furnishing the superconducting solenoid, and also for taking part in a discussion of the results of this work.

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TWO-PARTICLE INELASTIC REACTIONS AND MOVING BRANCH POINTS IN ANGULAR-MOMENTUM PLANE

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Submitted 9 March 1967
ZhETF Pis'ma 5, No. 11, 404-406 (1 June 1967)

Two-particle inelastic reactions at high energy (charge exchange, resonance production, etc.) are usually described with the aid of Regge poles in a crossing channel, with appropriate quantum numbers [1]. This approach leads, as it were, to a fair agreement with experiment in the region of small momentum transfers, the only region for which experimental data are presently available. On the other hand, a pure pole analysis is not satisfactory from the theoretical point of view, for it disregards the presence of branching in the angular-momentum plane.

We have previously calculated [2] the contribution of Mandelstam branch points to the asumptotic elastic-scattering amplitude at large momentum transfers. Since the branch points were in this case the singularities farthest to the right in the angular-momentum plane, they were decisive in the determination of the asymptotic behavior of the amplitude. A similar phenomenon takes place in the case of two-particle inelastic processes, too. We consider in this note the contribution of branch points to the asymptotic amplitudes of such reactions, at large momentum transfers $|t| >> m^2$, but at $|t| << s$ (s is the square of the total c.m.s. energy and m is the nucleon mass).

The method described in [3] makes it easy to obtain the branch-point trajectories arising upon exchange in the t-channel of one pole $\beta(t_1)$ with the required quantum numbers and n vacuum poles $\alpha(t_2)(\alpha(o) = 1)$:

$$j_{n+1}(t) = \beta(t_1) + n\alpha(t_2) - n. \qquad (1)$$

The "reggeon masses" t_1 and t_2 are determined from the equations

$$\sqrt{t_1}\,\beta'(t_1) = \sqrt{t_2}\,\alpha'(t_2), \quad \sqrt{t_1} + n\sqrt{t_2} = \sqrt{t}. \qquad (2)$$

At values of the number $n >> \sqrt{t/m^2}$ we have

$$\sqrt{t_2} \approx \frac{\sqrt{t}}{n}, \quad \sqrt{t_1} \approx \frac{\sqrt{t}}{n}\,\frac{\alpha'(o)}{\beta'(o)}, \qquad (3)$$

and the equations of the trajectories take the form

$$j_{n+1} \approx \beta(o) + \frac{\alpha'(o)\,t}{n}. \qquad (4)$$

Thus, the point $j = \beta(o)$ turns out to be a point of branch-point condensation*, just as the point j = 1 is a condensation point for Mandelstam branch points. If, for specified quantum numbers (isotopic spin, strangeness, etc.), $\beta(t)$ is the singularity farthest to the right in the j-plane for the case of small t, then for large $|t|$ the branchings (4) remain on the extreme right, although the pole $\beta(t)$ can go off far to the left. Therefore the singularities (4) are precisely those which determine the asymptotic behavior of the amplitude at large $s \gg |t|$.

It is shown in [2] that the use of the approximation (4) for the calculation of the elastic amplitude (when $\beta(o) = \alpha(o) = 1$) is valid only if $\ln(s/m^2) \gg 1$. On the other hand, if $\ln(s/m^2) \geq 1$, as is the case for present-day energies, then it is necessary to use the exact equations for the trajectories in order to calculate the contribution of the Mandelstam branch points. It can be shown that in our case, even when $\ln(s/m^2) \sim 1$, it is legitimate to use the following approximation for the trajectories j_{n+1}:

$$j_{n+1} = \beta(o) + n\alpha(t/n^2) - n. \tag{5}$$

Formula (5) is obtained from (1) by making the substitutions $t_1 = 0$ and $t_2 = t/n^2$. The difference between (5) and (1) influences only the factor preceding the exponential in the amplitude, a factor of no importance to us.

The asymptotic elastic-scattering amplitude was calculated in [2] by summing the contributions of the Mandelstam branch points. We made essential use of the alternation of the signs of the jumps at the branch points with neighboring numbers n. A similar alternation of signs takes place also in our case. Inasmuch as expression (5) for the trajectories j_{n+1} differs from the Mandelstam branch-point trajectories $n\alpha(t/n^2) - n + 1$ only by the constant number $\beta(o) - 1$, and the properties of the jumps at the discontinuities at the singularities (5) are the same as of the jumps at the Mandelstam branch points, all the calculations of [2] can be used in our case without change. The only difference will be the appearance of an additional factor in the amplitude, $s^{\beta(o)-1}$, corresponding to a shift of the condensation of the branch points in the j-plane by an amount $\beta(o) - 1$. For the cross section of any two-particle process we obtain the expression

$$\frac{d\sigma}{dt} = \frac{F(\xi,\tau)}{s^{2(1-\beta(o))}} \exp\{-2\sqrt{2\pi\tau\xi \operatorname{ctg}\phi/2}\,\Phi_1(\xi)\}(1 + \lambda\cos(2\sqrt{2\pi\tau\xi \operatorname{tg}\frac{\phi}{2}}\,\Phi_2(\xi) + \chi)), \quad \xi = \ln(s/m^2), \tag{6}$$

$$\tau = -\alpha'(o)t.$$

Here F, λ, χ, and ϕ are slowly varying functions of ξ and τ; the parameter ϕ is connected with the decrease of the jumps at the branch points with increasing number n. The functions $\Phi_1(\xi)$ and $\Phi_2(\xi)$ become equal to unity when $\xi \to \infty$ and are determined only by the trajectory of the vacuum pole and the value of ϕ. In principle, the parameter ϕ should be different for different processes. However, the case of particular interest from the theoretical point of view is $\phi = \pi/2$, which corresponds to not too rapid a drop in the jumps with increasing n and to the appearance of an exponential dependence of the cross section on the momentum transfer only owing to the alternation of the signs of the jumps at the neighboring branch points. If $\phi = \pi/2$

then the argument of the exponential in (6) turns out to be universal for all elastic and inelastic processes. An experimental confirmation of such a possibility would be highly interesting. The value $\phi = \pi/2$ agrees well, as shown in [2], with p-p scattering data.

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[2] A. A. Ansel'm and I. T. Dyatlov, YaF 6, 1967, in press.

[3] V. N. Gribov, I. Ya. Pomeranchuk, and K. A. Ter-Martirosyan, YaF 2, 361 (1965), Soviet JNP 2, 258 (1966).

*This condensation was pointed out, in connection with [3], by Gribov, Pomeranchuk, and Ter-Martirosyan back in 1964.

UNIVERSAL INTERACTION OF REGGE PARTICLES WITH VECTOR AND SCALAR CURRENTS

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Submitted 16 March 1967
ZhETF Pis'ma 5, No. 11, 407-409 (1 June 1967)

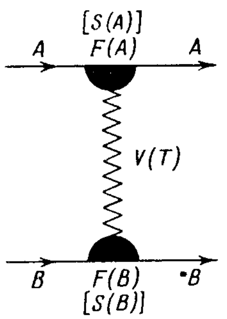
Recently Cabibbo, Horwitz, and Ne'eman [1] advanced the hypothesis that Regge particles corresponding to vector and tensor mesons (V and T reggeons) interact with vector and scalar currents:

$$F_i(A) = \langle A | \int d^3x\, \bar{q}(x) \gamma_0 \lambda_i q(x) | A \rangle, \tag{1}$$

$$S_i(A) = \langle A | \int d^3x\, \bar{q}(x) \lambda_i q(x) | A \rangle. \tag{2}$$

Here $F_i(A)$ and $S_i(A)$ are the vertices for the emission of a reggeon V_i or T_i by a particle A; the 4-momentum of the reggeon is assumed equal to zero; $i = 0, 1, \ldots, 8$ is the unitary index; $q(x) = (p, n, \lambda)$ is the quark field operator; γ_0 is a Dirac matrix; λ_i are Gell-Mann matrices ($\lambda_0 = \sqrt{2/3}\, \ell$); the matrix element in (2) is chosen, by definition, in the rest frame of particle A. The normalization of $|A\rangle$ is such that for quarks $F_i(q) = S_i(q) = \lambda_i$.

The contribution of the Regge pole to the scattering amplitude of particles A and B corresponds to the diagram shown in the figure, and the corresponding contribution to the total



The scattered particles exchange a V or T reggeon (wavy line). The reggeon emission vertices F(A), F(B) or S(A), S(B) are blackened.

cross section σ^t_{AB} from the V and T reggeons is, as is well known [2],

$$\sigma^t_{AB}(s) = s^{-1} \sum_{i=0}^{8} S_i(A) S_i(B) (s/\nu_i^s)^{\alpha_i^s} - F_i(A) F_i(B) (s/\nu_i^v)^{\alpha_i^v}. \tag{3}$$

Here s is the invariant square of the cms energy; α_i^V and α_i^S are the values of the trajectory