internal vertices by attaching a fictitious zero-momentum external line to each of the internal vertices. The resulting graph then belongs to type I above.

Thus it suffices to discuss the reducibility of a φ^n vertex belonging to an unsaturated $\mathcal{G}_n'$. The main difference in treating the reducibility for an unsaturated $\mathcal{G}_n'$ versus that for $\mathcal{G}_n$ is as follows: In $\mathcal{G}_n$, the perfect symmetry makes it immaterial as to which vertex one chooses to set up a system of base vectors for the internal lines. Now for an unsaturated $\mathcal{G}_n'$, one must pick the vertex which has the maximal complexity. Consider now the vertex V with φ^n coupling [1 external and $(n-1)$ internal lines]. The procedure stated in paragraph (4) can be modified to show that these $(n-1)$ internal lines emerging from the vertex V cannot be linearly independent in a ν-dimensional space for $n>\nu+1$. This is all that is really necessary to force a reduction. The Hankel transform method which we have used is just one formal technique to handle this reduction.

It should be emphasized that for our present purpose it is sufficient to establish the reducibility of the φ^n vertex (hence the graphs). We shall not discuss here the end products of the reduction.[5]

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U(12) AND TIME-REVERSAL ASYMMETRY IN THE WEAK INTERACTIONS*

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It has been indicated that the symmetry U(12) may be useful in the study of strong interactions.[1] The subgroups SU(6) and U(6)⊗U(6) have been investigated in some detail with encouraging results.[1,2] In this note we postulate that two U(8) subgroups of U(12) are pertinent symmetries for the weak interactions (WI). This, coupled with the approximate $|\Delta I|=\frac{1}{2}$ rule, leads "naturally" to a violation of time-reversal (T) invariance in the WI with a branching ratio

$$R=\frac{\Gamma(K_L^{\,0}\to\pi^+ +\pi^-)}{\Gamma(K_S^{\,0}\to\pi^+ +\pi^-)}\sim 2\times10^{-6}, \quad (1)$$

as observed by Christenson et al.[3] (L and S stand for the long- and short-lived K^0 components). The theory also requires that (1) there be no T violation in leptonic decays, (2) there be T violation only in the parity-nonconserving piece of the WI, (3) there be T-violating effects of order 4% in strangeness-conserving weak nuclear processes, and of order 0.2% in nonleptonic hyperon decays.

Experimental studies of leptonic decays leave little doubt that the leptons l enter the WI only through the combination $j_\mu^{\,l}=\bar{l}\gamma_\mu(1+i\gamma_5)\nu_l$. Theoretical arguments have been presented which make such a world quite "natural." Feynman and Gell-Mann,[4] for example, assume that the most "fundamental" description of a fermion is given by a two-component object which satisfies a second-order differential equation. They further assume that it is precisely this two-component spinor that enters the WI, thereby obtaining the V-A theory. Feynman and Gell-Mann argue in this manner not only for leptons, but also for the "bare" hadrons (strongly interacting particles). This exclusion of scalar (s), pseudoscalar (p), and tensor (t) pieces from the hadronic part of the WI has not been experimentally verified. In fact, the argument against s, p, and t in nonleptonic decays rests solely in our belief that there exist "fundamental" or "bare" hadrons. While leptons may be "fundamental," we have every reason to believe that hadrons are not. Consequently, we postulate (i) the existence of s, p, and t in addition to the conventional v, a terms in the hadronic part of the WI; the leptonic piece remains the same.

Although there is no a priori reason to believe that the WI are in any way related to the strong-interaction symmetries, the conserved-vector-current theory[4] and its experimental verification[5] indicate that there may be some connection. Previously, this connection has been made by assuming that the hadronic piece $j_\mu{}^h$ of the WI current J_μ was equal to a linear combination of generator densities of $U(3)\otimes U(3)$[6,7]:

$$j_\mu{}^h=(v_{1\mu}'+iv_{2\mu}')+(a_{1\mu}'+ia_{2\mu}'), \tag{2}$$

where

$$v_{1\mu}'=v_{1\mu}\cos\theta+v_{4\mu}\sin\theta,$$
$$v_{2\mu}'=v_{2\mu}\cos\theta+v_{5\mu}\sin\theta \quad (\mu=0,\cdots,3), \tag{3}$$

and similarly for the axial currents $a_{k\mu}'$. The WI was then written as

$$\mathfrak{L}_{\rm WI}=g(j_\mu{}^l+j_\mu{}^h)X_\mu+\text{H.c.}=gJ_\mu X_\mu+\text{H.c.}, \tag{4}$$

where X_μ is the intermediate vector boson. Although the introduction of an intermediate vector boson is not necessary, we will find it notationally convenient. In the proposed theory, we connect the strong and weak interactions with the assumption (ii) that

$$\mathfrak{L}_{\rm WI}=g(aJX+J_\mu X_\mu+bJ_{\mu\nu}X_{\mu\nu})+\text{H.c.}, \tag{5}$$

where a, b are presently phenomenological parameters presumably of order 1; where $X, X_{\mu\nu}$ are scalar, tensor intermediate bosons; where the "generalized current" densities $J, J_\mu, J_{\mu\nu}$ contain s and p, v_μ and a_μ, $t_{\mu\nu}$ and $\bar{t}_{\mu\nu} \equiv \frac{1}{2}\epsilon_{\mu\nu\lambda\sigma}t_{\lambda\sigma}$ (tensor and dual tensor terms), respectively; and where s, p, v_μ, a_μ, and $t_{\mu\nu}$ are linear combinations of the Hermitian generator densities of U(12). Since the s, p, and t pieces vanish for leptons (e.g., by the Feynman–Gell-Mann argument),

$$\mathfrak{L}_{\rm WI}=g[aj^hX+(j_\mu{}^l+j_\mu{}^h)X_\mu+bj_{\mu\nu}{}^hX_{\mu\nu}]+\text{H.c.} \tag{6}$$

Our final assumption comes in specifying precisely what combination of U(12) generator densities should be chosen for j^h, $j_\mu{}^h$, and $j_{\mu\nu}{}^h$. To clarify matters, it is helpful to first write down the generator densities in a quark or ace model.[8] They are

$$S_k=\bar{A}\tfrac{1}{2}\lambda_k A,$$
$$P_k=\bar{A}\tfrac{1}{2}\lambda_k\gamma_5 A,$$
$$V_{k\mu}=\bar{A}\tfrac{1}{2}\lambda_k\gamma_\mu A,$$
$$A_{k\mu}=i\bar{A}\tfrac{1}{2}\lambda_k\gamma_\mu\gamma_5 A,$$
$$T_{k\mu\nu}=\bar{A}\tfrac{1}{2}\lambda_k\sigma_{\mu\nu}A \tag{7}$$
$$(\mu,\nu=0,\cdots,3;\ k=0,\cdots,8),$$

where we use the Feynman notation for the Dirac matrices, and A is a 12-component ace field. Note that all 144 generators thus defined are Hermitian.

From the point of view of the strong interactions, an important subgroup of U(12) would be that obtained by letting the index k range from 1 to 3, and adding $(\frac{2}{3})^{1/2}\times$ ($k=0$ generator density) $+(\frac{1}{3})^{1/2}\times$ ($k=8$ generator density). This subgroup, designated by $U_I(8)$, is the strong-interaction symmetry group we would deal with in a world without strangeness; it is a direct generalization of the familiar isospin. U(12) possesses, in addition, the subgroups $U_U(8)$ and $U_V(8)$, which are the generalizations of the U spin and V spin.[9] These subgroups will be of importance for the weak interactions.

We select from $U_V(8)$ the familiar v, a terms, and from $U_U(8)$ the s, p, t objects. The generator densities of $U_V(8)$ cannot be used directly in the definition of the WI, for we would have no explanation of the nonstrangeness changing ($\Delta S=0$) leptonic decays. However, we may introduce $\Delta S=0$ currents by picking them from the generators of the group obtained by rotating $U_V(8)$ either about the v_{60} or v_{70} axis through an angle $\pi-2\theta$. The U(8) indices 6 and 7 are chosen to introduce $\Delta S=0$, while v with the space-time subscript 0 is taken to guarantee that, under $\pi-2\theta$ rotation, Lorentz transformation properties of the generator densities are left invariant. Rotation about v_{60} or v_{70} changes $U_V(8)$ into the group we will call $U_V{}^6(8)$ or $U_V{}^7(8)$. Similarly, rotation of $U_U(8)$ leads to $U_U{}^6(8)$ or $U_U{}^7(8)$.

If we were to treat $U_V(8)$ and $U_U(8)$ in a symmetrical fashion but use rotations only about one axis to obtain a WI theory, we would find that $|\Delta S| = 2$ nonleptonic decays are allowed (in contradiction to the measured value of the $K_S{}^0$–$K_L{}^0$ mass difference). Both the rotations about v_{60} and v_{70} must come into play.

Explicitly, we find that under rotation about v_{70},

$$v_{4\mu} \to v_{1\mu}{}^7 = v_{1\mu}\cos\theta + v_{4\mu}\sin\theta,$$

$$v_{5\mu} \to v_{2\mu}{}^7 = v_{2\mu}\cos\theta + v_{5\mu}\sin\theta, \tag{8}$$

$$s_6 \to s_1{}^7 = -s_3\cos\theta\sin\theta + s_6(1-2\cos^2\theta) + \sqrt{3}s_8\cos\theta\sin\theta,$$

$$s_7 \to s_2{}^7 = s_7, \tag{9}$$

while under rotations about v_{60},

$$v_{4\mu} \to v_{1\mu}{}^6 = v_{2\mu}\cos\theta + v_{4\mu}\sin\theta,$$

$$v_{5\mu} \to v_{2\mu}{}^6 = -v_{1\mu}\cos\theta + v_{5\mu}\sin\theta, \tag{10}$$

$$s_6 \to s_1{}^6 = s_6,$$

$$s_7 \to s_2{}^6 = s_3\cos\theta\sin\theta + s_7(1-2\cos^2\theta) - \sqrt{3}s_8\cos\theta\sin\theta, \tag{11}$$

with identical results for $a_{k\mu}$ replacing $v_{k\mu}$ and $p_k, t_{k\mu\nu}$ replacing s_k. If we impose the condition $|\Delta S| \neq 2$ and treat $U_V(8)$ and $U_U(8)$ symmetrically, we are forced to consider only the combination

$$2^{-1/2}[(s_1{}^6 + is_2{}^6) + (s_1{}^7 + is_2{}^7)] \equiv s'; \tag{12}$$

similarly for p, v_μ, a_μ, and $t_{\mu\nu}$. This leads us to our final assumption (iii) that

$$j^h = s' + cp',$$

$$j_\mu{}^h = v_\mu{}' + da_\mu{}',$$

$$j_{\mu\nu}{}^h = t_{\mu\nu}{}' + e\bar{t}_{\mu\nu}{}', \tag{13}$$

where c, d, e are real parameters, presumably of order 1, to be determined by experiment or a future theory. It is usually assumed, for example, that the constant d is precisely equal to one; however, experimentally we only know that the matrix element of the axial strangeness-conserving current between neutron and proton has an effective strength of about 1.2. We therefore prefer to leave the parameters unspecified for the present. It is important, however, to make it clear that the predictions we shall make do not depend on assuming particular values for them. Note that there are eight charged v, a, and eight neutral s, p, t "currents."

In detail,

$$j^h = -(s_3 + cp_3)\cos\theta\sin\theta + \sqrt{2}[s_6 + cp_6 + i(s_7 + cp_7)]\sin^2\theta + \sqrt{3}(s_8 + cp_8)\cos\theta\sin\theta,$$

$$j_\mu{}^h = [v_{1\mu} + da_{1\mu} + i(v_{2\mu} + da_{2\mu})]\cos\theta + \sqrt{2}[v_{4\mu} + da_{4\mu} + i(v_{5\mu} + da_{5\mu})]\sin\theta. \tag{14}$$

The result for $j_{\mu\nu}{}^h$ is similar to that for j^h. We have absorbed an irrelevant phase $\exp(-i\pi/4)$ into all objects that do not carry strangeness. The angle θ which we use here is different from that introduced by Cabibbo,[6] because of the $\sqrt{2}$ multiplying the strangeness-carrying term in $j_\mu{}^h$.

Note that the SU(3) combinations occurring in the s, p, t "currents" are all contained in the U-spin subgroup. These "currents" therefore carry one unit of U spin. The s, p, t part of the effective weak Hamiltonian contains symmetrized products of these currents; hence, it only has U spin zero and two. The SU(3) content of the s, p, t contribution to the strangeness-changing part of the weak Hamiltonian is then pure $\underline{27}$; the strangeness-conserving part, on the other hand, contains in addition $\underline{1}$ and $\underline{8}$.

In the real world, where aces probably do not exist, Eq. (7) no longer holds. However, we may still keep the third assumption along with the corresponding Eqs. (8)-(14) to obtain a formulation of the WI.

We now discuss some consequences of the theory.

(1) T invariance is violated only in the parity-changing part of nonleptonic decays. The violation comes in through (sp) and $(t\bar{t})$ interference, just as P and C violation come from (va) interference. This would imply T-invariance violation in $K^0 \to 2\pi$, but not in $K^0 \to 3\pi$. T conservation in leptonic decays is consistent

with results obtained in neutron β decay and may be further checked by measuring the μ polarization out of the plane of decay in $K^+ \to \mu^+ + \nu_\mu + \pi^0$.

(2) The kaon branching ratio is

$$R = \frac{\Gamma(K_L^0 \to \pi^+ + \pi^-)}{\Gamma(K_S^0 \to \pi^+ + \pi^-)} \sim 2\times 10^{-6}.$$

To obtain this branching ratio, we follow the notation and results of Wu and Yang[10]; we write

$$R = |\eta_{+-}|^2, \tag{15}$$

where

$$\eta_{+-} = \tfrac{1}{2}[\epsilon + 2iF\,\mathrm{Im}A_2/A_0] + O(\lambda^2). \tag{16}$$

Here, A_0 and A_2 are the $I=0, 2$ $K_{\pi 2}$-decay amplitudes with A_0 chosen to be real and positive, while

$$F = \exp[i(\delta_2 - \delta_0)], \tag{17}$$

where δ_0 and δ_2 are the $\pi\pi$ s-wave scattering phase shifts for the $I=0, 2$ states at the rest mass of K^0. Next, we write

$$\epsilon = \frac{-M_i + i(y_l + y_{3\pi})}{A_0^{\,2} + i(m_S - m_L)} + O(\lambda^2), \tag{18}$$

where $m_S - m_L$ is the K_S^0-K_L^0 mass difference, and M_i is the imaginary part of the K-$\bar{K}$ mass operator. Finally, $\lambda = |A_2/A_0|$ and $O(\lambda^2)$ indicates the presence of terms of order λ^2 (λ measures the validity of the $|\Delta I| = \frac{1}{2}$ rule).[12] We wish to estimate η_{+-}. We find, first, that

$$A_2 = A_0[O(\lambda) + iO(\lambda \sin^2\theta)]. \tag{19}$$

The factor $\sin^2\theta$ in the imaginary part is due to the fact that T violation occurs as a result of sp interference, and from Eq. (14) this is seen to be suppressed by $\sin^2\theta$ relative to the T-conserving part of the interaction. The factor λ is present in Eq. (19) because A_2 is a $|\Delta I| = \frac{3}{2}$ amplitude.[13]

Next, let us turn to an estimate of ϵ. As can be seen from Eq. (18), we need to know M_i, y_l, and $y_{3\pi}$. Now in our theory, all leptonic decays conserve T, and $K^0 \to 3\pi$ conserves T as well since it is purely parity conserving. Thus, y_l and $y_{3\pi}$ are nonzero only because of our choice of phase of the K_0 state in requiring A_0 to be real; $y_l/A_0^{\,2}$ and $y_{3\pi}/A_0^{\,2}$ are very small, and may be neglected. For M_i, we may write the expression

$$M_i = \frac{1}{\pi}\sum_n \int_{M_n^{\,2}}^\infty \frac{ds}{s - m_K^{\,2}} |A_n(s)|^2 \sin 2\varphi_n(s), \tag{20}$$

where $A_n(s) = |A_n(s)| \exp i\varphi_n(s)$ is the off-energy-shell amplitude for the transition $K_0 \to n$, and M_n is the total mass of n. That is, $A_n(s)$ is the amplitude with which a K_0 with mass $\sqrt{s}$ would decay into the state n. The only states in the summation over n which could possibly be important are 2π in $I=0$ and 3π.

Because of the specific form which the T violation takes, as can be seen from Eq. (14) and from the fact that the $|\Delta S| = 1$ T violation is pure $\underline{27}$, we know the phase is of the order $\lambda \sin^2\theta$ for all s. For the 2π state with $I=0$, we can then estimate that $M_i \sim A_0^{\,2}\lambda \sin^2\theta$.

The 3π state will give a smaller contribution. The phase $\varphi_{3\pi}(s)$ (which in our theory happens to be energy independent since $K \to 3\pi$ conserves CP) is of order $\lambda \sin^2\theta$, as in the $2\pi, I=0$ case. However, we can expect $|A_{3\pi}(s)|^2/A_0^{\,2}$ to be small over much of the important energy range in the integration. At $s = m_K^{\,2}$, for example, we know experimentally that the ratio $|A_{3\pi}(s)|^2/A_0^{\,2} \sim 10^{-3}$; for smaller values of s, it will be even smaller. Above $s = m_K^{\,2}$, the ratio will grow, but it is nevertheless probable that the over-all 3π contribution to $M_i/A_0^{\,2}$ will not rise above a fraction of $\lambda \sin^2\theta$.

Altogether, then, we find $\epsilon \approx O(\lambda \sin^2\theta)$. Consequently, we expect both the ϵ and $\mathrm{Im}A_2/A_0$ terms to be comparable in Eq. (16), yielding

$$\begin{aligned} \eta_{+-} &\approx O(\lambda \sin^2\theta), \\ R &\approx O(\lambda^2 \sin^4\theta) = O\left[\frac{1}{650} \times (0.184)^4\right] \\ &= O(2\times 10^{-6}). \end{aligned} \tag{21}$$

Here, λ^2 is estimated from the experimental ratio

$$\frac{\Gamma(K^+ \to 2\pi)}{\Gamma(K_S^0 \to 2\pi)} \approx \frac{1}{650}, \tag{22}$$

while $\tan\theta$ is essentially given by the ratio of the $|\Delta S| = 1$ to $|\Delta S| = 0$ weak-interaction strength:

$$\frac{\Gamma(K^+ \to \mu\nu)}{\Gamma(\pi^+ \to \mu\nu)} = 2\tan^2\theta \frac{m_K}{m_\pi}\left[\frac{1 - (m_\mu/m_K)^2}{1 - (m_\mu/m_\pi)^2}\right]^2. \tag{23}$$

This yields $\theta \approx 0.19$. Our estimate of R is to be compared with the experimental value[3]

$$R_{\text{expt}} = (3.3 \pm 0.6) \times 10^{-6}. \tag{24}$$

(3) Hyperon nonleptonic decays should exhibit a 0.2% T-nonconserving effect. This is evident since the strangeness-changing T-invariance violation may proceed through the $\underline{27}$ piece of $s'p' \sim s_3 p_6 \cos\theta \sin^2\theta + \cdots$, which is a factor $\lambda \sin^2\theta \sim 1/500$ smaller than the usual T-conserving interaction. The $\Delta S = 0$ weak nuclear processes, on the other hand, should exhibit a T-violating effect of order 4%, since here there is an octet part to the weak T-violating Hamiltonian, and thus we may expect the T-violating term to be only $\sin^2\theta \sim 1/25$ smaller than the T-conserving one.

(4) The T-conserving $\Delta S = 0$ WI amplitudes satisfy an approximate $\Delta I = 0$ rule. Explicitly,

$$\frac{|\Delta I| = 1}{|\Delta I| = 0} \sim \sin^2\theta.$$

For the $\Delta S = 0$ T-nonconserving amplitudes, however,

$$\frac{|\Delta I| = 1}{|\Delta I| = 0} \sim 1.$$

Consequently, in $\Delta S = 0$, $|\Delta I| = 1$ transitions, T violation and conservation are comparable.

(5) Finally, the vector coupling constant for strangeness-conserving β decays is not G, but $G \cos\theta = 0.98G$. (Remember that our θ is smaller than Cabibbo's angle by a factor of $\sim\sqrt{2}$.) This is to be compared with the experimental value of $(0.978 \pm 0.0015)G$.[11]

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[9]$U_U(8)$ is generated by $\{(k=6), (k=7), -\frac{1}{2}(k=3)+\frac{1}{2}\sqrt{3}(k=8), (\frac{2}{3})^{1/2}(k=0)-\frac{1}{2}(k=3)-\frac{1}{2}\times 3^{-1/2}(k=8)\}$. $U_V(8)$ is generated by $\{(k=4), (k=5), \frac{1}{2}(k=3)+\frac{1}{2}\sqrt{3}(k=8), (\frac{2}{3})^{1/2}(k=0)+\frac{1}{2}(k=3)-\frac{1}{2}\times 3^{-1/2}(k=8)\}$. See S. Meshkov, C. A. Levinson, and H. J. Lipkin, Phys. Rev. Letters 10, 361 (1963).

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RELATIVISTIC TREATMENT OF SPIN INDEPENDENCE

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There have been several interesting attempts[1-4] made recently to combine within one relativistic framework the internal symmetries of the strong interactions with their spin independence. The purpose of this Letter is to present for rapid evaluation and assimilation a somewhat different treatment of the spin aspects of the problem. A brief survey of the difficulties involved in the usual approach to the problem is given in order to clarify the subsequent derivation of the new technique, but detailed applications to the problem of combining spin and internal symmetries will be reserved for presentation at greater length in some more suitable publication.

In most usual treatments the infinitesimal