dynamical philosophy of duality. The approach of Schwarz is fairly simple, while the other three models are rather complex. However, the properties common to all are (i) crossing symmetry, (ii) the lack of $I=2$ resonances, (iii) the presence of zeros below threshold,[41] and (iv) the ρ meson with more-or-less correct parameters. The treatment of unitarity varies widely.

It seems therefore that as far as a_0 and a_2 are concerned, the essential ingredients in these models are not the particular dynamical philosophies (interference or duality) nor the explicit satisfaction of unitarity, but rather the presence of the four features listed above which are the essential characteristics of the present model. We expect them to be common to all successful descriptions of low-energy $\pi\pi$ scattering.

[41] The zeros are not in the same place in each model—also they arise from apparently rather different (but presumably deeply connected) sources.

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Behavior of Baryon-Baryon and Baryon-Antibaryon Total Cross Sections at High Energy*

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The energy-dependent parts of baryon-baryon and baryon-antibaryon total cross sections at high energies are obtained from (i) factorization of the residues of the leading meson trajectories, and (ii) the connection between the absence of resonances in "exotic" channels and the flatness of meson-meson and meson-baryon total cross sections in these channels. (Such a connection leads to exchange degeneracies, all of which are consistent with experiment.) The following is proven: (a) Nucleon-nucleon (and certain other $B=2$) total cross sections are flat whether or not trajectory couplings obey $SU(3)$; (b) the remaining $B=2$ cross sections are flat if tensor and vector exchanges couple via $SU(3)$, with equal F/D ratios at the baryon vertex; (c) non-Pomeranchuk contributions to baryon-antibaryon total cross sections persist in the **10**, $\overline{\mathbf{10}}$, and **27** channels despite the relaxation of assumptions regarding $B=2$ systems. [Results similar to (a) and (b) have been *assumed* to hold in previous approaches.]

I. INTRODUCTION

In the Regge-pole description of high-energy nucleon-nucleon interactions, the large real parts of the NN forward amplitude[1] and the rapid decrease of $\sigma_T(N\bar{N})$ with increasing energy[2] indicate that sizable non-Pomeranchuk contributions are present above 8 GeV/c. The relative flatness of $\sigma_T(NN)$ in this energy range indicates, however, that the contributions of + and − signature non-Pomeranchuk trajectories to the imaginary part of the forward spin-averaged NN amplitude almost cancel. This cancellation has been termed exchange degeneracy; it entails a particular "mixture of mesons" appropriate for describing the NN interaction in the combination of helicity states contributing to σ_T.[3]

The usual explanation of the near degeneracy of + and − signature contributions to $\sigma_T(BB)$ assumes the comparative weakness of $B=2$ s-channel forces as compared with presumably much stronger $B=0$ u-channel forces. From this assumption about the forces,

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one can easily derive exchange degeneracy and the flatness of the $B=2$ total cross sections.

In this paper we derive the flatness of baryon-baryon total cross sections directly from knowledge of the high-energy behavior of σ_T in meson-meson and meson-baryon channels. A primary assumption is the connection between flat total cross sections and the absence of resonances. Recently the validity of such constraints for the scattering of particles with high mass (such as BB or $B\bar{B}$) has been called into doubt.[4] Our results are clearly insensitive to such considerations since we have been careful *not* to apply constraints from $B=2$ or baryon-antibaryon channels. We use constraints only from meson-meson and meson-baryon channels, which lead to experimentally verified exchange degeneracies. No assumption about the baryon-baryon force is made, in contrast to previous approaches.[5–8] We find that the nucleon-nucleon (and certain other $B=2$) total cross sections are flat as a consequence of factorization of the leading meson trajectories, and that the remaining $B=2$ cross sections are flat if the ϕ and f' decouple from the nucleon.

The results are derived in Sec. II. We use only constraints arising from elastic meson-meson and meson-baryon reactions, but assume $\alpha_\phi<\alpha_\omega$ and $\alpha_{f'}<\alpha_f$. The results for both $SU(2)$- and $SU(3)$-invariant couplings are given. We also confirm previously obtained results for baryon-antibaryon total cross sections[6–8] without the need for assumptions regarding $B=2$ systems.[7,9]

The results of Sec. II are rederived in Sec. III under a different set of assumptions. We consider the constraints from both elastic and inelastic MM and MB reactions, but do not assume $\alpha_\phi<\alpha_\omega$ or $\alpha_{f'}<\alpha_f$. The resulting set of equations is identical to that of Sec. II for $SU(3)$-invariant couplings.

In Sec. IV, we discuss the validity of the constraint equations and state our conclusions.

II. NON-POMERANCHUK CONTRIBUTIONS TO TOTAL CROSS SECTIONS: TRAJECTORIES SPLIT

The dominant non-Pomeranchuk contributions to meson-meson and meson-baryon total cross sections come from the exchange of the vector (1^-) and tensor (2^+) meson Regge trajectories. We assume that these Regge poles have factorizable residues. Then the total cross section for the scattering of particle A on B is given by

$$\sigma_T(AB)=\sigma_T(AB)_{\text{pom}}+\text{Im}A/p_{\text{lab}}\,, \tag{1}$$

where ImA is the non-Pomeranchuk amplitude to which we shall be referring. It is given by

$$\text{Im}A=\sum_E g(EA\bar{A})g(EB\bar{B})\tau_E(\nu/\nu_0)^{\alpha_E}. \tag{2}$$

In Eq. (2) the sum is over the Regge exchanges E. The signature factor τ_E is $+1$ for the tensor exchanges (f, A_2^0, and f') and is -1 for the vector exchanges (ρ, ω, and ϕ).[10] The $g(EA\bar{A})$ are factorizable couplings, ν_0 is a common scale factor, and $\nu\equiv\frac{1}{2}(s-u)$.

Our main assumption is that total cross sections are flat[11] in exotic meson-meson and meson-baryon channels. Exotic channels are defined to be ones with little resonant activity. They are **10**, $\overline{\mathbf{10}}$, and **27** for meson-meson scattering, $\overline{\mathbf{10}}$ and **27** for meson-baryon scattering, and **27** and **35** for meson-decimet scattering. Where σ_T is measured, it is found to be flat in these channels. We assume that this correlation is true in general for meson-meson and meson-baryon systems.

To illustrate our method, we consider $\sigma_T(K^+K^+)$, $\sigma_T(K^+K^0)$, $\sigma_T(K^+p)$, and $\sigma_T(K^+n)$. Because these σ_T are flat regardless of the nucleon's or K's isospin, the net $I=0$ and $I=1$ non-Pomeranchuk contributions must vanish separately. If we assume for simplicity that $\alpha_{f'}<\alpha_f$ and $\alpha_\phi<\alpha_\omega$, we then obtain $\alpha_f=\alpha_\omega$, $\alpha_{f'}=\alpha_\phi$, and $\alpha_{A_2}=\alpha_\rho$; the coupling relations are $g(fN\bar{N})=g(\omega N\bar{N})$, $g(f'N\bar{N})=g(\phi N\bar{N})$, and $g(A_2N\bar{N})=g(\rho N\bar{N})$.[5,12] These conditions guarantee that $\sigma_T(pp)$ and $\sigma_T(pn)$ will be flat.

TABLE I. Nonet couplings $g(EA\bar{A})$ consistent with flatness of σ_T in exotic $PP\to PP$, $PB\to PB$, and $PD\to PD$ channels. Coefficients refer to the $SU(3)$ universality limit (in which $a=b=c=d=f=\cdots$ and $e=0$). The couplings of Y^{*+} and Ξ^{*0} are similar in form to those of Σ^+ and Ξ^0. Other couplings are related by charge conjugation or isospin to those of this table. The couplings of η and Λ are unconstrained. In this and subsequent tables, one must have $\alpha_{P'}=\alpha_{A_2}=\alpha_\omega=\alpha_\rho$ and $\alpha_{f'}=\alpha_\phi$.

A \ E	f	f'	A_2	ω	ϕ	ρ
π^+	$2a$	0	0	0	0	$2a$
K^+	b	$c\sqrt{2}$	b	b	$-c\sqrt{2}$	b
p	$3d$	e	f	$3d$	$-e$	f
Σ^+	$2g$	$h\sqrt{2}$	$2j$	$2j$	$h\sqrt{2}$	$2g$
Ξ^0	m	$2n\sqrt{2}$	p	p	$2n\sqrt{2}$	m
Δ^{++}	$3q$	0	$3q$	$3q$	0	$3q$
Ω^-	0	$3s\sqrt{2}$	0	0	$3s\sqrt{2}$	0

[4] J. Mandula, J. Weyers, and G. Zweig, Phys. Rev. Letters **23**, 266 (1969). A partial earlier statement of this result is given by R. H. Capps, Ref. 21.

[5] C. B. Chiu and J. Finkelstein, Phys. Letters **27B**, 510 (1968).

[6] J. Rosner, Phys. Rev. Letters **21**, 950 (1968).

[7] D. P. Roy and M. Suzuki, Phys. Letters **28B**, 576 (1969).

[8] H. J. Lipkin, Nucl. Phys. **B9**, 349 (1969).

[9] The experimental difficulty of obtaining a complete pattern of $B=2$ activity makes this especially important.

[10] This may be illustrated by comparing $\sigma_T(pp)$ with $\sigma_T(\bar{p}p)$. The latter displays sizable non-Pomeranchuk parts whereas the former does not. In the t channel, the former process is $\bar{p}p\to p\bar{p}$, for which the + and − signature poles contribute with opposite sign, while the latter is $pp\to\bar{p}\bar{p}$, for which they add.

[11] In general, we shall speak of a total cross section as "flat" if it lacks non-Pomeranchuk contributions, while realizing that this term may be an oversimplification if the Pomeranchuk singularity is in fact a cut giving rise to slow logarithmic variation of σ_T. The motivation for assuming flatness combines the well-documented scarcity of exotic meson-meson and meson-baryon resonances [e.g., A. H. Rosenfeld, in *Meson Spectroscopy* (W. A. Benjamin, Inc., New York, 1968), p. 455] with the conjecture that this scarcity is associated with flatness of σ_T. See P. G. O. Freund, Phys. Rev. Letters **20**, 235 (1968); Haim Harari, *ibid.* **30**, 1395 (1968).

[12] An over-all sign can be absorbed into the definition of the vector nonet and will not change our result.

TABLE II. Couplings of Table I satisfying $SU(3)$ symmetry. P couplings are to be multiplied by a factor γ_P, B couplings by γ_B, and D couplings by γ_D.

$A \diagdown E$	f	f'	A_2	ω	ϕ	ρ
π^+	2	0	0	0	0	2
K^+	1	$\sqrt{2}$	1	1	$-\sqrt{2}$	1
η	$\frac{2}{3}$	$4\sqrt{2}/3$	0	0	0	0
p	$2F_T+2F_V-1$	$(F_V-F_T)\sqrt{2}$	1	$2F_T+2F_V-1$	$(F_T-F_V)\sqrt{2}$	1
Σ^+	$2F_V$	$(F_T+F_V-1)\sqrt{2}$	$2F_T$	$2F_T$	$(F_T+F_V-1)\sqrt{2}$	$2F_V$
Λ	$2F_V+\frac{4}{3}F_T-\frac{4}{3}$	$[F_V+\frac{1}{3}(1-F_T)]\sqrt{2}$	0	$2F_T+\frac{4}{3}F_V-\frac{4}{3}$	$[F_T+\frac{1}{3}(1-F_V)]\sqrt{2}$	0
Ξ^0	$2F_V-1$	$(F_T+F_V)\sqrt{2}$	$2F_T-1$	$2F_T-1$	$(F_T+F_V)\sqrt{2}$	$2F_V-1$
Δ^{++}	3	0	3	3	0	3
Y^{*+}	2	$\sqrt{2}$	2	2	$\sqrt{2}$	2
Ξ^{*0}	1	$2\sqrt{2}$	1	1	$2\sqrt{2}$	1
Ω^-	0	$3\sqrt{2}$	0	0	$3\sqrt{2}$	0

Assumptions about meson-meson and meson-baryon scattering are not enough to guarantee that all the $B=2$ total cross sections are flat. The nucleon-nucleon system, belonging to the $\overline{\mathbf{10}}$ ($I=0$) and **27** ($I=1$) of $SU(3)$, is a special case. In general, the other baryon-baryon cross sections will not be flat unless the f' and ϕ decouple from nucleons. We now show this result.

Let us first assume $\alpha_\phi<\alpha_\omega$ and $\alpha_{f'}<\alpha_f$, and only $SU(2)$ symmetry for couplings. Then the couplings in Table I follow from assuming flatness of σ_T in exotic PP, PB, and PD systems.[13] Because of the constraints $\alpha_\phi<\alpha_\omega$ and $\alpha_{f'}<\alpha_f$, the contributions of the f' and the ϕ to the total cross section must cancel separately from the other contributions.[5]

These couplings lead to flatness of σ_T in all the BB, BD, and DD systems that cannot couple to the **1**, **8**, and **10** of $SU(3)$. For example, $\sigma_T(pn)$, $\sigma_T(pp)$, and $\sigma_T(\Sigma^+\Sigma^+)$ are predicted to be flat, while $\sigma_T(\Sigma^-p)$ and $\sigma_T(\Lambda p)$ are not, as is seen from Table I and Eq. (2).

TABLE III. S-channel contributions to ImA in PP, PB, and PD scattering. In this and subsequent tables we neglect $SU(3)$ breaking due to nondegenerate trajectories. We omit an over-all $\gamma_P{}^2(\nu/\nu_0)^\alpha$ for PP scattering, $\gamma_P\gamma_B(\nu/\nu_0)^\alpha$ for PB, and $\gamma_P\gamma_D(\nu/\nu_0)^\alpha$ for PD. Normalizations are those of Ref. 17.

PP: $\mathrm{Im}A\,(\mathbf{27})=\mathrm{Im}A\,(\mathbf{10})=\mathrm{Im}A\,(\overline{\mathbf{10}})\equiv0$
$\mathrm{Im}A\,(\mathbf{8}_{ss})=20/3$
$\mathrm{Im}A\,(\mathbf{8}_{aa})=12$
$\mathrm{Im}A\,(\mathbf{1})=64/3$
$\mathrm{Im}A\,(\mathbf{8}_{sa})=\mathrm{Im}A\,(\mathbf{8}_{as})=0$

PB: $\mathrm{Im}A\,(\mathbf{27})=\mathrm{Im}A\,(\overline{\mathbf{10}})\equiv0$
$\mathrm{Im}A\,(\mathbf{10})=8(F_V+F_T-1)$
$\mathrm{Im}A\,(\mathbf{8}_{ss})=10F_V+(10/3)F_T-10/3$
$\mathrm{Im}A\,(\mathbf{8}_{aa})=10F_V-2F_T+2$
$\mathrm{Im}A\,(\mathbf{1})=16F_V-(16/3)F_T+16/3$
$\mathrm{Im}A\,(\mathbf{8}_{sa})=\mathrm{Im}A\,(\mathbf{8}_{as})=(2\sqrt{5})(F_T-F_V+1)$

PD: $\mathrm{Im}A\,(\mathbf{35})=\mathrm{Im}A\,(\mathbf{27})\equiv0$
$\mathrm{Im}A\,(\mathbf{10})=16$
$\mathrm{Im}A\,(\mathbf{8})=20$

[13] We use the familiar notation $B=\frac{1}{2}^+$ octet, $D=\frac{3}{2}^+$ decimet, $P=0^-$ octet, $V=1^-$ nonet, $T=2^+$ nonet, and so on.

When we impose $SU(3)$ for the couplings[14] we find the results listed in Table II. Note that for $F_T=F_V$, the f' and ϕ decouple from nucleons. Both experiment[15] and the quark model[16] favor such a decoupling.

The non-Pomeranchuk parts of the forward spin-averaged amplitude for meson-meson, meson-baryon, and meson-decimet scattering are given in Table III. The use of $\equiv$, as in $\mathrm{Im}A_{PP}(\mathbf{27})\equiv0$, indicates an assumed constraint. The rest of the tables are derived from these constraints. Table IV lists the imaginary parts of the elastic BB, BD, and DD amplitudes, while those of baryon-antibaryon amplitudes are given in Table V.

The patterns of flat and decreasing total cross sections are essentially those found previously under more restrictive assumptions.[6] As mentioned, the condition $F_T=F_V$ which decouples f' and ϕ from nucleons also guarantees that all baryon-baryon total cross sections will be flat. Here we do not have to assume the absence of activity in $B=2$ channels.

TABLE IV. S-channel $B=2$ contributions to ImA. We omit an over-all $\gamma_B{}^2(\nu/\nu_0)^\alpha$, $\gamma_B\gamma_D(\nu/\nu_0)^\alpha$, or $\gamma_D{}^2(\nu/\nu_0)^\alpha$, respectively.

BB: $\mathrm{Im}A\,(\mathbf{27})=\mathrm{Im}A\,(\overline{\mathbf{10}})=0$
$\mathrm{Im}A\,(\mathbf{10})=16(F_V-F_T)(F_T+F_V-1)$
$\mathrm{Im}A\,(\mathbf{8}_{ss})=(20/3)(F_V-F_T)(2F_T+2F_V-1)$
$\mathrm{Im}A\,(\mathbf{8}_{aa})=4(F_V-F_T)(2F_T+2F_V+1)$
$\mathrm{Im}A\,(\mathbf{1})=(32/3)(F_V-F_T)(F_T+F_V+1)$
$\mathrm{Im}A\,(\mathbf{8}_{sa})=\mathrm{Im}A\,(\mathbf{8}_{as})=0$

BD: $\mathrm{Im}A\,(\mathbf{35})=\mathrm{Im}A\,(\mathbf{27})=0$
$\mathrm{Im}A\,(\mathbf{10})=16(F_V-F_T)$
$\mathrm{Im}A\,(\mathbf{8})=20(F_V-F_T)$

DD: $\mathrm{Im}A\,(\mathbf{35})=\mathrm{Im}A\,(\mathbf{28})=\mathrm{Im}A\,(\mathbf{27})=\mathrm{Im}A\,(\overline{\mathbf{10}})=0$

[14] V. Barger, M. Olsson, and K. V. L. Sarma, Phys. Rev. **147**, 1115 (1966).

[15] See, for example, V. Barger, in *Proceedings of the Topical Conference on High Energy Collisions of Hadrons, CERN, 1968* (Scientific Information Service, Geneva, 1968).

[16] G. Zweig, CERN Report No. CERN th-402, 1964 (unpublished).

TABLE V. S-channel baryon-antibaryon contributions to ImA. Normalizations as in Table IV.

$B\bar{B}$:	$\mathrm{Im}A(\mathbf{27})=4(F_T+F_V-1)^2$
	$\mathrm{Im}A(\mathbf{10})=\mathrm{Im}A(\overline{\mathbf{10}})=4(F_T+F_V+1)(F_T+F_V-1)$
	$\mathrm{Im}A(\mathbf{8}_{ss})=(32/3)(F_T^2+F_V^2)+8F_TF_V-(4/3)(F_T+F_V)-8/3$
	$\mathrm{Im}A(\mathbf{8}_{aa})=16(F_T^2+F_V^2)+8F_TF_V-12(F_T+F_V)+8$
	$\mathrm{Im}A(\mathbf{1})=(76/3)(F_T^2+F_V^2)+8F_TF_V-(56/3)(F_T+F_V)+44/3$
	$\mathrm{Im}A(\mathbf{8}_{sa})=\mathrm{Im}A(\mathbf{8}_{as})=(4\sqrt{5})(F_T^2+F_V^2-F_T-F_V)$
$B\bar{D}$:	$\mathrm{Im}A(\overline{\mathbf{35}})=0$
	$\mathrm{Im}A(\mathbf{27})=16(F_T+F_V-1)$
	$\mathrm{Im}A(\overline{\mathbf{10}})=8(F_T+F_V+1)$
	$\mathrm{Im}A(\mathbf{8})=16(F_T+F_V)+4$
$D\bar{D}$:	$\mathrm{Im}A(\mathbf{64})=0$
	$\mathrm{Im}A(\mathbf{27})=28$
	$\mathrm{Im}A(\mathbf{8})=48$
	$\mathrm{Im}A(\mathbf{1})=60$

The choice $F_T+F_V=1$ eliminates activity in the $\mathbf{10}$, $\overline{\mathbf{10}}$, and $\mathbf{27}$ $B\bar{B}$ channels. However, this gives unacceptable fits to high-energy data.[14,17] For example, experiment shows that [18]

$$\frac{\sigma_T(\pi^+p)-\sigma_T(\pi N)_\infty}{\sigma_T(\pi^-p)-\sigma_T(\pi N)_\infty}>\frac{1}{2}, \tag{3}$$

indicating that $F_T+F_V>2$, since the ratio in Eq. (3) is predicted to be $(F_T+F_V-1)/(F_T+F_V)$. If F_T+F_V were equal to 1, then $\sigma_T(\pi^+p)$ would be as flat as $\sigma_T(K^+p)$ or $\sigma_T(K^+n)$, which is certainly not the case. Independent information about F_T and F_V may be obtained, for example, from combinations of $\sigma_T(K^-p)$, $\sigma_T(K^-n)$, $\sigma_T(K^+n)$, and $\sigma_T(K^+p)$, from which we find that both F_T and F_V are at least 1.

In the elastic $B\bar{D}$ case, the choice $F_T+F_V=1$ eliminates the $\mathbf{27}$ activity, while the choice $F_T+F_V=-1$ eliminates the $\overline{\mathbf{10}}$ activity.[19] No choice can eliminate both.[20] The activity in the $\mathbf{27}$ of $D\bar{D}$ cannot be eliminated.[21]

The presence of activity in exotic $B\bar{B}$ channels implies large non-Pomeranchuk contributions to such total cross sections as $\sigma_T(\bar{\Sigma}^+p)$. In the limit $F_T=F_V$ one predicts, for example, that

$$\frac{\sigma_T(\bar{\Sigma}^+p)-\sigma_T(\Sigma^-p)}{\sigma_T(\bar{p}n)-\sigma_T(pn)}=\frac{1}{2}. \tag{4}$$

The denominator is almost 20 mb at 6 GeV/c.[2]

III. NON-POMERANCHUK CONTRIBUTIONS TO TOTAL CROSS SECTIONS: TRAJECTORIES DEGENERATE

In the limit of degenerate trajectories we appear to lose information because the negative-signature $SU(3)$ singlet trajectory must decouple from the pseudoscalar octet P merely by virtue of charge-conjugation invariance. Hence the channels considered in Sec. II will not tell us anything about the coupling of this trajectory to baryons. The lost information is recovered from the additional process

$$MM\to M'M',\quad MB\to M'B,\quad \text{and}\quad MD\to M'D, \tag{5}$$

where M and M' are two appropriately chosen *different* octets. The results are again those of Sec. II.

It is convenient to use the following method when starting with exact $SU(3)$ for both couplings and intercepts.[22] Consider the scattering of two pseudoscalar octets. The contribution of the leading crossed-channel Regge trajectories to the imaginary part of the amplitude is

$$\mathrm{Im}A_s(N)=\sum_{N'}(X_{st})_{NN'}m_{\tau N'}^2(N')(s/s_0)^{\alpha(t)}, \tag{6a}$$

for large s, and for large u, it is

$$\mathrm{Im}A_u(N)=\sum_{N'}(X_{ut})_{NN'}\tau_{N'}m_{\tau N'}^2(N')(u/u_0)^{\alpha(t)}, \tag{6b}$$

where N and N' are $SU(3)$ labels, $(X)_{NN'}$ is an element of the $SU(3)$ crossing matrix,[23] $m_{\tau_N}(N)$ is a factorized coupling of a trajectory of signature τ_N to the external octet, and $\alpha(t)$ is the common trajectory function.

The leading Regge trajectories which couple to the external pseudoscalar octets are a singlet and octet of positive signature, and an octet of negative signature. The contributions of these trajectories to Im$A(\mathbf{10})$, Im$A(\overline{\mathbf{10}})$, and Im$A(\mathbf{27})$ cancel if

$$\tfrac{1}{8}m_+^2(1)=\tfrac{2}{5}m_+^2(8_s)=(2/9)m_-^2(8_a). \tag{7}$$

Let us apply the same procedure to meson-baryon scattering. The vanishing of Im$A(\overline{\mathbf{10}})$ and Im$A(\mathbf{27})$ for $PB\to PB$ implies

$$\tfrac{1}{8}b_+(1)m_+(1)+(\sqrt{\tfrac{1}{5}})b_+(8_a)m_+(8_s)-(\sqrt{\tfrac{1}{5}})b_-(8_s)m_-(8_a)-\tfrac{2}{5}b_+(8_s)m_+(8_s)=0, \tag{8}$$

$$\tfrac{1}{8}b_+(1)m_+(1)+\tfrac{1}{5}b_+(8_s)m_+(8_s)-\tfrac{1}{3}b_-(8_a)m_-(8_a)=0, \tag{9}$$

where $b_\tau(N)$ is the coupling of the Regge trajectory to the external $B\bar{B}$ states. The conditions Im$A(\mathbf{27})$ $=$Im$A(\mathbf{35})=0$ in PD elastic channels give

$$(1/4\sqrt{5})d_+(1)m_+(1)+(3/5\sqrt{2})d_+(8)m_+(8_s)+(1/3\sqrt{10})d_-(8)m_-(8_a)=0, \tag{10}$$

$$(1/4\sqrt{5})d_+(1)m_+(1)-(1/5\sqrt{2})d_+(8)m_+(8_s)-(1/\sqrt{10})d_-(8)m_-(8_a)=0, \tag{11}$$

where $d_\tau(N)$ is the factorized coupling to external $D\bar{D}$ states.

[17] G. Renninger and K. V. L. Sarma, Phys. Rev. Letters **20**, 399 (1968).

[18] K. J. Foley, R. S. Jones, S. J. Lindenbaum, W. A. Love, S. Ozaki, E. D. Platner, C. A. Quarles, and E. H. Willen, Phys. Rev. Letters **19**, 330 (1967).

[19] This last result differs from one obtained in Ref. 7.

[20] M. Kugler, Phys. Rev. **180**, 1538 (1969).

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[22] J. Mandula, C. Rebbi, R. Slansky, J. Weyers, and G. Zweig, Phys. Rev. Letters **22**, 1147 (1969).

[23] See, for example, C. Rebbi and R. Slansky, Rev. Mod. Phys. **42**, 68 (1970).

The above equations do not determine $b_-(1)$ or $d_-(1)$. These are obtained by choosing M and M', in Eq. (5), to have charge conjugation opposite to that of a similar PP state.[24] Then the octet trajectory of positive signature couples via an antisymmetric coupling denoted by $\tilde{m}_+(8_a)$, the octet trajectory of negative signature couples with a symmetric coupling $\tilde{m}_-(8_s)$, the singlet positive-signature trajectory decouples $[\tilde{m}_+(1)=0]$, and a singlet trajectory of negative signature is exchanged $[\tilde{m}_-(1)\neq 0]$.

By requiring the cancellation of the imaginary parts in the exotic channels of $MM \to M'M'$, $MB \to M'B$, and $MD \to M'D$, we obtain equations identical to Eqs. (7)–(11), except for the replacement everywhere of m by $\tilde{m}$ and the interchange of $+$ and $-$ signature labels. Such equations involve the couplings $b_-(1)$ and $d_-(1)$, and thus give new information.

Equations (7)–(11) and their analog for reactions (5) are easily solved. A parametrization that we find convenient[25] is

$$\begin{aligned} m_+(1)&=-(8/3)\sqrt{3}\gamma_P\,,\\ m_+(8_s)&=\tfrac{2}{3}(\sqrt{15})\gamma_P\,,\\ m_-(8_a)&=-2\sqrt{3}\gamma_P\,; \end{aligned} \tag{12}$$

$$\begin{aligned} \tilde{m}_+(1)&=-(8/3)\sqrt{3}\gamma_M\,,\\ \tilde{m}_-(8_s)&=\tfrac{2}{3}(\sqrt{15})\gamma_M\,,\\ \tilde{m}_+(8_a)&=-2\sqrt{3}\gamma_M\,; \end{aligned} \tag{13}$$

$$\begin{aligned} b_+(1)&=[\tfrac{4}{3}\sqrt{3}(1-F_T)-4\sqrt{3}F_V]\gamma_B\,,\\ b_-(1)&=[\tfrac{4}{3}\sqrt{3}(1-F_V)-4\sqrt{3}F_T]\gamma_B\,,\\ b_+(8_s)&=\tfrac{2}{3}(\sqrt{15})(1-F_T)\gamma_B\,,\\ b_+(8_a)&=-2\sqrt{3}F_T\gamma_B\,,\\ b_-(8_s)&=\tfrac{2}{5}(\sqrt{15})(1-F_V)\gamma_B\,,\\ b_-(8_a)&=-2\sqrt{3}F_V\gamma_B\,; \end{aligned} \tag{14}$$

$$\begin{aligned} d_+(1)&=d_-(1)=2(\sqrt{15})\gamma_D\,,\\ d_+(8)&=d_-(8)=(\sqrt{30})\gamma_D\,. \end{aligned} \tag{15}$$

The couplings (12), (14), and (15) can be used in Eq. (6) to compute the imaginary parts of the remaining amplitudes. The results are identical to those of Sec. II, given in Tables III–V.

Table II can be reconstructed from Eqs. (12), (14), and (15) using the vector-coupling coefficients for $SU(3)$[26] and the usual "ideal" mixing[5]:

$$\left|\begin{matrix}f\\ \omega\end{matrix}\right\rangle=(\sqrt{\tfrac{2}{3}})|1,\tau=\pm\rangle+(\sqrt{\tfrac{1}{3}})|8,\tau=\pm\rangle\,, \tag{16}$$

$$\left|\begin{matrix}f'\\ \phi\end{matrix}\right\rangle=(\sqrt{\tfrac{1}{3}})|1,\tau=\pm\rangle-(\sqrt{\tfrac{2}{3}})|8,\tau=\pm\rangle\,. \tag{17}$$

IV. DISCUSSION AND CONCLUSIONS

We should stress that the relative lack of information on BB resonances has led some authors[7] to speculate that activity in exotic $B\bar{B}$ channels could be eliminated at the price of introducing it in BB channels. Our result shows this not to be the case. In particular (as is clear from Table IV), we do obtain activity in certain baryon-baryon channels when F_T and F_V are not exactly equal. [As the discussion below Eq. (3) indicates, this is irrelevant to the activity in exotic baryon-antibaryon channels, which we show to be present in any case.]

The factorization hypothesis is crucial to our discussion. Suppose there existed trajectories degenerate with the leading meson trajectories but coupling only to B and D. It would then appear possible to achieve nearly any pattern of total cross sections. However, unless the residues of the additional trajectories were to change sign between the t-channel resonance region and $t=0$, they would contribute more to baryon-antibaryon total cross sections than to baryon-baryon total cross sections, and would therefore not help to make exotic baryon-antibaryon total cross sections flat.

Our results should be considered as predictions of the shapes of the baryon-baryon and baryon-antibaryon total cross sections. We do not insist on any particular direct-channel mechanism (whether resonances or annihilation effects) for producing the non-Pomeranchuk imaginary parts of amplitudes. Since there is good evidence for approximately $SU(3)$-invariant couplings of Regge trajectories with $F_T \simeq F_V$,[14,18] a measurement of the energy-dependence of $\sigma_T(\Sigma^- p)$ or $\sigma_T(\Lambda p)$ would be most useful. We expect these last two cross sections to be as flat as $\sigma_T(pp)$ and $\sigma_T(pn)$ if indeed $F_T \simeq F_V$.

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[24] Combinations such as $A_s(K^-p \to K^{*-}p) \pm A_s(K^{*+}p \to K^+p)$ will involve trajectories of definite C parity.

[25] A choice of sign is involved here. As in Sec. II, this choice does not affect the contributions to imaginary parts listed in Tables III, IV, or V.

[26] J. J. de Swart, Rev. Mod. Phys. **35**, 916 (1963); P. NcNamee and F. Chilton, *ibid.* **36**, 1005 (1964).