Effective Lagrangian for Nonleptonic Hyperon Decays with $SU(3)\otimes SU(3)$-Symmetry Breaking

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(Received 23 September 1969)

The hadronic decays of hyperons are discussed in the context of broken chiral symmetry, the breaking being introduced according to the prescription of Gell-Mann, Oakes, and Renner. The corrections due to symmetry breaking lead to contributions to the p-wave amplitudes, giving better agreement with the experimental data, whereas the s-wave amplitudes remain almost unchanged.

NONLEPTONIC hyperon decays have been discussed from the point of view of chiral-invariant effective Lagrangians by Lee[1] and by Schechter,[2] who have shown that the method reproduces essentially the current-algebra results. These current-algebra results obtained especially by Sugawara[3] and by Suzuki[4] can be summarized by saying that the parity-violating hyperon decays are described very well while the predicted values of the parity-conserving decay amplitudes do not agree at all with experiments. In fact, Brown and Sommerfield[5] have demonstrated that their current-algebra results, as well as those of Hara, Nambu, and Schechter,[6] yield p-wave amplitudes half as small as experimental results.

In view of this large discrepancy, Kumar and Pati[7] have proposed a model for the hyperon decays in which the s-wave amplitudes remain essentially unchanged while the p-wave amplitudes are significantly altered, leading to better agreement with experiments. The model is designed to obtain corrections due to the mass splitting within the baryon octet and it is found that these corrections, which were neglected by Brown and Sommerfield, contribute significantly to the parity-conserving decays.

The success of the Kumar-Pati calculation encourages one to believe that a similar situation will also obtain in the context of the effective Lagrangian theory, provided that symmetry breaking is introduced in a systematic way. The construction of the $SU(3)\otimes SU(3)$ chiral-invariant Lagrangian has been described by many authors, and we shall here follow closely the method given by Zumino.[8] We shall then introduce the $SU(3)\otimes SU(3)$-symmetry breaking of the type proposed by Gell-Mann, Oakes, and Renner,[9] using for this purpose the elegant method of Yoshida.[10] This gives us a systematic way of introducing symmetry-breaking effects which can be exhibited in a transparent manner, unlike the model-dependent calculation of Kumar and Pati. We shall first write down the relevant Lagrangians and then give the expressions for the decay amplitudes. These will then be compared with the presently available experimental data as well as with the predictions of the model by Kumar and Pati.

Since the basic tool in all calculations of nonleptonic (NL) hyperon decays is still the time-honored pole model of Feldman, Matthews, and Salam,[11] let us first write down the strong meson-baryon interaction that we need for the strong-vertex part. A simple $SU(3)\otimes SU(3)$ chiral-invariant Lagrangian describing the interaction of the pseudoscalar meson octet with the baryon octet is[8]

$$L_{\rm inv}'=\mathrm{Tr}\{-(1/2a^2)p_\mu{}^2+i\bar{B}(\gamma_\mu\partial_\mu+M)B+\bar{B}\gamma_\mu[v_\mu,B]+\bar{B}\gamma_\mu\gamma_5(b_1p_\mu B+b_2Bp_\mu)\}\,, \quad (1)$$

where

$$p_\mu=\partial_\mu\xi+\cdots \quad (2)$$

and

$$v_\mu=\tfrac{1}{2}i[\xi,\partial_\mu\xi]+\cdots. \quad (3)$$

From the first term, it follows that the normalized pseudoscalar fields are given by

$$\mathbf{P}=(1/a)\xi \quad (4)$$

and hence that the relevant meson-baryon interaction Lagrangian comes out to be

$$L_{\rm inv}=a\,\mathrm{Tr}\{\bar{B}\gamma_\mu\gamma_5(b_1\partial_\mu PB+b_2B\partial_\mu P)\} \quad (5)$$

together with a contact interaction term, which, however, does not contribute in the hyperon decays. The partially conserved axial-vector current (PCAC) condition is written in this construction as $\partial_\mu\mathbf{A}_\mu=(m_\pi{}^2/2a)\boldsymbol{\pi}$ so that $a=\frac{1}{2}f_\pi$, where f_π is the usual pion decay constant known to be about 100 MeV.

As regards the nonleptonic Lagrangian, we shall follow the method of Lee,[12] who has shown that if

[1] B. W. Lee, Phys. Rev. **170**, 1359 (1968).
[2] J. Schechter, Phys. Rev. **174**, 1829 (1968).
[3] H. Sugawara, Phys. Rev. Letters **15**, 870 (1965); **15**, 997 (1965).
[4] M. Suzuki, Phys. Rev. Letters **15**, 986 (1965); G. Murtaza and P. J. O'Donnell, Can. J. Phys. **45**, 2375 (1967).
[5] L. S. Brown and C. M. Sommerfield, Phys. Rev. Letters **16**, 751 (1966).
[6] Y. Hara, Y. Nambu, and J. Schechter, Phys. Rev. Letters **16**, 380 (1966).
[7] A. Kumar and J. C. Pati, Phys. Rev. Letters **18**, 1230 (1967).
[8] B. Zumino, in *Proceedings of the Fifth Coral Gables Conference on Symmetry Principles at High Energy, University of Miami, 1968*, edited by A. Perlmutter and B. Kursunoğlu (W. A. Benjamin, Inc., New York, 1968).
[9] M. Gell-Mann, R. Oakes, and B. Renner, Phys. Rev. **175**, 2195 (1968).
[10] K. Yoshida, University of Durham Report, 1969 (unpublished).
[11] G. Feldman, P. T. Matthews, and Abdus Salam, Phys. Rev. **121**, 302 (1961).
[12] See, e.g., B. W. Lee, in Proceedings of the Argonne International Conference on Weak Interactions, p. 421 [Argonne National Laboratory Report No. ANL-7130, 1965 (unpublished)]; see also Ref. 1.

(a) the weak interaction Lagrangian is of the current ×current type and if (b) the currents are of the form proposed by Cabibbo,[13] then CP invariance of the theory determines that the nonleptonic weak interaction Lagrangian must transform like Gell-Mann's λ_6. Furthermore, in the spirit of Zumino's as well as Weinberg's[14] methods of construction of chiral-invariant effective Lagrangians, we shall adopt here the point of view that we have only derivative couplings in the theory. This then gives us the nonleptonic Lagrangian

$$\begin{aligned}L_{\rm LN}=&\tfrac{1}{2}(d'-f')\,\mathrm{Tr}\{\bar{B}\gamma\cdot\overleftrightarrow{\partial}B\lambda_6+(i/2f_\pi)\bar{B}\gamma\cdot\overleftrightarrow{\partial}B[\lambda_6,P]\}\\&+\tfrac{1}{2}(d'+f')\,\mathrm{Tr}\{\bar{B}\lambda_6\gamma\cdot\overleftrightarrow{\partial}B+(i/2f_\pi)\bar{B}[\lambda_6,P]\gamma\cdot\overleftrightarrow{\partial}B\}\\&+\tfrac{1}{2}(d''-f'')\,\mathrm{Tr}\{\bar{B}\gamma_5\gamma\cdot\overleftrightarrow{\partial}B\lambda_6\\&+(i/2f_\pi)\bar{B}\gamma_5\gamma\cdot\overleftrightarrow{\partial}B[\lambda_6,P]\}+\tfrac{1}{2}(d''+f'')\\&\quad\times\mathrm{Tr}\{\bar{B}\lambda_6\gamma_5\gamma\cdot\overleftrightarrow{\partial}B+(i/2f_\pi)\bar{B}[\lambda_6,P]\gamma_5\gamma\cdot\overleftrightarrow{\partial}B\}\,.\end{aligned}\tag{6}$$

There are other terms involving derivatives of meson fields, some of which vanish in the $SU(3)$ limit, and some are ruled out by current algebra.[1] We shall not consider them here.

We must now introduce the symmetry-breaking term. Recently, Gell-Mann, Oakes, and Renner[9] have proposed a definite way of breaking chiral $SU(3)\otimes SU(3)$ symmetry in current algebra and Macfarlane and Weisz,[15] as well as Yoshida,[10] have shown how to construct a parallel theory in terms of chiral Lagrangians. Following Gell-Mann *et al.*, we write the symmetry-breaking (SB) Hamiltonian as

$$H_{\rm SB}=v_0+cv_8\,,\tag{7}$$

which transforms as $(3,3^*)+(3^*,3)$ representation of chiral $SU(3)\otimes SU(3)$. We obtain the v's from the expression (given by Yoshida)

$$\begin{bmatrix}u_0\\ \vdots\\ u_8\\ v_0\\ \vdots\\ v_8\end{bmatrix}=e^{iQ_5\cdot\xi}\begin{bmatrix}X_0\\ X_1\\ \vdots\\ X_8\\ 0\\ \vdots\\ 0\end{bmatrix},\tag{8}$$

where X_0 and X_i $(i=1,\ldots,8)$ are bilinear functions of the baryon fields transforming as singlet and octet under $SU(3)$,

$$\begin{aligned}X_i&=m_0\alpha d_{ijk}\bar{B}_jB_k+m_0\beta(-if_{ijk})\bar{B}_jB_k\,,\\X_0&=m_0\gamma\sum_i\bar{B}_iB_i\,,\end{aligned}\tag{9}$$

and $Q_5{}^\alpha$ $(\alpha=1,\ldots,8)$ are the 18×18 generators of $(3,3^*)+(3^*,3)$:

$$Q_5\cdot\xi=i\begin{pmatrix}0 & D'\cdot\xi\\ -D'\cdot\xi & 0\end{pmatrix},\tag{10}$$

$$D'\cdot\xi=\begin{pmatrix}0 & (\sqrt{\tfrac{2}{3}})\xi\\ (\sqrt{\tfrac{2}{3}})\xi & D\cdot\xi\end{pmatrix},\quad D\cdot\xi=d_{ijk}\xi_i.\tag{11}$$

For further details of this method of construction we refer to Yoshida's paper.[10] For our purposes we simply pick up terms corresponding to the process $B'\to B+P$, so that

$$L_{\rm SB}=\kappa(v_0+cv_8)\,,\tag{12}$$

where κ is a scale factor and

$$\begin{aligned}v_0=&(\sqrt{\tfrac{2}{3}})\xi_i[m_0\alpha d_{ijk}\bar{B}_j\gamma_5B_k+m_0\beta(-i)f_{ijk}\bar{B}_j\gamma_5B_k]\,,\\v_8=&(\sqrt{\tfrac{2}{3}})m_0\gamma\sum_\alpha\bar{B}_\alpha\gamma_5B_\alpha\xi_8+d_{8jk}\xi_k[m_0\alpha d_{jlm}\bar{B}_l\gamma_5B_m\\&+m_0\beta(-i)f_{jlm}\bar{B}_l\gamma_5B_m]\,.\end{aligned}\tag{13}$$

The complete Lagrangian is then given by

$$L=L_{\rm inv}+L_{\rm NL}+L_{\rm SB}\,.\tag{14}$$

We are now in a position to write down the decay amplitudes. In standard notation, we have the p-wave amplitudes

$$\begin{aligned}P(\Lambda_-{}^0)=&-\frac{d'+3f'}{2\sqrt{6}}(\Lambda+N)\left[-ab_1\frac{\Lambda+N}{\Lambda-N}+\kappa\frac{m_0(\alpha+\beta)}{\sqrt{3}f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Lambda-N}\right]\\&+\tfrac{1}{2}(d'-f')(\Sigma+N)\left[\frac{a}{\sqrt{6}}(b_1+b_2)\frac{\Lambda+N}{\Sigma-N}-\kappa\frac{2}{3\sqrt{2}}\frac{m_0\alpha}{f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Sigma-N}\right]-\frac{d''+3f''}{4\sqrt{6}}\frac{\Lambda-N}{f_\pi}\,,\end{aligned}\tag{15}$$

$$\begin{aligned}P(\Xi_-{}^-)=&-\tfrac{1}{2}(d'+f')(\Xi+\Sigma)\left[-\frac{a(b_1+b_2)}{\sqrt{6}}\frac{\Xi+\Lambda}{\Xi-\Sigma}+\kappa\frac{2}{3\sqrt{2}}\frac{m_0\alpha}{f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Xi-\Sigma}\right]\\&+\frac{-d'+3f'}{2\sqrt{6}}(\Xi+\Lambda)\left[-ab_2\frac{\Xi+\Lambda}{\Xi-\Lambda}+\kappa\frac{1}{\sqrt{3}}\frac{m_0(\alpha-\beta)}{f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Xi-\Lambda}\right]+\frac{d''-3f''}{4\sqrt{6}}\frac{\Xi-\Lambda}{f_\pi}\,,\end{aligned}\tag{16}$$

[13] N. Cabibbo, Phys. Rev. Letters **10**, 531 (1963).
[14] S. Weinberg, Phys. Rev. Letters **18**, 188 (1967); Phys. Rev. **166**, 1568 (1968).
[15] A. J. Macfarlane and P. H. Weisz, Cambridge Report No. 69/20, 1969 (unpublished).

$$P(\Sigma_-^-)=-\frac{d'+3f'}{2\sqrt{6}}(\Lambda+N)\left[\frac{a}{\sqrt{6}}(b_1+b_2)\frac{\Sigma+N}{\Lambda-N}-\kappa\frac{2}{3\sqrt{2}}\frac{m_0\alpha}{f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Lambda-N}\right]$$
$$-\frac{d'-f'}{2\sqrt{2}}(\Sigma+N)\left[\frac{a}{\sqrt{2}}(b_1-b_2)\frac{\Sigma+N}{\Sigma-N}-\kappa\frac{2}{\sqrt{6}}\frac{m_0\beta}{f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Sigma-N}\right]+\tfrac{1}{4}(d''-f'')\frac{\Sigma-N}{f_\pi}, \quad (17)$$

$$P(\Sigma_+^+)=\tfrac{1}{2}(d'-f')(\Sigma+N)\left[-ab_1\frac{\Sigma+N}{\Sigma-N}+\kappa\frac{1}{\sqrt{3}}\frac{m_0(\alpha+\beta)}{f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Sigma-N}\right]$$
$$-\frac{d'+3f'}{2\sqrt{6}}(\Lambda+N)\left[\frac{a}{\sqrt{6}}(b_1+b_2)\frac{\Sigma+N}{\Lambda-N}-\kappa\frac{2}{3\sqrt{2}}\frac{m_0\alpha}{f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Lambda-N}\right]$$
$$-\frac{d'-f'}{2\sqrt{2}}(\Sigma+N)\left[\frac{a(b_2-b_1)}{\sqrt{2}}\frac{\Sigma+N}{\Sigma-N}+\kappa\frac{2}{\sqrt{6}}\frac{m_0\beta}{f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Sigma-N}\right], \quad (18)$$

and the s-wave amplitudes

$$S(\Lambda_-^0)=\frac{d''+3f''}{2\sqrt{6}}(\Lambda-N)\left[ab_1\frac{\Lambda-N}{\Lambda+N}+\kappa\frac{m_0(\alpha+\beta)}{\sqrt{3}f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Lambda+N}\right]$$
$$+\tfrac{1}{2}(d''-f'')(\Sigma-N)\left[\frac{a}{\sqrt{6}}(b_1+b_2)\frac{\Lambda-N}{\Sigma+N}-\kappa\frac{2m_0\alpha}{3\sqrt{2}f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Sigma+N}\right]-\frac{d'+3f'}{4\sqrt{6}}\frac{\Lambda+N}{f_\pi}, \quad (19)$$

$$S(\Xi_-^-)=\tfrac{1}{2}(d''+f'')(\Xi-\Sigma)\left[\frac{a}{\sqrt{6}}(b_1+b_2)\frac{\Xi-\Lambda}{\Xi+\Sigma}+\kappa\frac{2m_0\alpha}{3\sqrt{2}f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Xi+\Sigma}\right]$$
$$+\frac{-d''+3f''}{2\sqrt{6}}(\Xi-\Lambda)\left[-ab_2\frac{\Xi-\Lambda}{\Xi+\Lambda}+\kappa\frac{m_0(\alpha-\beta)}{\sqrt{3}f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Xi+\Lambda}\right]+\frac{d'-3f'}{4\sqrt{6}}\frac{\Xi+\Lambda}{f_\pi}, \quad (20)$$

$$S(\Sigma_-^-)=-\frac{d''+3f''}{2\sqrt{6}}(\Lambda-N)\left[\frac{a}{\sqrt{6}}(b_1+b_2)\frac{\Sigma-N}{\Lambda+N}-\kappa\frac{2m_0\alpha}{3\sqrt{2}f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Lambda+N}\right]$$
$$-\frac{d''-f''}{2\sqrt{2}}(\Sigma-N)\left[\frac{a}{\sqrt{2}}(b_1-b_2)\frac{\Sigma-N}{\Sigma+N}-\kappa\frac{2m_0\beta}{(\sqrt{6})f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Sigma+N}\right]+\tfrac{1}{4}(d'-f')\frac{\Sigma+N}{f_\pi}, \quad (21)$$

$$S(\Sigma_+^+)=\tfrac{1}{2}(d''-f'')(\Sigma-N)\left[-ab_1\frac{\Sigma-N}{\Sigma+N}-\kappa\frac{m_0(\alpha+\beta)}{\sqrt{3}f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Sigma+N}\right]$$
$$-\frac{d''+3f''}{2\sqrt{6}}(\Lambda-N)\left[\frac{a}{\sqrt{6}}(b_1+b_2)\frac{\Sigma-N}{\Lambda+N}-\kappa\frac{2m_0\alpha}{3\sqrt{2}f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Lambda+N}\right]$$
$$-\frac{d''-f''}{2\sqrt{2}}(\Sigma-N)\left[\frac{a}{\sqrt{2}}(b_2-b_1)\frac{\Sigma-N}{\Sigma+N}+\kappa\frac{2m_0\beta}{(\sqrt{6})f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Sigma+N}\right]. \quad (22)$$

The strong-coupling F/D ratio is denoted by f, so that $ab_1=g_A/2f_\pi$ and $a(b_1+b_2)=(1-f)g_A/f_\pi$. To facilitate comparison with the calculations of Kumar and Pati, although it is by no means essential for the analysis, we further introduce $d'+f'=\bar{g}'/\sqrt{2}$ and $f'/(d'+f')=\bar{f}'$. Notice that, since we have used pseudovector coupling, our Born terms B^j incorporate the $\Delta M/2M$ mass corrections. However, because the ratios of the sums of baryon masses are not $SU(3)$ symmetric, the use of $SU(3)$ for the pseudovector couplings leads to a fit different from that of Kumar and Pati. Since we are interested in showing the close correspondence of the present work and the model of these authors, we rewrite our amplitudes using the Goldberger-Treiman relation

$$g_{B_iB_jP_k}/(m_i+m_j)=g_A/2f_\pi, \quad (23)$$

and absorb some over-all factors in a redefinition of the

TABLE I. Decay amplitudes.

	Born term: Kumar and Pati	Born term: Present work	Symmetry breaking term: Kumar and Pati	Symmetry breaking term: Present work[a]	Total: Kumar and Pati[b]	Present work	Experiment[c]
$10^6 P(\Lambda_-^0)$	4.9−3.8	5.34−3.41	0.88	−0.36	1.98 (+0.08=2.1)	1.57	2.267±0.071
$10^6 P(\Xi_-^-)$	−0.7+2.4	−0.71+2.19	−0.22	−0.42	1.48 (−0.02=1.46)	1.06	1.611±0.141
$10^6 P(\Sigma_-^-)$	−2.6+2.4	−2.43+2.16	−0.05	+0.34	−0.25 (−0.04=−0.3)	0.07	−0.119±0.013
$10^6 P(\Sigma_+^+)$	7.0−5.0	8.05−4.59	1.4	−0.08	3.4	3.37	4.143±0.076
$[P(\Lambda_-)+2P(\Xi_-^-)]\times[\sqrt{3}P(\Sigma_0^+)]^{-1}$					1.1	1.09	

	Born+SB terms: Kumar and Pati	Born+SB terms: Present work	Nonpole term: Kumar and Pati	Nonpole term: Present work	Total: Kumar and Pati	Present work	Experiment[c]
$10^6 S(\Lambda_-^0)$	−0.12+0.03+0.1 =0.01	0.013+0.002−0.114 =−0.099	−0.27	−0.32	−0.26	−0.42	−(0.33±0.004)
$10^6 S(\Xi_-^-)$	0.08+0.01−0.1 =−0.01	0.0069−0.0005−0.016 =−0.01	0.49	0.56	0.48	0.55	0.405±0.007
$10^6 S(\Sigma_-^-)$	−0.07−0.02+0.08 =−0.01	−0.007−0.0019−0.0066 =−0.015	−0.57	−0.66	−0.58	−0.67	−(0.406±0.007)
$10^6 S(\Sigma_+^+)$	0.5−0.05−0.02 =−0.02	−0.0072−0.0013−0.075 =−0.08	0	0	−0.02	−0.08	0.004±0.009
$[S(\Lambda_-)+2S(\Xi_-^-)]\times[\sqrt{3}S(\Sigma_0^+)]^{-1}$					1.02	1.1	

[a] For $\kappa=100$. [b] Equal-time commutator contribution is given in parentheses. [c] H. Filthuth, CERN Report No. 69–7 (unpublished).

weak-coupling constant $\bar{g}'$. We thus get, e.g.,

$$P(\Lambda_-^0)=-g'(1+2\bar{f}')(\sqrt{\tfrac{2}{3}})\times\left[-\frac{g}{\Lambda-N}\left(1+\frac{\Lambda-N}{2N}\right)+\kappa\frac{m_0(\alpha+\beta)}{\sqrt{3}f_\pi}\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Lambda-N}\right]+2g'(1-2\bar{f}')\left[\frac{2g(1-f)}{(\sqrt{6})(\Sigma-N)}\left(1+\frac{N-\Sigma}{\Lambda+\Sigma}\right)-\kappa\frac{2m_0\alpha}{3\sqrt{2}f_\pi}\times\left(1+\frac{c}{\sqrt{2}}\right)\frac{1}{\Sigma-N}\right]-gf_Km_K[(\sqrt{12})f_\pi]^{-1}(1+2f)\,,\quad(24)$$

where we take

$$g'=-6\times10^{-6}\ \text{MeV}\,,\quad \bar{f}'=6 \quad\text{and}\quad f=0.34\,.\quad(25)$$

We must emphasize that this redefinition of the coupling constants is done only for the purpose of comparison with the calculations of Kumar and Pati and is in no way essential for the model here considered. To reduce the number of parameters further, we assume that d'' and f'' are proportional to strong interaction d and f, respectively, and, introducing as before $\bar{g}''$ and $\bar{f}''$, we identify $\bar{g}''$ with $f_Km_K/\sqrt{2}$ of Kumar and Pati. Thus we are able to rewrite our amplitudes in the form, e.g.,

$$S(\Lambda_-^0)=gf_Km_K\left[\frac{g}{\sqrt{3}}\frac{1}{\Lambda+N}\frac{\Lambda-N}{2N}+\kappa\left(1+\frac{c}{\sqrt{2}}\right)\frac{m_0(\alpha+\beta)}{3f_\pi}\frac{1}{\Lambda+N}\right](1+2f)+gf_Km_K\left[\frac{2g}{\sqrt{3}}\frac{(1-f)}{\Sigma+N}\frac{\Lambda-N}{\Lambda+\Sigma}-\kappa\left(1+\frac{c}{\sqrt{2}}\right)\frac{2m_0\alpha}{3\sqrt{2}f_\pi}\frac{1}{\Sigma+N}\right](1-2f)-\frac{g'(1+2\bar{f}')}{(\sqrt{6})f_\pi}\,,\quad\text{etc.,}\quad(26)$$

where we use $f_Km_K=1.4\times10^{-6}$ MeV. As regards the parameters α and β, we use Yoshida's estimate from the Gell-Mann–Okubo mass formula

$$m_0c\alpha/\sqrt{3}=39\ \text{MeV}\,,\quad \tfrac{1}{2}m_0\sqrt{3}c\beta=190\ \text{MeV}\,,\quad(27)$$

with the value of $c=-1.26$.

The numerical results for the amplitudes are given in Table I. The close correspondence of the present work and the model of Kumar and Pati is clearly exhibited.

In this work we have not attempted to obtain an exact numerical agreement with experimental data since, in view of the number of parameters of the theory, such an agreement would not be very significant. Our purpose in this note has been to demonstrate that the effective Lagrangian theory incorporating Gell-Mann-type symmetry breaking is well adapted to reproduce the current-algebra results of Kumar and Pati.

ACKNOWLEDGMENTS

The authors thank Professor Abdus Salam and Professor P. Budini as well as the International Atomic Energy Agency for hospitality at the International Centre for Theoretical Physics, Trieste. One author (AMHR) is grateful to the Ford Foundation for making possible his Associateship at the ICTP.