New origin of CP violation in $SU(2) \times U(1)$ gauge models

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It is proposed that CP violation originates in a Yukawa interaction of a single, charged scalar meson, not a Higgs particle, with bilinear forms which contain a light quark and a heavy quark. These Yukawa couplings are the only source of CP violation and the only direct interaction between light quarks and heavy quarks. Detailed results are worked out for a local $SU(2) \times U(1)$ gauge model with spontaneous parity violation. It is required that $\Delta S = 2$ CP-conserving processes are dominated by W exchange and $\Delta S = 2$ CP-violating process are dominated by exchange of the charged scalar meson, called χ. The $\Delta S = 1$ CP violations and electric dipole of moments of quarks are shown to be extremely small for any values of the CP-violating phases of the χ Yukawa coupling constants. Quasistable, heavy hadrons are a special feature of the model presented, and their properties are summarized. Distinctive features of such hadrons are long lifetimes (greater than 10^{-9} sec), absence of nonleptonic decay modes with only light hadrons, and the likelihood of large CP-violation effects as compared to those of the K-meson system.

I. INTRODUCTION

The steadily growing literature on gauge-theory explanations of CP violation can be loosely categorized as either gauge-boson-mediated[1] or Higgs-particle-mediated.[2] We pursue a different line in this paper and introduce new Yukawa interactions between the usual light quarks, new heavy quarks, and a single, charged scalar boson, which we call χ, not a Higgs particle. Only these new interactions produce vertices with both light quarks[3] and heavy quarks. Their Yukawa coupling constants are allowed to be complex and the phases cannot all be absorbed by redefinition of quark fields. On the other hand, all Higgs-particle Yukawa couplings, gauge-boson couplings, and the couplings in the potential of the spin-zero fields can be chosen real and CP-conserving by proper phase redefinitions. Thus, except in the interactions of the charged, scalar boson χ, all interactions are CP-conserving because it is impossible to write down a renormalizable Lagrangian which violates CP. This condition is enforced by choice of the representation content of the fermions and Higgs fields[4] and by extra symmetries, among them parity conservation and a global chiral U(1) symmetry which can be used to ensure that neutrinos are massless in the tree approximation, imposed on the Lagrangian.

The spontaneous breaking of parity has an important implication for the problem of naturally suppressing strong P and T noninvariance in the presence of instantons. This problem develops because the instanton contributions to the vacuum-to-vacuum transition amplitude give rise to a term in the effective Lagrangian of the form[5]

$$\mathcal{L}_{\text{eff}}(\theta_{\text{QCD}}) = \frac{\theta_{\text{QCD}}}{64\pi^2} g_{\text{QCD}}{}^2 \epsilon^{\mu\nu\lambda\sigma} \operatorname{Tr}(G_{\mu\nu} G_{\lambda\sigma}) ,$$

where G is the covariant curl of the color-gluon field matrix. This manifestly P- and T-violating term, though formally a total divergence, cannot be neglected because the surface terms in the action integral do not vanish. For a theory where P is broken spontaneously, one imposes $\theta_{\text{QCD}} = 0$ as a symmetry requirement. The global, chiral U(1) symmetry mentioned above is sufficient to diagonalize the tree-approximation mass matrix, including counterterms. The corresponding axial-vector current is anomaly-free and, therefore, the tree-level mass diagonalization does not shift the value $\theta_{\text{QCD}} = 0$. Then those radiative corrections which introduce phases which are not removable by the U(1) chiral transformation are finite, and in our model the lowest order to which they could appear is the sixth order.

If the mass of the charged scalar χ is larger than the W-boson mass, we can expect two consequences. One consequence will be a suppression of CP-violating processes in the interactions among hadrons composed of light quarks, regardless of the size of Yukawa-coupling phases. This is because a virtual χ exchange and a virtual heavy-quark exchange must be involved. The experimental documentation of this suppression and its isolation in the neutral-kaon system is the cause of the steadily growing literature referred to in the opening sentence. The other consequence will be a stability, or quasistability, of all of those hadrons which contain a heavy quark.[6] The differences between such stable hadrons, whose mass spectrum should start at about half the Υ mass,[7] and hadrons which contain the usual t and b quarks, which have gauge-boson interactions with light quarks, are quite dramatic[6] and easy to distinguish experimentally.

Our gauge-theory assumptions will be the following: (1) SU(2)×U(1) is the local gauge group of the weak and electromagnetic interactions, and color SU(3) is the local gauge theory of the strong interactions. (2) Parity is spontaneously broken. (3) Neutral-gauge-boson and neutral-Higgs-particle couplings preserve all flavors naturally.[8] (4) Other than the new interactions, it is impossible to write down a renormalizable Lagrangian which violates CP. This set of assumptions puts our work into the category of theories with natural suppression of CP violation.

In the following section we outline a specific model which implements our ideas. We then turn to a computation of the $\Delta S=2$ CP-violating effective Langrangian and the expression for the parameter ϵ in K decays in Sec. III. In Sec. IV we show that $\Delta S=1$ CP violation is guaranteed to be negligible if the χ interactions produce the $\Delta S=2$ effects. An upper bound on the electric dipole moment of the d quark is computed and found to have the astonishingly small value $\simeq 10^{-35}$ e cm,[9] ten orders of magnitude smaller than the present experimental bound for the neutron. Some general features of the quasistable hadrons which are predicted in our scheme are discussed in Sec. V, and concluding remarks are made in Sec. VI.

II. CP-VIOLATING QUARK YUKAWA INTERACTIONS

Let us develop the idea that CP violation is due to the interaction between light quarks and heavy quarks by means of a single, charged spin-zero field by outlining a specific SU(2)×U(1) gauge model and pursuing its implications. Parity is assumed to be broken spontaneously, and we have two Higgs-particle doublets ϕ_L and ϕ_R which transform into each other under parity, and which carry a global phase transformation $\phi_{L,R} \to e^{\mp i\delta}\phi_{L,R}$ in addition to the local gauge transformations. We need at least two doublets and four singlets of quarks in order to ensure that all quarks gain a mass by vacuum symmetry breaking and that the usual $\Delta S=1$, neutral-current suppression mechanism occurs.[10]

The Yukawa couplings are written in terms of the doublets

$$D^u = \begin{pmatrix} u_L + u_R' \\ d_L(\theta_C) + d_R'(\theta_C) \end{pmatrix},$$

$$D^c = \begin{pmatrix} c_L + c_R' \\ s_L(\theta_C) + s_R'(\theta_C) \end{pmatrix},$$

where $d(\theta_C)$ and $s(\theta_C)$ are the standard Cabbibo-mixed d and s quark combinations. Singlets are designated by $u=u_R+u_L'$, $c=c_R+c_L'$, $d=d_R+d_L'$, and $s=s_R+s_L'$. The primes denote heavy quarks. The Yukawa couplings, whose neutral-Higgs-particle interactions naturally conserve flavor,[11] are expressed as

$$\mathcal{L}_{\text{Yukawa}} = \frac{m_u}{\sigma_L}\,\bar{u}_R(\phi_L^c)^\dagger D_L^u + \frac{m_c}{\sigma_L}\,\bar{c}_R(\phi_L^c)^\dagger D_L^c + \frac{(m_s \sin\theta_C\,\bar{s}_R + m_d\cos\theta_C\,\bar{d}_R)}{\sigma_L}\phi_L^\dagger D_L^u$$
$$+(m_s\cos\theta_C\,\bar{s}_R - m_d\sin\theta_C\,\bar{d}_R)\phi_L^\dagger D_L^c + \text{H.c.} + (L \leftrightarrow R,\, q \to q',\, m_q \to M_{q'}), \qquad (1)$$

where ϕ_L^c are the charge-conjugate isospinors and where

$$\begin{pmatrix} 0 \\ \sigma_L \end{pmatrix} = \langle \phi_L \rangle$$

and

$$\begin{pmatrix} 0 \\ \sigma_R \end{pmatrix} = \langle \phi_R \rangle .$$

The physical charged Higgs boson will be designated by $H^{(+)}$.

The gauge-boson couplings of the light (unprimed) quarks are those of the standard model,[12] while those of the heavy (primed) quarks have the sign of γ_5 reversed with respect to the standard ones. Neutral-Higgs-particle, flavor-changing interactions are suppressed naturally in both light- and heavy-quark sectors. The identification of Υ (Ref. 13) as, for example, a $\bar{u}'u$ composite requires that $\sigma_L/\sigma_R = M_u/M_{u'} \simeq M_\rho/M_\Upsilon = 0.075$, where the constituent quark picture[14] indicates that the masses of the quarks can be considered to be roughly half of the mass of the corresponding $q\bar{q}$ vector-meson state. As written in Eq. (1), the Lagrangian possesses a global γ_5 invariance under the transformations

$$\phi_{L,R} \to e^{\mp i\delta}\phi_{L,R}, \quad D^{u,c} \to e^{i\alpha\gamma_5}D^{u,c},$$
$$q_{Q=2/3} \to e^{-i(\alpha+\delta)\gamma_5} q_{Q=2/3}, \qquad (2)$$
$$q_{Q=-1/3} \to e^{-i(\alpha-\delta)\gamma_5} q_{Q=-1/3},$$

where q are quark singlets and Q is expressed in positron charge units. The γ_5 invariance prevents mixing between the light and heavy quarks. The Lagrangian, including the Higgs potential, is naturally CP conserving. The Yukawa couplings in Eq. (1), for example, can always be rendered real by proper phase choice of quark fields. The global-U(1)-invariant Higgs-particle couplings are real by Hermiticity, and bilinear terms which might be added to break this symmetry[15] can be made real by redefinition of ϕ_L and ϕ_R. The gauge-boson couplings, as in the standard model, are also naturally CP conserving. At this stage, there is no term in the Lagrangian density which couples heavy quarks and light quarks, nor is there CP violation. The simplest way which we can think of to include CP violation is by coupling heavy-quark singlets to light-quark singlets by means of a single, charged, scalar field with arbitrary, complex Yukawa couplings. The interaction terms are as follows:

$$\mathcal{L}_{\chi q q'} = g_{ud}(\bar{u}_R d'_L + \bar{u}'_L d_R)\chi + g_{us}(\bar{u}_R s'_L + \bar{u}'_L s_R)\chi + g_{cd}(\bar{c}_R d'_L + \bar{c}'_L d_R)\chi + g_{cs}(\bar{c}_R s'_L + \bar{c}'_L s_R)\chi + \text{H.c.} \tag{3}$$

The chiral phase of χ is zero, and the phase α in Eq. (2) must be zero in order to admit the terms in Eq. (3). However, the combination of Eq. (1) and Eq. (3) is invariant under a discrete symmetry where

$$\chi \to -\chi, \quad u \to \exp\left(-i\,\frac{\pi}{2}\,\gamma_5\right)u, \quad c \to \exp\left(-i\,\frac{\pi}{2}\,\gamma_5\right)c,$$

$$d \to \exp\left(-i\,\frac{\pi}{2}\,\gamma_5\right)d, \quad s \to \exp\left(-i\,\frac{\pi}{2}\,\gamma_5\right)s,$$

$$D^{uc} \to \exp\left(i\,\frac{\pi}{2}\,\gamma_5\right)D^{u,c}, \quad \phi_{L,R} \to \phi_{L,R}\,.$$

Mass terms of the form $\bar{D}D$ do not respect this discrete symmetry and should not be included. Potential-energy terms of the form $\chi\phi_L^\dagger\phi_R^c$ are likewise prohibited by this discrete symmetry. This symmetry ensures that such interactions do not arise in higher order. Consequently, a selection rule prohibits decay of a state which contains a heavy-quark composite into one which contains nothing but light-quark composites.

The arbitrariness in phase of the quark fields has been completely eliminated by making the Yukawa couplings in Eq. (1) all real. We are not, then, at liberty to trivially redefine all of the χ couplings in Eq. (3) to be real. One phase could be eliminated by redefinition of χ, but this allows all the *relative* phases among g_{ud}, g_{us}, g_{cd}, and g_{cs} to be nonzero. All of the CP violation resides in the χ interactions, which introduce these CP-violating phases.

III. $\Delta S=2$ EFFECTIVE LAGRANGIAN AND CP VIOLATION

The relative phases among the coupling coefficients in Eq. (2) can lead to CP violation, and we limit the mass of the χ and its couplings to quarks, g_{ud} etc., by the requirement that $\Delta S=2$ CP violation be accounted for by fourth-order χ-exchange diagrams,[16] Fig. 1. The effective Lagrangian, to leading order in $M_{q'}{}^2/M_\chi{}^2$, is

$$\mathcal{L}_{\text{eff}}^{\Delta S=2}(\chi\text{ exchange}) = \frac{i(g_{ud}\bar{g}_{us} + g_{cd}\bar{g}_{cs})^2}{64\pi^2 M_\chi{}^2} \times \bar{u}_s\gamma_\mu\left(\frac{1+\gamma_5}{2}\right)u_d\,\bar{u}_s\gamma^\mu\left(\frac{1+\gamma_5}{2}\right)u_d\,. \tag{4}$$

The CP-violation parameter ϵ can be evaluated in the free-quark approximation from the expression[17]

$$|\epsilon| = 2^{-3/2}\left|\frac{\langle d\bar{s}|S_-|s\bar{d}\rangle}{\langle d\bar{s}|S_+|s\bar{d}\rangle}\right|,$$

where the S-matrix components S_- and S_+ are odd and even, respectively, under the time-reversal operation. We can now use Eq. (4) and evaluate the imaginary part of $(g_{ud}\bar{g}_{us} + g_{dc}\bar{g}_{cs})^2$ by the equation

$$|\epsilon| = 2\times10^{-3} = \frac{1}{2\sqrt{2}}\;\frac{\dfrac{\text{Im}(g_{ud}\bar{g}_{us} + g_{dc}\bar{g}_{cs})^2}{32\pi^2 M_\chi{}^2}}{\dfrac{G_F}{\sqrt{2}}\,\dfrac{\alpha}{4\pi}\,\dfrac{m_c{}^2}{M_W{}^2}\,\dfrac{\cos^2\theta_C\sin^2\theta_C}{\sin^2\theta_W}}\,, \tag{5}$$

where the denominator is the W-boson contribution to the $\Delta S=2$ neutral kaon mass difference, assumed to be the dominant mechanism,[16,18] and ϵ is the CP-violation parameter. In order to get an idea of the magnitude of $\text{Im}(g_{ud}\bar{g}_{us} + g_{dc}\bar{g}_{cs})^2$ one can use $M_\chi \simeq M_{H^+} \simeq 10M_W$ which is obtained by requiring that $\Delta S=2$ CP-conserving amplitudes be determined by W-meson exchange, and one finds that $\text{Im}(g_{ud}\bar{g}_{us} + g_{cd}\bar{g}_{cs})^2 \simeq 2\times10^{-7}$. This is the same order of magnitude as that of the fourth power of the (real) mass-determining Higgs coupling coefficients in expression (1). As we show later, the above constraint on $\text{Im}(g_{ud}\bar{g}_{us} + g_{cd}\bar{g}_{cs})^2$ leads to extremely small $\Delta S=1$ neutral-current CP violation and electric dipole moments of the quarks. We can, therefore, say that our model has natural suppression of CP violation in the sense advocated by B. W. Lee.[4]

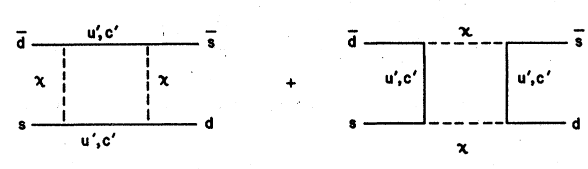

FIG. 1. Box diagrams with χ exchange which produce the $\Delta S=2$ CP violation.

IV. ESTIMATES OF CP VIOLATION IN $\Delta S=1$ NEUTRAL CURRENT AND ELECTRIC DIPOLE MOMENT OF QUARK

CP violation can arise in $\Delta S=1$ processes (or flavor-changing processes in general) by the finite second-order corrections to the dsZ and $ds\gamma$ vertices as are shown in Fig. 2. Neutral-Higgs-particle $\Delta S=1$ effects turn out to be negligible compared to Z and γ effects.

In these graphs, as in those relevant to the electric dipole moments of the quarks, one notes that the CP-violating phase originates from the relative phases between χ coupling constants defined in Eq. (3). In a given graph, there must be present two different quark flavors of the same charge in order to produce a CP-violating phase. An example of a class of flavor-preserving graphs which do *not* produce CP-violating phases is shown in Fig. 3.

The contributions to the effective γds vertex and the Zds vertex from the graphs shown in Fig. 2 will produce CP-violating effects. The effective photon vertex is composed of a transition magnetic moment term

$$\bar{s}(p)\left(\frac{i\sigma_{\mu\nu}k^{\nu}}{m_s+m_d}\right)M_{ds}d(q), \quad k=p-q$$

and a transition charge radius,

$$\bar{s}(p)\gamma^{\nu}\left(g_{\mu\nu}k^2-k_{\mu}k_{\nu}\right)\frac{\langle r^2\rangle_{ds}}{6}\left(\frac{1-\gamma_5}{2}\right)d(q)\,.$$

The CP-violating χ contributions to M_{ds} and $\frac{1}{6}\langle r^2\rangle_{ds}$ from calculations of graphs of Fig. 2 in 't Hooft-Feynman gauge are

$$M_{ds}\simeq e\,\frac{g_{sc}\bar{g}_{cd}}{32\pi^2}\,\frac{(m_s+m_d)^2}{M_\chi^{\,2}}\times\left(1+\gamma_5\,\frac{m_s-m_d}{m_s+m_d}\right)\left(\frac{5}{36}+2\ln\frac{M_\chi^{\,2}}{M_{c'}^{\,2}}\right) \qquad (6)$$

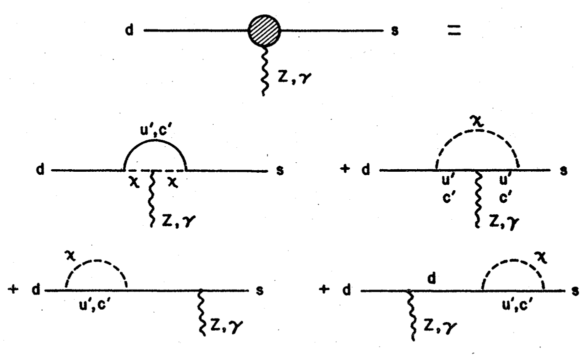

FIG. 2. Diagrams with internal χ lines which produce CP violation in the effective dsZ and $ds\gamma$ vertices.

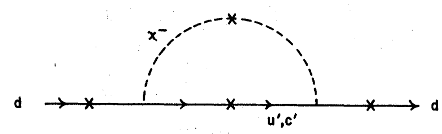

FIG. 3. Illustration of graphs in which no CP-violating phase can appear. The crosses represent flavor preserving vertex insertions.

and

$$\frac{\langle r^2\rangle_{ds}}{6}\simeq e\,\frac{g_{sc}\bar{g}_{cd}}{16\pi^2}\,\frac{1}{M_\chi^{\,2}}\left(\frac{1}{6}-2\ln\frac{M_\chi^{\,2}}{M_{c'}^{\,2}}\right), \qquad (7)$$

respectively. The CP-violating $\Delta S=1$ effects which arise from the matrix elements in Eqs. (6) and (7), denoted $\mathfrak{M}$, with a photon exchange of momentum k will be typically of a strength

$$\frac{e}{k^2}\,\mathfrak{M}_{sd\gamma}\simeq e^2\,\frac{g_{cd}\bar{g}_{cs}}{16\pi^2}\,\frac{1}{M_\chi^{\,2}}\ln\frac{M_\chi^{\,2}}{M_{c'}^{\,2}} \simeq 4\times10^{-8}G_F\,, \qquad (8)$$

where the $M_{c'}$, M_χ, and g_{cd} values are those suggested following Eq. (5).[19]

For the Zds effective vertex, it turns out that the dominant contribution to $\Delta S=1$ CP violation can be evaluated from the coupling of the weak-isospin-triplet component of Z in the second graph of Fig. 2.[20] The resulting "transition charge" vertex is

$$\bar{s}(p)\gamma_\mu\,\frac{(1+\gamma_5)}{2}\,d(q)\left[g\bar{g}_{cd}g_{cs}\,\frac{M_{c'}^{\,2}}{M_\chi^{\,2}}\left(-1+\ln\frac{M_\chi^{\,2}}{M_{c'}^{\,2}}\right)\right]. \qquad (9)$$

The CP-violating $\Delta S=1$ effects which come from the matrix element in Eq. (9) when a Z is exchanged will be of strength

$$\frac{g}{M_z^{\,2}}\,\mathfrak{M}_{sdz}\simeq\frac{g^2}{M_z^{\,2}}\,\frac{g_{cd}\bar{g}_{cs}}{16\pi^2}\,\frac{M_{c'}^{\,2}}{M_\chi^{\,2}}\ln\frac{M_\chi^{\,2}}{M_{c'}^{\,2}} \simeq 9\times10^{-8}G_F\,. \qquad (10)$$

From Eqs. (8) and (10) we see that the criterion for natural suppression of $\Delta S=1$ CP-violating processes is satisfied, since $\Delta S=1$ CP violation is negligible for any value of the CP-violating phases in the model.

Finally, we turn to the estimate of the electric dipole moments of the quarks, which tell us approximately what to expect for the order of magnitude of the electric dipole moment of the neutron.

As we mentioned earlier, there must be a change in flavor for CP-violating phases to appear, and the electric-dipole moment appears only in the *fourth-order* correction to the photon vertex. There are three classes of graphs which contribute the electric dipole moment, and they are shown in Figs. 4(a), 4(b), and 4(c) for the d quark. These leading contributions to the electric dipole moment are so small that we can simply put upper bounds on the loop integrations which indicate the order of magnitude one might expect.[21] Our point about the distinctive smallness of the electric dipole moment can be amply made in this way. Taking one set of graphs at a time, we can summarize our results as follows: The electric dipole moment of a quark, D_E (the d quark is chosen for illustration) is identified as the $q^2=0$ coefficient of the CP-violating form $\bar{d}(p')(\sigma_{\lambda\nu}\gamma_5)d(p)q^\nu$ which arises from each set of graphs separately.

$$|D_E(a)| < \left| \frac{7}{3} \frac{g^2 \sin\theta_C \cos\theta_C}{(4\pi)^4} \operatorname{Im}(\bar{g}_{ud} g_{uc}) \frac{m_c^3 m_u m_d}{M_W^4 M_\chi^2} \left(2 \ln \frac{M_W^2}{m_c^2} - 3 + \ln \frac{M_\chi^2}{m_u m_c}\right) \right| = 1.7\times 10^{-35}\, e\,\text{cm}\,, \quad \text{(11a)}$$

$$|D_E(b)| < \left| \frac{1}{18} \frac{\operatorname{Im}(\bar{g}_{ud} g_{cs} g_{us} \bar{g}_{cd})}{(4\pi)^4} \frac{m_d M_{c'}^2}{M_\chi^4} \ln \frac{M_\chi^2}{M_{c'}^2} \right| = 3.2\times 10^{-35}\, e\,\text{cm}\,, \quad \text{(11b)}$$

$$|D_E(c)| < \left| \frac{1}{9} \frac{M_{c'}^3 M_{s'}^2 M_{u'} m_d}{M_\chi^4 M_W^4} \operatorname{Im}(\bar{g}_{ud} g_{dc}) g^2 \cos\theta_C \sin\theta_C \left(\ln \frac{M_W^2}{M_{c'}^2} - 2\right) \right| = 7\times 10^{-38}\, e\,\text{cm} \ll D_E(a,b)\,. \quad \text{(11c)}$$

The graphs (4c) are much more convergent than the others and, as indicated by our bound in Eq. (11c) which was obtained from the first of graphs of Fig. 4(c), their contributions are typically two orders of magnitude smaller than the others, even though it is heavy-quark masses which are involved in the Glashow-Iliopoulos-Maiani (GIM)[10,12] suppression. The scalar exchange mechanism has the curious (though essentially untestable) feature that the dipole moments of baryons are predicted to be five orders of magnitude smaller than even the superweak or microweak predictions, which hover about 10^{-30} e cm. The present experimental upper bound for the neutron is[8]

$$(D_E)_{\text{expt}} < 4\times 10^{-25}\, e\,\text{cm}\,.$$

Contributions to the electric dipole moments of the quarks which arise from complex phases in the mass matrix could enter at most in sixth order,[22] and we expect the above estimates to be the dominant ones.

V. FEATURES OF THE STABLE, HEAVY HADRONS

Guided by the above considerations on $\Delta S=2$ processes, we assume that the mass of the χ field is large ($\simeq 10 M_W$). There is no decay $q' \to q\chi$ in this case. Since none of the interactions discussed so far admits a decay of heavy quarks into lighter ones, and because there is no mixing between the χ field and the charged member of the Higgs doublet (the $\phi_L^\dagger \phi_R^c \chi$ interaction is forbidden by the discrete symmetry $\chi \to -\chi$ and $\phi_{L,R} \to \phi_{L,R}$), the lightest hadron state which contains a primed quark can only decay if we include χ-lepton couplings.[23] In the latter situation, hadron stability depends upon the masses of lepton pairs which interact with χ. In any case, lifetimes of the lightest of these hadrons will be much longer than charmed hadrons and clearly distinguishable from hadrons composed of t or b quarks,[24] sequential analogs of u, c and d, s,

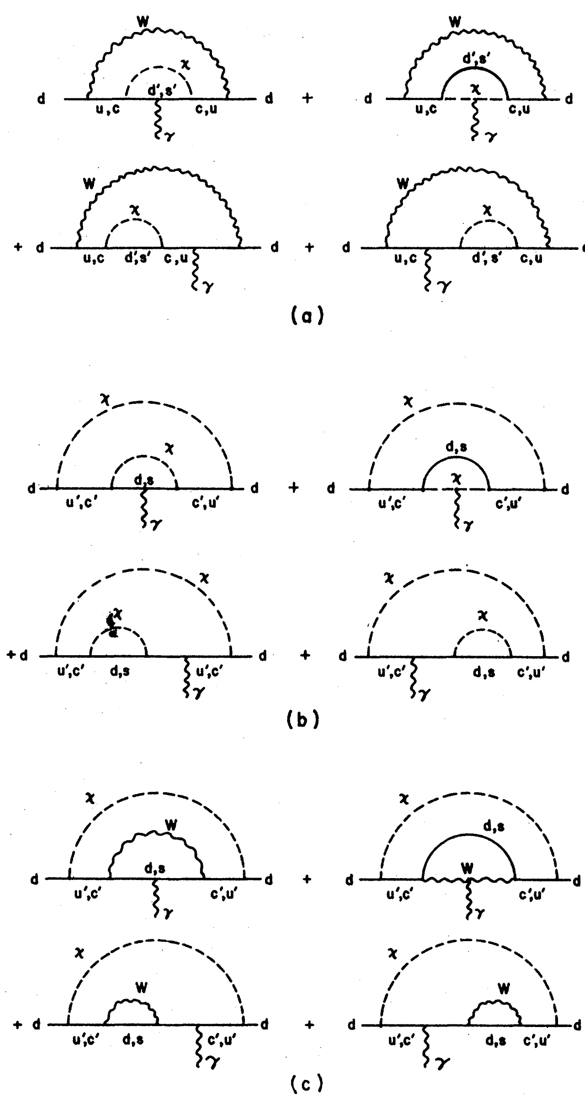

FIG. 4. (a)–(c) The graphs which produce fourth-order contributions to the electric-dipole moment of the d quark.

respectively. A new feature of our picture of CP violation which is also independent of the details of the model is the possibility of large CP-violating mixing between neutral mesons which have a heavy, primed-quark constituent. These latter mixings arise from the graphs of the type shown in Fig. 5. The CP-violating parts of these processes arise from the relative phases between χ couplings to different quarks. These phases are not necessarily small, and the CP-violating amplitudes could be of the same order of magnitude as the CP-conserving ones.

If χ is *not* coupled to leptons, we expect a stable meson, $M(qq')$, probably neutral since $M_{u'} < M_d'$ in tree approximation, at about 5 GeV. A meson with one unit difference in charge and about a 10 MeV $\simeq M_d - M_u$ heavier mass would decay slowly ($\tau \simeq 1$ sec, as estimated from naive extension of the neutron β-decay calculation) into the stable meson plus $e+\nu_e$.

A more interesting situation is the one where χ couples to leptons. Examples are not difficult to construct,[15] and one can arrange to couple χ with light lepton pairs, heavy-light pairs, and heavy pairs. An example of the first situation is one in which e shares a neutrino with μ or τ in the Konopinski-Mahmoud scheme.[25,15] A doublet-doublet χ interaction gives rise to a term in the Lagrangian

$$g_{\chi e\nu}\bar{\nu}\left(\frac{1+\gamma_5}{2}\right)e\chi + g_{\chi\mu\nu}\bar{\nu}\left(\frac{1-\gamma_5}{2}\right)\mu\chi + \text{H.c.}$$

If the couplings satisfy $g_{\chi e\nu} \simeq g_{ud} \simeq 2\times 10^{-2}$, then a free-quark decay estimate yields

$$\frac{\tau(q' \to q+e+\nu)}{\tau(\mu \to e+\nu+\bar{\nu})} = \frac{M_\mu{}^5}{M_q{}^5}\left(\frac{g}{g_{\chi e\nu}}\right)^4\left(\frac{M_\chi}{M_W}\right)^4 \simeq 25 \qquad (12)$$

or $\tau(q') \simeq 10^{-5}$ sec which can be used as a rough indication of the lifetime of the neutral low-mass quasistable mesons. The charged mesons can decay purely leptonically, and the lifetime can be written

$$\Gamma(M' \to l+\nu) = \frac{M'}{\pi}\,\frac{|g_{e\nu\chi}g_{qq'\chi}|^2}{M_\chi{}^4}\,|f'|^2\,, \qquad (13)$$

where M' on the right-hand side of Eq. (13) stands for the mass of a charged quasistable meson and

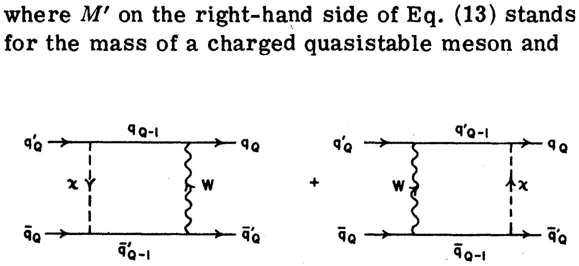

FIG. 5. Graphs which can give rise to mixing between neutral, heavy mesons. The real and imaginary parts of the mass matrix can be comparable in general.

f' is the counterpart of the pion decay constant.[26] With the values $M' \simeq 5$ GeV, $g_{e\nu\chi} \simeq g_{qq'\chi} \simeq 0.02$, $M_\chi \simeq 10 M_W$, and $M_q M_{q'} < |f'| < M'^2$, we find 10^{-9} sec $< \tau(M'^{\pm} \to 1^{\pm}+\nu) < 6\times 10^{-7}$ sec. This estimate indicates the lifetimes of charged heavy mesons which one expects. These estimates of $\tau(q')$ and $\tau(M')$ are much longer than the corresponding ones for mesons composed of t and b quarks,[22] whose lifetime estimates are typically $\tau(b\bar{q}) \lesssim 10^{-13}$ sec, where q denotes a light quark. The decays which occur via χ exchange are fast enough to permit observation of the decay products when the heavy mesons are produced near threshold.

We have re-emphasized here the clear distinction between lifetimes of mesons composed of the b and t quarks which arise in straightforward generalization of the standard model and the (long) lifetimes of quarks which are prohibited by a quantum number or a selection rule from decaying nonleptonically into light quarks and prohibited from decaying semileptonically by W exchange into light quarks and leptons.[6] *The complete absence of nonleptonic decays into known mesons should provide, with the lifetimes, a distinctive signature for the presence of such stable hadrons.* We will present a detailed study of lifetime, branching ratios, and CP-violating effects of such states in a separate publication.

VI. SUMMARY AND CONCLUSIONS

We have presented CP-violation results in a spontaneous-parity-violating local SU(2)×U(1) gauge model of weak and electromagnetic interactions. The CP violation is due to the Yukawa interaction of a single, charged scalar meson, not a Higgs particle, with bilinear forms which contain a light quark and a heavy quark. These Yukawa couplings provide the only possible sources of CP violation in the model as well as the only direct interaction between light quarks and heavy quarks. We required that $\Delta S=2$ CP-conserving processes are dominated by the GIM-suppressed W-exchange mechanism and that the χ exchange describes the $\Delta S=2$ CP-violating processes. We then showed that the $\Delta S=1$ CP violations and the electric dipole moments of quarks are extremely small, well below experimental bounds for any values of CP-violating phases in the χ Yukawa couplings. The violations of CP occur in the quark couplings. The violations of CP occur in the quark mass matrix and in dipole-moment corrections at sixth and fourth order, respectively, because change in flavor *and* change in momentum in light-quark or heavy-quark lines must both be present in the relevant processes in order for CP violation to occur.

There is a selection rule which inhibits decay of heavy quarks into final states which contain light quarks. The primary signals that heavy mesons are composed of such quarks rather than sequential generalizations of u, c, d, and s, generally called t and b, are (1) the lifetimes should be much longer ($10^{-5}-10^{-9}$ sec at least as compared to $\sim 10^{-13}$ for the latter case) and (2) the purely nonleptonic decay into normal hadrons is prohibited. It is likely that CP-violation effects in such stable, heavy hadrons would be much larger than in the K^0 system.

We argued that the strong CP-violation effects due to instanton contributions to the effective action can be naturally suppressed, and we feel that the dipole-moment calculations correctly represent the dominant CP violation in flavor-preserving amplitudes.

In conclusion, we believe that the new χ interaction which we have introduced in order to incorporate CP violation in an SU(2)×U(1) gauge scheme gives a simple and realistic account of known weak-interaction phenomenology. Calculations are relatively straightforward, the Higgs sector is the simplest required for spontaneous parity violation and effects of the charged Higgs particle are readily assessed, and the χ-related phenomena which involve heavy, stable hadrons are distinctive and testable. It is also plausible that strong CP violation is suppressed to the required degree. The χ particle has no role in gauge- or Higgs-boson mechanisms within the SU(2)×U(1) framework, but it might be viewed as a Higgs remnant of a larger, unified gauge theory.

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[11]The general features of the present paper which we wish to stress—spontaneous P-violation, the new CP-violating interaction, quasistable heavy hadrons, and natural suppression of CP violation—do not require some of the special features which occur in the models used for illustration. Among these special features are (1) flavor-conserving neutral-Higgs-particle interactions, (2) $m_q/m_{q'}=\epsilon$ and its implication that the charge $\frac{2}{3}$ u' quark should be lighter than the rest of the primed quarks, (3) vector coupling of the neutral current to the electron [D. W. McKay and H. J. Munczek, Phys. Rev. D 19, 985 (1979)]. For example, consider a quark Yukawa Lagrangian

$$\begin{aligned}\mathcal{L}_Y=&(m_1/\sigma_R)(\bar d_R^{(1)}\phi_L^\dagger D_L^{(1)}+\bar d_L^{(1)}\phi_R^\dagger D_R^{(1)})\\&+(m_2/\sigma_R)(\bar d_R^{(2)}\phi_L^\dagger D_L^{(2)}+\bar d_L^{(2)}\phi_R^\dagger D_R^{(2)})\\&+(m_3/\sigma_R)(d_R^{(1)}\phi_R^\dagger D_L^{(2)}+d_L^{(1)}\phi_L^\dagger D_L^{(2)})\\&+\text{terms } d^{(i)}\to u^{(i)} \text{ and } \phi_{L,R}\to\phi_{L,R}^{(c)},\end{aligned}$$

where d's denote singlet charge $-\frac{1}{3}$ quarks and u's denote $\frac{2}{3}$ charge and $D^{(i)}$ denote doublets. The chiral phase symmetry assignments are

$$d^{(2)}\to e^{i\gamma_5\beta}d^{(2)},\quad d^{(1)}\to e^{-i\gamma_5(2\delta-\beta)}d^{(1)},$$

$$D^{(2)}\to e^{i\gamma_5\delta}D^{(2)},\quad D^{(1)}\to e^{i\gamma_5(3\delta-\beta)}D^{(1)},$$

$$u^{(2)}\to e^{-i\gamma_5(2\delta-\beta)}u^{(2)},\quad u^{(1)}\to e^{i\gamma_5(4\delta-\beta)}u^{(1)}.$$

L_y does not have natural flavor conservation of neutral-Higgs-particle couplings, though gauge-boson couplings are flavor conserving. The masses of u' and c' are unrelated to those of d' and s'; the current quark mass values for u and d can be chosen without inducing extremely large Higgs couplings; and a parallel development in the lepton sector reproduces the standard model of leptons (Ref. 12).

[12]The standard model refers to the Weinberg-Salam theory for leptons extended to include two left-handed quark doublets and four right-handed singlets. The mechanism of Glashow, Iliopoulos, and Maiani, Ref. 10, is used to suppress $\Delta S=1$ neutral-current transitions.

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[19]$\Delta S=1$ processes in which two gluons instead of a photon are exchanged between a quark line in which flavor changes and one in which flavor is preserved can be roughly estimated to be

$$\frac{eM_{sd\gamma}}{(p-q)^2}\frac{e^2}{8\pi^2}\frac{\alpha_s}{\alpha},$$

where α_s is the gluon "fine-structure constant." If $\alpha_s \approx 1$, then this two-gluon amplitude is smaller than corresponding one-photon amplitude by a factor $1/2\pi$ Clearly this is no danger.

[20]We simplify the Z-vertex calculations to this order by writing

$$J_\mu^Z(\theta_W)=-\sin\theta_W J_\mu^B+\cos\theta_W J_\mu^3 ,$$

and

$$J_\mu^\gamma(\theta_W)=\cos\theta_W J_\mu^B+\sin\theta_W J_\mu^3 .$$

Since $\partial^\mu J_\mu^\gamma(\theta_W)=0$,

$$(p_s-p_d)^\mu \langle s\,|\,J_\mu^\gamma(0)\,|\,d\rangle=0$$
$$=\bar{U}_s(p_s)(A^\gamma+B^\gamma\gamma_5)|(\not p_s-\not p_d)\,U_d(p_d)$$
$$\Rightarrow A^\gamma=0,\ B^\gamma=0.$$

Now

$$A^\gamma=0=\cos\theta_W A^B+\sin\theta_W A^3 \Rightarrow A^Z=A^3/\cos\theta_W,$$

where

$$A^Z=-\sin\theta_W A^B+\cos\theta_W A^3.$$

The "transition-charge" form-factor contribution in the matrix elements of J_μ^Z is the dominant one, and the above argument shows that only the isospin triplet component of the Z coupling need be computed. This eliminates the need to consider graphs with Z coupled to a χ line, for example.

[21]An example of the procedure which we use to place upper bounds on the two-loop integrals is the following: A finite interior-loop integral is reduced to an integral over a Feynman parameter x, to give for example

$$(I_{\rm int})_\sigma=\not k\, k_\sigma\int_0^1\frac{x\,dx}{k^2-M^{(2)}(x)},$$
$$M^2(x)\equiv\frac{M_\chi^2}{1-x}-\frac{M_{d'}^2}{x},$$

k is the outer-loop momentum and σ is a Lorentz index It is understood that leading powers of external momentum relevant to the magnetic moment have been extracted. The outer-loop integrals are exemplified by

$$I_{\rm out}=\int_0^1 x\,dx\int_0^\infty\frac{K^3dK}{(K+M_W^2)^2(K+M_u^2)^2(K+M_c^2)^2(K+M^2(x))},$$

and the Lorentz index has been shifted to an external momentum or γ matrix. Now $M^2(x)>(M_\chi^2+M_{d'}^2)\simeq M_\chi^2$ and $K^2/(K+M_u^2)^2>1$, so

$$I_{\rm out}<\int_0^\infty\frac{KdK}{(K+M_W^2)^2+(K+M_c^2)^2(K+M_\chi^2)}$$
$$\simeq\frac{2}{M_W^4M_\chi^2}\left(\ln\frac{M_\chi^2}{M_cM_W}-1\right).$$

[22]Unless internal flavor changes occur in a graph which has the same flavor of quark in and out and unless a momentum change occurs in a light-quark or heavy-quark internal-line flow, no CP-violating dependence on χ couplings can occur. These are the reasons why CP-violating corrections to the electromagnetic vertex occur only in fourth order and mass corrections only in sixth order.

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[26]As a reminder of the difference in space-time structure between the π decay and the one discussed here, we note that the χ is spin zero so that the two-lepton mode is not suppressed by helicity conservation and the decay constant has dimensions of mass squared rather than mass.