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# $\Delta I = \frac{1}{2}$ Rule in Nonleptonic Weak Interactions

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We investigate whether the $\Delta I = \frac{1}{2}$ rule for weak nonleptonic interactions can be embedded into unified theories of the weak and electromagnetic interactions. In such theories weak interactions are mediated not only by (gauge) vector-boson exchange, as usually envisaged, but also through exchange of Higgs scalar bosons. Although the former contributions cannot (it seems) be arranged to have a pure $\Delta I = \frac{1}{2}$ structure, we discuss a model in which the latter *can* be so arranged. Owing to strong-interaction complexities the effective strengths of these two sources of weak interactions cannot easily be assessed. The discussion nevertheless emphasizes the possibility that Higgs-exchange effects may have a significant role for the $\Delta I = \frac{1}{2}$ question.

## I. INTRODUCTION

Among the various regularities and selection rules which are suggested by the phenomenology of weak interactions, the $\Delta I = \frac{1}{2}$ rule for nonleptonic processes uniquely occupies an anomalous theoretical status. With respect to most aspects of hyperon and $K$ decays the $\Delta I = \frac{1}{2}$ rule appears to enjoy ample experimental support, to within corrections which are at the few-percent level and which can perhaps be attributed to the intervention of electromagnetic effects. On present experimental evidence, substantially larger deviations arise only for the slope parameters in $K \to 3\pi$ decays.[1] Nevertheless, even when these latter discrepancies are provisionally ignored, the straightforward conclusion that the relevant weak-interaction Lagrangian must have an essentially pure $\Delta I = \frac{1}{2}$ character is not universally accepted as a principle of model building. In the "conventional" theoretical picture, the weak interactions are described as having a current × current structure, such as would arise (in lowest order) from exchange of *charged* vector bosons coupled to leptonic and hadronic currents. This has the elegant feature that the nonleptonic interactions are built up solely out of the charged hadronic currents that figure into semileptonic interactions. But it also entails the existence of $\Delta I = \frac{3}{2}$ as well as $\Delta I = \frac{1}{2}$ terms in the $|\Delta S| = 1$ nonleptonic sector. On this scheme it is necessary to appeal to strong-interaction dynamics for a mechanism that somehow selectively enhances the $\Delta I = \frac{1}{2}$ (or suppresses the $\Delta I = \frac{3}{2}$) contribution to physical amplitudes.

In order to achieve a pure $\Delta I = \frac{1}{2}$ character for the effective (lowest-order) interaction Lagrangian in the above framework, it is necessary to introduce neutral intermediate vector bosons, coupled with appropriate strength to appropriate neutral hadron currents. A pair of neutral bosons, represented by a complex field and its conjugate, is required in order to avoid effective interactions in lowest order which give rise to $|\Delta S| = 2$ transitions. Moreover, neutral-boson couplings to neutral lepton currents formed from the known light leptons ($e$, $\mu$, $\nu_e$, and $\nu_\mu$) would have to be much weaker, if such couplings exist at all, than the corresponding couplings to neutral hadron currents. This asymmetry follows from the fact that processes such as $K \to \pi + l + \bar{l}$ are known to have, at most, very tiny branching ratios (typically $\lesssim 10^{-6}$). All of these requirements for the $\Delta I = \frac{1}{2}$ rule, or its SU(3) generalization – the octet rule – can of course be met, and have indeed been elaborated in the literature. In our present state of quantitative ignorance about strong-interaction dynamics, commitment to the simpler conventional picture seems to be a matter of taste – and faith in strong-interaction octet dominance. For that matter, there is no reliable, empirically derived evidence whatsoever that weak nonleptonic processes in fact arise predominantly from an effective current×current interaction, of either sort described above. For either picture the size of observed nonleptonic decay amplitudes seems somewhat too large, at least on the basis of naive estimates (although it may be said that this reflects the workings of octet *enhancement*). Perhaps the main argument for a current×current structure of the nonleptonic interactions rests on the idea that the usual semileptonic and leptonic interactions are mediated by charged vector bosons. The introduction of neutral vector bosons, as described above, is not similarly tied to known semileptonic and leptonic phenomena and therefore represents an *ad hoc* supplement designed to achieve the $\Delta I = \frac{1}{2}$ rule. But in this spirit, additional interactions of a wholly different sort can be imagined for the nonleptonic sector: for example, interactions generated by exchange of spinless bosons coupled to scalar and pseudoscalar hadron densities. On this approach one can contemplate the possibility that these additional interactions predominate for nonleptonic decays and that they have a pure $\Delta I = \frac{1}{2}$ character.

At this qualitative level the options are of course wide open. The question is whether there exists any theoretical framework in which the new interactions have a natural place. One such framework is embodied in the notion of spontaneously broken gauge symmetry as a basis for unified theories of the weak and electromagnetic interactions.[2] The idea here is to embed lowest-order phenomenology into a renormalizeable field theory, where higher-order corrections are not only finite but small. Intermediate vector bosons arise here as the quanta corresponding to the gauge fields of the theory. Although there is considerable flexibility with respect to choice of gauge and choice of group representations, however, no model has so far been discussed in which one encounters the neutral vector-boson pair needed, along with the "usual" charged vector boson, to complete the $\Delta I = \frac{1}{2}$ rule through vector exchanges alone. Indeed, owing to the various constraints discussed earlier, it seems impossible to arrange for this selection rule in the interactions generated by gauge-vector-boson exchange. It might be imagined that additional *nongauge* vector bosons could be introduced, along with appropriate couplings to neutral hadron currents, all of this adjusted in such a way as to make up the $\Delta I = \frac{1}{2}$ rule.[3] However, in order to preserve renormalizability these nongauge vector bosons must couple to conserved currents, and this is something that it does not seem possible to arrange. (See also the note added in proof below.)

For the question under discussion here a hopeful new element of broken gauge symmetry theories is that weak interactions are mediated not only by gauge vector bosons but also by spinless bosons. These particles correspond to fields whose nonvanishing vacuum expectation values are responsible for breakdown of the gauge symmetry, according to the mechanism first described by Higgs.[4] An intriguing possibility, and this is the essential observation of the present paper, is that the bulk of $|\Delta S| = 1$ nonleptonic weak amplitudes may arise from exchange of Higgs particles. There are two issues here: (1) Can it be arranged that the lowest-order interactions so generated have a pure $\Delta I = \frac{1}{2}$ character? (2) Can it be arranged, through adjustment of parameters, that these interactions predominate for nonleptonic processes over those generated by vector-boson exchange, without at the same time running afoul of other aspects of phenomenology?

In at least one scheme (the eight-quark version of the Georgi-Glashow model[5]) the $\Delta I = \frac{1}{2}$ (and octet) rule in fact emerges naturally for the interactions generated by Higgs boson exchange. We shall discuss the situation in some detail, not because we seriously think this model has been chosen by nature, but in order to illustrate the possibility of assigning a creative role to the Higgs particles in connection with the $\Delta I = \frac{1}{2}$ rule. The quantitative question whether scalar exchanges dominate over vector exchanges for nonleptonic decays is of course much harder to deal with. This depends not only on the parameters of the model but also on strong-interaction dynamics. One relevant parameter is the mass of the single physical Higgs boson of the model under discussion. For broken gauge symmetry theories in general there is no known principle which relates the Higgs boson masses to other masses; and for the particular model being considered there are no observational effects which at present preclude the possibility of a rather small mass for the Higgs particle. It is clear that scalar exchange effects will be overwhelmed by vector exchange if this mass is made very large relative to other masses in the model. Unfortunately, the converse does not necessarily follow[6] if the Higgs mass is made very small, but we shall entertain this hope.

## II. NONLEPTONIC TRANSITIONS GENERATED BY HIGGS PARTICLE EXCHANGE

In a previous paper[7] we have pointed out various difficulties with the five-quark version of the Georgi-Glashow model.[5] An eight-quark version (we are concerned here only with the hadron sector of the model) is based on fundamental spin-$\frac{1}{2}$ fermions grouped into two SO(3) triplets and two singlets:

$$\psi_{\mathfrak{N}}=\begin{pmatrix}\mathcal{P}\\(\mathfrak{N}_C\sin\beta+q\cos\beta)_L+q_R\\r\end{pmatrix},$$

$$\psi_{\lambda}=\begin{pmatrix}\mathcal{P}'\\(\lambda_C\sin\beta+q'\cos\beta)_L+q'_R\\r'\end{pmatrix}, \tag{1}$$

$$S_{\mathfrak{N}}=(\mathfrak{N}_C\cos\beta-q\sin\beta)_L+\mathfrak{N}_R,$$

$$S_{\lambda}=(\lambda_C\cos\beta-q'\sin\beta)_L+\lambda_R,$$

where

$$\mathfrak{N}_C=\mathfrak{N}\cos\theta+\lambda\sin\theta,$$

$$\lambda_C=-\mathfrak{N}\sin\theta+\lambda\cos\theta.$$

The (integral) quark charges are

$$Q=\begin{cases}1, & \mathcal{P},\ \mathcal{P}',\\ 0, & \mathfrak{N},\ \lambda,\ q,\ q',\\ -1, & r,\ r'.\end{cases} \tag{2}$$

The subscripts $L$ and $R$ refer respectively to left- and right-handed chiral projections, $\theta$ is the Cabibbo angle, and $\beta$ is a parameter of the model. With respect to ordinary isospin we suppose that $\mathcal{P}$ and $\mathfrak{N}$ form a doublet; with respect to SU(3), that $\mathcal{P}$, $\mathfrak{N}$, and $\lambda$ form a triplet. The other fermions are singlets. For simplicity we shall assume SU(3) invariance, so that the masses of $\mathcal{P}$, $\mathfrak{N}$, and $\lambda$ are equal: $m(\mathcal{P})=m(\mathfrak{N})=m(\lambda)\equiv m$.

The strong interactions are presumed to be mediated by a single neutral vector gluon, so that in addition to charge, hypercharge, and baryon number, there are various kinds of charm quantum numbers that are additively conserved. We may take all the quarks to have unit baryon number. The particle $q$ is taken to be uncharmed and to have zero hypercharge (actually, in the following discussion $q$ may be replaced by any linear combination of $q$ and $q'$ with only minor changes, but we forgo this). The remaining particles bear charm quantum numbers. The physical mesons and baryons are of course uncharmed. The assignments are such that the physical $\pi^+$ meson, for example, may be thought of as $\mathcal{P}\bar{\mathfrak{N}}$ and any number of uncharmed, SU(3)-singlet pairs $(\bar{\mathcal{P}}\mathcal{P}+\bar{\mathfrak{N}}\mathfrak{N}+\bar{\lambda}\lambda)$, $\bar{q}q$, $\bar{q}'q'$, $\bar{r}r$, and $\bar{r}'r'$. Similarly the physical $K^+$ is $\mathcal{P}\bar{\lambda}$ plus pairs; the physical proton is $\mathcal{P}\bar{\lambda}q$ plus pairs; etc.

In addition to the quark fields, one introduces a triplet $\vec{W}$ of vector gauge fields (whose neutral member describes the photon) and a triplet $\vec{\phi}$ of Higgs scalar fields (whose neutral member describes the physical Higgs particle). The Lagrangian of the model has the form

$$\mathcal{L}=\bar{\psi}\gamma\cdot(i\partial-e\vec{W}\cdot\vec{T})\psi-\psi^\dagger\beta(M_0+\vec{\Gamma}\cdot\vec{\phi})\psi+\mathcal{L}_{W,\phi}, \tag{3}$$

where $\mathcal{L}_{W,\phi}$ does not involve the fermion fields. It is convenient to choose a basis in which the quark mass matrix is diagonal. The matrix $\vec{T}$ is a reducible representation of the O(3) generators, and $\psi$ represents all eight fermions. The matrices $\beta M_0$ and $\beta\vec{\Gamma}$ transform under O(3) like a singlet and triplet, respectively. The full mass matrix is given by

$$\begin{aligned}M&=M_0+\vec{\Gamma}\cdot\langle\vec{\phi}\rangle_0\\&=M_0+\frac{M_W}{e}\Gamma_0,\end{aligned} \tag{4}$$

where $M_W=e\langle\phi_0\rangle_0$ is the mass of the charged vector bosons and $\langle\phi_0\rangle_0$ is the vacuum expectation value of the neutral Higgs field. It follows from Eq. (4) that $M$ transforms like a mixture of $J=0$ and $J=1$ objects under O(3). The absence of a $J=2$ term implies the "zeroth-order" sum rules[8]:

$$\begin{aligned}m(\mathcal{P})+m(r)&=2m(q)\cos\beta,\\m(\mathcal{P}')+m(r')&=2m(q')\cos\beta.\end{aligned} \tag{5}$$

The coupling of the fermion and Higgs fields is given by

$$\mathcal{L}_{\psi,\phi}=-\psi^\dagger\beta\Gamma_0\psi\phi,$$

where $\phi=\phi_0-\langle\phi_0\rangle_0$ has zero vacuum expectation value. The Higgs coupling matrix $\Gamma_0$ can be worked out in terms of the quark mass matrix. We note that

$$\begin{aligned}\beta\Gamma_0&=[T_+,\beta\Gamma_-]\\&=\frac{e}{M_W}[T_+,[T_-,\beta M]].\end{aligned} \tag{7}$$

For the most part $\Gamma_0$ is diagonal in the fermion fields. The only nondiagonal terms are $(\bar{q}\lambda)$, $(\bar{q}\mathfrak{N})$, $(\bar{q}'\lambda)$, and $(q'\mathfrak{N})$ and their conjugates; and similar pseudoscalar densities. An important point, as noted in the Appendix of Ref. 7, is that there are no $(\bar{\lambda}\mathfrak{N})$ or $(\bar{\lambda}\gamma_5\mathfrak{N})$ terms or their conjugates. In more detail, the off-diagonal terms in the coupling between fermions and the Higgs field are given by

$$\frac{e\cos\beta\sin\beta}{M_W}\,\phi\left\{(\bar{\mathfrak{N}}\cos\theta+\bar{\lambda}\sin\theta)\left[m\left(\frac{1-\gamma_5}{2}\right)-m(q)\left(\frac{1+\gamma_5}{2}\right)\right]q + (-\bar{\mathfrak{N}}\sin\theta+\bar{\lambda}\cos\theta)\left[m\left(\frac{1-\gamma_5}{2}\right)-m(q')\left(\frac{1+\gamma_5}{2}\right)\right]q'+\text{H.c.}\right\}. \tag{8}$$

Recall that we have set $m(\mathcal{P})=m(\mathfrak{N})=m(\lambda)=m$.

Charm-conserving $\Delta S=0$ and $|\Delta S|=1$ nonleptonic interactions are seen to be generated already in lowest order through exchange of the Higgs particle; for the $|\Delta S|=1$ sector it is evident that the effective interaction has a pure $\Delta I=\frac{1}{2}$, indeed a pure octet, character. In order to estimate the strength and structure of this interaction, we first consider single Higgs exchange contributions to the quark scattering process $q+\mathfrak{N}\to q+\lambda$, $q'+\mathfrak{N}\to q'+\lambda$, provisionally neglecting all strong-interaction effects and computing, moreover, in the limit of small external momenta. We then assemble the results into an effective Lagrangian, whose matrix elements between physical hadron states determine the amplitudes for physical transitions. In this small-momentum limit the Higgs propagator function is simply $m_\phi^{-2}$ where $m_\phi$ is the Higgs particle mass. Insofar as the wave functions of physical hadrons favor small momenta for the quark constituents (small compared to $m_\phi$), this effective-Lagrangian approach is perhaps not too unreasonable; and the strength of the effective coupling then grows with decreasing $m_\phi$. However, we surely cannot increase the strength at will in this way, since the approximation becomes misleading for small enough $m_\phi$.[6] Nevertheless, this is the approximation we shall adopt, if only for a rough estimate of the state of affairs. For the charm-conserving $|\Delta S|=1$ sector we find the effective Lagrangian

$$\frac{e^2\cos^2\beta\sin^2\beta}{M_W{}^2m_\phi{}^2}\cos\theta\sin\theta\left\{m^2\bar{\mathfrak{N}}\left(\frac{1-\gamma_5}{2}\right)q\,\bar{q}\left(\frac{1+\gamma_5}{2}\right)\lambda-mm(q)\left[\bar{\mathfrak{N}}\left(\frac{1-\gamma_5}{2}\right)q\,\bar{q}\left(\frac{1-\gamma_5}{2}\right)\lambda+\bar{\mathfrak{N}}\left(\frac{1+\gamma_5}{2}\right)q\,\bar{q}\left(\frac{1+\gamma_5}{2}\right)\lambda\right] + m^2(q)\bar{\mathfrak{N}}\left(\frac{1+\gamma_5}{2}\right)q\,\bar{q}\left(\frac{1-\gamma_5}{2}\right)\lambda-(q\longleftrightarrow q')+\text{H.c.}\right\}. \tag{9}$$

We shall now assume that $m\ll m(q),\ m(q')$. There seems to be no theoretical or observational argument against this supposition. In this case Eq. (9) simplifies considerably and takes the form, after the performance of a Fierz transformation,

$$\mathcal{L}_{|\Delta S|=1}=-\frac{1}{4}\frac{G_F}{\sqrt{2}}\cos^2\beta\cos\theta\sin\theta[\bar{\mathfrak{N}}\gamma^\mu(1-\gamma_5)\lambda+\bar{\lambda}\gamma^\mu(1-\gamma_5)\mathfrak{N}]\left[\left(\frac{m(q)}{m_\phi}\right)^2\bar{q}\gamma_\mu(1+\gamma_5)q-\left(\frac{m(q')}{m_\phi}\right)^2\bar{q}'\gamma_\mu(1+\gamma_5)q'\right]. \tag{10}$$

With the proviso noted above, we see that this Higgs exchange interaction may be stronger than the usual vector exchange interaction, provided $m_\phi$ is small compared to $m(q)$, $m(q')$.

In a simple quark model in which hadrons are built up out of valence quarks and a sea of unitary singlet pairs of quarks, nonleptonic decays of mesons must proceed via the interaction of a valence quark $(\mathfrak{N},\lambda,\bar{\mathfrak{N}},\bar{\lambda})$ with sea quarks $(q,q',\bar{q},\bar{q}')$; in baryon decays the valence $q$ (or $\bar{q}$) quarks may participate directly.

It should be noted that the standard results on nonleptonic decays, as obtained from current algebra and partial conservation of axial-vector current (PCAC), remain unaltered with the interactions of Eq. (10). This is because the latter couples a $V-A$ octet current to a $V+A$ singlet current, so that one recovers the usually assumed result[9]

$$[Q_i^5,\mathcal{L}_{|\Delta S|=1}]=[Q_i,\mathcal{L}_{|\Delta S|=1}], \tag{11}$$

where $Q_i$ and $Q_i^5$ are the generators of SU(3)×SU(3). The familiar results then follow for $K\to3\pi$ decays[10] and for $s$-wave hyperon decays[11]; e.g.,

$$S(\Sigma_+^+)=0\,,\qquad 2S(\Xi_-^-)+S(\Lambda_-^0)-\sqrt{3}\,S(\Sigma_0^+)=0\,. \tag{12}$$

Here $S(\Lambda_-^0)$, for example, is the $s$-wave amplitude for the decay $\Lambda\to p+\pi^-$. Notice however that we achieved this (and have achieved the conditions for substantial parity violation in nonleptonic decays) by requiring that $m \ll m(q)$, $m(q')$. It was this assumption that permitted us to approximate the interaction of Eq. (9) by the simpler one of Eq. (10). However, this exposes the quantitative problem why the low-lying physical hadrons are uncharmed. The problem is a general one for models involving charmed quarks and we can offer no convincing resolution.

### III. SUPPRESSION OF "DANGEROUS" PROCESSES

In a model in which nonleptonic weak interactions are principally mediated by a neutral scalar boson as in this model, there is a potential danger that $\Delta S=2$ transitions (such as $K^0 \longleftrightarrow \bar{K}^0$) and $|\Delta S|=1$ induced neutral current effects (such as $K_L \to \mu\bar{\mu}$) may become intolerably large. Fortunately in this model two-scalar meson contributions to these processes are suppressed for two reasons which we shall elaborate. We have shown elsewhere[7] that the two-$W$-exchange mechanism is suppressed for these processes in the eight-quark version we are discussing.

In making estimates of higher-order corrections, we observe that the main contribution comes from large internal momentum, and under such circumstances it is perhaps not unreasonable to assume that the relevant operator products of hadronic currents and densities at short distances are not affected by strong-interaction effects. The first reason why the above processes are suppressed in the present model is that the interaction of the Higgs scalar to the neutral fermions $\mathfrak{N}$, $\lambda$, $q$, and $q'$ possesses the permutation symmetry [see Eq. (8)]

$$q \longleftrightarrow q'$$
$$\mathfrak{N}_C \longleftrightarrow \lambda_C, \tag{13}$$

where

$$\mathfrak{N}_C=\mathfrak{N}\cos\theta+\lambda\sin\theta, \quad \lambda_C=\lambda\cos\theta-\mathfrak{N}\sin\theta,$$

in the limit $m(q)=m(q')$; $m(\mathfrak{N})=m(\lambda)$. This implies that the process $\mathfrak{N}\bar{\lambda}\to 2\phi$ is forbidden in this limit, as can be seen from the effective interaction for this process with the symmetry of Eq. (13):

$$(\bar{\mathfrak{N}}_C\mathfrak{N}_C+\bar{\lambda}_C\lambda_C)\phi^2=(\bar{\mathfrak{N}}\mathfrak{N}+\bar{\lambda}\lambda)\phi^2. \tag{14}$$

The second reason is a little more subtle and is kinematical in origin. Let us consider the process $\mathfrak{N}\bar{\lambda}\to\lambda\bar{\mathfrak{N}}$. There are four diagrams for this process. They are shown in Fig. 1. The contribution of Fig. 1(a) is of the order of

$$T^{(a)}(\mathfrak{N}\bar{\lambda}\to\lambda\bar{\mathfrak{N}}) \sim \frac{G_F\,\alpha\cos^2\beta\sin^2\beta\sin\theta}{48\pi\sqrt{2}}\left(\frac{\Delta m(q)}{M_W}\right)^2$$
$$\times\bar{\lambda}\gamma_\mu(1-\gamma_5)\mathfrak{N}\,\bar{\mathfrak{N}}\gamma^\mu(1-\gamma_5)\lambda, \tag{15}$$

where the suppression factor $[\Delta m(q)/M_W]^2$, $\Delta m(q)=m(q)-m(q')$ is evident. When we neglect the external momenta compared to the internal ones, the contribution of Fig. 1(b) is precisely $-T^{(a)}$. Likewise the contributions of Figs. 1(c) and 1(d) cancel. Thus, we find that the entire $2\phi$-exchange contribution to the $K^0 \longleftrightarrow \bar{K}^0$ transition vanishes to order

$$G_F\alpha\sin^2\theta[\Delta m(q)/53\text{ GeV}]^2 f_K{}^2 m_K{}^2,$$

where we have used the PCAC-inspired estimate

$$\langle K^0|\bar{\lambda}\gamma_\mu(1-\gamma_5)\mathfrak{N}\,\bar{\mathfrak{N}}\gamma^\mu(1-\gamma_5)\lambda|\bar{K}^0\rangle \simeq f_K{}^2 m_K{}^2.$$

The next leading term to the estimate of Eq. (15) is expected to be of order

$$\frac{\Delta m(K)}{m(K)} \sim G_F\,\alpha f_K{}^2\left(\frac{\Delta m(q)}{53\text{ GeV}}\right)^2\left(\frac{m(K)}{m(q)}\right)^2,$$

where $m(K)$ is the kaon mass and $\Delta m(K)=m(K_1)-m(K_2)$, and to be well within the experimental value $7\times10^{-15}$.

Next, let us consider the process $\mathfrak{N}\bar{\lambda}$ (or $\lambda\bar{\mathfrak{N}}$) $\to\bar{\mu}\mu$. The two diagrams in Fig. 2 are opposite in sign:

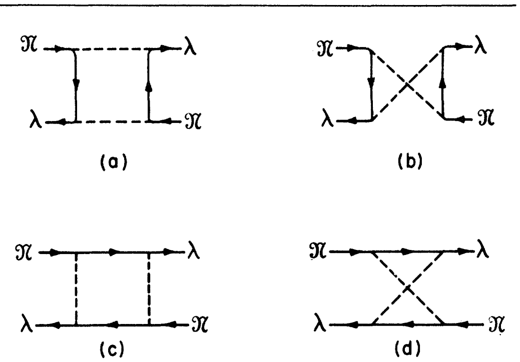


FIG. 1. Diagrams for the process $\mathfrak{N}\bar{\lambda}\to\lambda\bar{\mathfrak{N}}$. The dotted lines are the $\phi$ meson lines. The internal quark lines stand for both $q$ and $q'$.

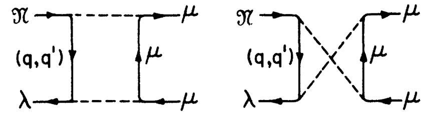


FIG. 2. Diagrams for the process $\mathfrak{N}\bar{\lambda}\to\mu\bar{\mu}$. The internal quark lines stand for both $q$ and $q'$.

$$T^a(\mathfrak{N}\bar{\lambda}\to\mu\bar{\mu}) = -T^b(\mathfrak{N}\bar{\lambda}\to\mu\bar{\mu})$$
$$\sim G_F\alpha\cos^2\beta\sin^2\beta\sin\theta\,\frac{m(q)\Delta m(q)}{M_W^{\,2}}\left(\frac{m(Y^0)}{m(q)}\right)^2\ln[m(q)/m_\phi]\,\bar{\lambda}\gamma_\mu(1-\gamma_5)\mathfrak{N}\,\bar{\mu}\gamma^\mu\mu\,. \qquad (16)$$

In addition to the cancellation of the two diagrams, there is another reason why the $2\phi$ contribution to $K_L\to 2\mu$ vanishes. It is that

$$\langle 0|\,\bar{\lambda}\gamma_\mu(1-\gamma_5)\mathfrak{N}\,|K_L\rangle\bar{\mu}\gamma^\mu\mu = -if_K P_\mu\bar{\mu}\gamma^\mu\mu = 0\,,$$

where $P_\mu$ is the momentum of the decaying particle.

*Note added in proof.* In our general discussion, before specializing to the Georgi-Glashow model, we contemplated situations involving the existence of only a single pair of charged gauge bosons $W^\pm$, where e.g., $W^+$ must then couple to both $\Delta S=0$ and $\Delta S=1$ hadron currents. We asserted that no arrangements with respect to possible neutral gauge bosons could be achieved to complete the $\Delta I=\frac{1}{2}$ rule. However, one can imagine a gauge group involving, in addition to neutral gauge bosons, two pairs $W_1^\pm, W_2^\pm$ of charged gauge bosons, such that $W_1^+$, say, couples to $\Delta S=0$, $W_2^+$ to $\Delta S=1$ hadron currents. In this case $\Delta S=1$ nonleptonic interactions do not arise at all, in lowest order, from charged gauge boson exchange. This, in turn, leaves open the possibility that $\Delta S=1$ nonleptonic interactions, with a pure $\Delta I=\frac{1}{2}$ character, can be generated by exchange of neutral gauge bosons. Just such a model has been produced recently by Pais.[12] It is based on broken O(4) gauge symmetry and incorporates *CP* violation as well as the $\Delta I=\frac{1}{2}$ rule.

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