were to make the principal contribution to the total cross section, then the elastic process would be due essentially to the exchange of two vacuum reggeons. Such an elastic process would not have the usual Regge properties, and in particular the corresponding partial amplitude in the crossed channel would, according to [8], have a branch point in the l-plane.

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