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S-WAVE NON-LEPTONIC DECAYS AND GAUGE FIELDS†

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Abstract: S-wave non-leptonic decay amplitudes are obtained from a gauge-field model and the results are also interpreted from the point of view of mixing of particle fields. The model gives different effective enhancements for K- and hyperon decays as required by experiment.

The current-current theory of weak interactions which works so well when leptonic currents are involved gives, for purely hadronic processes, an interaction Lagrangian

$$\mathcal{L}_{\mathrm{NL}} = \frac{G}{\sqrt{2}} J^{\dagger}_{\mu} J_{\mu} \,, \tag{1}$$

in terms of the Fermi constant $G = 1.01 \times 10^{-5} m_{\mathrm{p}}^{-2}$ and the Cabibbo current

$$J_{\mu} = J_{\mu}^{1+i2} \cos\theta + J_{\mu}^{4+i5} \sin\theta \,, \qquad J_{\mu}^{\alpha} = (j_{\mu}^{\alpha} + j_{5\mu}^{\alpha}) \,. \tag{2}$$

To account for the $|\Delta I| = \frac{1}{2}$ rule one can either assume dynamical octet enhancement or the existence of neutral hadronic currents. In the latter (simplest) case the strangeness changing part of the interaction reads

$$\mathcal{L}_{|\Delta s|=1} = \frac{G}{\sqrt{2}} 2 d_{6\alpha\beta} J_{\mu}^{\alpha} J_{\mu}^{\beta} \cos\theta \sin\theta \,. \tag{3}$$

The main trouble with this interaction is that at first sight it seems to lead to amplitudes which are about a factor $\sin\theta$ too small. Such a result is obtained [2] by first applying vector-meson dominance and PCAC to the parity violating part of eq. (3) in order to get the couplings

$$\sqrt{2} G \sin\theta \cos\theta \left[\frac{m_{\mathrm{K}^*}^2}{f_{\mathrm{K}^*}} C_{\pi} (K_{\mu}^{*+} \partial_{\mu} \pi^- - \frac{1}{\sqrt{2}} K_{\mu}^{*0} \partial_{\mu} \pi^0) \right.$$
$$\left. + \frac{m_{\rho}^2}{f_{\rho}} C_{\mathrm{K}} (\rho_{\mu}^{+} \partial_{\mu} K^- - \frac{1}{\sqrt{2}} \rho^0 \partial_{\mu} K^0) + \mathrm{h.c.} \right] , \tag{4}$$

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and then calculating the amplitudes from diagrams with intermediate K* and ρ [2, 3] ‡.

In the present work it will be shown that when the full Lagrangian for the currents involved is taken into account one obtains a result much closer to experiment.

The hadronic currents seem to satisfy SU(2) ⊗ SU(2) symmetry very well, the chiral invariance being (spontaneously) broken by the presence of the almost massless pion acting as a Nambu-Goldstone boson. The most elegant presentation of this state of affairs was provided by the gauge-fields of Lee, Weinberg and Zumino [4]. Their original model contained only the isovector and its axial partner but the other vector currents can also be incorporated without any problem provided strong symmetry breaking and mixing are part of the kinetic terms of the Lagrangian [5, 6]. Since our main interest here is in the parity-violating decays which involve currents transforming like K^* and A_1, we explicitly write the part of the Lagrangian containing just such (gauge) fields only:

$$\mathcal{L}_{\mathrm{g}} = -\tfrac{1}{4}\,\tilde{\boldsymbol{F}}_{\mu\nu}\;K\boldsymbol{F}_{\mu\nu} - \tfrac{1}{2}\,\tilde{\boldsymbol{\Phi}}_{\mu} M_{\mathrm{o}}^{2}\,\boldsymbol{\Phi}_{\mu}\;, \tag{5}$$

where

$$\tilde{\boldsymbol{\Phi}}_{\mu} \equiv (a_{\mu}^{1}\ldots a_{\mu}^{3}\;v_{\mu}^{4}\ldots v_{\mu}^{7})\;; \qquad \tilde{\boldsymbol{F}}_{\mu\nu} \equiv (A_{\mu\nu}^{1}\ldots A_{\mu\nu}^{3}\;V_{\mu\nu}^{4}\ldots V_{\mu\nu}^{7})\;, \tag{6}$$

with

$$V_{\mu\nu}^{\alpha} = \partial_{\mu} v_{\nu}^{\alpha} - \partial_{\nu} v_{\mu}^{\alpha} - \tfrac{1}{2} f_{\mathrm{o}}\, f_{\alpha\beta\gamma}\,\left([v_{\mu}^{\beta}, v_{\nu}^{\gamma}]_{+} + [a_{\mu}^{\beta}, a_{\nu}^{\gamma}]_{+}\right)\,,$$

$$A_{\mu\nu}^{\alpha} = \partial_{\mu} v_{\nu}^{\alpha} - \partial_{\nu} v_{\mu}^{\alpha} - \tfrac{1}{2} f_{\mathrm{o}}\, f_{\alpha\beta\gamma}\,\left([v_{\mu}^{\beta}, a_{\nu}^{\gamma}]_{+} + [a_{\mu}^{\beta}, v_{\nu}^{\gamma}]_{+}\right)\,. \tag{7}$$

The kinetic and mass matrices in eq. (5) are given by

$$K_{\alpha\beta} = \delta_{\alpha\beta} + Dd_{8\alpha\beta}\,; \qquad (M_{\mathrm{o}}^{2})_{\alpha\beta} = m^{2}\left[\delta_{\alpha\beta} - 2d_{6\alpha\beta}\left(\frac{m}{f_{\mathrm{o}}}\right)^{2}\sqrt{2}G\,\sin\theta\,\cos\theta\right], \tag{8}$$

where, as we pointed out already, strong symmetry breaking are contained in the matrix K. The weak interaction appears in M_{o}^{2} with the factor $(m^{2}/f_{\mathrm{o}})^{2}$ since as shown in ref. [4] the hadronic currents are related to the gauge fields ϕ_{μ} by ‡‡

$$j_{\mu}^{\alpha} = \frac{m^{2}}{f_{\mathrm{o}}}\,\Phi_{\mu}^{\alpha}\,. \tag{9}$$

Using the canonical rules for the fields Φ_{μ} one can derive [6] for the spin one and zero spectral functions of the currents the sum rules

‡ Of course the virtual ρ can only contribute to the off-shell $K_{2\pi}$ amplitudes.

‡‡ The indices α, β, γ, λ run from 1 to 8. Except for eqs. (2), (29) and (30) (where we explicitly differentiate between vector and axial vector currents) a current j_{μ}^{α} (as well as a gauge field ϕ_{μ}^{α}) is a vector if $\alpha = 4, \ldots, 7$ and an axial vector for $\alpha = 1, \ldots, 3$. j_{μ}^{8} is not needed in the following.

$$\int d\mu^2 \left(\frac{\rho^{(1)}}{\mu^2} + \rho^{(0)}\right)_{\alpha\beta} = \left(\frac{m^2}{f_0}\right)^2 (M_0^{-2})_{\alpha\beta} , \tag{10}$$

$$\int d\mu^2 \rho^{(1)}_{\alpha\beta} = \left(\frac{m^2}{f_0}\right)^2 (K^{-1})_{\alpha\beta} . \tag{11}$$

To proceed now we have to assume single-particle saturation of the spectral functions. The matrix elements contributing are

$$\frac{1}{\sqrt{2}}\langle 0 | j_{5\mu}^{(1-i2)};\, j_{\mu}^{(4-i5)} | \pi^+(q)\rangle = q_\mu C_\pi(1;\nu) ,$$

$$\langle 0 | j_{5\mu}^{3};\, j_{\mu}^{6} | \pi^0\rangle = q_\mu C_\pi(1;\nu) ,$$

$$\frac{1}{\sqrt{2}}\langle 0 | j_{5\mu}^{(1-i2)};\, j_{\mu}^{(4-i5)} | A_1^+(\lambda)\rangle = \epsilon_\mu^{(\lambda)} \frac{m_{A_1}^2}{f_{A_1}} (1;\alpha\nu) ,$$

$$\langle 0 | j_{5\mu}^{3};\, j_{\mu}^{6} | A_1^0\rangle = \epsilon_\mu \frac{m_{A_1}^2}{f_{A_1}} (1;\alpha\nu) ,$$

$$\frac{1}{\sqrt{2}}\langle 0 | j_{\mu}^{(4-i5)};\, j_{5\mu}^{(1-i2)} | K^*\rangle = \epsilon_\mu \frac{m_{K^*}^2}{f_{K^*}} (1;\beta) ,$$

$$\frac{1}{\sqrt{2}}\langle 0 | j_{\mu}^{6};\, j_{5\mu}^{3} | K^{*0} + \overline{K^{*0}}\rangle = \frac{m_{K^*}^2}{f_{K^*}} \epsilon_\mu(1;\beta) . \tag{12}$$

Including the ρ-like current, which we did not write explicitly, the diagonal elements of (10) and (11) lead to ‡

$$2C_\pi^2 = \frac{m_\rho^2}{f_\rho^2} = \frac{m_{K^*}^2}{f_{K^*}^2} = 2\,\frac{m_{A_1}^2}{f_{A_1}^2} = \left(\frac{m}{f_0}\right)^2 , \qquad m_{A_1}^2 = 2m_\rho^2 , \tag{13}$$

while from the non-diagonal elements we get

$$\beta = -\alpha\nu \frac{m_\rho^2}{m_{K^*}^2} , \tag{14}$$

‡ Actually the first equality in eq. (13), the famous KSRF relation (ref. [7]) is an input rather than a result from the sum rules.

$$\nu = \left[\tfrac{1}{2} - \alpha\left(\frac{m_\rho^2}{m_{K^*}^2} - \tfrac{1}{2}\right)\right]^{-1} \left(\frac{m}{f_0}\right)^2 \sqrt{2}G \sin\theta \cos\theta \,. \tag{15}$$

Now that we have all the couplings in eqs. (12) (in terms of the constant α) we can proceed to calculate the decay amplitudes. Besides $\mathcal{L}_g$ we have to take into account terms coming from the matter Lagrangian (the B are the baryon fields)

$$\mathcal{L}_m = f_0 f_{\alpha\beta\lambda} \bar{B}^\beta \gamma_\mu B^\alpha v_\mu^\lambda + \ldots = \left(\frac{f_0}{m}\right)^2 f_{\alpha\beta\lambda} \bar{B}^\beta \gamma_\mu B^\alpha j_\mu^\lambda + \ldots \,. \tag{16}$$

Thus, neglecting final-state interactions, we have the effective weak couplings

$$g_{B^\beta B^\alpha \pi^\gamma} \bar{u}^\beta u^\alpha = \langle B^\beta \pi^\gamma | \mathcal{L}_g + \mathcal{L}_m | B^\alpha \rangle \,, \tag{17}$$

which immediately give the parity-violating amplitudes

$$A(B_\gamma^\alpha) = \left[\tfrac{1}{2} - \alpha\left(\frac{m_\rho^2}{m_{K^*}^2} - \tfrac{1}{2}\right)\right]^{-1} (m_\alpha - m_\beta) \left(\sum_{\lambda=4}^{7} 2 f_{\alpha\beta\gamma} d_{6\lambda\gamma}\right) C_\pi \sqrt{2}G \sin\theta \cos\theta \,, \tag{18}$$

defined by

$$\Gamma(B^\alpha \to B^\beta \pi^\gamma)_{\text{s-wave}} = \frac{|q_{\text{c.m.}}|}{8\pi m_\alpha^2} \left[(m_\alpha + m_\beta)^2 - m_\pi^2\right] |A(B_\gamma^\alpha)|^2 \,. \tag{19}$$

In order to fix α we would need, besides the sum-rules some extra condition. A convenient one is that the matrix element of $\mathcal{L}_g$ between π (or A_1) and K^* be zero. By using eq. (12) these matrix elements can be reduced to †

$$\langle \pi^b (A_1^b) | \mathcal{L}_g | K^{*a} \rangle \sim (M_0^2)_{\alpha\beta} \langle \pi^b (A_1^b) | j_\mu^\alpha | 0 \rangle \langle 0 | j_\mu^\beta | K^{*a} \rangle$$

$$\sim d_{6ab} \left[\nu + \beta - \sqrt{2}G \sin\theta \cos\theta \left(\frac{m}{f_0}\right)^2\right] \left(d_{6ab} \left[\alpha\nu + \beta - \sqrt{2}G \sin\theta \cos\theta \left(\frac{m}{f_0}\right)^2\right]\right) . \tag{20}$$

When eqs. (14) and (15) are replaced into eq. (20) the square brackets vanish if $\alpha = 1$. Thus, for α equal (or close to) unity we have effectively eliminated from our model all bilinear couplings. The equivalence between the above procedure and diagonalization will be further substantiated shortly,

† In and following eq. (20) the indices a, $a' = 4, \ldots, 7$ and $b, c = 1, \ldots, 3$. In eq. (20) we only write explicitly the mass terms of $\mathcal{L}_g$ since the matrix element of the kinetic part between π and K^* is, of course, zero and it is easy to show that it also vanishes between A_1 and K^* as long as eq. (14) holds.

when we will translate our results in more phenomenological terms. Instead of computing amplitudes with effective weak bilinear couplings as in pole models, since (the relevant part of) the Lagrangian is assumed to be known, we prefer to diagonalize the bilinear couplings away and use the resulting trilinear couplings.

Let us try next to interpret our results in the more conventional language of particle fields for which we write the Lagrangian

$$\mathcal{L} = -\tfrac{1}{2} m_{\mathrm{K}^*}^2 (\mathrm{K}_\mu^{*\prime a})^2 - \tfrac{1}{2} m_{\mathrm{A}_1}^2 (A_{1\mu}^{\prime b} \cos\varphi + \pi_\mu^{\prime b} \sin\varphi)^2$$

$$- \tfrac{1}{2} m_{\mathrm{A}_1}^2 (A_{1\mu}^{\prime b} \sin\varphi - \pi_\mu^{\prime b} \cos\varphi)^2$$

$$+ 2d_{6ab} \sqrt{2} G \sin\theta \cos\theta \frac{m_{\mathrm{K}^*}^2}{f_{\mathrm{K}^*}} \frac{m_{\mathrm{A}_1}^2}{f_{\mathrm{A}_1}} (A_{1\mu}^{\prime b} + \pi_\mu^{\prime b}) K_\mu^{*\prime a} + \text{kinetic terms}$$

$$+ f_{\mathrm{K}^*} f_{a\beta\lambda} \bar{B}^\lambda \gamma_\mu B^\beta K_\mu^{*\prime a} + \dots \quad (21)$$

The meaning of the primes is that, for instance, $K_\mu^{*\prime}$ is mainly the physical K* field but due to the weak interactions it also contains a small amount of π and A_1. With π_μ defined by

$$\pi_\mu^b \equiv m_{\mathrm{A}_1}^{-1} \partial_\mu \pi^b , \quad (22)$$

eqs. (12) tell us that the weak interaction induces the transformation

$$K^{*a} \to K^{*\prime a} = K^{*a} + 2d_{6ab}\, \nu \frac{m_\rho}{m_{\mathrm{K}^*}} (\alpha A_{1\mu}^b + m_{\mathrm{A}_1}^{-1} \partial_\mu \pi^b) , \quad (23)$$

$$(A_{1\mu}^b + \pi_\mu^b) \to (A_{1\mu}^{\prime b} + \pi_\mu^{\prime b}) = (A_{1\mu}^b + m_{\mathrm{A}_1}^{-1} \partial_\mu \pi_b) + 2d_{6ab} \frac{m_{\mathrm{K}^*}}{m_\rho} \beta K_\mu^{*a} , \quad (24)$$

while the angle φ in eq. (21) is chosen in such a way that

$$(A_{1\mu}^b \sin\varphi - m_{\mathrm{A}_1}^{-1} \partial_\mu \pi^b \cos\varphi) \to (A_\mu^{\prime b} \sin\varphi - \pi_\mu^{\prime b} \cos\varphi)$$

$$= (A_{1\mu}^b \sin\varphi - m_{\mathrm{A}_1}^{-1} \partial_\mu \pi^b \cos\varphi) . \quad (25)$$

Thus, if we accept that α and β are related as in eq. (14) we can either eliminate the $K^*\pi$ bilinear coupling from eq. (21) with

$$\nu = \nu' = \frac{\sqrt{2} G \sin\theta \cos\theta \, (m/f_0)^2}{1 - \alpha (m_\rho^2/m_{\mathrm{K}^*}^2)\, 2 \tan\varphi \, (1 + \tan\varphi)^{-1}} , \quad (26)$$

or we can eliminate the K^*A_1 coupling with ‡

$$\nu = \nu'' = \frac{\sqrt{2}G \sin\theta \cos\theta \, (m/f_0)^2}{\alpha\left(1 - \frac{m_\rho^2}{m_{K^*}^2} \frac{2}{1+\tan\varphi}\right)} . \tag{27}$$

Notice that for

$$\tan(\varphi - \tfrac{1}{4}\pi) = \frac{1}{2} \frac{1-\alpha}{\alpha} \frac{m_{K^*}^2}{m_\rho^2} , \tag{28}$$

both $K^*\pi$ and K^*A_1 bilinear terms can be simultaneously transformed away and eqs. (26) and (27) coincide with eq. (15) obtained previously.

According to eq. (28), $\alpha = 1$ corresponds to $\varphi = \frac{1}{4}\pi$ for which $A_{1\mu}$ and π_μ (which share the axial-isovector current in equal proportions), acquire, once the weak interaction is switched on, an equal amount of K^* each. In this case what we have done is equivalent to taking Lagrangian (21) with $\varphi = \frac{1}{4}\pi$, ignore the third term (which does not take part in the mixing with K^*) and simply get rid of bilinear couplings by diagonalization.

Going back to Lagrangian (5), if we add the ρ-like current we get the equalties

$$\int d\mu^2 \mu^{-2} \rho^{(1)}(j^a, j^b) = \int d\mu^2 [\mu^{-2}\rho^{(1)}(j^a, j_5^b) + \rho^{(0)}(j^a, j_5^b)] = \ldots = 2d_{6ab} A , \tag{29}$$

$$\int d\mu^2 \rho^{(1)}(j^a, j^b) = \int d\mu^2 \rho^{(1)}(j^a, j_5^b) = \int d\mu^2 \rho^{(1)}(j_5^a, j_5^b) = 2d_{6ab} B . \tag{30}$$

These sum rules (which can also be justified from an asymptotic-symmetries point of view) could, alternatively be taken as a starting point for the calculation of the decay amplitudes since A and B can be obtained in the following way: from the sum rules and just assuming a baryon-current interaction (like eq. (16)) one gets the weak vertices $\Lambda \to p\rho^-$ and $n \to pK^{*-}$ as a function of A and B. But these vertices also follow from diagrams with virtual K^* and ρ and a ρ-K^* coupling taken directly from vector-meson dominance of the currents in eq. (3). In this way one obtains (the same as from Lagrangian (5))

$$2d_{6ab} A = (M_0^2)_{ab} , \qquad B = 0 . \tag{31}$$

Having justified (we hope) our model we go on to calculate the K_S decay. Again neglecting final-state interactions we have the couplings

‡ The A_1K^* bilinear coupling in the kinetic term is also eliminated if we use eq. (14) and assume that under the weak interaction $(\partial_\mu A_{1\nu} - \partial_\nu A_{1\mu}) = [\partial_\mu(A_{1\nu} + \pi_\nu) - \partial_\nu(A_{1\mu} + \pi_\mu)] \to [\partial_\mu(A'_{1\nu} + \pi'_\nu) - \partial_\nu(A'_{1\mu} + \pi'_\mu)]$.

$$g_{\mathrm{K}^a\pi^b\pi^c} = \langle \pi^b \pi^c | \mathcal{L}_{\mathrm{g}} | \mathrm{K}^a \rangle = \left(\frac{f_0}{m^2}\right)^2 (M_0^2)_{\lambda\gamma} \langle \pi^b | j_\mu^\lambda | \mathrm{K}^a \rangle \langle \pi^c | j_\mu^\gamma | 0 \rangle$$

$$= \sum_{a',b'} \langle \pi^b(q) | \left[\frac{f_{\mathrm{K}^*}^2}{m_{\mathrm{K}^*}^2} j_\mu^{a'} - \sqrt{2} G \sin\theta \cos\theta \, (2 d_{6a'b'}) j_{5\mu}^{b'} \right] | 0 \rangle \langle \pi^c(k) | j_\mu^{a'} | \mathrm{K}^a(p) \rangle$$

$$+ \frac{1}{C_\pi} q_\mu \langle \pi^c(k) | j_{5\mu}^b | \mathrm{K}^a(p) \rangle . \quad (32)$$

We do not know how to calculate the last term but if we assume it can be neglected then

$$g_{\mathrm{K}^a\pi^b\pi^c} = \sum_{a'} 2 f_{aa'c} \, d_{6a'b} \left[\left(\frac{f_0}{m}\right)^2 \nu - \sqrt{2} G \sin\theta \cos\theta \; C_\pi \right] (m_{\mathrm{K}}^2 - m_\pi^2) \; , \quad (33)$$

which leads to the decay amplitude

$$A(\mathrm{K}_1 \to \pi^+\pi^-) = \left[\frac{1}{\frac{1}{2} - \alpha(m_\rho^2 / m_{\mathrm{K}^*}^2 - \frac{1}{2})} - 1 \right] (m_{\mathrm{K}^0} - m_\pi) C_\pi \sqrt{2} G \sin\theta \cos\theta \; , \quad (34)$$

defined by

$$\Gamma(\mathrm{K}_1 \to \pi^+\pi^-) = \frac{|q_{\mathrm{c.m.}}|}{8\pi} \left(1 + \frac{m_\pi^2}{m_{\mathrm{K}}^2}\right) |A(\mathrm{K}_1 \to \pi^+\pi^-)|^2 \; . \quad (35)$$

Comparing eqs. (18) and (34) we see that the effective couplings responsible for hyperon and K-decay are in the ratio (between square brackets)

$$\left[\frac{1}{2} + \alpha \left(\frac{m_\rho^2}{m_{\mathrm{K}^*}^2} - \frac{1}{2} \right) \right]^{-1} \xrightarrow[\alpha \to 1]{} \frac{m_{\mathrm{K}^*}^2}{m_\rho^2} = 1.36 \; . \quad (36)$$

If the last term in eq. (32) which can be written in terms of form factors as

$$-2 d_{6a'c} f_{baa'} \frac{1}{C_\pi} (p_\mu - k_\mu) [(p_\mu + k_\mu) G_+(-m_\pi^2) + (p_\mu - k_\mu) G_-(-m_\pi^2)] \quad (37)$$

is indeed negligible ‡ compared to the rest, then we have found the explanation of another mystery: K_1 non-leptonic decays require an effective coupling $\approx 30\%$ lower than hyperon decays ‡‡.

In a recent paper Nishijima and Sato [8] also obtain an enhancement factor ($\approx 40\%$ higher than ours) from a chiral phenomenological Lagrangian. Since these authors do not consider the part of the Lagrangian containing baryons it would seem that their final result does not take into account all

‡ Both form factors cannot vanish since, for instance current algebra in the soft-K limit gives $G_-(-m_\pi^2) - G_+(-m_\pi^2) = \nu \; C_\pi / C_{\mathrm{K}}$.

‡‡ See for instance ref. [2].

Table 1
Parity-violating non-leptonic amplitudes in units of 10^{-7}.

Amplitude	Experiment	Cabibbo cur.-cur. (eq. (3))		Tomozawa, $\alpha = 1$	
		$\alpha = 1$	$\alpha = 1.48$	octet	nonet
$\lvert A(\Lambda_-)\rvert$	3.30 ± 0.04	1.98	3.41	3.41	2.78
$\lvert A(\Sigma^-_-)\rvert$	4.07 ± 0.07	2.34	4.04	4.04	3.30
$\lvert A(\Sigma^+_+)\rvert$	0.044 ± 0.087	0	0	0	0
$\lvert A(\Sigma^+_0)\rvert$	3.33 ± 0.03 2.50 ± 0.04	1.62	2.79	2.79	2.28
$\lvert A(\Xi^-_-)\rvert$	4.50 ± 0.07	2.29	3.95	3.95	3.22
$\lvert A(\Xi^0_0)\rvert$	3.33 ± 0.11	1.57	2.72	2.72	2.22
$\lvert A(K_S \to \pi^+\pi^-)\rvert$	6.21	3.38	6.72	5.86	4.78

possible contributions. But while the authors of ref. [8] claim their model gives octet enhancement (over the 27-plet) without assuming it as input, we have to put octet dominance by hand (by adding the neutral currents). Another distinguishing feature between the model of ref. [8] and ours is that we get a different enhancement for hyperon and K_1 decays.

We have listed our results in table 1. The first two columns after the experimental values‡ contain the amplitudes (in units of 10^{-7}) which follow from the hadronic current-current interaction (3) with $\theta = 0.235$ (± 0.006) [9]. We have also calculated the amplitudes by using the non-leptonic interaction from octet and nonet intermediate W-meson models recently proposed by Tomozawa [10]. In these models the hadronic weak interaction comes out as in eq. (3) (up to a phase) but multiplied by $\sqrt{3}$ and $\sqrt{2}$ respectively.

Following tradition we have listed in table 1 most of the hyperon amplitudes even though the fit to their ratios is more general than our model. The same hyperon amplitude ratios were obtained, for instance, in the intermediate K* model of ref. [2].

The largest discrepancies between theoretical and experimental hyperon amplitudes are in the Ξ decays. On the other hand a comparison between $A(\Xi^-_-)$ and $A(\Xi^0_0)$ shows a sizable $\Delta I = \frac{1}{2}$ violating effect at work. One might

‡ We have taken the experimental hyperon amplitudes from ref. [9] and changed the normalization as to make them non-dimensional. The experimental $A(K_S \to \pi^+\pi^-)$ corresponds to $\Gamma(K_S \to \pi^+\pi^-) = (0.797 \pm 0.009) \cdot 10^{10}\ \mathrm{sec}^{-1}$ listed in the Particle Property Tables, Jan. 1970.

then be hopeful that if this effect (which is outside the scope of the present paper) were taken into account some improvement (specially in $A(\Xi^0_0)$) would result.

The sum rules (10) and (11), from which the amplitude (18) was derived, can be more general than our gauge fields model while our conclusion from eq. (20) that $\alpha = 1$ might not be on as firm grounds. That is why we also give the amplitudes for a different value of α (= 1.48). Nevertheless, since we do have arguments in favor of $\alpha = 1$ we feel justified in concluding that our model with Tomozawa's current-current interaction (from his W-octet model) provides a good description of s-wave non-leptonic decays ‡.

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‡ The enhancement of the non-leptonic amplitudes is also being considered, from a different point of view, by Tomozawa [11].

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