for $\vec{k}_\alpha = -30$ MeV/c is 2.2 at $\theta_{c.m.} = 61°$ (Ref. 1). This ratio is easily understood from the full width at half-maximum of about 60 MeV/c [see R. Hagelberg, E. L. Haase, and Y. Sakamoto, Nucl. Phys. A207, 366 (1973)] for the d-α intercluster momentum distribution, $|F(k_\alpha)|^2$, found by analyzing the data from the reactions ^{6}Li(x, xd) and ^{6}Li(x, $x\alpha$). In Fig. 1 these values are used to determine the relative height between the two groups of points.

[5]M. L. Goldberger and K. M. Watson, *Collision Theory* (Wiley, New York, 1964).

[6]A. C. L. Barnard, C. M. Jones, and J. L. Weil, Nucl. Phys. 50, 604 (1964).

[7]The Watson-Migdal formula does not give the correct description for FSI effects in the higher $E_{p-\alpha}$ region where the factor $\mathfrak{F}_{FSI}$ calculated with this formula falls to zero. The Jost function (Ref. 5) has to be used for a more correct treatment for the enhancement of the cross sections.

[8]The c.m.'s of the d and α clusters are well separated from each other because of the small d-α separation energy of ^{6}Li. The asymptotic part of the d-α intercluster wave function dominates $\mathfrak{F}(k_\alpha)$ for small k_α. The antisymmetrization effects of particles between the clusters do not play an important role for QFR involving small k_α, although these effects may give rise to the $\vec{k}_\alpha$ dependence of the shrinkage of the clusters for larger values of k_α.

CP Nonconservation and the (J,ψ) Particles*

S. Pakvasa and S. F. Tuan
Department of Physics and Astronomy, University of Hawaii, Honolulu, Hawaii 97822
(Received 6 January 1975)

We consider in a schizon-model framework the suggestion of Nieh, Wu, and Yang that $J^0(3.105)$ [or $\psi^0(3.695)$?] exhibits CP nonconservation. Constraints are delineated for consistency with known CP nonconservation in $K_L{}^0 \to 2\pi$. We examine also the consequences of a different avenue for CP nonconservation should the J particles (with pairwise strong interactions to hadrons) be identified with the a particles of Lee.

Recently Nieh, Wu, and Yang[1] made the enormously interesting proposal that (i) the $J(3.105)$ particle,[2] as part of an isotopic doublet set (J^+,J^0) and ($J^-,\bar{J}^0$), is very likely to involve CP nonconservation analogous to that of the $W_a{}^0$-$W_b{}^0$ complex of the old schizon model,[3] and (ii) the doublets have pairwise strong interactions with hadrons (e.g., $J^+J^-, J^0\bar{J}^0 \leftrightarrow$ hadrons) and hence an additional additive quantum number t ($=\pm1$) conserved in strong and electromagnetic interactions but changed in weak interactions. Here we consider these propositions and their possible consequences for CP (or T via the CPT theorem) nonconservation given the currently known constraints in both CP-conserving and -nonconserving interactions.

Schizon model for (J, ψ).—We shall henceforth use the generic symbol X to denote $J(3.105)$ for convenience of notation. For the most part our estimates will be made on the assumption that $X = J(3.105)$ though they can be readily adapted for $\psi(3.695)$.[4] As pointed out by Nieh, Wu, and Yang,[1] $X^+(\bar{\nu}_l l)$ coupling has to be very small, hence $X^\pm$ cannot be the mediator $W^\pm$ for β decay. We shall assume that $X^\pm$ *is not coupled to the charge-changing lepton current at all.* Similarly the $X^0(\nu\bar{\nu})$ *coupling also has to be negligibly small,* hence we write the total interaction Lagrangian density in the following form:

$$\mathfrak{L} = \mathfrak{L}_{\text{strong}} + \mathfrak{L}_\gamma + \mathfrak{L}_{X^0\text{-}l} + \mathfrak{L}_{X\text{-}J} + \mathfrak{L}_{X\text{-}S}, \qquad (1)$$

where $\mathfrak{L}_\gamma$ denotes the electromagnetic interaction and

$$\mathfrak{L}_{X^0\text{-}l} = ig[\bar{e}\gamma_\lambda(1+\gamma_5)e + (e \to \mu)]X^0 + \text{H.c.}, \qquad (2)$$

$$\mathfrak{L}_{X\text{-}J} = gb[JX^* + J_a{}^0X_a{}^0 + J^*X] + g'J_b{}^0X_b{}^0, \qquad (3)$$

$$\mathfrak{L}_{X\text{-}S} = ga[SX^* + S^0X^{0*}] + \text{H.c.} \qquad (4)$$

Here $g^2/4\pi \sim 2\times10^{-6}$ (for $V-A$ lepton current) and $\sim 4\times10^{-6}$ (for pure V lepton current), consistent with the "medium-weak" leptonic decay properties of the $X^0(3.105)$ particle.[1] The combinations $X_a{}^0$ and $X_b{}^0$ are[5]

$$X_a{}^0 = 2^{-1/2}(-X^0 - X^{0*}), \quad X_b{}^0 = 2^{-1/2}i(X^0 - X^{0*}), \qquad (5)$$

while ($J, J_a{}^0, J^*$) and $J_b{}^0$ represent isovector and scalar hadronic (V, A) currents, respectively, for which[5] $\Delta N = \Delta S = 0$. For Eq. (4), S, S^0 is an isotopic spin-doublet hadronic (V, A) current for which $\Delta N = 0$, $\Delta S = -1$. The parameters b, a, and g' are to be constrained below by experiments involving both CP-conserving and -nonconserving decay interactions. *In particular, for* $g' \neq 0$ *and* *simple possible forms for* J_b^0 [*e.g.*, $\bar{p}p + \bar{n}n$ *or* $2(\Lambda\bar{\Lambda}) - (\bar{p}p) - (\bar{n}n)$], *violation of time-reversal invariance would occur from* $\mathfrak{L}_{X\text{-}S} + \mathfrak{L}_{X\text{-}J}$ *and hence for* $|\Delta S| = 1$ *nonleptonic processes.* It should also be stressed that though Eqs. (1)–(5) with X instead of the usual W intermediary do not represent "conventional" weak interactions (i.e., are not necessarily related to the square root of usual weak interactions), they do preserve the $|\Delta \vec{I}| = \frac{1}{2}$ rule for strangeness-changing decays of particles and suppression of $\Delta S = \pm 2$ interactions.

We next estimate b, a, and g' from decay constraints. For a full-fledged (V,A) doublet current for S and S^0, the $K_L{}^0 \to \mu^+\mu^-$ decay requires that

$$g^2 a/m_X{}^2 \lesssim G_F\alpha^2 \text{ or } a \lesssim 1.6\times 10^{-4}, \tag{6}$$

if we identify m_X with 3.105 GeV. From Eqs. (2), (3), and (5), it is evident that $X_a{}^0 = J(3.105)$. Hence, assuming that $\Gamma(X_a{}^0 \to \text{hadrons}) \lesssim 100$ keV, we have crudely $g^2b^2m_X \lesssim 100$ keV and, within reasonable limits,

$$1 < b < 10. \tag{7}$$

It seems reasonable that the near-degenerate orthogonal state $X_b{}^0(3.105)$ (which has only hadronic couplings) should have a substantially smaller hadronic width than $\Gamma(X_a{}^0)$. For definiteness, we take $\Gamma(X_b{}^0) \sim 20$ keV, hence

$$\Gamma(X_b{}^0) \sim 20 \text{ keV} \sim g'^2 m_X \text{ or } g' \sim 1/400. \tag{8}$$

The Fitch-Cronin CP-nonconservation parameter $|\eta_{+-}|$ in $K_L{}^0 \to 2\pi$ is given by

$$|\eta_{+-}| \sim (gg'a/m_X{}^2)(G_F/\sqrt{2})^{-1}. \tag{9}$$

For parameter a given by Eq. (6), the CP effects due to the present model are *far too small* to account for the observed value $|\eta_{+-}| \sim 2\times 10^{-3}$. We shall therefore make the futher assumption that the hadronic currents J, $J_a{}^0$, and J^* and S, S^0, S^*, and S^{0*} are *pure vector currents*[6]; likewise Eq. (2) is modified by dropping the axial-vector parts. It is important however to *retain (at least) the axial-vector piece for* $J_b{}^0$ since the decay $K_L{}^0 \to 2\pi$ is also parity nonconserving! The decay $K_L{}^0 \to 2\mu$ is no longer relevant to our discussion and the limit on parameter a is now determined from $K \to \pi l\bar{l}$ $(l = l, e)$. With $g^2/4\pi = 4\times 10^{-6}$, we have

$$g^2 a/m_X{}^2 < G_F\alpha, \text{ or } a < 1.5\alpha. \tag{10}$$

The estimate for $|\eta|$ from (9) and (10) is now compatible with the experimental number.

Such a model for CP (or T) nonconservation will of course preserve the $|\Delta \vec{I}| = \frac{1}{2}$ rule and hence $|\eta_{+-}| \sim |\eta_{00}|$ is expected. Except for $|\Delta S| = 1$ nonleptonic decays [with T-nonconserving amplitude $\sim 10^{-3}\times(T\text{-conserving amplitude})$], it would be difficult to find other evidence for nonconservation of CP. The neutron dipole moment[7] is of order $g^2a^2bgg'(m_N/m_X)^4 m_N{}^{-1}$ and is hence in the range $(10^{-26}-10^{-27})e$ cm. However establishment of $X_b{}^0$ of narrow hadronic width and the study of interference and regeneration effects in the decay of $X_a{}^0$ and $X_b{}^0$ analogous to the $K_S{}^0$-$K_L{}^0$ complex situation[8] will offer fairly convincing support for the model. Note that the decay angular distribution for $X_a{}^0 \to l\bar{l}$ $(l = e, \mu)$ is predicted to be $1 + \cos^2\theta$. Finally, the set of currents (J, S) can be taken to form an SU(3) octet.[9]

Modified chimeron model for (J, ψ).—Here we take advantage of the postulated pairwise strong interaction[1] $X\bar{X}$(hadrons) and identify the X or J particle with the a particle of Lee[10] analogous to the old $W \equiv a$ chimeron model.[11] We may retain the medium-weak structure given by Eqs. (2) to (7); however, we set $g' \equiv 0$. The source of CP nonconservation is now intrinsic to the properties attributed to the $X \equiv a$ particles themselves. In this case the $Y = S = 0$ isodoublets (X^+, X^0) and $(\bar{X}^0, X^-)$ have charges[1,10]

$$Q = Q_J + Q_K, \quad (Q_J, Q_K) = (I_z, t/2), \tag{11}$$

with the following quantum numbers (conserved in strong and electromagnetic interactions): for X^+,

$$(Q_J, Q_K) = (\tfrac{1}{2}, \tfrac{1}{2});$$

for $\bar{X}^0$,

$$(Q_J, Q_K) = (\tfrac{1}{2}, -\tfrac{1}{2});$$

for X^0,

$$(Q_J, Q_K) = (-\tfrac{1}{2}, \tfrac{1}{2});$$

for X^-,

$$(Q_J, Q_K) = (-\tfrac{1}{2}, -\tfrac{1}{2}).$$

They are therefore hybrid a particles in possessing both Q_K charges[12] and "ordinary" charges Q_J.[10,13] According to this model[14] CP nonconservation in weak (and strong) interactions involving just the usual hadrons is of order αf. Here f is the strong-coupling constant (~1) associated with the strong interaction of two neutral currents $\mathfrak{L}_{st} = fj_Xj_n$ (e.g., $j_X = \bar{X}\sigma_{\alpha\beta}\gamma_5 q_\beta X$, $j_n = \bar{N}\gamma_\alpha\gamma_5 N$); such an $\mathfrak{L}_{st}$ vertex is C-noninvariant. It has been pointed out[14,15] that the $W \equiv a$ model leads to the constraint that in the weak interaction CP nonconservation must take place in the first order in f and in the zero order in electric charge e, hence $f\sim 10^{-3}$. For the $X\equiv a$ model, the $X^{\pm}$ is not coupled to the charged lepton current and the medium-weak $\mathcal{L}_{X\text{-}S}$ vertex is drastically attenuated *vis-à-vis* $\mathcal{L}_{W\text{-}S}$ [cf. Eqs. (4) and (6)]. Thus in processes like $K_L{}^0\to 2\pi$, $\Lambda^0\to p+\pi^-$, $\nu+n\to\mu^-+p$, etc., the αf-type CP nonconservation (with $f\sim\frac{1}{10}$ to 1) can be the dominant contribution.

It has long been recognized[13] that CP-nonvariant vertices describing the emission of photons by ordinary hadrons must be necessarily isotopically invariant ($\Delta I=0$) independent of whether the a particles respect minimal electromagnetic interaction and whether these particles are isotopic singlets. Hence CP effects are absent in $\eta^0\to\pi^0+e^++e^-$ and $\Sigma^0\to\Lambda^0+e^++e^-$ and are drastically suppressed in $\eta^0\to\pi^+\pi^-\pi^0$. To obtain consistency with present limits on T nonconservation in the neutron dipole moment,[16] $e^-+N\to e^-+N^*(I=\frac{1}{2})$, and C nonconservation in $\eta^0\to\pi^+\pi^-\gamma$, it is best to assume that the $K_L{}^0\to 2\pi$ decay[17] represents "deci-electromagnetic" violation[14] (with $f\sim\frac{1}{10}$).

The $X,\overline{X}$ pair can be produced in strong interactions from $p+p\to X+\overline{X}+\dots$ or $p+\overline{p}\to X+\overline{X}+\dots$. The C-noninvariant effects must be of order f here and *large* asymmetry between the energy distribution of X and $\overline{X}$, say, is expected from the latter process [e.g., $dN(E_{X^+}=E_1,\ E_{X^-}=E_2)\neq dN(E_{X^+}=E_2,\ E_{X^-}=E_1)$. The decays of $X,\overline{X}$ are generally of medium-weak strength, satisfying the rule $|\Delta Q_J|=|\Delta Q_K|=\frac{1}{2}$; examples are $(X^{\pm},X_a{}^0,X_b{}^0)\to$ hadrons ($\Delta S=0$), $X_a{}^0\to l\overline{l}$ ($l=e,\mu$). Processes $(X^{\pm},X_a{}^0)\to(l\overline{l})+$hadrons ($\Delta S=0$) have strength fgb approximately semiweak.

Finally, Gell-Mann[18] has stressed the striking eightfold mass degeneracy obtained should X form bound states with a usual hadron. For instance, from the set $(X^+p, X^0p;\ \overline{X}{}^0p, X^-p)$ we generate, via CPT, the set $(X^-\overline{p},\overline{X}{}^0\overline{p};\ X^0\overline{p},X^+\overline{p})$ with medium-weak decays with $\Delta S=0$ to hadrons and leptons. Such spectacular events should be watched for in the neighborhood of m_J+m_p and/or $m_\psi+m_p$. Again if the decay $\psi^0(3.695)\to J^0(3.105)+(\pi\pi)$ is strong,[19] we would be tempted to assign similar isodoublet structures with $Q_K=t/2=\pm\frac{1}{2}$ for both ψ and J. This raises the further possibility that bootstrap between them might lead to an enriched mass spectrum replete with "nonexotic" $Q_K=0$ states ($\psi\overline{\psi}$, $\psi\overline{J}$, and $\overline{\psi}J$) coupled to normal hadrons and perhaps exotic states (ψJ, $\psi\psi$, and JJ) with $Q_K\neq 0$ (akin to the Z^* hyperon resonance).

We have considered also the nonhybrid ($Q_K=1$, $Q_J=0$) isotriplet $a\equiv X$ particles X^{++}, X^+, and X^0. However the resulting model is extremely unattractive.

Note added.—The present model with a vector $(X^+,X_a{}^0,X^-)$ would, however, predict a predominant decay of $X_a{}^0(3.105)$ into $2\pi^+2\pi^-$, $3\pi^+3\pi^-$, etc. If it turns out that this is not the case experimentally, we can modify our models in the following way: We choose $(J,J_a{}^0,J^*)$ to be axial-vector, hence $\Gamma(2\pi^+2\pi^-+3\pi^+3\pi^-+\dots)\approx 0$. Either the leptonic coupling of $X_a{}^0$ is now axial vector or else there will be parity-nonconserving effects in the resonance region. Though the $X_b{}^0$ has a significant decay into $2\pi^+2\pi^-$, $3\pi^+3\pi^-$, etc., it is not coupled to lepton pairs and hence can only be identified through production in hadronic-initiated reactions.

---

*Work supported in part by the U. S. Atomic Energy Commission under Contract No. AT(04-3)-511.

[1]H. T. Nieh, T. T. Wu, and C. N. Yang, Phys. Rev. Lett. 34, 49 (1975).

[2]J. J. Aubert *et al.*, Phys Rev. Lett. 33, 1404 (1974); J.-E. Augustin *et al.*, Phys. Rev. Lett. 33, 1406 (1974); C. Bacci *et al.*, Phys. Rev. Lett. 33, 1408 (1974).

[3]T. D. Lee and C. N. Yang, Phys. Rev. 119, 1410 (1960).

[4]G. S. Abrams *et al.*, Phys. Rev. Lett. 33, 1453 (1974); see also, however, J. J. Aubert *et al.*, Phys. Rev. Lett. 33, 1624 (1974).

[5]Unexplained notations are as given in Ref. 3.

[6]This assumption is needed at least for the $|\Delta S|=1$ currents. For simplicity we assume that it is also true for the set $(J,J_a{}^0,J^*)$.

[7]R. Gourishankar and A. N. Kamal, Can. J. Phys. 47, 693 (1969); S. Pakvasa, unpublished.

[8]V. L. Fitch, R. F. Roth, J. S. Russ, and W. Vernon, Phys. Rev. Lett. 15, 73 (1965).

[9]T. D. Lee and C. S. Wu, Annu. Rev. Nucl. Sci. 15, 381 (1965).

[10]T. D. Lee, Phys. Rev. 140, B959 (1965).

[11]S. Pakvasa, S. F. Tuan, and T. T. Wu, Phys. Rev. Lett. 20, 1546 (1968).

[12]Since Q_K charge is essential towards obtaining CP nonconservation here, it can be said to supply a physically meaningful quantum number in place of t.

[13]L. B. Okun, Phys. Lett. 23, 595 (1966); we regard here the doublet X particles as themselves *elementary* rather than as a "composite" of a normal hadron doublet and an X singlet.

[14]L. B. Okun and C. Rubbia, in *Proceedings of the International Conference on Elementary Particles, Heidelberg, Germany, 1967,* edited by H. Filthuth (North-Holland, Amsterdam, The Netherlands, 1968), p. 301.

[15]I. Yu. Kobsarev and L. B. Okun, Phys. Lett. 27B, 172 (1968).

[16]L. Wolfenstein, CERN Report No. TH. 1837, 1974 (unpublished).

[17]The αf-type CP nonconservation does not of course

guarantee $|\eta_{+-}| \sim |\eta_{00}|$, hence dynamical suppression of the $|\Delta \vec{I}| = \frac{3}{2}$ component must be invoked. It is not ruled out that the source of CP nonconservation in $K_L{}^0 \to 2\pi$ is indeed dominantly superweak with $|\eta_{+-}| = |\eta_{00}|$ and that the electromagnetic αf effects represent an *additional* source of CP nonconservation which contribute <10% of $|\eta|$ for $f \lesssim \frac{1}{20}$.

[18]M. Gell-Mann, in *Fundamental Interactions at High Energy I*, edited by T. Gudehus, G. Kaiser, and A. Perlmutter (Gordon and Breach, New York, 1969), p. 380; S. Pakvasa and S. F. Tuan, Nucl. Phys. B36, 173 (1972).

[19]In this connection the possiblity of pairwise strong interactions was also raised by G. Rajasekaran, private communication.

Theorem on the Scalar Form Factor in K_{l3} Decay

P. R. Auvil and P. L. Pritchett
Department of Physics, Northwestern University, Evanston, Illinois 60201
(Received 23 December 1974)

A current-algebra result is derived for the slope of the scalar form factor in K_{l3} decay which is valid at all momentum transfers. This theorem is compared with the previous results of Dashen and Weinstein and Dashen, Li, Pagels, and Weinstein. Our analysis indicates that there are large corrections to these previous results at $t = m_K{}^2 + m_\pi{}^2$.

The semileptonic decays of the K meson have provided an important testing ground for basic theoretical ideas regarding the symmetry properties and the dynamics of the strong interaction. In particular, considerable effort has been devoted to understanding the momentum-transfer dependence of the form factors in K_{l3} decays. A theorem regarding the slope of the scalar form factor in such decays was obtained some time ago by Dashen and Weinstein[1] and was later modified by Dashen, Li, Pagels, and Weinstein[2] (DLPW). In its amended form the theorem determined the slope at the unphysical point $t = m_K{}^2 + m_\pi{}^2$ to be $\frac{1}{2}(F_K/F_\pi - F_\pi/F_K) + O(\epsilon)$, where the parameter ϵ sets the scale of chiral symmetry breaking. The leading term, $\frac{1}{2}(F_K/F_\pi - F_\pi/F_K)$, is of order $\epsilon \ln \epsilon$. In this paper we generalize the DLPW analysis to yield a theorem valid for all t. We then make a simple estimate which indicates that the corrections to the DLPW result are of order 50% at $t = m_K{}^2 + m_\pi{}^2$. Finally, we give an exact evaluation of the theorem based on a previous analysis[3] we have made of three-point functions within the framework of the $(\underline{3}, \underline{3}^*)$ model of chiral symmetry breaking. We again find corrections on the order of 50% at $t = m_K{}^2 + m_\pi{}^2$. The reason the corrections are so large is that they are of order $2m_K{}^2/m_\kappa{}^2$ relative to the leading term.[4] The exact predictions of the theorem are in good agreement with the experimental data.[5]

First we derive the generalized theorem. We define the standard K_{l3} form factors $f_\pm(t)$ by

$$\langle M_a(q)| V_b{}^\mu(0)| M_c(k)\rangle = i f_{abc}[f_+(t)(k+q)^\mu + f_-(t)(k-q)^\mu], \tag{1}$$

where $t = (k-q)^2$. Our interest is in the function $D_{abc}(t)$ defined by

$$D_{abc}(t) \equiv i f_{abc}\tilde{D}_{abc}(t) \equiv i\langle M_a(q)|\partial_\mu V_b{}^\mu(0)| M_c(k)\rangle = i f_{abc}[(m_c{}^2 - m_a{}^2)f_+(t) + t f_-(t)]. \tag{2}$$

We begin by studying the function

$$S(q^2, k^2, t') \equiv \int d^4x\, d^4y\, e^{iqx} e^{-iky}\langle 0| T[\partial_\mu A_a{}^\mu(x)\partial_\sigma V_b{}^\sigma(0)\partial_\nu A_c{}^\nu(y)]|0\rangle, \tag{3}$$

which we consider as a function of the variables q^2, k^2, and $t' \equiv (k-q)^2 - \sigma(k^2+q^2)$. σ is an arbitrary constant. For $\sigma = 1$ we reproduce the DLPW analysis. It is useful to isolate the meson poles appearing in Eq. (3) since this will enable us to identify the on-shell amplitude D_{abc}. Thus, following DLPW, we introduce currents $\hat{A}^\mu$ in which the meson poles have been removed,

$$\partial_\mu \hat{A}^\mu(p) = \partial_\mu A^\mu(p) - (p^2 - m^2)^{-1} \lim_{p^2 \to m^2} (p^2 - m^2)\partial_\mu A^\mu(p). \tag{4}$$