PHYSICAL REVIEW VOLUME 186, NUMBER 5 25 OCTOBER 1969

Broken-Symmetry Couplings for $S_8 \rightarrow P_8P_8$†

HSIUNG CHIA TZE
The Enrico Fermi Institute, and the Department of Physics, The University of Chicago, Chicago, Illinois 61801
(Received 26 May 1969)

The existence of an isosinglet η'(1070 MeV) and an isotriplet π'(1016 MeV) of a presumed scalar octet has been established recently. A simple model, essentially combining a set of triangle Feynman graphs and the Gell-Mann–Okubo mass formula, allows the calculation of the broken-symmetry coupling constants for $S_8 \rightarrow P_8P_8$. In particular, the anomalous value observed in the branching ratio $R_{\eta'}(\pi\pi/K\bar{K})$ is accounted for.

I. INTRODUCTION

ALTHOUGH the possible existence of scalar mesons has come up repeatedly in the theoretical literature,[1,2] such $J^P=0^+$ hadronic states have been observed only recently. An isosinglet η'(1070 MeV) and an isotriplet π'(1016 MeV) have been identified as peaks in the $K\bar{K}$ amplitude near threshold.[3] Presumably, they are members of a scalar octet with the remaining isodoublet K'(1100 MeV) still awaiting confirmation.[4]

If π', K', and η' do indeed form an octet, there is an interesting feature which needs explanation. On the basis of an $SU(3)$-symmetric interaction and insertion of actual masses in the phase-space factor, the branching ratio $R_{\eta'}(\pi\pi/K\bar{K})$ should be 8.8. However, experiments yield the anomalously low value $R_{\eta'} \lesssim 2.3$.[3]

At first one might think of using the scalar singlet-octet mixing to explain this large effect of symmetry breaking. However, as Evans and Fulton[5] argue, the nearly degenerate 0^+ spectrum implies either a singlet of very different mass, which would induce very little mixing, or one of the same mass as the octet, which is not observed. To account for the relative suppression of the $\pi\pi$ mode, they put forth an effective Hamiltonian

$$H_{\eta'} = f\eta'[\eta\eta + \alpha K\bar{K} - (2-\alpha)\pi\cdot\pi], \qquad (1)$$

where $\alpha=1$ corresponds to the $SU(3)$ limit. η' is then identified as a tadpole responsible for the discrepancies in the Gell-Mann–Okubo (GMO) mass formula for the pseudoscalars. Such a procedure allows a numerical determination of the parameter α.

In the phenomenological approach of Evans and Fulton, the $SU(3)$ deviation is chosen to be symmetric for the $K\bar{K}$ and the $\pi\pi$ couplings. We find this to be an *ad hoc* assumption. The present paper attempts to show how a simple model based on customary dynamical inputs can provide a successful and manifest symmetry-breaking mechanism for the scalar mesons.

Our model is motivated by the work of Wali and Warnock[6] and the subsequent paper of Johnson and McCliment[7] on decuplet decays into baryons and pseudoscalars. They view the coupling shifts as due solely to the propagation of the GMO rule in the masses. For the case of the branching ratio $R_{Y_1^*}(\Sigma\pi/\Lambda\pi)$, where the $SU(3)$ value disagrees with experiments, the above treatments give corrections of the right order of magnitude. While our approach stems from one of Cutkosky's bootstrap equations, it draws upon the successful features of both of the above models for its simplifying approximations.

In Sec. II of this paper, our model is described and the numerical analysis is carried out. The model has as its basis the Cutkosky's bootstrap equation for the coupling constants.[8] Here the S_8 states are considered as P_8P_8 composites bound by exchanges of S_8 and V_8, where S_8, P_8, and V_8 denote, respectively, the 0^+, 0^-, and 1^- octets. The general expression of the broken-symmetry couplings is then formulated. It is done via a first-order expansion in the mass differences about the $SU(3)$-symmetric solution of the Cutkosky equation. The mass shifts are required to obey the GMO formula. Yet the couplings are kept $SU(3)$-symmetric at the vertices. The triangle graphs involved are approximated by the Feynman diagrams of standard perturbation theory. In the context of these approximations, our model makes contact with that of Johnson and McCliment.

In Sec. III, our results for the broken-symmetry couplings tabulated in terms of the Singh-Gupta coefficients are discussed.

II. MODEL

The existence of scalar mesons is indeed a compelling possibility in terms of our current ideas about the composite structure of mesons. In fact, when the boson-

† This work supported in part by the U. S. Atomic Energy Commission.

[1] S. Coleman and S. Glashow, Phys. Rev. **134**, B671 (1964); and references cited therein.

[2] H. Miyazawa, in *Proceedings of the 1965 First Pacific International School in Physics* (Gordon and Breach, Science Publishers, Inc., New York, 1966).

[3] A. H. Rosenfeld *et al.*, University of California Lawrence Radiation Laboratory Report No. UCRL-8030, 1969 (unpublished).

[4] T. G. Trippe, C. Y. Chien, E. Malamud, J. Mellama, P. E. Schlein, W. E. Slater, D. H. Stork, and H. K. Ticho, Phys. Letters **28B**, 203 (1968).

[5] L. E. Evans and T. Fulton, Phys. Rev. **168**, 1706 (1968).

[6] K. C. Wali and R. L. Warnock, Phys. Rev. **135**, B1358 (1964).

[7] E. M. Johnson and E. McCliment, Phys. Rev. **139**, B951 (1965).

[8] R. E. Cutkosky, in *Proceedings of the International Trieste Seminar in High-Energy Physics and Elementary Particles, Trieste, 1965* (International Atomic Energy Agency, Vienna, 1965).

boson forces are assumed dominant over their baryon-antibaryon counterparts, strong attraction in scalar states results for almost all simple models of mesons examined.[9] This assumption forms part of Cutkosky's bootstrap rule I.[8] It is then natural to inquire specifically whether, when applied to the 0^+ mesons, this very rule could account for the large symmetry-breaking effect observed in $R_{\eta'}$.

In the spirit of bootstrap rule I, a meson a of S_8 is viewed as a bound state of two others, b and c, of P_8. α and β, two other constituents of a, could exchange a boson γ before turning into b and c. Now, bootstrap rule II states that the vertex abc so obtained and its lowest-order modification are equal (Fig. 1). Thus we have the self-consistent equation for the $S_a P_b P_c$ coupling constant $g_{bc}{}^a$

$$g_{bc}{}^a = \Delta_{\alpha\beta\gamma}{}^{abc} g_{\alpha\beta}{}^a g_{\alpha\gamma}{}^b g_{\beta\gamma}{}^c , \tag{2}$$

where the Δ's are dynamical factors depending on the masses and the spins of the particles indicated. We approximate the Δ's as being given by all the Feynman triangle graphs allowed by the Lagrangians

$$\mathcal{L}_{SPP} = g_S m_{\eta'} d_{bc}{}^a P^b P^c S_a , \tag{3}$$

$$\mathcal{L}_{VPP} = \tfrac{1}{2} g_V f_{bc}{}^a P^b P \overleftrightarrow{\partial}_\mu{}^c V_a{}^\mu . \tag{4}$$

g_S and g_V are dimensionless constants. The d's and f's are identified as the familiar D- and F-type coupling coefficients in the $SU(3)$ limit, i.e., when the masses are taken to be degenerate within the multiplets. S_i, P_i, and $V_i{}^\mu$, with $i=1,2,\cdots,8$, are the fields associated with the respective octets.

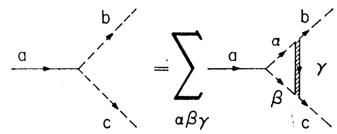

Fig. 1. Self-consistent equation for the coupling constant $g_{bc}{}^a$. Dashed, solid, and hatched lines correspond to pseudoscalar, scalar, and vector or scalar mesons, respectively.

In restricting ourselves to only scalar and vector exchanges, we are guided by the ensuing considerations. In their study of scattering of 0^- mesons, Hong Mo Chan *et al.*[10] found a scalar singlet and octet to be possible bound states through exchange of both the 0^+ and the 1^- octets. Indeed, scalar exchange alone would give rise to a weak repulsion in the direct 0^+ channels. Dynamical self-consistency requires the inclusion of vector-meson exchanges which in general play so dominant a role in strong-interaction physics. Furthermore, by the customary dynamical argument, both scalars and vectors should be considered here since they are mutually degenerate in mass at the $SU(3)$ limit and thus correspond to forces of approximately the same range.

We denote by $\mathcal{S}_i = m_i{}^2$ the mass squared of particle i, by p_i the four-momentum of the particles except for that of the exchange boson γ whose momentum is labeled q. By use of the Feynman rules[11] and Lagrangians (3) and (4), Eq. (2) takes the form

$$g_S d_{bc}{}^a = g_S d_{\alpha\beta}{}^a [d_{\beta c}{}^\gamma d_{\alpha b}{}^\gamma (g_S{}^2/4\pi)\Delta^{(S)} + f_{\beta c}{}^\gamma f_{\alpha b}{}^\gamma (g_V{}^2/4\pi)\Delta^{(V)}] , \tag{5}$$

with

$$\Delta^{(S)} = \frac{i}{2^2\pi^3}\int d^4q \frac{1}{[(p_b+q)^2 - \mathcal{S}_\alpha + i\epsilon][(p_c-q)^2 - \mathcal{S}_\beta + i\epsilon][q^2 - \mathcal{S}_\gamma + i\epsilon]} \tag{6}$$

and

$$\Delta^{(V)} = \frac{-i}{2^2\pi^3}\int d^4q \frac{(p_\beta+p_c)\cdot(p_\alpha+p_b)}{[(p_b+q)^2 - \mathcal{S}_\alpha + i\epsilon][(p_c-q)^2 - \mathcal{S}_\beta + i\epsilon][q^2 - \mathcal{S}_\gamma + i\epsilon]} , \tag{7}$$

where $A^{(S)}$ and $\Delta^{(V)}$ correspond to the Feynman triangle graphs without coupling constants. In the expression for $\Delta^{(V)}$, the omission of a term proportional to $q_\mu q_\nu / \mathcal{S}_\gamma$ in the intergrand should be noted. Since our later calculations involve only values of $\Delta^{(V)}$ and $\partial\Delta^{(V)}/\partial \mathcal{S}_i$ at the $SU(3)$ limit, such a term clearly does not contribute.

We apply the standard Feynman α_i parametrization and integration techniques[12] to Eqs. (6) and (7), to obtain

$$\Delta^{(S)} = \frac{\mathcal{S}_{\eta'}}{4\pi}\int_0^1 \prod_{i=1}^3 d\alpha_i \delta(1-\sum_i \alpha_i) N(\alpha_i, \mathcal{S}_i) , \tag{8}$$

with

$$N^{-1} = [-\alpha_1\alpha_2 \mathcal{S}_a + (\alpha_1{}^2 + \alpha_1\alpha_2 - \alpha_1)\mathcal{S}_b + (\alpha_2{}^2 + \alpha_1\alpha_2 - \alpha_2)\mathcal{S}_c + \alpha_1 \mathcal{S}_\alpha + \alpha_2 \mathcal{S}_\beta + \alpha_3 \mathcal{S}_\gamma] \tag{9}$$

and

$$\Delta^{(V)} = \frac{1}{4\pi}\int_0^1 \prod_{i=1}^3 d\alpha_i \delta(1-\sum_i \alpha_i) \tilde{N}(\alpha_i, \mathcal{S}_i) , \tag{10}$$

with

$$\tilde{N} = N[(\alpha_1+\alpha_2)\mathcal{S}_a - (3\alpha_1+\alpha_2)\mathcal{S}_b - (3\alpha_2+\alpha_1)\mathcal{S}_c - 2(\mathcal{S}_a - \mathcal{S}_b - \mathcal{S}_c) + K^2] - \tfrac{1}{2}(-NL^2) + \tfrac{3}{4} , \tag{11}$$

[9] R. E. Cutkosky, in *Proceedings of the Oxford International Conference on Elementary Particles, 1965* (Rutherford High-Energy Laboratory, Chilton, Berkshire, England, 1966).

[10] Hong-Mo Chan, P. C. DeCelles, and J. E. Paton, Nuovo Cimento 33, 70 (1964).

[11] J. Bjorken and S. Drell, *Relativistic Quantum Mechanics* (McGraw-Hill Book Co., New York, 1964), Appendix B.

[12] A. I. Akhiezer and V. B. Berestetskii, *Quantum Electrodynamics* (Wiley-Interscience, Inc., New York, 1965).

TABLE I. Results for $\chi_{bc}{}^a$ terms.

Resonance	Component	Threshold (MeV)	$\chi_{bc}{}^a$
π'	$\pi\eta$	690	+0.055
	KK	1000	+0.082
K'	πK	640	+0.600
	$K\eta$	1050	+0.070
η'	$\pi\pi$	280	+0.280
	$K\bar{K}$	1000	−0.555
	$\eta\eta$	1100	−0.846

where

$$K^2=(\alpha_2{}^2+\alpha_1\alpha_2)\,\mathcal{S}_c+(\alpha_1{}^2+\alpha_1\alpha_2)\,\mathcal{S}_b-\alpha_1\alpha_2\,\mathcal{S}_a$$

and L^2 is a momentum cutoff.

Equation (5) is known[8] to have a $SU(3)$-symmetric solution for $\Delta^{(S)}$ and $\Delta^{(V)}$ evaluated at the central masses $(\mathcal{S}_0)_{S,P,V}$ of the octets. $(\mathcal{S}_0)_i$ is given by the GMO formula

$$m^2=\mathcal{S}_0\{1+b[I(I+1)-\tfrac{1}{4}Y^2]\}\,, \tag{12}$$

i.e.,

$$(\mathcal{S}_0)_{S,P,V}=m_{\eta'}{}^2,m_\eta{}^2,m_\phi{}^2\,.$$

Thus we have

$$g_SD_{bc}{}^a=g_SD_{\alpha\beta}{}^a[D_{\beta c}{}^\gamma D_{\alpha b}{}^\gamma(g_S{}^2/4\pi)\Delta_0{}^{(S)}+F_{\beta c}{}^\gamma F_{\alpha b}{}^\gamma(g_V{}^2/4\pi)\Delta_0{}^{(V)}]\,, \tag{13}$$

where the D's and F's are now the coupling coefficients of $SU(3)$. $\Delta_0{}^{(S,V)}$ is the $SU(3)$-symmetric value of $\Delta^{(S,V)}$ and $g_V=g_{\rho\pi\pi}$, with the latter defined through

$$H_{\rho\pi\pi}=g_{\rho\pi\pi}\boldsymbol{\rho}_\mu\cdot(\boldsymbol{\pi}\times\partial_\mu\boldsymbol{\pi})\,. \tag{14}$$

Moreover, it can be shown[10] that if

$$D_{\alpha\beta}{}^aD_{\alpha\beta}{}^b=(5/3)\delta_{ab}\,,\quad F_{\alpha\beta}{}^aF_{\alpha\beta}{}^b=3\delta_{ab}\,, \tag{15}$$

then

$$D_{\alpha\beta}{}^aD_{\beta c}{}^\gamma D_{\alpha b}{}^\gamma=-\tfrac{1}{2}D_{bc}{}^a \tag{16}$$

and

$$D_{\alpha\beta}{}^aF_{\beta c}{}^\gamma F_{\alpha b}{}^\gamma=\tfrac{3}{2}D_{bc}{}^a\,. \tag{17}$$

Hence, using Eqs. (16) and (17), we obtain from Eq. (13)

$$\tfrac{3}{2}(g_V{}^2/4\pi)\Delta_0{}^{(V)}-\tfrac{1}{2}(g_S{}^2/4\pi)\Delta_0{}^{(S)}=1\,. \tag{18}$$

This relation corresponds to a renormalization condition on the symmetric coupling constant $g_SD_{bc}{}^a$.

If the masses within the octets are now taken to be nondegenerate, the coupling constants are shifted from their symmetric values by amounts calculable from Eq. (5). We consider only a first-order expansion of Eq. (5) about the symmetric solution in the mass differences $\Delta\mathcal{S}_i=\mathcal{S}_i-\mathcal{S}_{0i}$. All external and internal masses associated with the triangle graphs corresponding to $\Delta^{(S,V)}$ are perturbed from their symmetric values. However, the vertices are kept $SU(3)$-symmetric. In so doing, we get the general form for $g_{bc}{}^a$:

$$g_{bc}{}^a=g_SD_{bc}{}^a\chi_{bc}{}^a\,, \tag{19}$$

with

$$\chi_{bc}{}^a=\Bigg(1-\frac{1}{2}\frac{g_S{}^2}{4\pi}\sum_i^{a,b,c}\frac{\partial\Delta^{(S)}}{\partial\mathcal{S}_i}\Delta\mathcal{S}_i+\frac{D_{\alpha\beta}{}^aD_{\beta c}{}^\gamma D_{\alpha b}{}^\gamma}{D_{bc}{}^a}\frac{g_S{}^2}{4\pi}\sum_i^{\alpha,\beta,\gamma}\frac{\partial\Delta^{(S)}}{\partial\mathcal{S}_i}\Delta\mathcal{S}_i+\frac{3}{2}\frac{g_V{}^2}{4\pi}\sum_i^{a,b,c}\frac{\partial\Delta^{(V)}}{\partial\mathcal{S}_i}\Delta\mathcal{S}_i+\frac{D_{\alpha\beta}{}^aF_{\beta c}{}^\gamma F_{\alpha b}{}^\gamma}{D_{bc}{}^a}\frac{g_V{}^2}{4\pi}\sum_i^{\alpha,\beta,\gamma}\frac{\partial\Delta^{(V)}}{\partial\mathcal{S}_i}\Delta\mathcal{S}_i\Bigg)\,, \tag{20}$$

where $\chi_{bc}{}^a$ is the ratio of the coupling constant to its degenerate limit. This ratio corresponds to a renormalized Singh-Gupta coefficient as defined by Wali and Warnock. To simplify the notation in Eq. (20), we delete the fact that the $\partial\Delta^{(S,V)}/\partial\mathcal{S}_i$ terms are evaluated at the $SU(3)$-symmetric limit. We note that the second and fourth terms get their simple forms through the use of Eqs. (16) and (17). Equation (18) is used along with the above-mentioned identities to obtain the first term in Eq. (20). It is through this renormalization condition that the cutoff-dependent term $\Delta_0{}^{(V)}$ is absorbed into the symmetric coupling. The resulting expression for $\chi_{bc}{}^a$, Eq. (20), involves only $\partial\Delta^{(V)}/\partial\mathcal{S}_i$ terms which are cutoff-free. This particular feature insures a better chance of quantitative success in the calculation of coupling and mass shifts than in the computation of absolute couplings and masses obtainable via bootstrap constraint relations such as Eq. (18).

By inspection of Eqs. (8) and (10), it is clear that

$$\frac{\partial\Delta^{(S,V)}}{\partial\mathcal{S}_b}=\frac{\partial\Delta^{(S,V)}}{\partial\mathcal{S}_c}\,,\quad\frac{\partial\Delta^{(S,V)}}{\partial\mathcal{S}_\beta}=\frac{\partial\Delta^{(S,V)}}{\partial\mathcal{S}_\alpha}\,. \tag{21}$$

It is also found that $\partial\Delta^{(S,V)}/\partial\mathcal{S}_i$ can be written as sum of integrals of the form

$$\int_0^1dx\,A(x,\mathcal{S}_{0_i})\int_0^x dy\,N_0{}^{-r}B(x,y,\mathcal{S}_{0_i})\,, \tag{22}$$

where x and y are related to the α_i's by

$$\alpha_1=y\,,\quad\alpha_2=x-y\,,\quad\alpha_3=1-x\,.$$

N_0 denotes the $SU(3)$-symmetric value of the N already defined in Eq. (9) and $r=1,2$. A and B are simple polynomial functions in x, y, and $\mathcal{S}_{0_i}$. While the y integrations are done analytically, the x integrations require numerical evaluation (Guass method).

III. RESULTS AND DISCUSSION

Our results for the $\chi_{bc}{}^a$'s are presented in Table I. In these calculations, we chose $g_V{}^2/4\pi=2.15$, which corresponds to $\Gamma_{\rho\pi\pi}=110$ MeV. $g_S{}^2/4\pi$ is determined from the observed width $\Gamma_{\eta' K_0 \bar{K}_0}=80$ MeV and the formula

$$\Gamma_{\eta' K_0 \bar{K}_0}=\tfrac{1}{6}(g_S{}^2/4\pi)(\chi_{K\bar{K}}{}^{\eta'})^2 P\,, \qquad (23)$$

with

$$P=\tfrac{1}{2}(\mathcal{S}_{\eta'}-4\,\mathcal{S}_K)^{1/2}.$$

The resulting cubic equation for $g_S/\sqrt{(4\pi)}$ has a single real root giving $g_S{}^2/4\pi=9$.

From Table I, we obtain $R_{\eta'}=2.25$, in agreement with the experimental upper limit $R_{\eta'}\leqslant 2.3$. A choice of $g^2/4\pi>2.15$ results in an even stronger suppression of the $\pi\pi$ mode of η'. As to the remaining $S_8P_8P_8$ couplings, no precise experimental data are yet available. Qualitatively, however, in the case of π', the predicted ratio $(\chi_{KK}{}^{\pi'}/\chi_{\pi\eta}{}^{\pi'})^2=2.2\times10^2$ concurs with the fact that $K^\pm K^0$ is the only mode seen.

A priori, our model does give a rather simplified dynamical picture of the 0^+ mesons. Thus, it could be expected to offer little more than a suppression mechanism for the η' decay, and it is successful in this respect. But, in addition, we find remarkably close agreement with experiments.

However, the validity of our results is sensitive to at least two kinds of dynamical effects not included in our simple model. First, the large deviations in the coupling shifts indicate that we should consider not just the mass shifts but also the coupling-shift feedback terms, certainly in the scalar-exchange part of the model. Second, the higher-order effects of nonlinear mass differences could either significantly alter our results or leave them intact through the same compensating mechanism which saves the decuplet equal-spacing rule.[13] The second alternative may possibly explain why the GMO rule in particular works so well in accounting for experimental facts. The study of such phenomena should be part of any future effort to improve on the present model.

ACKNOWLEDGMENT

The author is grateful to Professor K. C. Wali for suggesting this problem and for his unfailing interest and advice throughout the course of the work.

[13] F. J. Ernst, R. L. Warnock, and K. C. Wali, Phys. Rev. **141**, B1354 (1966).

PHYSICAL REVIEW VOLUME 186, NUMBER 5 25 OCTOBER 1969

Calculation of Re ϵ for the K^0-$\bar{K}^0$ System

A. N. Kamal* and B. G. Kenny†
Rutherford High-Energy Laboratory, Chilton, Didcot, Berkshire, England
(Received 6 May 1969)

An attempt is made to estimate the contribution to Reϵ for the K^0-$\bar{K}^0$ system from the existence of CP nonconservation in the $|\Delta I|>\frac{1}{2}$ $K^0\to 2\pi$ amplitude. In doing so, we have used as input three parametrizations of the S-wave $I=2$ π-π scattering phase shift δ_2. These, in turn, were used to calculate the self-energy contribution to the complex-mass matrix for the K^0-$\bar{K}^0$ system. It is found that the contribution of the above amplitude to Reϵ, although an order of magnitude larger than previous estimates by Truong and Barshay, is smaller than experimental measurements of Reϵ by at least a factor of 2.

I. INTRODUCTION

THE discovery[1] of $K_L{}^0\to\pi^+\pi^-$ established CP nonconservation. When the mode $K_L{}^0\to\pi^0\pi^0$ was first detected,[2] the ratio $\Gamma(K_L{}^0\to\pi^0\pi^0)/\Gamma(K_L{}^0\to\pi^+\pi^-)$ strongly suggested that the source of CP nonconservation was in the $|\Delta I|>\frac{1}{2}$, $K^0\to2\pi$ mode.[3] Such a model of CP nonconservation was first put forward by Truong.[4] Some of the consequences of this model were worked out by Truong,[4] and subsequently by Barshay.[5] One of the predictions of this model was that the ratio $\Gamma(K_L{}^0\to\pi^0\pi^0)/\Gamma(K_L{}^0\to\pi^+\pi^-)$ be close to 2. This was borne out by the original experiments[2] on $K_L{}^0\to\pi^0\pi^0$, but more recent experiments[6] suggest that this ratio is somewhat less than 2, but still large enough to indicate the presence of CP nonconservation

* On leave of absence from the University of Alberta, Edmonton, Canada. Present address: University of Alberta.

† Present address: Australian National University, Canberra, Australia.

[1] J. H. Christenson, J. W. Cronin, V. L. Fitch, and R. Turlay, Phys. Rev. Letters **13**, 138 (1964).

[2] J. M. Gaillard, F. Krienan, W. Galbraith, A. Hussri, M. R. Jane, N. H. Lipman, G. Manning, T. Ratcliffe, P. Day, A. G. Parham, B. T. Payne, A. C. Sherwood, H. Faissner, and H. Reithler, Phys. Rev. Letters **18**, 20 (1967); J. W. Cronin, P. F. Kunz, W. S. Risk, and P. C. Wheeler, *ibid.* **18**, 25 (1967).

[3] By this we mean that if we follow the notation of T. T. Wu and C. N. Yang [Phys. Rev. Letters **13**, 380 (1964)] and choose the phase of K^0 so that the amplitude for $K^0\to2\pi$ $(I=0)$ is real, apart from the phase δ_0 due to the final-state interaction, then the amplitude for $\bar{K}^0\to2\pi$ $(I=2)$ is complex, apart from the phase δ_2 due to the final-state interaction.

[4] T. N. Truong, Phys. Rev. Letters **13**, 358a (1964).

[5] S. Barshay, Phys. Rev. **149**, 1229 (1966). Note that this paper contains corrections to Ref. 4.

[6] For example, see J. W. Cronin, Rapporteur's talk in *Proceedings of the Fourteenth International Conference on High-Energy Physics, Vienna, 1968* (CERN, Geneva, 1968), p. 281.