Implication of scalar-pseudoscalar mixing on $\epsilon'/\epsilon$ in SUSY models

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Abstract

We study the effects of scalar-pseudoscalar mixing induced from quantum loop on $\epsilon'/\epsilon$ in SUSY models with $\tan \beta \sim m_t/m_b$. We find that even the non-universal soft $A_d$ term and Yukawa matrix, $Y^d$, are hermitian, the predicted value of $|\epsilon'/\epsilon|$ can be consistent with the measured results of NA48 and KTeV. And also the EDMs are compatible with experimental bounds.
Since its discovery in the neutral kaon decays in 1964 [1], the origin of CP violation (CPV) still puzzles the physicist. Although recently another time-dependent CP asymmetry (CPA), $\sin 2\phi_1$, in the decay of $B \rightarrow J/\Psi K_s$ is observed by BARBAR [2] and BELLE [3], our understanding of CPV is not much better than before. What we are certain at present is only that it is necessary to exist unrotatable phase and it is believable that such kind of phase is associated with weak interactions, usually called weak phase. In the case of standard model (SM), the unique source of CPV is from the Cabibbo-Kobayashi-Maskawa (CKM) matrix [4] induced from the three-generation quark mixings and described by the three angles $\alpha$, $\beta$ and $\gamma$ or $\phi_2$, $\phi_1$ and $\phi_3$.

Even though the SM prediction on the indirect CP violating parameter $\epsilon$ in the kaon system can be fitted well with current experimental data, due to the large uncertainties from hadronic matrix elements, so far it has not been settled yet whether the result in the SM can explain the observed value of the direct CP violating parameter $\epsilon'$ measured by NA48 [5] and KTeV [6]. And also, because of the unitarity in the CKM matrix, the predicted electric dipole moments (EDMs) of neutron and lepton are quite small and unreachable experimentally. Moreover, the requirement of Higgs boson with the mass being less than 60 GeV to solve the problem of baryogenesis is ruled out by the LEP experiment. Therefore, it becomes important to search the possibility of existing other CP violating sources for explaining all CP phenomena.

One of reliable models in the extension of SM is supersymmetric (SUSY) model. SUSY theories not only supply an elegant mechanism for the breaking of the electroweak symmetry and a solution to the hierarchy problem, but also guarantee the unification of gauge couplings at GUTs scale [7]. In addition, SUSY possesses abundant flavor structures, such as upper and down type squark mixing matrices, and CP violating phases, which are arisen from the trilinear and bilinear SUSY soft breaking $A$ and $B$ terms, the $\mu$ parameter for the scalar mixing etc.. Unfortunately, one can check easily that those phases are severely bounded by electric dipole moments (EDMs) [8] so that the effects on $\epsilon$ and $\epsilon'$ cannot enough explain the current experimental values. In order to handle the small CP phase problem, it has been suggested to use the non-universal soft $A$ terms instead of universal ones [9]. Furthermore, to evade any fine-tuning on the specific phases which contribute to EDM of neutron, it is proposed that SUSY soft-breaking $A^q$ and Yukawa, $Y^q$, matrices are hermitian [10, 11]. Consequently, the CP phases of $O(1)$ can exist naturally even considering the contributions of EDM. And also, it implies that the CPA in hyperon decays could reach the value of $O(10^{-4})$.
proposed by the experiment E871 at Fermilab\cite{13}. However, the effect on \( \epsilon' \), dominated by gluon-penguin with gluino as the internal particle in the one-loop, will be suppressed due to \( (\delta_{12}^d)_{LR} \approx (\delta_{12}^d)_{RL} \) (the definition is shown below). It is interesting to ask whether \( \epsilon' \), based on the gluino mechanism, can be satisfied in the framework of hermitian \( A^d \) and \( Y^d \) matrices by considering other effects which are available after including the constraints from experimental measurements.

According to the analysis in \cite{14}, we find that because of the enhancement of large \( A^q \), \( \mu \) and \( \tan \beta \), the coupling \( N^0 - \tilde{q}_L - \tilde{q}_R \), with \( N^0 \) and \( \tilde{q} \) being scalar (or pseudoscalar) and the squark of corresponding to the \( q \)-quark respectively, will be enhanced. Thus, it is easy to conjecture that the vanished \( \epsilon' \) problem in hermitian case could be solved if we consider the induced interactions \( C_{f N^0} N^0 \bar{f}_R f_{Lj} \) in which \( C_{f N^0} \) is the effective coupling and \( f \) could be upper or down type quarks and the indicies \( i \) and \( j \) denote the possible flavour. From the induced effective vertices, if \( C_{f N^0} \) is complex, we see clearly that \( \epsilon \) will also contribute. If so, we will face the strict constraints from it. In our following analysis, the induced effective coupling will be taken as real. But, the CPV is generated by the scalar-pseudoscalar mixing.

It is known that if the scalar-pseudoscalar mixing comes from spontaneous CPV, such as realized by radiative corrections in the MSSM \cite{15}, the predicted pseudoscalar mass, \( m_{A^0} \), is far below the current experimental limit and excluded \cite{16}. However, it is found that if CP is broken at the tree level by the SUSY soft breaking sector, the large mixing between CP-even and CP-odd boson can be obtained through radiative corrections without the limit of small mass on the pseudoscalar boson \cite{17}. As also been shown in \cite{18}, the radiative CP effects will modify the couplings of Z-boson to Higgs boson such that the mass of lightest Higgs boson with \( 60-70 \) GeV can escape from the bound of LEP2. With above conclusions, in this paper, we will show their implications on \( \epsilon' \).

We start by writing the effective interactions between the relevant neutral Higgs bosons or gluino and squarks as

\[
\mathcal{L}_{\text{eff}} = \mathcal{L}_{N f_L f_R} + \mathcal{L}_{\tilde{q} \tilde{f} f},
\]

\[
= \frac{g}{2 M_W} \left[ h^0 \left( (M^2_{\tilde{q}})_{LR} \tilde{q}^\dagger \tilde{q} + \mu^* m_{\tilde{q}} \tilde{\omega}^2 \right) \tilde{q}_R 
- \frac{i A^0 \tilde{q}^\dagger_i \left( (M^2_{\tilde{q}})_{LR} Z^a_{\beta} - \mu^* m_{\tilde{q}} \right) \tilde{q}_R 
- \sqrt{2} g_s \left[ \bar{q} P_R \tilde{g}^a T^a \tilde{q}_L - \bar{q} P_L \tilde{g}^a T^a \tilde{q}_R \right] + h.c. \right]
\]  

(1)

where \( P_{L(R)} = (1 \mp \gamma_5)/2 \), \( h^0 \) stands for the lightest scalar particle while \( A^0 \) is for pseudoscalar boson. The bold \( \mathbf{q} \) and \( \tilde{\mathbf{q}} \) denote three generation quarks and the corresponding squarks,
respectively. The generators of $SU(3)_c$ are normalized by $tr(T^a T^b) = 1/2 \delta^{ab}$. $\tilde{\omega}^{u(d)}_1 = \cos \alpha / \sin \beta (- \sin \alpha / \cos \beta)$ and $\tilde{\omega}^{u(d)}_2 = - \sin \alpha / \sin \beta (\cos \alpha / \cos \beta)$ in which angle $\alpha$ describes the mixing between two CP-even Higgs particles. For simplicity, we only concentrate on the contributions from the lightest CP-even Higgs. Due to the mass suppression, we expect that the effects of heavier CP-even boson are smaller than those of light one. One the other hand, the squared squark mass matrices responsible for flavor change are described by

$$\mathcal{M}^2_\tilde{q} = \begin{pmatrix} (m^2_{\tilde{q}})^{LL} & (m^2_{\tilde{q}})^{LR} \\ (m^2_{\tilde{q}})^{LR} & (m^2_{\tilde{q}})^{RR} \end{pmatrix},$$

(2)

$$\begin{align*}
(m^2_{\tilde{q}})^{LL} &= (M^2_{\tilde{q}})^{LL} + m^2_q - M^2_Z \cos 2 \beta C^q_L \hat{1}, \\
(m^2_{\tilde{q}})^{LR} &= (M^2_{\tilde{q}})^{LR} - \mu^* Z^q \tilde{m}_q, \\
(m^2_{\tilde{q}})^{RR} &= (M^2_{\tilde{q}})^{RR} + m^2_q + M^2_Z \cos 2 \beta C^q_R \hat{1},
\end{align*}$$

(3)

where we have adopted the so-called super-CKM basis that the quarks have been the mass eigenstates so that $m_q$ is the diagonalized quark mass matrix. $q$ and $\tilde{q}$ stand for quark and its superpartner. They could be upper or down type quark. $Z^u(d) = \cot \beta (\tan \beta)$ and

$$\begin{align*}
C^q_L &= T^3_q - Q_q \sin^2 \theta_W, \\
C^q_R &= Q_q \sin^2 \theta_W
\end{align*}$$

(4)

with $T^3_q$ and $Q_q$ being the z-component of isospin $SU(2)_L$ for the squark $\tilde{q}$ and its charge, respectively. $\hat{1}$ denotes the $3 \times 3$ unit matrix. The definition angle $\beta$ is followed by $\tan \beta = v_u/v_d$ with $v_u$ and $v_d$ being the vacuum expectation values (VEVs) of Higgs fields $\Phi^u$ and $\Phi^d$ responsible for the masses of upper type quarks and down type quarks, respectively. $\mu$ is the mixing effects of $\Phi^u$ and $\Phi^d$. $(M^2_{\tilde{q}})^{LL(RR)}$ stand for the soft breaking masses for the corresponding squarks and $(M^2_{\tilde{q}})^{LR}$ describe the trilinear soft breaking couplings and are written as

$$\begin{align*}
(M^2_{\tilde{q}})^{LR} &= \frac{v_q}{\sqrt{2}} V_{ql} \tilde{A}^q V_{qr}^\dagger,
\end{align*}$$

(5)

where $V_{ql(r)}$ transform the left(right)-handed quarks from weak eigenstates to mass eigenstates and $\tilde{A}^q_{ij} = Y^q_{ij} A^q_{ij}$ with $Y^q_{ij}$ being the Yukawa matrix.

For convenience, we adopt the so-called mass-insertion approximation method [19] in which the basis for squark is chosen such that the gluino-squark-quark vertices involving
quarks are flavour diagonal (super-CKM basis) instead of diagonalizing the squark mass matrix itself. Hence, the squared squark mass matrices are regarded as effective couplings. If necessary, we can insert the proper effective couplings in the propagator of squark. According

\[ N \]

\[ q_{Ai} \]

\[ q_{Bj} \]

\[ \tilde{g} \]

\[ \tilde{q}_{Ai} \]

\[ \tilde{q}_{Bj} \]

\[ (a) \]

\[ N \]

\[ q_{Ai} \]

\[ q_{Bj} \]

\[ \tilde{g} \]

\[ \tilde{q}_{Ai} \]

\[ \tilde{q}_{Bj} \]

\[ (b) \]

Figure 1: Feynman diagrams for one-loop induced couplings of neutral Higgs to quarks: \( N^0 \) could be scalar or pseudoscalar boson. \( q \) could be upper or down type quarks, \( \tilde{q} \) is the corresponding superpartner, and the indices \( i,j,k \) stand for the possible flavour of quarks. \( A^{(i)} \) and \( B^{(i)} \) denote the chirality and \( B^{(i)} = R(L) \) while \( A^{(i)} = L(R) \).

to Eqs. (1) and (2), the interactions for \( N^0 \bar{d}_{BA} \) and \( N^0 \bar{q}_{B}q_{A} \), illustrated in Figure 1, can be derived as

\[
\mathcal{L} = (4\sqrt{2}G_F)^{1/2} \left[ (\tilde{N}^{d}_{12})_{BA} \tilde{d}_{BA} A N^0 + (\tilde{N}^{d}_{11})_{BA} \tilde{d}_{BA} d_{A} N^0 \right. \\
\left. + (\tilde{N}^{u}_{11})_{BA} \tilde{u}_{BA} u_{A} N^0 \right]
\]

(6)

where \( A \) and \( B \) denote the chiralities and they are always opposite to each other, \((\tilde{N}^q)_{BA} = (\tilde{H}^q)_{BA} ((\tilde{A}^q)_{BA}) \) if \( N^0 \) is scalar (pseudoscalar) Higgs and their expressions are written as

\[
(\tilde{N}^q)_{BA} = \frac{\alpha_s}{4\pi} \sqrt{x_q} C_F \left[ \hat{n}^{q}_{ij} I_1(x_q) \\
+ \frac{1}{3} (\delta^q_{kj}) B^B \hat{n}^{q}_{jk} (\delta^q_{ik}) A A' I_2(x_q) \right],
\]

(7)

\[
(\delta^q_{ij})_{AA'} = \frac{(M_{\tilde{q}_{ij}}^2)_{AA'}}{m_{\tilde{q}}^2}
\]

with \( x_q = m_q^2/m_{\tilde{q}}^2 \) and \( m_{\tilde{g}} \) and \( m_{\tilde{q}} \) being the average masses of gluino and \( q \) type squark.

The first term in Eq. (3) is from the lowest order contributions, illustrated in Figure 1(a), while the second term is generated by double mass-insertion, shown in Figure 1(b). The chirality \( A' \) in Eq. (3) can be \( L \) or \( R \), but chirality \( B' \) is opposite to \( A' \) so that beside \((\delta^q_{ij})_{LR(LR)} \), \((\delta^q_{ij})_{LL(RR)} \) will also contribute. \( \hat{n}^{q}_{A} = \hat{h}^{q}_{A}(\hat{a}^{q}_{A}) \) are related to the couplings of scalar (pseudoscalar) boson to the squark \( \tilde{q} \) in which the expressions are

\[
\hat{h}^{q}_{L} = \frac{(M_{\tilde{q}}^2)_{LR \tilde{q}}}{m_{\tilde{q}}} \hat{\omega}_1 + \frac{\mu^* m_{\tilde{q}}}{m_{\tilde{q}}} \hat{\omega}_2,
\]
\[ \hat{a}_L^q = -i \frac{(M_d^q)^{LR} Z_f^q}{m_{\tilde{q}}} + i \frac{\mu m_{\tilde{q}}}{m_{\tilde{q}}} \]

and \( \hat{n}_R = \hat{n}_L^q \).

\[ I_1(x) = \frac{1}{1-x} + \frac{x \ln(x)}{(1-x)^2}, \]
\[ I_2(x) = \frac{2 + 3x - (6-x)x^2}{2(1-x)^4} + \frac{3x \ln(x)}{(1-x)^4}. \]

We note that the relevant effects in the first term of Eq. (7) are \( \hat{a}_{L12}^q, \hat{a}_{L11}^q \), and \( \hat{a}_{L11}^u \). Via their definitions, we have

\[ \hat{h}_{L12}^d = m_d(\delta_{12}^q)_{LR}\tilde{\omega}_1, \]
\[ \hat{h}_{L11}^d = m_d(\delta_{11}^q)_{LR}\tilde{\omega}_1^d + \frac{\mu^* m_{d1}}{m_{\tilde{d}}} \tilde{\omega}_2^d, \]
\[ \hat{h}_{L11}^u = m_u(\delta_{11}^q)_{LR}\tilde{\omega}_1^u + \frac{\mu^* m_{u1}}{m_{\tilde{u}}} \tilde{\omega}_2^u \]

with \( m_{u1} \) and \( m_{d1} \) being the masses of u-quark and d-quark, respectively. According to the analysis in [21], the constraint on \( |Im(\delta_{12}^q)_{LR}| \) from \( \epsilon' \) is of order of \( 10^{-5} \). With the assumption of the phase of \( O(1) \), we expect that the \( |(\delta_{12}^q)_{LR}| \) has the similar order of magnitude so that its contribution is negligible. Although it is not necessary, to simplify our analysis, we adopt that the Yukawa matrices have the structure \( Y_u(d)_{ij} > Y_u(d)_{33} \) in which \( Y_u(d) \) is any entry except \( Y_{33} \) and each of them is proportional to \( \sqrt{m_i m_j / m_3} \) with \( m_3 \) being the top or bottom quark mass for the corresponding Yukawa matrix [21]. As a result, \( (\delta_{11}^q)_{LR} \) will be related to the first two generation quark masses or be suppressed by flavour mixing elements as defined in Eq. (5). In this paper, we also neglect their contributions. The similar situation is also applied to \( \hat{a}_L^q \). According to the above assumption, we see that the dominant effect in second term of Eq. (7) is from the component \( \hat{h}_{L33}^d \) which are expressed by

\[ \hat{h}_{L33}^q = m_q(\delta_{33}^q)_{LR}\tilde{\omega}_1^q + \frac{\mu^* m_3}{m_{\tilde{q}}} \tilde{\omega}_2^q, \]
\[ \hat{a}_{L33}^q = -im_q(\delta_{33}^q)_{LR}\tilde{Z}_\beta + i \frac{\mu^* m_3}{m_{\tilde{q}}}. \]

It is worth mentioning that due to \( \hat{h}_{L33}^d \) and \( \hat{a}_{L33}^d \) associated with \( 1 / \cos \beta \) and \( \tan \beta \) respectively, if considering large \( \tan \beta \) case, there is a large enhancement.

In terms of Eq. (3), we can immediately obtain the effective operators for \( |\Delta S| = 2 \) and \( |\Delta S| = 1 \) as

\[ \mathcal{L}_{|\Delta S|=2} = -4\sqrt{2} G_F \sum_{A,A'=L,R} \frac{(\hat{H}_{12}^d)^{BA}(\hat{H}_{12}^d)^{B'A'}}{m_h^2}, \]
the main effects on the FCNC decays, thus, due to $m_B$, we only concentrate on the contributions of $m_A$, respectively, $\Omega = \Omega_0 = m_0$ and $u, d$. $\Delta K$ denote the two-pion final states of $K_L \to \pi \pi$ in isospin $I = 0$ and $I = 2$, respectively, $\Omega = ReA_2/ReA_0 \approx 1/22$ and $ReA_0 \approx 2.7 \times 10^{-7}$ GeV. From Eqs. (10) and (14),

\begin{align}
\mathcal{L}_{\Delta S=1} &= -4\sqrt{2}G_F \sum_{A',A=L,R} \left[ \frac{(\tilde{H}_{12}^q)_{BA}(\tilde{A}_{1l}^q)_{BA'}}{m_h^2} + \frac{(\tilde{A}_{12}^q)_{BA}(\tilde{H}_{11}^q)_{BA'}}{m_h^2} \right] \Delta SP \bar{q} B_{q'A} \bar{d} B_{s'A},
\end{align}

where $q = u, d$ and $\Delta SP = M_{SP}/m_A^2$ describes the scalar-pseudoscalar mixing effect and its qualitative dependence can be described by \[17, 18\]

\begin{align}
\mathcal{M}_{SP}^2 &\approx \frac{m_q^2}{v^2} \frac{Im(A^t \mu)}{32\pi^2 M^2_{SUSY}} \left( 1, \frac{|A|^2}{M^2_{SUSY}}, \frac{|\mu|^2}{M^2_{SUSY}}, 2Re(A^t \mu) \right).
\end{align}

In terms of the results in \[17\], one can find that one loop radiative effects $M_{SP}$ could be of order of few hundred GeV, that is the CP mixing factor $\Delta SP$ could be $O(1)$. In order to study the contributions to $|\Delta S| = 2$ and 1 decays, such as $\Delta m_K$, $\epsilon$ and $\epsilon'$, the relevant hadronic matrix elements are estimated by vacuum saturation method and written as

\begin{align}
\mathcal{K}_1 &= \langle K^0|\bar{d}_R s_L \bar{d}_R s_L|K^0\rangle = -\frac{5}{24} \left( \frac{m_K}{m_s + m_d} \right)^2 m_K f_K^2, \\
\mathcal{K}_2 &= \langle K^0|d_L s_R \bar{d}_R s_L|K^0\rangle = \left[ \frac{1}{24} + \frac{1}{4} \left( \frac{m_K}{m_s + m_d} \right)^2 \right] m_K f_K^2
\end{align}

for $\Delta S = 2$ decays \[22\] and

\begin{align}
\mathcal{P}_1 &= \langle \pi^-|\bar{d}\gamma_5 u|0\rangle \langle \pi^+|\bar{u}s|K^0\rangle = -f_\pi B_0^2 \left( 1 + 2 \frac{m_d^2}{\Lambda^2} \right), \\
\mathcal{P}_2 &= \langle \pi^+ \pi^-|\bar{q}q|0\rangle \langle 0|\bar{d}\gamma_5 s|K^0\rangle = -f_\pi B_0^2 \left( 1 + 2 \frac{m_K^2}{\Lambda^2} \right)
\end{align}

with $B_0 = m_K^2/(m_s + m_d)$ and $\Lambda \approx 1$ GeV for $|\Delta S| = 1$ \[23\]. What we are concerned is only the main effects on the FCNC decays, thus, due to $m_h < m_A$, in the following calculations, we only concentrate on the contributions of $m_h$ in Eqs. (10).

It is known that the observable for $\epsilon'/\epsilon$ is expressed as

\begin{align}
Re \left( \frac{\epsilon'}{\epsilon} \right) = \frac{\Omega}{\sqrt{2} ReA_0 |\epsilon|} (\Omega^{-1} Im A_2 - Im A_0)
\end{align}

where $A_0$ and $A_2$ denote the two-pion final states of $K_L \to \pi \pi$ in isospin $I = 0$ and $I = 2$, respectively, $\Omega = ReA_2/ReA_0 \approx 1/22$ and $ReA_0 \approx 2.7 \times 10^{-7}$ GeV. From Eqs. (11) and (14),
we know that the new effects on Im\(A_2\) and Im\(A_0\) have similar value in magnitude so that due to the suppressed factor \(\Omega\) in isospin \(I = 0\), the final state with \(I = 2\) is dominant. Because the CP violating phases are arisen from soft SUSY breaking terms, in our considering case the \((\delta^d_{33})_{LL(RR)}\) are real. On the other hand, although \((\delta^d_{33})_{LR(RL)}\) \((i = 1, 2)\) could be complex, however, charge and color breaking (CCB) minima and the potential unbounded from below (UFB) will give strict constraints \([24]\) so that their contributions are small and negligible. In order to avoid the constraints from \(\epsilon\), we impose that \(Arg(\mu) = 0, A^d\) and \(Y^d\) are hermitian matrices so that \((\delta^d_{33})_{LR(RL)}\) is real, and \((\delta^g_{ij})_{LL} \approx (\delta^g_{ij})_{RR}\). As a result, \(\hat{h}^d_{A33}\) and \(\hat{a}^d_{A33}\) are real and purely imaginary, respectively. And then we get the identities \((\hat{H}^d_{12})_{LR} = (\hat{H}^d_{12})_{RL}\) and \((\hat{A}^d_{12})_{LR} = -(\hat{A}^d_{12})_{RL}\) such that the \(\epsilon\) relating effect, such as \((\hat{H}^d_{12})_{RL}(\hat{A}^d_{12})_{RL} + (\hat{H}^d_{12})_{LR}(\hat{A}^d_{12})_{LR}\) is vanished. Altogether, Eq. (10) can only contribute to \(\Delta m_K\). Moreover, from Eq. (12), we clearly see that the \(M^2_{SP}\) is related to \(Im(A^t\mu)\), therefore, \((\delta^u_{33})_{LR}\) in Eq. (3) should be complex. Nevertheless, in terms of Eq. (11) and due to purely imaginary \((\hat{A}^d_{12})_{BA}\), it is obvious that only real part has the contribution.

Combining the results of Eqs. (10)-(17), we get

\[
\Delta m_K = 2\text{Re}(K^0|\mathcal{H}|_{\Delta S} = 2|\bar{K}^0),
\]

\[
\sim 16\sqrt{2}G_F \frac{(\hat{H}^d_{12})_{RL}^2}{m_h^2}(\mathcal{K}_1 + \mathcal{K}_2),
\]

\[
\text{Re} \left( \frac{\epsilon'}{\epsilon} \right)_{l=2} \sim \frac{2\sqrt{2}G_F}{3m_h^2 \text{Re} A_0|\epsilon|} \left( 1 - \frac{1}{2N_c} \right) \mathcal{P}_1 \times [(\hat{H}^d_{12})_{RL}(\hat{A}^u_{11})_{RL} + (\hat{A}^d_{12})_{RL}(\hat{H}^u_{11})_{RL}] \hat{\Delta}_{SP} (16)
\]

where \(N_c = 3\) is the color number, for simplicity we have used \((\hat{H}^u_{ij})_{RL} \approx (\hat{H}^u_{ij})_{LR}\) and \((\hat{A}^d_{ij})_{RL} \approx -(\hat{A}^d_{ij})_{LR}\) for \(q = u, d\). Because the constraint on \((\delta^d_{13})_{LL}\) is stricter than that on \((\delta^u_{13})_{LL}\), in Eq. (17), we only show the contributions of \((\hat{H}^u_{11})_{AB}\) and \((\hat{A}^u_{11})_{AB}\).

To escape the constraint from \(\Delta m_K\) directly, we set the \(A^b, \mu\) and angle \(\alpha\) satisfy with \(\cot \alpha \approx m_q^2(\delta^d_{33})_{RL}/\mu m_b\) so that \((\hat{H}^d_{12})_{RL} \approx 0\). In order to obtain the measured value of \(\epsilon'\), the values of relevant parameters are taken as \(m_1 \approx m_3 \approx m_{\tilde{q}}\), \(x_d \approx x_u \approx x_q = 0.3\), \((\delta^u_{13})_{LR} \sim 0.1 m_{\tilde{q}}/500 GeV\), \((\delta^u_{13})_{LL} \sim 0.3 m_{\tilde{q}}/500 GeV\) \([23]\), \((\delta^d_{23})_{LL} \sim 0.4 (m_{\tilde{q}}/500 GeV)^2\), \((\delta^d_{31})_{LL} \sim 4.5 \times 10^{-2} (m_{\tilde{q}}/500 GeV)^2\) \([22]\), \(|A^t|/m_{\tilde{q}} = |A^b|/m_{\tilde{q}} \sim 2\), \(\mu/m_{\tilde{q}} \sim 1\), \((\delta^u_{33})_{RL} \approx A^{(b)} m_b/m_{\tilde{q}}^2, m_q \approx 800\) GeV and \(\tan \beta \sim m_t/m_b\). As mentioned before, the magnitude of squared mass arisen from radiative corrections for the mixing between CP-even and CP-odd pseudoscalar could be order of \((100 GeV)^2\), i.e., \(\hat{\Delta}_{SP}\) could be order of unity. As a consequence, from Eq. (17) and above taken values, the predicted direct CP violating parameter for \(K_L \rightarrow \pi\pi\)
is given by $|Re(e'/e)| \approx 1.83 \times 10^{-3}$ with $m_h = 120$ GeV and $\Delta_{SP} \approx 0.30$. The result is consistent with $(15.3 \pm 3.6) \times 10^{-4}$ and $(20.7 \pm 2.8) \times 10^{-4}$ given by NA48 [3] and KTeV [3], respectively.

Finally, we give the estimation on the electric dipole moments (EDMs) of electron and neutron. In our present considering case, the scalar-pseudoscalar mixing only depends on the complex $(\delta_{33}^u)_{LR}$. That is, even the mixing between CP-even and CP-odd boson is small, it still can introduce CP violating effects, such as EDMs of neutron and lepton. According to the results of [26], the neutral Higgs can contribute to EDMs of neutron and lepton via two-loop topologies. Due to $\Delta_{SP}$ being less than unity, the main effects should be from pseudoscalar exchange. Hence, following the results of [26], the EDM of fermion can be written as

$$
(d_f e)_{\gamma} = Q_f \frac{N_c \alpha_{em} \tan \beta m_f}{32 \pi^3} \xi_t \left[ F\left(\frac{m^2_{\tilde{t}_1}}{m_A^2}\right) - F\left(\frac{m^2_{\tilde{t}_2}}{m_A^2}\right) \right]
$$

where $\tilde{t}_1$ and $\tilde{t}_2$ are the mass eignestates of stop-quark, $Q_f$ denotes the corresponding fermion charge and $\xi_t = Z^u_\beta \mu_t \text{Im}(\delta_{33}^u)_{LR} / \sin \beta \cos \beta v^2$ with $v = \sqrt{v_u^2 + v_d^2}$, the definition of function $F$ can be found in [26]. By using naive valence quark mode and taking $g_s(\Lambda) = 4\pi/\sqrt{6}$, $\alpha_s(M_Z) = 0.12$, $m_u = 7$ MeV, $m_d = 10$ MeV, $m_A = 400$ GeV, the predicted EDMs of electron and neutron with QCD renormalization effects, described by

$$
\frac{d_N}{e} \approx \left( \frac{g_s(M_Z)}{g_s(\Lambda)} \right)^{32/23} \left[ \frac{4}{3} \left( \frac{d_d}{e} \right)_A - \frac{1}{3} \left( \frac{d_u}{e} \right)_A \right],
$$

are around $1.01 \times 10^{-27}$ cm and $3.62 \times 10^{-27}$ cm, respectively. Both are satisfied the current experimental limits given as $d_e/e < 4.3 \times 10^{-27}$ cm [27] and $d_N/e < 6.3 \times 10^{-26}$ cm [28].

In summary, we have studied the effects of scalar-pseudoscalar mixing on $e'/e$ in SUSY models with $\tan \beta \sim m_t/m_b$. We find that if the non-universal soft $\mathcal{A}^d$ term and Yukawa matrix, $Y^d$, are hermitian, and $(\delta_{ij}^d)_{LL} \approx (\delta_{ij}^d)_{RR}$, the predicted value of $e'/e$ is consistent with the measured results of NA48 and KTeV. And also the EDMs are compatible with experimental bounds.

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