that they obey our requirements.

We do not know how to include spin effects in performing quantitative calculations with the diagrams. We may use "quark-counting" assumptions or assume SU(3) invariance for the various vertices, but such requirements do not tell us much about the dynamics of the intermediate baryon and meson resonances. It should be interesting to find out whether the diagrams or some extension of them can be utilized for a more explicit understanding of the hadronic spectrum.

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[1]R. Dolen, D. Horn, and C. Schmid, Phys. Rev. 166, 1768 (1968).

[2]G. Veneziano, Nuovo Cimento 57A, 190 (1968).

[3]Exotic resonances are mesons which do not belong to SU(3) singlets or octets and baryons which do not belong to SU(3) singlets, octets, or decuplets.

[4]P. G. O. Freund, Phys. Rev. Letters 20, 235 (1968); C. Schmid, Phys. Rev. Letters 20, 628 (1968); H. Harari, Phys. Rev. Letters 20, 1395 (1968); C. Schmid and J. Yellin, Phys. Letters 27B, 19 (1968); M. Kugler, to be published; H. J. Lipkin, to be published.

[5]G. F. Chew and A. Pignotti, Phys. Rev. Letters 20, 1078 (1968); H. Harari, in Proceedings of the Fourteenth International Conference on High Energy Physics, Vienna, Austria, September, 1968 (CERN Scientific Information Service, Geneva, Switzerland, 1968), p. 195; K. Bardakci and H. Ruegg, to be published; M. A. Virasoro, to be published; Chan Hong-Mo, to be published.

[6]J. L. Rosner, Phys. Rev. Letters 21, 950 (1968). See also H. J. Lipkin, Ref. 4.

[7]A similar assumption was proposed by H. J. Lipkin [Phys. Rev. Letters 16, 1015 (1966)] on pure phenomenological grounds, without appealing to the duality concept or the resonance-dominance hypothesis. See also J. J. J. Kokkedee and L. Van Hove, CERN Report No. CERN 66/248/5-TH642, 1966 (unpublished).

[8]G. Alexander, H. J. Lipkin, and F. Scheck, Phys. Rev. Letters 17, 412 (1966).

[9](i) $\sigma_t(K^+p)=\sigma_t(K^+n)$ holds experimentally between 2 and 20 GeV and leads to $\mathrm{Im}(K^+n\rightarrow K^0p)=0$. (ii) The production of φ mesons in πN scattering is known to be extremely rare. (iii) The $\Lambda\rho$ and $\Lambda\omega$ production rates in K^-p scattering are roughly equal at 4.1 and 5.5 BeV [J. Mott et al., Phys. Rev. Letters 18, 355 (1967)]. At 2.24 BeV only the forward ρ and ω cross sections are equal, and the angular distributions are completely different [G. W. London et al., Phys. Rev. 143, 1034 (1966)].

MODEL FOR THE POMERANCHUK TERM*

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A model of the Pomeranchuk term as a sum over two-body intermediate states is constructed. Experimental tests involving resonance production are suggested.

The Pomeranchuk singularity (P) that drives elastic high-energy scattering is the least understood of all complex angular-momentum plane singularities. The nature of lower lying singularities is reasonably clear. A simple picture in terms of straight-line Regge trajectories is found to explain well both their effects in high-energy scattering and the observed particle spectrum. Particularly relevant for us are the two approximately exchange-degenerate nonets of trajectories (which we shall jointly label by R)[1] on which the vector and tensor mesons are located. Is there any chance of understanding the nature of P starting from these "usual" R trajectories? It has been suggested[2] that the P contribution to an elastic process $AB\rightarrow AB$ is generated by the sequence $AB\rightarrow AB+$[secondaries produced by multiperipheral R exchange (multi-Regge exchange)] $\rightarrow AB$, as shown in Fig. 1(a). The difficulty involved in this approach is that of mathematical complexity in treating many-particle intermediate states. We want to show that the model of Fig. 1(a) can be simplified to one including only two-body intermediate states. Using duality[3] one can successively "reduce" the many secondary lines in Fig. 1(a) while at the same time including higher and higher excitations of the "elastic lines" [Fig. 1(b)]. For a large class of diagrams

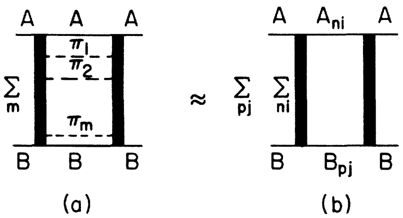

FIG. 1. (a) Multi-Regge diagrams: Heavy lines stand for usual R Regge poles; dotted lines stand for secondaries, say pions. (b) After repeated use of duality thus "reducing" all secondaries we obtain these "two-particle intermediate-state" diagrams. This diagram can be viewed as the box diagram with all four box lines being Regge poles.

such as that of Fig. 2(a) all secondaries can be reduced to one side [see Figs. 2(b) and 2(c)] without exciting the other side. The intermediate state is then $A_{ni}B$ rather than $A_{ni}B_{pj}$. We call such diagrams one sided. Diagrams of the type Fig. 2(d) cannot be fully reduced to either side [see Fig. 2(c)]. We call them two sided. When we then iterate these diagrams through unitarity (i.e., find the shadow they cast in the elastic process), diagrams with increasing number of secondaries will only contribute higher resonances and higher multiplicity decay modes of lower resonances in the intermediate state. As long as these intermediate-state resonances are sufficiently narrow, we may treat them as if they were stable and then the corresponding diagrams are precisely those of Fig. 1(b). To include both one-sided and two-sided diagrams, we have to sum over both ni and pj in Fig. 1(b).[4] We thus reduced the multi-Regge model to the mathematically much simpler "two-body intermediate-state" model of Fig. 1(b).

Before we go into more detail, let us make an important dynamical observation. By successively replacing R-trajectory exchanges by direct-channel resonances why do we stop at Fig. 1(b) rather than make one more step to get $AB \to \sum(AB)$ resonances $\to AB$? Were this possible, we would contradict the result that P is built completely of nonresonant s-channel background.[5] What prevents us from this last step is exchange degeneracy. Take for instance $\pi^+\pi^+$ scattering. Then at the level of Fig. 1(b) we have $\pi^+\pi^+ \to R^+R'^+$ (R^+, R'^+ = any meson resonance allowed by conservation laws) in which we can exchange only the ρ^0 and f trajectories. If these are exchange degenerate, then the $\pi^+\pi^+ \to R^+R'^+$ amplitude will be purely real and no further reduction is possible.

Now to the details of the model. For generality let us discuss the scattering of any two hadrons A and B. Consider the sequence of resonance pairs $A_{ni}B_{pj}$ such that the reaction $AB \to A_{ni}B_{pj}$ can proceed through usual R exchange. Let the resonance mass spectra be given by

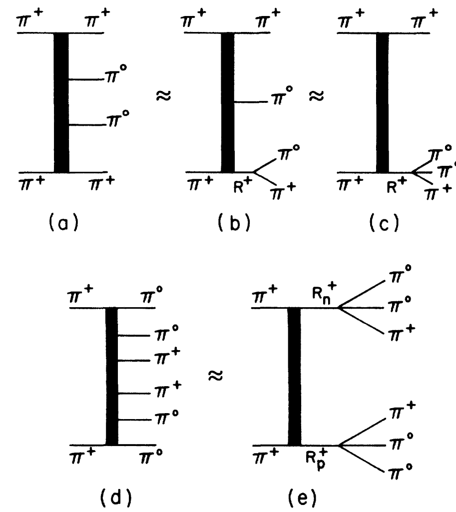

FIG. 2. (a) One-sided triperipheral Regge exchange, (b) reduced to biperipheral Regge exchange, (c) reduced to single Regge exchange in one-sided way. (d) Two-sided multiperipheral Regge exchange reducible only in the form (e) to single Regge exchange. Heavy lines stand again for usual R Regge poles.

$$m_{A_{ni}}^2 = m_{An}^2 = m_{A0}^2 + \frac{1}{\alpha'}n, \quad i = 1, \cdots, a_n;$$

$$m_{B_{pj}}^2 = m_{Bp}^2 = m_{B0}^2 + \frac{1}{\alpha'}n, \quad j = 1, \cdots, b_p, \tag{1}$$

so that at mass m_{An} (m_{Bp}) there are a_n (b_p) degenerate A_{ni}'s (B_{pj}'s) of various spins.

First of all we have to find the restrictions on the $A_{ni}B_{pj}$ intermediate state in order that it contribute appreciably [through diagrams of type Fig. 1(b)] to the Pomeranchuk amplitude $A_P(s,t)$. This contribution can be significant only if the $A_{ni}B_{pj}$ intermediate state can be produced with small momentum transfer $|t|_{\min} \approx 0$.

Now

$$|t|_{\min} \approx m_{An}^2 m_{Bp}^2 / s, \tag{2}$$

so that to the surprise of no one we have to require[4] that

$$m_{An}^2 m_{Bp}^2 \lesssim \mu^2 s, \tag{3}$$

where μ is some small constant mass. Then from Fig. 1(b)[6] (at $s \gg m_{A0}^2, m_{B0}^2$)

$$\mathrm{Im}A_P(s,0) \approx s^{2\alpha(0)-1} \sum_{n=0}^{\alpha' s} \sum_{p=0}^{\alpha'^2\mu^2 s/n} \sum_{i=1}^{a_n} \sum_{j=1}^{b_p} \beta_{ni,pj}(0), \quad (4)$$

where $\alpha(t)$ is the usual exchange-degenerate R trajectory. The $\beta_{ni,pj}$ are related to the residua of the usual R poles in $AB \to A_{ni}B_{pj}$ and can be determined experimentally. Indeed, let us set

$$\beta_{np}(0) \equiv \sum_{i,j} \beta_{ni,pj}(0) \approx \bar{\beta}(0)(np)^x \text{ for } np \gg 1. \quad (5)$$

Then[6]

$$\mathrm{Im}A_P(s,0) = \beta_P(0) s^{\alpha_P(0)}, \quad (6)$$

where $\beta_P(0)$ is related to $\bar{\beta}(0)$ and

$$\alpha_P(0) = 2\alpha(0) + x. \quad (7)$$

From $\alpha_P(0) \approx 1$, $\alpha(0) \approx \frac{1}{2}$ we predict

$$x \approx 0. \quad (8)$$

Consider now the process $AB \to A_{ni}B_{pj}$ and let $d\sigma_{ni,pj}/dt$ be the corresponding differential cross section. Define

$$\frac{d\sigma_{n,p}}{dt} = \sum_{i=1}^{a_n} \sum_{j=1}^{b_p} \frac{d\sigma_{ni,pj}}{dt}. \quad (9)$$

From Eqs. (5) and (9)[6,7]

$$\left.\frac{d\sigma_{n,p}}{dt}\right|_{t=0} = \mathrm{const} \times n^x p^x s^{2\alpha(0)-2}. \quad (10)$$

From the n and p dependence at fixed s of $d\sigma_{n,p}/dt|_{t=0}$ one can directly measure x and compare it with our prediction (8). Now let us say a few words about $t \neq 0$. Consider first the sum of one-sided diagrams [of type 1(b) with $pj = 00$]

$$A_P^{(0)}(s,t) = \sum_{n=0}^{\alpha' s} \sum_{i=1}^{a_n} A_{Pni,00}(s,t). \quad (11)$$

We write

$$\mathrm{Im}A_P^{(0)}(s,t) = \beta_P^{(0)}(t) s^{\alpha_P^{(0)}(t)}, \quad (12)$$

while for the full Pomeranchuk amplitude

$$\mathrm{Im}A_P(s,t) = \beta_P(t) s^{\alpha_P(t)}. \quad (13)$$

A straightforward specialization of the argument that led to Eq. (7) to one-sided graphs leads to

$$\alpha_P^{(0)}(0) = \alpha_P(0). \quad (14)$$

Furthermore in some small neighborhood $-|t_0| \leq t \leq 0$ all contributions to $\mathrm{Im}A_P(s,t)$ are coherent, so that in this neighborhood

$$\mathrm{Im}A_P(s,t) > \mathrm{Im}A_P^{(0)}(s,t) > 0. \quad (15)$$

If for small t we write

$$\alpha_P(t) \approx \alpha_P(0) + \alpha_P' t, \quad (16a)$$

$$\alpha_P^{(0)}(t) = \alpha_P^{(0)}(0) + \alpha_P^{(0)\prime} t, \quad (16b)$$

then from the relations (12)-(16) it follows that

$$\alpha_P' \leq \alpha_P^{(0)\prime}. \quad (17)$$

We shall now try to place an upper bound on $\alpha_P^{(0)\prime}$. From (11) and Fig. 1(b) we have[6]

$$\mathrm{Im}A_P^{(0)}(s,t) \approx \sum_{n=0}^{\alpha' s} \sum_{i=1}^{a_n} \beta_{ni,00}(t) s^{\alpha_n(t,s)}, \quad (18)$$

where

$$\alpha_n(t,s) = 2\alpha[\tfrac{1}{4}t(i - m_{An}^2/s)] - 1. \quad (19)$$

To proceed, we have to feed in some further information on $d\sigma_{n0}/dt$ [defined in Eq. (9)]. Indeed, assume that for $n \leq c\alpha' s$ ($c \leq 1$) we can write

$$d\sigma_{n0}/dt = \mathrm{const} \times n^{x-2y\alpha' t} s^{2\alpha(t)-2}. \quad (20)$$

From Eq. (20) we can determine

$$\beta_{ni}(t) = \sum_{i=1}^{a_n} \beta_{ni,00}(t)$$

and insert the result into the unitarity condition (18), whereupon we obtain a contribution to $\mathrm{Im}A_P^{(0)}(s,t)$ of the form[6] (12) with $\alpha_P^{(0)}(0)$ given by Eqs. (7) and (14) and the slope

$$\alpha_P^{(0)\prime} = \tfrac{1}{2}\alpha'(1-y)(1-c). \quad (21)$$

The Froissart bound ($x_P' \geq 0$) requires $y \leq 1$. Experimentally[8] the width of the forward peak in $d\sigma_{n0}/dt$ at fixed s appears to be either expanding or fixed but never shrinking as n increases. This means $y \geq 0$ in (20) [Eq. (20) actually predicts that the expansion of the peak is logarithmic in n]. For $0 \leq y \leq 1$ since $c \leq 1$, Eqs. (17) and (21) yield

$$\alpha_P' \leq \tfrac{1}{2}\alpha' \quad (22)$$

which seems to be obeyed in nature. We wish to re-emphasize here that the inequality (22) follows from our model only after making the additional assumptions (20) and $y \geq 0$.

The main point of this paper is that of reducing a complicated multiparticle problem to a much simpler two-body problem. Further work, especially at $t \neq 0$, is needed to obtain a full understanding of the Pomeranchuk singularity in this model.[9]

I wish to thank Professor R. Oehme and Professor R. Arnold for valuable discussions.

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*Work supported in part by U. S. Atomic Energy Commission.

[1]This should not be confused with the redundant label R attached sometimes to the A_2 trajectory.

[2]G. F. Chew and A. Pignotti, to be published; A. Pignotti, private communication.

[3]G. F. Chew and A. Pignotti, Phys. Rev. Letters 20, 1078 (1968).

[4]I wish to acknowledge an illuminating conversation with Dr. G. Veneziano on this point.

[5]P. G. O. Freund, Phys. Rev. Letters 20, 235 (1968); in this paper the result is stated in its full generality and not only for the case of $\pi\pi$ scattering [see especially Aachen-Berlin-CERN-London (I.C.)-Vienna Collaboration, Phys. Letters 27B, 336 (1968) and the discussion to which it refers]. A restatement of this result is contained in H. Harari, Phys. Rev. Letters 20, 1385 (1968).

[6]Throughout this paper we disregard factors like $(\ln s)^{\lambda_1}(\ln \ln s)^{\lambda_2}\cdots$ and concentrate only on powers of s and n. Logarithmic factors can always be included by suitably modifying the *Ansatz* (6).

[7]Assuming the width of the forward peak at fixed s due to R exchange to depend at most logarithmically on n.

[8]Aachen-Berlin-CERN-London (I.C.)-Vienna Collaboration, Ref. 5.

[9]The approach opposite to that taken in this paper and in Refs. 2 and 3 would be a model in which the Pomeranchuk term bootstraps by itself, independently of any lower Regge trajectories. A two-body intermediate-state reduction may prove useful even for such a model but the resonances ought to be of the type that can be produced by diffraction dissociation.

HIGH-ENERGY GAMMA RAYS, COSMIC-RAY ELECTRONS, AND THE FAR-INFRARED BACKGROUND RADIATION*

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It is shown that the observed flux of galactic gamma rays, which is one order of magnitude higher than the theoretical value, can be explained if one assumes the recently detected far-infrared radiation to be galactic or universal. Effects of the infrared radiation on cosmic rays are discussed.

Shivanandan, Houck, and Harwit[1] have recently detected a background infrared radiation with intensity 5×10^{-2} erg/cm^2 sec sr (within a factor of 2) in the spectral range from 0.4 to 1.3 mm. We wish to point out that inverse Compton scattering of such radiation, if it is universal, or at least galactic, would provide a natural explanation of the otherwise unaccountable gamma flux observed recently by Clark, Garmire, and Kraushaar.[2] Some secondary evidence supporting the existence of a background infrared radiation in the galaxy can also be found from the spectrum of cosmic electrons observed at Earth. The gamma-ray data obtained by Clark, Garmire, and Kraushaar can be summarized as follows. A diffuse photon flux of 2×10^{-4} (cm^2 sec rad)$^{-1}$ was detected in coincidence with the plane of the galaxy at energy $E \geq 100$ MeV (the instrumental limit). The intensity of the radiation in this plane has a broad maximum [5×10^{-4} (cm^2 sec rad)$^{-1}$] toward the galactic center. With less certainty an isotropic background intensity of 1×10^{-4} (cm^2 sec sr)$^{-1}$ was also observed.

Theoretically, the primary processes for generating high-energy photons in space are the decay of π^0 mesons produced by cosmic-ray collisions, and the bremsstrahlung and Compton scattering of relativistic electrons.[3] The predicted γ-ray intensity due to π^0 decay and bremsstrahlung are more than one order of magnitude less than the observed intensity. One should notice that $I_{\pi^0}(E_\gamma)$, the γ-ray flux due to π^0 decay, and to a lesser degree $I_B(E_\gamma)$, the bremsstrahlung flux, are closely related to the cosmic-ray positron intensity. They are all proportional to the interstellar mass density. Assuming equal pro-