Strong and weak CP in a model with a new gauged U(1) symmetry

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We discuss a natural solution to the strong CP problem with spontaneous CP breaking in the context of $\mathrm{SU}_C(3)\times\mathrm{SU}_L(2)\times\mathrm{U}_Y(1)\times\mathrm{U}_{\mathrm{new}}(1)$ with two vectorlike quarks. The gauged $\mathrm{U}_{\mathrm{new}}(1)$ symmetry is imposed such that $\bar{\theta}$ remains zero at the tree level when CP and $\mathrm{U}_{\mathrm{new}}(1)$ are broken by singlet Higgs scalars. Flavor-changing neutral currents at the tree level are naturally suppressed because only singlet Higgs scalars are introduced. In order to agree with the experimental measurement of ϵ and ϵ'/ϵ, we find that the CP-breaking scale and the masses of the new quarks are less than 2.0 TeV and 530 GeV, respectively.

I. INTRODUCTION

The *absence* of observable CP violation in strong interactions, and its *presence* at a strength 0.3% of the weak interaction, are empirical facts that suggest physics beyond the standard model of strong and electroweak forces. In quantum chromodynamics, the theory of strong interactions, instanton effects[1] induce P- and CP-violating interactions of the form

$$L_\theta=\frac{\theta}{32\pi^2}\mathrm{Tr}G_{\mu\nu}\tilde{G}^{\mu\nu}\ , \qquad (1)$$

where $G_{\mu\nu}$ is the strong gauge field tensor and $\tilde{G}^{\mu\nu}$ is its dual. Although this term can be rewritten as a total divergence, it has nonzero effects because of the existence of nontrivial topological field configurations. This CP-violating effect induces a nonvanishing neutron electric dipole moment, which is measured[2] to be less than 1.2×10^{-25} e cm. Hence, it requires $\bar{\theta}$ to be less than 2×10^{-10},[3] where $\bar{\theta}$, which is given by

$$\bar{\theta}=\theta+\mathrm{Arg}(\det\mathcal{M}_{\mathrm{up}}+\det\mathcal{M}_{\mathrm{down}})\ , \qquad (2)$$

is the effective θ parameter in Eq. (1). $\mathcal{M}_{\mathrm{up}}$ and $\mathcal{M}_{\mathrm{down}}$ are the up- and down-quark mass matrices, respectively. The extreme smallness of $\bar{\theta}$ is *unnatural* in the standard model; this is the strong CP problem.

A *global* chiral symmetry $\mathrm{U}_{\mathrm{PQ}}(1)$ was introduced by Peccei and Quinn[4] such that $\bar{\theta}$ can be rotated away by the symmetry. A physical pseudoscalar, axion, is inevitable when the global $\mathrm{U}_{\mathrm{PQ}}(1)$ symmetry is spontaneously broken. Other classes of models which solve the strong CP problem have been discussed by several authors,[5] in particular by Nelson and Barr in grand unified theories, some years ago. These models have a vanishing $\bar{\theta}$ at the tree level after CP is broken. $\bar{\theta}$ is then arranged to be small at one-loop corrections. Recently, Frampton and Kephart[6] proposed a simple model, the aspon model, in which the gauge group is $\mathrm{SU}_C(3)\times\mathrm{SU}_L(2)\times\mathrm{U}_Y(1)\times\mathrm{U}_{\mathrm{new}}(1)$ with an additional vectorlike quark doublet and two singlet Higgs scalars transforming nontrivially under the global $\mathrm{U}_{\mathrm{new}}(1)$. Vacuum expectation values (VEV's) of the Higgs singlet are responsible for $\mathrm{U}_{\mathrm{new}}(1)$ and CP breaking. $\mathcal{M}_{\mathrm{up}}$ and $\mathcal{M}_{\mathrm{down}}$ are complex but their determinants are real at the tree level. Therefore, $\bar{\theta}$ picks up a nonzero value only through radiative corrections. In order to agree with the experimental measurements of ϵ and ϵ'/ϵ, it is shown in this paper that the spontaneous CP-breaking scale and the new quark masses are bounded by 2.0 TeV and 530 GeV, respectively.

In Sec. II, the aspon model is described. In Sec. III, flavor-changing neutral currents at the tree level are considered. In Sec. IV, a calculation of $\bar{\theta}$ is presented. In Sec. V, weak CP violation parameters ϵ and ϵ'/ϵ are discussed. Finally, in Sec. VI, there are some concluding remarks.

II. DESCRIPTION OF MODEL

The gauge group of this model which is called the aspon model is $\mathrm{SU}_C(3)\times\mathrm{SU}_L(2)\times\mathrm{U}_Y(1)\times\mathrm{U}_{\mathrm{new}}(1)$. Although a global $\mathrm{U}_{\mathrm{new}}(1)$ was discussed in Ref. 6, since $\mathrm{U}_{\mathrm{new}}(1)$ is anomaly-free it may be *gauged* as here; it makes this solution of strong CP problem seem more appealing than solutions which involve the *necessarily global anomalous* $\mathrm{U}_{\mathrm{PQ}}(1)$. The particle assignments are given by

$$q_L^T=(u,d)_L:\ (3,2,\tfrac{1}{6},0)\ , \qquad (3)$$

$$u_R:\ (3,1,\tfrac{2}{3},0)\ , \qquad (4)$$

$$d_R:\ (3,1,-\tfrac{1}{3},0)\ , \qquad (5)$$

$$l_L^T=(\nu,e)_L:\ (1,2,-\tfrac{1}{2},0)\ , \qquad (6)$$

$$e_R:\ (1,1,-1,0)\ , \qquad (7)$$

$$\Phi_L^T=(\phi^+,\phi^0):\ (1,2,\tfrac{1}{2},0)\ , \qquad (8)$$

$$Q_L^T=(U,D)_L:\ (3,2,\tfrac{1}{6},-1)\ , \qquad (9)$$

$$Q_R^T=(U,D)_R:\ (3,2,\tfrac{1}{6},-1)\ , \qquad (10)$$

$$\chi_1,\chi_2:\ (1,1,0,1)\ , \qquad (11)$$

where quarks states are not necessarily mass eigenstates. Three generations are assumed in Eqs. (3)–(7). Equations (3)–(8) correspond to the standard model and Eqs. (9)–(11) are new particles. [The particle content is not unique; U and D can be *alternatively* assigned to be SU(2) singlets.] The Yukawa interactions are given by

$$-L_Y=\bar{q}_L\mathbf{m}_d d_R\left[\frac{\sqrt{2}}{v}\Phi\right]+\bar{q}_L\mathbf{m}_u u_R\left[\frac{\sqrt{2}}{v}\tilde{\Phi}\right]+\bar{l}_L\mathbf{m}_e e_R\left[\frac{\sqrt{2}}{v}\Phi\right]+\mathbf{h}^\alpha\bar{q}_L Q_R\chi_\alpha+\text{H.c.}\ ,\quad \alpha=1,2, \tag{12}$$

where $v/\sqrt{2}$ is defined as the VEV of ϕ^0 and $\tilde{\Phi}$ as $(\bar{\phi}^0,-\phi^-)^T$. The generation indices are implicit. Usual quarks and leptons acquire their masses through spontaneous symmetry breaking (SSB) induced by the VEV of the doublet Higgs scalar. The new quarks acquire their mass through a gauge-invariant mass of the form $M\bar{Q}_L Q_R$. Hence, U and D quarks are degenerate in mass. $\mathbf{m}_d$, $\mathbf{m}_u$, $\mathbf{m}_e$, v, $\mathbf{h}^{1,2}$, and M are real by the assumption of CP invariance. The VEV's of χ_1 and χ_2 are chosen to be

$$\langle\chi_1\rangle=\frac{1}{\sqrt{2}}\kappa_1 e^{i\theta}\ \text{ and }\ \langle\chi_2\rangle=\frac{1}{\sqrt{2}}\kappa_2\ . \tag{13}$$

Hence CP is broken spontaneously. [CP can be broken softly by $i(\chi^*{}_1\chi_2-\chi^*{}_2\chi_1)$.]

The up- and down-quark mass matrices linking from right-handed sector to left-handed sector are in the form

$$\mathcal{M}_{\text{up}}=\begin{bmatrix}\mathbf{m}_u & \mathbf{F}\\ 0 & M\end{bmatrix}\ \text{ and }\ \mathcal{M}_{\text{down}}=\begin{bmatrix}\mathbf{m}_d & \mathbf{F}\\ 0 & M\end{bmatrix}, \tag{14}$$

where

$$\mathbf{F}=\mathbf{h}^1\langle\chi_1\rangle+\mathbf{h}^2\langle\chi_2\rangle\ . \tag{15}$$

The Kobayashi-Maskawa (KM) matrices will be generalized to 4×4. From the constraint $|V_{ud}|^2+|V_{us}|^2+|V_{ub}|^2=0.9979\pm0.0021$, we obtain $|F_1|/M$ and $|F_2|/M$ to be less than 10^{-2} and 10^{-1}, respectively. Although $\mathbf{F}$ is a complex column matrix, the determinants of $\mathcal{M}_{\text{up}}$ and $\mathcal{M}_{\text{down}}$ are real. All entries become complex and, therefore, a nonvanishing value of $\bar{\theta}$ arises through radiative corrections. The calculation of $\bar{\theta}$ at the 1-loop level will be done in Sec. IV. In the next section, we will discuss how the flavor-changing neutral currents in the presence of new quarks are suppressed at the tree level.

III. FLAVOR-CHANGING NEUTRAL CURRENTS (FCNC's)

After introducing the new vectorlike quark doublet, we find that there are FCNC's induced by Z coupling because of the mismatch of the new and usual quarks in the right-handed sector. Therefore, the flavor-changing Z couplings are induced by the terms

$$L_Z^{\text{FCNC}}=(-\tfrac{1}{2})\frac{g_2}{\cos\theta_W}\bar{D}_R\gamma_\mu D_R Z^\mu+(U_R\text{ contributions})\ , \tag{16}$$

where the factor $-\frac{1}{2}$ is the isospin of D_R and g_2 is the SU(2) gauge coupling constant. Let us consider the down sector first. Without losing any generality, we assume the down-quark mass matrix is in the partially diagonalized form

$$\mathcal{M}_{\text{down}}=\begin{bmatrix}m_d & 0 & 0 & F_1\\ 0 & m_s & 0 & F_2\\ 0 & 0 & m_b & F_3\\ 0 & 0 & 0 & M\end{bmatrix}. \tag{17}$$

This mass matrix can be diagonalized by a biunitarity transformation, $\mathbf{K}_L^\dagger\mathcal{M}_{\text{down}}\mathbf{K}_R$. The transformation matrices[7] are given, to second order in the x_i, by

$$\mathbf{K}_L=\begin{bmatrix}1-\frac{1}{2}|x_1|^2 & x_1x_2^*\dfrac{m_s^2}{m_d^2-m_s^2} & x_1x_3^*\dfrac{m_b^2}{m_d^2-m_b^2} & x_1\\ x_1^*x_2\dfrac{m_d^2}{m_s^2-m_d^2} & 1-\frac{1}{2}|x_2|^2 & x_2x_3^*\dfrac{m_b^2}{m_s^2-m_b^2} & x_2\\ x_1^*x_3\dfrac{m_d^2}{m_b^2-m_d^2} & x_2^*x_3\dfrac{m_s^2}{m_b^2-m_s^2} & 1-\frac{1}{2}|x_3|^2 & x_3\\ -x_1^* & -x_2^* & -x_3^* & 1-\frac{1}{2}\sum_{j=1}^{3}|x_j|^2\end{bmatrix}, \tag{18}$$

$$\mathbf{K}_R=\begin{bmatrix}1 & x_1x_2^*\dfrac{m_dm_s}{m_d^2-m_s^2} & x_1x_3^*\dfrac{m_dm_b}{m_d^2-m_b^2} & \dfrac{m_d}{M}x_1\\ x_1^*x_2\dfrac{m_dm_s}{m_s^2-m_d^2} & 1 & x_2x_3^*\dfrac{m_sm_b}{m_s^2-m_b^2} & \dfrac{m_s}{M}x_2\\ x_1^*x_3\dfrac{m_dm_b}{m_b^2-m_d^2} & x_2^*x_3\dfrac{m_sm_b}{m_b^2-m_s^2} & 1 & \dfrac{m_b}{M}x_3\\ -\dfrac{m_d}{M}x_1^* & -\dfrac{m_s}{M}x_2^* & -\dfrac{m_b}{M}x_3^* & 1\end{bmatrix}, \tag{19}$$

where $x_i = F_i/M$. Thus Eq. (16) can be rewritten in terms of mass eigenstates d'^i as

$$L_Z^{\mathrm{FCNC}}(\mathrm{down}) = \beta_{ij}\bar{d}'^i_R \gamma_\mu d'^j_R Z^\mu \quad \text{for } i \neq j \ , \tag{20}$$

where

$$\beta_{ij} = (-\tfrac{1}{2})\frac{g_2}{\cos\theta_W}(K_R)^*_{4i}(K_R)_{4j}$$
$$= (-\tfrac{1}{2})\frac{g_2}{\cos\theta_W}\frac{m_{d_i}m_{d_j}}{MM}x_i x_j^* \ . \tag{21}$$

Therefore, the FCNC induced by Z coupling is highly suppressed by the small mass ratio of usual to new quarks. It is because the mixings of right-handed quarks require a helicity flip of the usual quarks. For example, $\beta_{12} \approx 5.8 \times 10^{-8} x_1 x_2^* < 10^{-10}$ for $M = 100$ GeV, while the experimental limits on FCNC's require only that $\beta_{12} < 10^{-6}$.

FCNC's can also be induced by aspon (A) couplings, which are given by

$$L_A^{\mathrm{FCNC}}(\mathrm{down}) = \alpha_{ij}\bar{d}'^i_L \gamma_\mu d'^j_L A^\mu \quad \text{for } i \neq j \ , \tag{22}$$

where

$$\alpha_{ij} = -g_A x_i x_j^* \ . \tag{23}$$

Therefore, FCNC's induced by A will be important if A is not too heavy compared to Z. Consider the K^0-$\bar{K}^0$ mixing matrix element M_{12}. $\mathrm{Re}(M_{12})$ is expected to be dominated by standard $2W$-exchange box diagrams, while $\mathrm{Im}(M_{12})$ receives its largest contribution from A exchange shown in Fig. 1. We obtain

$$\mathrm{Im}(M_{12}) = \frac{f_K^2 m_K}{6}\frac{1}{\kappa^2}\mathrm{Im}(x_1 x_2^*)^2 \ , \tag{24}$$

where $\kappa^2 = \kappa_1^2 + \kappa_2^2$. The color factor has been taken into account in Eq. (24). $\mathrm{Im}(M_{12})$ receives contributions from the new $2W$-exchange box diagrams shown in Fig. 2, but these contributions are negligible. We will consider the CP-violating parameters, such as $\mathrm{Im}(M_{12})/\Delta M_K$, in more detail in Sec. V.

Next we consider the up sector for completeness. We can also choose the states such that $\mathcal{M}_{\mathrm{up}}$ in Eq. (14) is replaced by

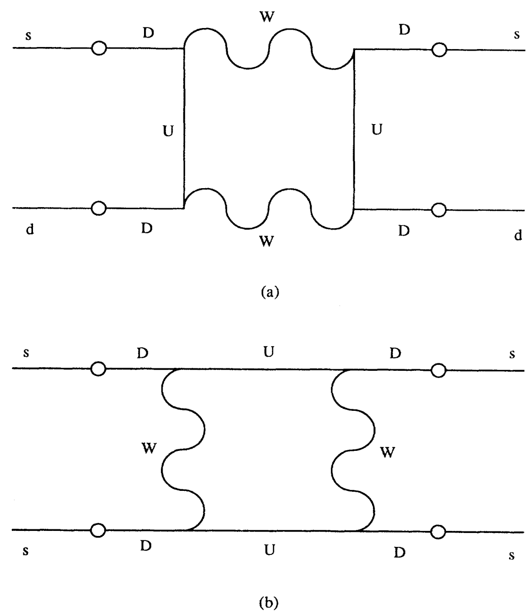

FIG. 2. Contributions to $\mathrm{Im}(\Delta M_{12})$ by new quarks and two-W box diagrams. (Circles mean mixings.)

$$\mathcal{M}_{\mathrm{up}} = \begin{bmatrix} m_u & 0 & 0 & \tilde{F}_1 \\ 0 & m_c & 0 & \tilde{F}_2 \\ 0 & 0 & m_t & \tilde{F}_3 \\ 0 & 0 & 0 & M \end{bmatrix} , \tag{25}$$

with

$$\tilde{F}_i = C_{ij}F_j \ , \tag{26}$$

where $\mathbf{C}$ is the real standard 3×3 KM matrix. The transformation matrices $\mathbf{J}_L$ and $\mathbf{J}_R$ that diagonalize $\mathcal{M}_{\mathrm{up}}$ in Eq. (25) can be related to $\mathbf{K}_L$ and $\mathbf{K}_R$ by changing x_i into $\tilde{x}_i(=\tilde{F}_i/M)$ and m_d, m_s, and m_b into m_u, m_c, and m_t. The generalized 4×4 KM matrix is given by

$$\mathbf{V}_{\mathrm{KM}}^{(4)} = \mathbf{J}_L^\dagger \begin{bmatrix} \mathbf{C} & 0 \\ 0 & 1 \end{bmatrix} \mathbf{K}_L \ . \tag{27}$$

Before leaving this section, we note that flavor-changing Z coupling can be induced by the one-loop diagram shown in Fig. 3. By naive dimensional arguments, the effective coupling coefficient β_{ij} is

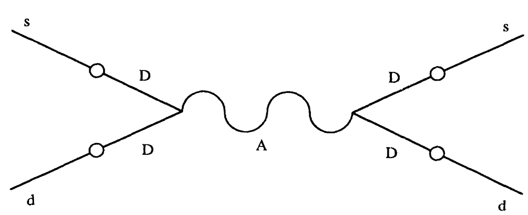

FIG. 1. Contributions to $\mathrm{Im}(\Delta M_{12})$ by aspon exchange. (Circles mean mixings.)

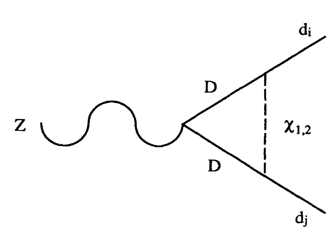

FIG. 3. Flavor-changing Z coupling induced by new quarks at one loop.

$$(-\tfrac{1}{2})\frac{g_2}{\cos\theta_W}\frac{h_i h_j}{16\pi^2}\left[\frac{M}{m_\chi}\right]^2 .$$

Using $h_{1,2,3}\approx 0.01$ and $(M/m_\chi)^2\approx 0.1$, we conservatively estimate β_{ij} to be less than 10^{-7}. Therefore, we expect these FCNC's to be smaller than those in the standard model.

IV. CALCULATION OF $\bar{\theta}$

Consider now the one-loop corrections to $\bar{\theta}$. Although the mass matrices in Eq. (14) are complex, their determinants are real. Therefore, $\bar{\theta}$ defined in Eq. (2) is zero at the tree level. $\bar{\theta}$ will be nonzero when the mass matrices receive radiative corrections. For example, the contributions to $\bar{\theta}$ from the up sector are given by

$$\begin{aligned}\bar{\theta}(\text{up}) &= \text{Arg}[\det(\mathcal{M}_{\text{up}}+\delta\mathcal{M}_{\text{up}})]\\ &= \text{Im}\,\text{Tr}\,\ln[\mathcal{M}_{\text{up}}(1+\mathcal{M}_{\text{up}}^{-1}\delta\mathcal{M}_{\text{up}})]\\ &\approx \text{Im}\,\text{Tr}(\mathcal{M}_{\text{up}}^{-1}\delta\mathcal{M}_{\text{up}}) ,\end{aligned} \tag{28}$$

where we have used the fact that $\mathcal{M}_{\text{up}}$ is real and that the corresponding radiative corrections $\delta\mathcal{M}_{\text{up}}$ are small. The last line in Eq. (28) is valid at one-loop order. Defining the one-loop corrections $\delta\mathcal{M}_{\text{up}}$ by

$$\delta\mathcal{M}_{\text{up}}=\begin{bmatrix}\delta\mathbf{m}_u & \delta\mathbf{F}\\ \delta\mathbf{m}_{Uu} & \delta M\end{bmatrix} , \tag{29}$$

and combining with Eq. (28), we obtain

$$\bar{\theta}(\text{up})=\text{Im}\,\text{Tr}(\mathbf{m}_u^{-1}\delta\mathbf{m}_u-\mathbf{m}_u^{-1}\mathbf{F}M^{-1}\delta\mathbf{m}_{Uu}+M^{-1}\delta M) . \tag{30}$$

Notice that δF will not contribute to $\bar{\theta}$ at one-loop order because of the structure of $\mathcal{M}_{\text{up}}$. The expression for $\bar{\theta}$(down) is strictly analogous to Eq. (30).

First we consider δM. The diagrams are shown in Figs. 4(a)–4(c). The integrals are real because the gauge coupling constants and M are real. Therefore, $\text{Im}\,\text{Tr}(M^{-1}\delta M)$ is equal to zero. Next we consider $\delta\mathbf{m}_{Uu}$, the diagrams of which are given in Fig. 4(d). $\delta\mathbf{m}_{Uu}$ is found to be proportional to $M\mathbf{F}^\dagger\mathbf{m}_u$. It is complex, but $\text{Tr}(\mathbf{m}_u^{-1}\mathbf{F}M^{-1}M\mathbf{F}^\dagger\mathbf{m}_u)$ is real. Finally, we consider $\delta\mathbf{m}_u$, the diagrams of which are given in Figs. 4(e) and 4(f). The contributions from Figs. 4(e) and 4(f) are real. However, the contributions from Fig. 4(g) are proportional to $\mathbf{h}_1 e^{i\theta}\mathbf{F}^\dagger\mathbf{m}_u$ for χ_1 and $\mathbf{h}_2\mathbf{F}^\dagger\mathbf{m}_u$ for χ_2. Therefore, $\text{Tr}(\mathbf{m}_u^{-1}\mathbf{h}_1 e^{i\theta}\mathbf{F}^\dagger\mathbf{m}_u)$ and $\text{Tr}(\mathbf{m}_u^{-1}\mathbf{h}_2\mathbf{F}^\dagger\mathbf{m}_u)$ yield a nonvanishing $\bar{\theta}$. By dimensional arguments, $\bar{\theta}$ is estimated to be

$$\bar{\theta}=\frac{h_i^1 h_i^2}{16\pi^2}\left[\frac{m_{u_i}M}{v^2}\right]\ln\left[\frac{M}{m_{u_i}}\right]^2+\bar{\theta}(\text{down}) . \tag{31}$$

Using $\bar{\theta}<2\times 10^{-10}$, $v=250$ GeV, and $M=150$ GeV (as illustration), we obtain $|h_1^\alpha|\le 10^{-2}$, $|h_2^\alpha|\le 10^{-3}$, and $|h_3^\alpha|\le 10^{-4}$ (for $m_t\approx 100$ GeV) for $\alpha=1,2$. Notice that if $\chi_1=\chi_2$, we find that $\text{Tr}(\mathbf{m}_u^{-1}\mathbf{h}_1 e^{i\theta}\mathbf{F}^\dagger\mathbf{m}_u)+\text{Tr}(\mathbf{m}_u^{-1}\mathbf{h}_2\mathbf{F}^\dagger\mathbf{m}_u)$ is real as expected. Therefore, $\bar{\theta}$ is

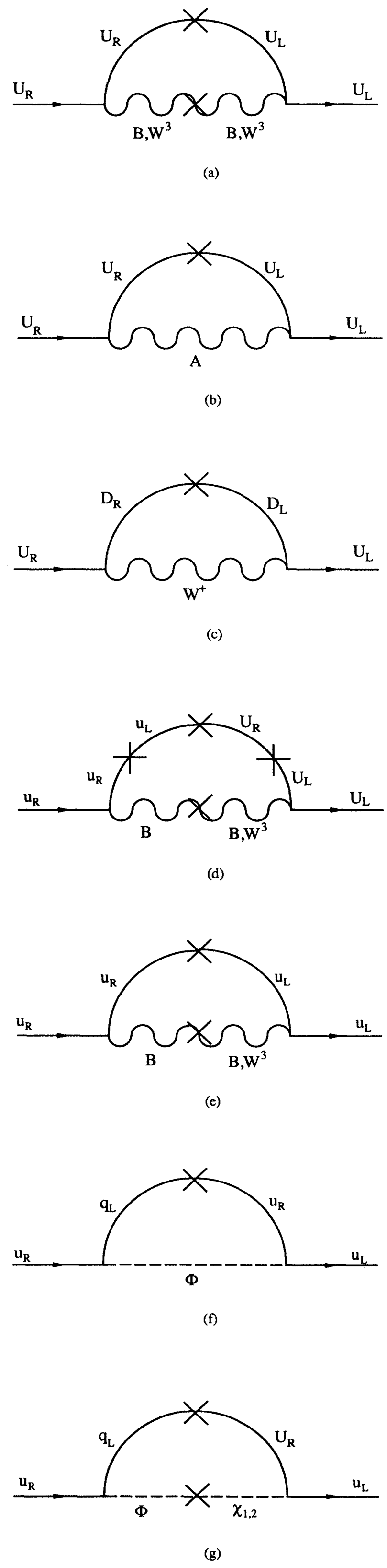

FIG. 4. One-loop corrections to $\mathcal{M}_{\text{up}}$. (Crosses mean mass insertions.)

calculable in this model as well as the models proposed in Ref. 5. However, this model is much simpler and gives a connection between strong and weak *CP*.

V. WEAK *CP*-VIOLATION PARAMETERS ϵ AND ϵ'/ϵ

The model-independent *CP*-violating parameter ϵ can always be fixed to agree with experiment. However, ϵ'/ϵ, which is model dependent, cannot be made arbitrary. For example, $|\epsilon'/\epsilon|$ was estimated[8] to be 0.048 in Weinberg's spontaneous *CP*-violation model,[9] which is at least ten times greater than the experimental measurement. The challenge is to make $\bar{\theta}$ small for values of ϵ and ϵ'/ϵ which are consistent with experiments.

Since flavor-changing Z couplings are negligible, the weak *CP*-violating effects are similar to the standard model with four generations. We first consider the $\Delta S=1$ effective Hamiltonian[10]

$$H^{\Delta S=1}=\frac{G_F}{\sqrt{2}}\sum_r (z_r\lambda_u - y_r\lambda_t - y_r'\lambda_U)Q_r \quad (r=1,2,3,5,6,7,8)\ , \tag{32}$$

where the definitions of z_r, y_r, and Q_r follow Buchalla *et al.*, Ref. 10. y_r' is the Wilson coefficient for the new heavy quark U. λ_q is defined as $V_{qs}^* V_{qd}$. The imaginary part of the amplitude A_0 of $K^0\to 2\pi(I=0)$ is dominated by the contribution from Q_6:[10]

$$\mathrm{Im}\,A_0 \approx -\frac{G_F}{\sqrt{2}} y_6' \mathrm{Im}\lambda_U \langle 2\pi^2(I=0)|Q_6|K^0\rangle\ , \tag{33}$$

where $\mathrm{Im}\lambda_U\ [=O(x_i^2)] \gg \mathrm{Im}\lambda_u, \mathrm{Im}\lambda_t\ [=O(x_i^4)]$ is used. Using $\mathrm{Re}\,A_0=4\times10^{-7}$ GeV, $y_6'\approx 0.08$ ($y_6'\approx y_6$ and $M_U\approx M_t\approx 150$ GeV are assumed), and $\langle 2\pi^2(I=0)|Q_6|K^0\rangle \approx 1.6\ \mathrm{GeV}^3$, we obtain the magnitude of the ratio $\delta_0=\mathrm{Im}\,A_0/\mathrm{Re}\,A_0$ to be

$$|\delta_0| = \left|\frac{\mathrm{Im}\,A_0}{\mathrm{Re}\,A_0}\right| \approx 1.9|\mathrm{Im}(x_1 x_2^*)|\ . \tag{34}$$

On the other hand, the magnitude of the ratio $\epsilon_m=\mathrm{Im}M_{12}/\Delta M_K$ is given by

$$|\epsilon_m| = \left|\frac{\mathrm{Im}M_{12}}{\Delta M_K}\right| \approx 6.1\times10^5 \left|\frac{1\ \mathrm{TeV}}{\kappa}\right|^2 |\mathrm{Im}(x_1x_2^*)^2|\ , \tag{35}$$

where $\Delta M_K = 3.5\times10^{-15}$ GeV, $M_K=0.5$ GeV, and $f_K^2=0.16$ GeV are used.

Next we consider ϵ'/ϵ. ϵ'/ϵ is defined as

$$\left|\frac{\epsilon'}{\epsilon}\right| \approx \frac{1}{\sqrt{2}}\left|\frac{1}{\epsilon}\right| \left|\frac{\mathrm{Re}\,A_2}{\mathrm{Re}\,A_0}\right| \left|\frac{\mathrm{Im}\,A_0}{\mathrm{Re}\,A_0}\right| \tag{36}$$

$$\approx \left|\frac{\delta_0}{\epsilon_m+\delta_0}\right| \left|\frac{\mathrm{Re}\,A_2}{\mathrm{Re}\,A_0}\right|\ , \tag{37}$$

where $|\epsilon|$, defined as $(1/\sqrt{2})|\epsilon_m+\delta_m|$, is used in Eq. (37). Comparing with experimental data,[11,12]

$$\mathrm{Re}\left[\frac{\epsilon'}{\epsilon}\right] = \begin{cases} (3.3\pm1.1)\times10^{-3} & \mathrm{CERN(NA31)}\ , \quad (38) \\ (-0.5\pm1.5)\times10^{-3} & \mathrm{Fermilab(E731)}\ , \quad (39) \end{cases}$$

we find that ϵ_m is much greater than δ_0 as $\mathrm{Re}\,A_2/\mathrm{Re}\,A_0$ is approximately equal to $\frac{1}{22}$. Therefore, ϵ is essentially equal to $\epsilon_m/\sqrt{2}$. Using the experimental value[13] for ϵ, 2.258×10^{-3}, Eq. (35) yields

$$\left|\frac{1\ \mathrm{TeV}}{\kappa}\right|^2 |\mathrm{Im}(x_1x_2^*)^2| \approx \left|\frac{1\ \mathrm{TeV}}{\kappa}\right|^2 |x_1x_2^*|^2 = 5.2\times10^{-9}\ . \tag{40}$$

Although the values in Eqs. (38) and (39) tend to exclude each other, we can make a conservative bound for the magnitude of ϵ'/ϵ being less than 4×10^{-3}. Then we obtain

$$\left|\frac{\kappa}{1\ \mathrm{TeV}}\right|^2 |x_1x_2^*|^{-1} < 2.8\times10^4\ . \tag{41}$$

Combining Eqs. (40) and (41), we obtain

$$|x_1x_2^*| < 1.5\times10^{-4}\ , \tag{42}$$

and hence one of our principal results which is

$$\kappa < 2.0\ \mathrm{TeV}\ . \tag{43}$$

This shows that for consistency the SSB of *CP* symmetry must occur at relatively low energy, below 2.0 TeV. In the early universe, this SSB of a discrete symmetry will generate domain walls (unless *CP* is softly broken) which if they persisted would dominate the present energy density of the universe. Thus we required cosmic inflation to occur at a temperature *below* 2 TeV and then baryogenesis must occur at a temperature below that at which inflation occurs. Furthermore, h_1 and h_2 are expected to be less than 10^{-2} and 10^{-3} from $\bar{\theta}$. Therefore Eq. (41) yields

$$M < 530\ \mathrm{GeV}\ . \tag{44}$$

The *aspon mass,* which is given by $g_{\mathrm{new}}\kappa$, *is less than* 600 GeV ($g_{\mathrm{new}}=e$ is assumed). To get the feel of how κ and M depend on ϵ'/ϵ, we tabulated the result in Table I for several values of ϵ'/ϵ. We find that κ depends linear in ϵ'/ϵ, while M depends on the square root of ϵ'/ϵ.

It is worth remarking that $\mathrm{SU}_C(3)\times\mathrm{SU}_L(2)\times\mathrm{U}_Y(1)$ singlet scalars that one could identify with our χ occur naturally in the string model of Ref. 14.

We have predicted that the aspon and the new quark masses are bounded to be 600 and 530 GeV, respectively.

TABLE I. The bounds for the breaking scale, κ, and the mass of the new quarks, M, for various values of ϵ'/ϵ.

ϵ'/ϵ	κ (TeV)	M (GeV)
$<4\times10^{-3}$	<2.0	<530
3×10^{-3}	1.5	<460
1×10^{-3}	0.5	<270

They can thus be produced in the CERN Large Hadron Collider (LHC) and the Superconducting Super Collider (SSC), for example, by gluon fusions, such as $gg \to U\bar{U}(D\bar{D})$ by $U(D)$ t-channel exchange and $gg \to A+g(\gamma)$ at one loop with U and D in the internal loop. If the aspon is lighter than 200 GeV, it can be produced by two-W fusions, or e^+e^- collision (if there exists a vectorlike lepton doublet), in LEP II as early as 1993.

VI. CONCLUSION

In this paper, we have discussed a gauge model of $SU_C(3)\times SU_L(2)\times U_Y(1)\times U_{\rm new}(1)$. A vectorlike quark doublet and two singlet Higgs scalars were introduced. *CP* invariance is assumed to be broken spontaneously by the VEV's of the Higgs singlet. The model was arranged such that the determinants of quark mass matrices are real at the tree level. Therefore, $\bar{\theta}$ is nonzero only through one-loop radiative corrections. Flavor-changing neutral currents are naturally suppressed in the model. Using $\epsilon = 2.258\times 10^{-3}$ and $|\epsilon'/\epsilon| < 4\times 10^{-3}$, we obtained that the *CP*-breaking scale and the new quark masses are bounded to be less than 2.0 TeV and 530 GeV, respectively. Therefore, the aspon and new quarks can be produced in LHC and SSC.

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