SU(4) × U(1) gauge theory. II. *CP* nonconservation

Nilendra G. Deshpande, Rudolph C. Hwa, and Philip D. Mannheim
Institute of Theoretical Science, University of Oregon, Eugene, Oregon 97403
(Received 6 July 1977)

We exploit the higher symmetry inherent in an SU(4) × U(1) gauge theory to construct a spontaneously broken theory of *CP* nonconservation. Higgs multiplets in the adjoint representation of SU(4) contain both even and odd *CP* fields; thus, requiring the simultaneous nonvanishing of the vacuum expectation values of these fields leads to *CP* noninvariance of the vacuum. We find that all the *CP*-nonconserving effects are mediated in our theory by the superheavy gauge bosons of the broken SU(4) × U(1) symmetry. In fact, the very existence of *CP* violation sets an upper limit on the masses of these bosons. In our model the dominant *CP* effect lies in the neutral kaon system and is found to arise through a direct ($\Delta S = 2$) K_1-K_2 transition. The model has all the features of a superweak theory, with a neutron electric dipole moment substantially smaller than 10^{-24} e cm.

I. INTRODUCTION

We have recently proposed[1] a weak-interaction model based on the gauge group SU(4) × U(1). In such a model the four leptons (ν_e, e^-, μ^-, ν_μ) and the four quarks (u, d, s, c) form fundamental representations of the gauge group. The model incorporates the SU(2) × U(1) gauge group of Weinberg and Salam[2] as an approximate subgroup, while the additional gauge bosons present have high masses. These gauge bosons allow many exotic interactions, for example, muon-number-nonconserving processes. These have been treated fully in paper I. In this paper we show that the model also allows for the spontaneous violation of *CP* invariance. Such *CP*-nonconserving effects are, however, limited to the interactions involving the heavy gauge bosons only. A small K_1-K_2 mixing arises due to these interactions, leading naturally to a superweak theory of *CP* violation.

Most models of *CP* nonconservation in the context of gauge theories are of the milliweak variety. Various models based on the presence of right-handed currents have been proposed.[3] In an alternate scheme[4] it has been shown that the presence of three or more Higgs doublets in the SU(2) × U(1) model leads to *CP* nonconservation. Both these approaches expect the neutron electric dipole moment to be of the order of 10^{-24} e cm, a value very close to the present experimental limit.[5] The value expected in our model is much smaller; consequently, a small improvement in the experimental limit will serve to distinguish the various models.

We present our model in Sec. II, and discuss its phenomenology in Sec. III. In Sec. IV we discuss the relation between *CP* nonconservation and the Cabibbo angle; our conclusions are contained in Sec. V.

II. THE MODEL

The basic model has already been described in paper I. Here we present only the specific aspects which are essential for understanding *CP* nonconservation. There are 16 gauge bosons in the model, and the 4 quarks and 4 leptons belong to fundamental representations of SU(4). As before, we break the symmetry down to SU(2) × U(1) by introducing two adjoint multiplets of Higgs bosons, ϕ and ψ. In the tensor basis defined by

$$\phi_{ab} = \sum_i \lambda^i_{ab}\phi_i \,, \tag{2.1}$$

the vacuum expectation value of ϕ_{ab} is chosen, as in paper I:

$$\langle\phi_{ab}\rangle = \eta_0\begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & -1 & 0 \\ 0 & 0 & 0 & -1 \end{pmatrix}. \tag{2.2}$$

However, for ψ_{ab} we now choose

$$\langle\psi_{ab}\rangle = \xi_0\begin{pmatrix} 0 & 0 & 0 & e^{-i\delta} \\ 0 & 0 & e^{-i\delta} & 0 \\ 0 & e^{i\delta} & 0 & 0 \\ e^{i\delta} & 0 & 0 & 0 \end{pmatrix}. \tag{2.3}$$

This choice for $\langle\psi_{ab}\rangle$ differs from our previous one in paper I only in that the parameter δ is nonvanishing. Equations (2.2) and (2.3) imply in the canonical basis

$$\langle\phi_8\rangle = \sqrt{2}\langle\phi_{15}\rangle = 2\eta_0\sqrt{3}\,, \tag{2.4a}$$

$$\langle\psi_6\rangle = \langle\psi_9\rangle = \xi_0\cos\delta\,, \tag{2.4b}$$

$$\langle\psi_7\rangle = \langle\psi_{10}\rangle = \xi_0\sin\delta\,. \tag{2.4c}$$

In paper I we have shown that the parameters in the Higgs potential can be chosen so as to produce a minimum of the potential consistent with Eqs. (2.2) and (2.3). The value of δ itself is, however, arbitrary and implications of the nonvanishing of δ were not explored in paper I. We see now that the presence of the angle δ leads to CP nonconservation because ψ_6 and ψ_7 possess opposite CP quantum numbers, as do ψ_9 and ψ_{10}. Nonvanishing of the quantities in Eqs. (2.4b) and (2.4c) for arbitrary δ implies that the vacuum is no longer CP invariant. As explained in paper I, the choice of the vacuum expectation values given above leaves unbroken the usual Weinberg-Salam subgroup generated by

$$G \equiv \{F_1 + F_{13}, F_2 - F_{14}, F_3 + \frac{1}{\sqrt{3}}F_8 - (\tfrac{2}{3})^{1/2}F_{15}, F_0\}. \tag{2.5}$$

The mass matrix of the heavy gauge bosons can easily be found using formulas presented in paper I. We find it convenient to introduce a new basis

$$\begin{pmatrix} \pm i\tilde{U}_\pm \\ \tilde{W}_6 \\ \tilde{W}_9 \end{pmatrix} = \cos\delta \begin{pmatrix} \pm iU_\pm \\ W_6 \\ W_9 \end{pmatrix} + \sin\delta \begin{pmatrix} V_\pm \\ W_7 \\ W_{10} \end{pmatrix}, \tag{2.6}$$

$$\begin{pmatrix} \tilde{V}_\pm \\ \tilde{W}_7 \\ \tilde{W}_{10} \end{pmatrix} = -\sin\delta \begin{pmatrix} \pm iU_\pm \\ W_6 \\ W_9 \end{pmatrix} + \cos\delta \begin{pmatrix} V_\pm \\ W_7 \\ W_{10} \end{pmatrix}. \tag{2.7}$$

The fields $V_\pm$, $U_\pm$ above, as well as $X_\pm$, $W_\pm$, S, T, Z, A to appear presently, have been defined in paper I, Eqs. (2.3) and (3.5). The mass term in the Lagrangian is found to be

$$\mathcal{L}_{\rm mass} = \tfrac{1}{2}g^2\eta_0^2(2\tilde{U}_+\tilde{U}_- + \tilde{W}_6^2 + \tilde{W}_9^2 + 2\tilde{V}_+\tilde{V}_- + \tilde{W}_7^2 + \tilde{W}_{10}^2) + \tfrac{1}{2}g^2\zeta_0^2(2X_+X_- + S^2 + T^2 + 2\tilde{V}_+\tilde{V}_- + \tilde{W}_7^2 + \tilde{W}_{10}^2). \tag{2.8}$$

From Eq. (2.8) we obtain the following important mass relations:

$$\begin{aligned} &M^2(\tilde{W}_7) = M^2(\tilde{W}_{10}) = M^2(\tilde{V}_\pm), \\ &M^2(\tilde{W}_6) = M^2(\tilde{W}_9) = M^2(\tilde{U}_\pm), \\ &M^2(\tilde{W}_7) - M^2(\tilde{W}_6) = M^2(S) = M^2(T) = M^2(X_\pm). \end{aligned} \tag{2.9}$$

It should be noted that rotation through δ leaves invariant the SU(2) $\times$ U(2) $\times$ U(1) subgroup associated with gauge bosons $W_\pm$, Z, $X_\pm$, S, T, A. Consequently, their interactions with the fermions are as in paper I, and the associated phenomenology presented there is unaltered. However, the heavier states $U_\pm$, $V_\pm$, W_6, W_7, W_9, and W_{10} have undergone mixing leading to transitions between currents of opposite CP properties. For our purposes here, we present only the interactions of the quarks and gauge bosons (couplings to leptons can be inferred by analogy). In terms of the fields given in Eqs. (2.6) and (2.7) and the Cabibbo-rotated quark fields

$$\begin{pmatrix} u' \\ c' \end{pmatrix} = \begin{pmatrix} \cos\theta_C & -\sin\theta_C \\ \sin\theta_C & \cos\theta_C \end{pmatrix} \begin{pmatrix} u \\ c \end{pmatrix}, \tag{2.10}$$

the Lagrangian of paper I, Eq. (3.7) is reexpressed as

$$\begin{aligned} \frac{\sqrt{2}}{g}\mathcal{L}_{\rm int} = &-\sqrt{2}\sin\theta_W(A + \tan\theta_W Z)J^{\rm em} + (\sqrt{2}\cos\theta_W)^{-1}Z(\bar{u}u + \bar{c}c - \bar{d}d - \bar{s}s) + W_+(\bar{u}'d + \bar{c}'s) + X_+(\bar{u}'d - \bar{c}'s) \\ &+ S(\bar{d}d - \bar{s}s) + T(\bar{u}'u' - \bar{c}'c') + \tilde{W}_6[\cos\delta(\bar{d}s + \bar{s}d) + i\sin\delta(\bar{s}d - \bar{d}s)] \\ &+ \tilde{W}_7[i\cos\delta(\bar{s}d - \bar{d}s) - \sin\delta(\bar{d}s + \bar{s}d)] + \tilde{W}_9[\cos\delta(\bar{u}'c' + \bar{c}'u') + i\sin\delta(\bar{c}u - \bar{u}c)] \\ &+ \tilde{W}_{10}[i\cos\delta(\bar{c}u - \bar{u}c) - \sin\delta(\bar{u}'c' + \bar{c}'u')] + \tilde{V}_+[\cos\delta(\bar{u}'s - \bar{c}'d) + i\sin\delta(\bar{u}'s + \bar{c}'d)] \\ &+ \tilde{U}_+[\cos\delta(\bar{u}'s + \bar{c}'d) + i\sin\delta(\bar{u}'s - \bar{c}'d)] + \text{H.c. of the charged sector.} \end{aligned} \tag{2.11}$$

Here $\bar{a}b$ stands for $\bar{a}\gamma_\lambda\frac{1}{2}(1-\gamma_5)b$. The masses of $W_\pm$ and Z are generated as in paper I by the introduction of additional Higgs bosons, χ_i^s, in the fundamental representation of SU(4). This breaks the SU(2) $\times$ U(1) symmetry and causes some small mixing of the light bosons with the heavy ones as well as mixing among the heavy bosons. This effect, though crucial for processes such as $\mu \to e + \gamma$, is completely insignificant as far as CP nonconservation is concerned, and we shall neglect it here.

III. PHENOMENOLOGY OF *CP* NONCONSERVATION

In this section we shall estimate the CP-violation effects. If we ignore the heavy gauge bosons for the moment, the $K_1 - K_2$ decay matrix is diagonal and can be represented as

$$\begin{pmatrix} m_1 - i\Gamma_1/2 & 0 \\ 0 & m_2 - i\Gamma_2/2 \end{pmatrix}. \tag{3.1}$$

Here Γ_1 and Γ_2 are the lifetimes of K_1 and K_2, and m_1 and m_2 are their respective masses including weak corrections, but not including CP-violation effects.

Inspection of the interactions in Eq. (2.11) reveals two different mechanisms for CP violation. The dominant mechanism is the direct $\Delta S = 2$ transition from K_1 to K_2 via $\tilde{W}_6$ and $\tilde{W}_7$ exchange. These contributions do not cancel because $\tilde{W}_6$ and $\tilde{W}_7$ are not degenerate. In fact, it is the very nonvanishing of $\langle\psi_6\rangle$, $\langle\psi_7\rangle$, and $\langle\phi_8\rangle$ which forces this nondegeneracy, as can be seen from Eq. (2.8). A much smaller $\Delta S = 1$ contribution arises through $\tilde{U}_\pm$ and $\tilde{V}_\pm$ exchange. We first evaluate the $\Delta S = 2$ contribution. We define

$$\langle 0 | \tfrac{1}{2}[\bar{d}\gamma_\mu(1-\gamma_5)s + \bar{s}\gamma_\mu(1-\gamma_5)d] | K_i\rangle = iF_K k_\mu \delta_{i1}, \tag{3.2}$$

$$\langle 0 | i\tfrac{1}{2}[\bar{d}\gamma_\mu(1-\gamma_5)s - \bar{s}\gamma_\mu(1-\gamma_5)d] | K_i\rangle = iF_K k_\mu \delta_{i2}, \tag{3.3}$$

where $i = 1$ or 2, and F_K is the kaon decay constant ($F_K = 1.28F_\pi \approx 120$ MeV). We then obtain

$$\Delta_{12} \equiv \langle K_1 | i\mathcal{L}_{\text{int}} | K_2\rangle = \sqrt{2}G_F \sin 2\delta\, F_K{}^2 m_K{}^2 \left[\frac{M_W{}^2}{M^2(\tilde{W}_6)} - \frac{M_W{}^2}{M^2(\tilde{W}_7)}\right], \tag{3.4}$$

where $G_F\sqrt{2} = g^2/8M_W{}^2$. We now calculate the CP-violating parameter ϵ,

$$\epsilon \equiv \frac{A(K_L \to 2\pi)}{A(K_S \to 2\pi)} = \frac{\Delta_{12}}{2m_K(m_L - m_S - i\Gamma_S/2)}, \tag{3.5}$$

where we have inserted the K_S pole in the $K_L \to 2\pi$ amplitude. The small mass difference in the denominator justifies this pole dominance, and enhances this mechanism for ϵ. Since $m_L - m_S \approx \Gamma_S/2$ experimentally, we have

$$\epsilon = \frac{\Delta_{12} e^{i\phi}}{2\sqrt{2}m_K(m_L - m_S)}, \quad \phi \approx \pi/4. \tag{3.6}$$

The experimental values of $|\epsilon| \sim 2 \times 10^{-3}$ and $(m_L - m_S)/m_K \sim 7 \times 10^{-15}$ lead to the following useful relation:

$$\left| \sin 2\delta \left[\frac{M_W{}^2}{M^2(\tilde{W}_6)} - \frac{M_W{}^2}{M^2(\tilde{W}_7)}\right] \right| = 1.8 \times 10^{-10}. \tag{3.7}$$

To estimate the contribution from $\tilde{U}_\pm$ and $\tilde{V}_\pm$ exchange to ϵ, we note that the effective CP-nonconserving interaction due to their exchange is obtained from Eq. (2.11) as

$$\mathcal{L}_{\text{eff}} = \frac{G_F}{\sqrt{2}}\left[\frac{M_W{}^2}{M^2(\tilde{U}_+)} - \frac{M_W{}^2}{M^2(\tilde{V}_+)}\right] \sin\theta_C \cos\theta_C e^{-2i\delta}[\bar{s}\gamma^\lambda(1-\gamma_5)u\bar{u}\gamma_\lambda(1-\gamma_5)d - \bar{s}\gamma^\lambda(1-\gamma_5)c\bar{c}\gamma_\lambda(1-\gamma_5)d] + \text{H.c} \tag{3.8}$$

The contribution to ϵ arising from $\mathcal{L}_{\text{eff}}$ is of order

$$\left[\frac{M_W{}^2}{M^2(\tilde{U}_+)} - \frac{M_W{}^2}{M^2(\tilde{V}_+)}\right]\sin 2\delta. \tag{3.9}$$

Since $M^2(\tilde{U}_+) = M^2(\tilde{W}_6)$ and $M^2(\tilde{V}_+) = M^2(\tilde{W}_7)$ we see from Eq. (3.7) that the quantity in Eq. (3.9) is only 1.8×10^{-10}, and is therefore negligible. It is because of the nearby pole that the $\Delta S = 2$ contribution dominates over the $\Delta S = 1$ exchange in our model.

From Eq. (3.7) we obtain the inequality

$$\frac{M_W{}^2}{M^2(\tilde{W}_6)} > \frac{M_W{}^2}{M^2(\tilde{W}_7)} + 1.8 \times 10^{-10} \tag{3.10}$$

or

$$M(\tilde{W}_6) < 7 \times 10^4 M_W. \tag{3.11}$$

We thus find that the heavy gauge bosons of our theory cannot be made indefinitely massive. To obtain a further constraint we calculate the contribution to the K_L-K_S mass difference due to $\tilde{W}_6$ and $\tilde{W}_7$ exchange. We have

$$\Delta_{11} \equiv \langle K_1 | i\mathcal{L}_{\text{int}} | K_1\rangle = -\frac{g^2}{2}F_K{}^2 m_K{}^2\left[\frac{\cos^2\delta}{M^2(\tilde{W}_6)} + \frac{\sin^2\delta}{M^2(\tilde{W}_7)}\right], \tag{3.12}$$

$$\Delta_{22} \equiv \langle K_2 | i\mathcal{L}_{\text{int}} | K_2\rangle = -\frac{g^2}{2}F_K{}^2 m_K{}^2\left[\frac{\cos^2\delta}{M^2(\tilde{W}_7)} + \frac{\sin^2\delta}{M^2(\tilde{W}_6)}\right]. \tag{3.13}$$

From the above equations and from Eq. (3.4) we obtain the result

$$\frac{m_L - m_S}{m_K} = \sqrt{2}G_F F_K{}^2\left[\frac{M_W{}^2}{M^2(\tilde{W}_6)} - \frac{M_W{}^2}{M^2(\tilde{W}_7)}\right], \tag{3.14}$$

which is independent of δ in leading order. This contribution should be less than the experimental value, hence we have the constraint

$$\left|\frac{M_W{}^2}{M^2(\tilde{W}_6)} - \frac{M_W{}^2}{M^2(\tilde{W}_7)}\right| < 3 \times 10^{-8}. \tag{3.15}$$

Using Eqs. (3.15) and (3.7) we find that

$$\sin 2\delta > 6 \times 10^{-3}. \tag{3.16}$$

Thus the smallness of the right-hand side of Eq. (3.7) cannot arise from the smallness of $\sin\delta$, but rather arises from the large masses of the gauge bosons. Indeed, δ can be as large as $\pi/4$, which corresponds to maximal CP nonconservation (i.e., $\langle\psi_6\rangle = \langle\psi_7\rangle$); this case will be discussed below.

At this point we recognize the essence of our approach to be a subtle interplay between the breakings of global and local symmetries. By allowing $\langle\psi_7\rangle \neq 0$ we break a discrete CP invariance and also a continuous global SU(4) × U(1) invariance. The local extension of the latter then gives heavy masses to the gauge bosons making the CP-nonconservation effects small.

We consider next the $K_{L,S} \to \mu e$ decays. Both $\tilde{W}_6$ and $\tilde{W}_7$ contribute, and calculation of the decay rates gives the ratios

$$\frac{\Gamma(K_L\to\mu e)}{\Gamma(K^+\to\mu\nu)}=\frac{2M_W^4}{\sin^2\theta_C}\left[\frac{\cos^2\delta}{M^4(\tilde W_7)}+\frac{\sin^2\delta}{M^4(\tilde W_6)}\right] \quad (3.17)$$

and

$$\frac{\Gamma(K_S\to\mu e)}{\Gamma(K^+\to\mu\nu)}=\frac{2M_W^4}{\sin^2\theta_C}\left[\frac{\sin^2\delta}{M^4(\tilde W_7)}+\frac{\cos^2\delta}{M^4(\tilde W_6)}\right], \quad (3.18)$$

where for simplicity we have assumed no mixing in the lepton mass matrix, so that the $\bar\mu e$ couplings are identical to the $\bar s d$ couplings in Eq. (2.11). By adding the above ratios and using Eq. (3.11) we obtain

$$\frac{\Gamma(K_L\to\mu e)+\Gamma(K_S\to\mu e)}{\Gamma(K^+\to\mu\nu)}>1.3\times10^{-18}. \quad (3.19)$$

While this is an extremely small number, its significance lies in the fact that Eq. (3.19) is a lower bound.

We consider now the case of maximal CP nonconservation, i.e., $\delta=\pi/4$. From Eq. (3.7) we then obtain

$$M^2(\tilde W_6)[M^2(\tilde W_6)+M^2(X_+)]=5\times10^9M^2(X_+)\,, \quad (3.20)$$

where we have used the mass sum rule of Eq. (2.9). Now in paper I we obtained a lower limit of $14M_W$ for $M(X_+)$ using data from the $K^+\to\mu\nu_\mu$ and $K^+\to e\nu_e$ decays. As we have shown in paper I the mass of X_+ determines the rate of interesting processes such as $\mu+N\to e+N$. These processes are maximal for a light X_+. Should X_+ be lighter than $\tilde W_6$ and $\tilde W_7$, we find from Eq. (3.20) that

$$M(\tilde W_6),\, M(\tilde W_7)>10^3M_W. \quad (3.21)$$

This implies from Eq. (3.17) that the branching ratio for $K_L\to\mu e$ is less than 10^{-10}, an order of magnitude below the present bound.

We have also investigated the contribution to the neutron dipole moment that can arise from Eq. (2.11). All CP-nonconserving phases disappear in the one-loop graphs involving any of the heavy gauge bosons. There are effects at the two-loop level, which we do not calculate. However, these contributions are many orders of magnitude smaller than those of milliweak theories; thus the observation of the dipole moment at the level of 10^{-24} -10^{-25} e cm, would definitely rule out the mechanism proposed here.

IV. *CP* NONCONSERVATION AND THE CABIBBO ANGLE

Having now discussed the main phenomenological applications of our model, we would like to make a short theoretical remark on the nature of the CP nonconservation mechanism we have employed. Our discussion of Eq. (2.11) followed by identifying the states K_1 and K_2 as

$$K_1=\bar d s+\bar s d\,, \quad (4.1)$$

$$K_2=i(\bar s d-\bar d s)\,. \quad (4.2)$$

These states then mix through $\tilde W_6$ and $\tilde W_7$ to give K_S and K_L. Suppose, however, that we define a new state

$$\hat s=e^{-i\delta}s\,, \quad (4.3)$$

which is just a change of phase. There is no *a priori* reason why we should not define K_1 and K_2 in terms of the quarks d and $\hat s$. This is not a change in the physics, and yet now $\tilde W_6$ would only couple to K_1, and $\tilde W_7$ only to K_2, and thus apparently remove the CP nonconservation effect we have analyzed. We note, however, that now there is a complex phase in the coupling of W_+ given by the term

$$W_+(\bar u'd+\bar c'\hat s e^{i\delta})\,. \quad (4.4)$$

Because of the Cabibbo angle we cannot absorb this phase in c' without a simultaneous change of phase on u', since the change in phase on the u and c quarks is the same. Then a phase change in u' would require a phase change in d, recovering our original interaction. Thus the phase cannot be removed by redefinition of the quark fields. There is therefore still CP violation in the theory, and the calculation of the preceding section remains valid.

Note that if the Cabibbo angle were zero, the redefinition of the phase of c would not affect the u field. We then obtain the surprising result that in this theory there cannot be any CP nonconservation without Cabibbo mixing.[6] This point and its generalization will be discussed in more detail in paper III.

V. CONCLUSIONS

We have shown how the simultaneous requirement of $\langle\psi_6\rangle\neq0$ and $\langle\psi_7\rangle\neq0$ breaks the CP invariance of the vacuum and leads to a theory of CP violation. The CP-violating interactions are all mediated by the superheavy gauge bosons of the model. Further, the dominant effect occurs in the K_1-K_2 system because the direct contribution $K_2\to K_1$ is enhanced by the small mass difference of K_L and K_S. All other effects turn out to be very small. Thus if the experimental limit on the neutron dipole moment were pushed below a level of, say, 10^{-25} e cm, most of the milliweak theories would be eliminated while ours would still be a viable theory.

Finally, we remark that the breaking of CP invariance is natural in our model because SU(4) [unlike SU(2)] possesses explicit CP-odd operators in its adjoint representations. In paper III we shall show how the phenomenological angle δ can actually arise through minimizing the complete SU(4) × U(1) invariant Higgs potential.

ACKNOWLEDGMENTS

This research was supported in part by the U. S. Energy Research and Development Administration under Contract No. AT(45-1)-2230.

[1] N. G. Deshpande, R. C. Hwa, and P. D. Mannheim, preceding paper, Phys. Rev. D 19, 2686 (1979), hereafter referred to as I; following paper, *ibid.* 19, 2708 (1979), hereafter referred to as III.

[2] S. Weinberg, Phys. Rev. Lett. 19, 1264 (1967); Phys. Rev. D 5, 1412 (1972); A. Salam, in *Elementary Particle Theory: Relativistic Groups and Analyticity (Nobel Symposium No. 8)*, edited by N. Svartholm (Almqvist and Wiksell, Stockholm, 1968), p. 367.

[3] R. N. Mohapatra, Phys. Rev. D 6, 2023 (1972); other theories are reviewed in R. N. Mohapatra, *Particles and Fields—1974*, proceedings of the 1974 Williamsburg Meeting of the Division of Particles and Fields of the American Physical Society, edited by C. E. Carlson, (N. Y., 1974), p. 127.

[4] S. Weinberg, Phys. Rev. Lett. 37, 657 (1976); N. G. Deshpande and E. Ma, Phys. Rev. D 16, 1513 (1977); 16, 1583 (1977).

[5] W. B. Dress, P. D. Miller, J. M. Pendlebury, P. Perrin, and N. F. Ramsey, Phys. Rev. D 15, 9 (1977).

[6] Of course, without a Cabibbo angle we would not be able to see the main CP-violating transition $K_L \to 2\pi$. Our point here is that without a Cabibbo angle there would be no CP violation anywhere in the theory.