ELSEVIER
30 August 2001
PHYSICS LETTERS B
Physics Letters B 515 (2001) 387–394
www.elsevier.com/locate/npe

Chargino contributions to ε and ε'

Shaaban Khalil [a,b], Oleg Lebedev [a]

[a] *Centre for Theoretical Physics, University of Sussex, Brighton BN1 9QJ, UK*
[b] *Ain Shams University, Faculty of Science, Cairo 11566, Egypt*

Received 2 June 2001; received in revised form 5 July 2001; accepted 6 July 2001
Editor: P.V. Landshoff

Abstract

We analyze the chargino contributions to the K–$\overline{K}$ mixing and ε' in the mass insertion approximation and derive the corresponding bounds on the mass insertion parameters. We find that the chargino contributions can significantly enlarge the regions of the parameter space where CP violation can be fully supersymmetric. In principle, the observed values of ε and ε' may be entirely due to the chargino–up-squark loops.

A convenient way to parameterize SUSY contributions to the flavor changing processes is to employ the so called mass insertion approximation [1]. The advantage of this approach is that it allows to treat such contributions in a model independent way without resorting to specific assumptions about the SUSY flavor structures (a technical definition of this approximation will be given below).

The gluino contributions to the kaon observables in the mass insertion approximation have been studied in detail [2–4], but the chargino contributions have not received similar attention. The latter have been considered either in the context of minimal flavor violation, that is in SUSY models with the flavor mixing given by the CKM matrix [5,6] or as contributing to the K–$\overline{K}$ mixing only [7]. In general, the flavor structure in the squark sector may be very complicated. In particular, flavor patterns in the up and down sectors can be entirely different, which may result in the dominance of the chargino contributions to the K and B observables. On the other hand, the neutralino contributions involve the same mass insertions as the gluino ones (i.e., down type mass insertions) and thus cannot qualitatively change the picture. In this Letter, we study the chargino contributions to the K–$\overline{K}$ mixing and ε' using the mass insertion approximation and derive the corresponding bounds on the mass insertion parameters.

Let us first consider the K–$\overline{K}$ mixing. The two observables of primary interest are the K_L–K_S mass difference and indirect CP violation in $K \to \pi\pi$ decays:

$$\Delta M_K = M_{K_L} - M_{K_S}, \qquad \varepsilon = \frac{A(K_L \to \pi\pi)}{A(K_S \to \pi\pi)}. \tag{1}$$

The experimental values for these parameters are $\Delta M_K \simeq 3.489 \times 10^{-15}$ GeV and $\varepsilon \simeq 2.28 \times 10^{-3}$. The Standard Model predictions for them lie in the ballpark of the measured values, however, a precise prediction cannot be made due to the hadronic and CKM uncertainties.

E-mail address: khalil@mail.cern.ch (S. Khalil).

0370-2693/01 © 2001 Elsevier Science B.V. Open access under CC BY license.
PII: S0370-2693(01)00880-2

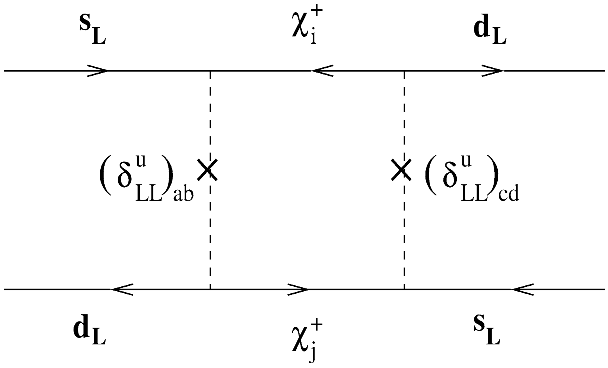

Fig. 1. Leading chargino–up-squark contribution to K–$\overline{K}$ mixing.

Generally, ΔM_K and ε can be calculated via

$$\Delta M_K = 2\,\mathrm{Re}\langle K^0|H_{\mathrm{eff}}^{\Delta S=2}|\overline{K}^0\rangle, \qquad \varepsilon = \frac{1}{\sqrt{2}\Delta M_K}\,\mathrm{Im}\langle K^0|H_{\mathrm{eff}}^{\Delta S=2}|\overline{K}^0\rangle. \tag{2}$$

Here $H_{\mathrm{eff}}^{\Delta S=2}$ is the effective Hamiltonian for the $\Delta S = 2$ transition. It can be expressed via the Operator Product Expansion (OPE) as

$$H_{\mathrm{eff}}^{\Delta S=2} = \sum_i C_i(\mu) Q_i, \tag{3}$$

where $C_i(\mu)$ are the Wilson coefficients and Q_i are the relevant local operators. The main uncertainty in this calculation arises from the matrix elements of Q_i, whereas the Wilson coefficients can be reliably calculated at high energies and evolved down to low energies via the Renormalization Group (RG) running.

In supersymmetric extensions of the SM, the dominant chargino contribution to $H_{\mathrm{eff}}^{\Delta S=2}$ comes from the "super-box" diagram in Fig. 1. We perform our calculations in the super CKM basis, i.e., the basis in which the gluino–quark–squark vertices are flavor-diagonal. In this basis, the chargino–left-quark–left-squark vertices involve the usual CKM matrix:

$$\Delta\mathcal{L} = -g\sum_k\sum_{a,b} V_{k1}K_{ba}^{*}\, d_L^{a\dagger} i\sigma_2 \big(\tilde{\chi}_{kL}^{+}\big)^{*}\tilde{u}_L^{b}, \tag{4}$$

where K is the CKM matrix, a, b are the flavor indices, $k = 1, 2$ labels the chargino mass eigenstates, and V, U are the chargino mixing matrices defined by

$$M_{\chi^+} = \begin{pmatrix} M_2 & \sqrt{2}\,M_W\sin\beta \\ \sqrt{2}\,M_W\cos\beta & \mu \end{pmatrix}, \qquad U^{*}M_{\chi^+}V^{-1} = \mathrm{diag}(m_{\chi_1^+}, m_{\chi_2^+}). \tag{5}$$

Only the gaugino components of the charginos lead to significant contributions to the K–$\overline{K}$ mixing since the higgsino couplings are suppressed by the quark masses (except for the stop coupling) and are not important even at large $\tan\beta$. The stop loop contribution is suppressed by the CKM mixing at the vertices: each vertex involving the stop is suppressed by λ^2 or λ^3 with λ being the Cabibbo mixing, whereas we will be working in $\mathcal{O}(\lambda)$ order. The super-box involving higgsino interactions with the stops depends on the left–right mass insertions and, as will be clear later, does not lead to useful constraints on the SUSY flavor structures.

Due to the gaugino dominance, chargino–squark loops will generate a significant contribution to only one operator

$$Q_1 = \bar{s}_L^{\alpha}\gamma^{\mu}d_L^{\alpha}\bar{s}_L^{\beta}\gamma_{\mu}d_L^{\beta}, \tag{6}$$

similarly to the Standard Model (α, β are the color indices). The corresponding Wilson coefficient is given by the sum of the Standard Model and the chargino contributions:

$$C_1(\mu) = C_1(\mu)^{\text{SM}} + C_1(\mu)^{\tilde{\chi}^+}. \tag{7}$$

Generally, there are additional contributions from gluinos and the Higgs sector, but they are not correlated with the chargino contributions and are unimportant for the present study. We calculate $C_1(\mu)^{\tilde{\chi}^+}$ using the mass insertion approximation. That is, we express the left–left squark propagator as

$$\langle \tilde{u}_L^a \tilde{u}_L^{b*} \rangle = \mathrm{i}\left(k^2\mathbf{1} - m^2\mathbf{1} - \delta m^2\right)^{-1}_{ab} \simeq \frac{\mathrm{i}\,\delta_{ab}}{k^2 - m^2} + \frac{\mathrm{i}\,(\delta m^2)_{ab}}{(k^2 - m^2)^2}, \tag{8}$$

where $\mathbf{1}$ is the unit matrix and m is the average up-squark mass. The SUSY contributions are parameterized in terms of the dimensionless parameters $(\delta^u_{LL})_{ab} \equiv (\delta m^2)_{ab}/m^2$. The corresponding Wilson coefficient is calculated to be

$$C_1(M_W)^{\tilde{\chi}^+} = \frac{g^4}{768\pi^2 m^2}\left(\sum_{a,b} K^*_{a2}\left(\delta^u_{LL}\right)_{ab} K_{b1}\right)^2 \sum_{i,j} |V_{i1}|^2 |V_{j1}|^2 \frac{x_i h(x_i) - x_j h(x_j)}{x_i - x_j}, \tag{9}$$

where $x_i \equiv m^2_{\tilde{\chi}^+_i}/m^2$ and

$$h(x) = \frac{2 + 5x - x^2}{(1-x)^3} + \frac{6x\ln x}{(1-x)^4}. \tag{10}$$

It is interesting to note that "flavor-conserving" mass insertions $(\delta^u_{LL})_{aa}$ contribute to $C_1(M_W)$, unlike for the gluino case. Such mass insertions arise from nondegeneracy of the squark masses and are proportional to the difference of the average squark mass squared and the diagonal matrix elements of the squark mass matrix. If the diagonal elements are equal, the "flavor-conserving" mass insertions drop out of the sum due to the GIM cancellations.

The flavor structure appearing in Eq. (9) can be expanded in powers of λ:

$$\sum_{a,b} K^*_{a2}\left(\delta^u_{LL}\right)_{ab} K_{b1} = \left(\delta^u_{LL}\right)_{21} + \lambda\left[\left(\delta^u_{LL}\right)_{11} - \left(\delta^u_{LL}\right)_{22}\right] + \mathcal{O}\left(\lambda^2\right). \tag{11}$$

Assuming the presence of *one type* of the mass insertions at a time in Eq. (11) at each order in λ, one can derive constraints on $(\delta^u_{LL})_{21}$ and $\delta \equiv (\delta^u_{LL})_{11} - (\delta^u_{LL})_{22}$ imposed by ΔM_K and ε. A much weaker constraint on $(\delta^u_{LL})_{31}$ can also be obtained if we are to keep $\mathcal{O}(\lambda^2)$ terms in Eq. (11).

To derive constraints on the mass insertions, one has to take into account the RG evolution of the Wilson coefficients. In our numerical numerical analysis, we use the NLO QCD result $C_1(\mu)^{\tilde{\chi}^+} \simeq 0.8\, C_1(M_W)^{\tilde{\chi}^+}$ with $\mu = 2$ GeV [4]. The matrix element of Q_1 is computed via $\langle K^0|Q_1|\overline{K}^0\rangle = \frac{1}{3} M_K f_K^2 B_1(\mu)$ with the lattice value $B_1(\mu) = 0.61$ [4]. In addition, the SM contribution should be taken into account. Its detailed discussion can be found in Ref. [8]. In our analysis, we assume a zero CKM phase which corresponds to a conservative bound on the mass insertion. The Wolfenstein parameters are set to $A = 0.847$ and $\rho = 0.4$. The other relevant constants are $M_K = 0.498$ GeV and $f_K = 0.16$ GeV.

The resulting bounds on $(\delta^u_{LL})_{21}$ and δ as functions of M_2 and the average squark mass m are presented in Tables 1 and 2. We find that these bounds are largely insensitive to $\tan\beta$ in the range 3–40 and to μ in the range 200–500 GeV. This can be understood since these parameters do not significantly affect the gaugino components of the charginos and their couplings. Note that δ is real due to the hermiticity of the squark mass matrix and therefore does not contribute to ε. The presented bounds on the real part of $(\delta^u_{LL})_{21}$ are a bit stronger than those derived from the gluino contribution to the D–$\overline{D}$ mixing [3], whereas the imaginary part of $(\delta^u_{LL})_{21}$ is not constrained by any other FCNC processes.

Table 1
Bounds on $\sqrt{|\mathrm{Re}[(\delta^u_{LL})_{21}]^2|}$ from ΔM_K (assuming a zero CKM phase). To obtain the corresponding bounds on δ, these entries are to be multiplied by 4.6. These bounds are largely insensitive to $\tan\beta$ in the range 3–40 and to μ in the range 200–500 GeV

$M_2 \setminus m$	300	500	700	900
150	0.04	0.06	0.08	0.09
250	0.07	0.08	0.09	0.11
350	0.09	0.10	0.11	0.12
450	0.12	0.12	0.13	0.14

Table 2
Bounds on $\sqrt{|\mathrm{Im}[(\delta^u_{LL})_{21}]^2|}$ from ε. These bounds are largely insensitive to $\tan\beta$ in the range 3–40 and to μ in the range 200–500 GeV

$M_2 \setminus m$	300	500	700	900
150	5.3×10^{-3}	7.2×10^{-3}	9.1×10^{-3}	1.1×10^{-2}
250	7.8×10^{-3}	9.2×10^{-3}	1.1×10^{-2}	1.3×10^{-2}
350	1.1×10^{-2}	1.2×10^{-2}	1.3×10^{-2}	1.5×10^{-2}
450	1.5×10^{-2}	1.5×10^{-2}	1.6×10^{-2}	1.7×10^{-2}

In principle, it is possible to constrain the $(\delta^u_{LL})_{31}$ mass insertion as well. At order λ^2, there are two contributions in Eq. (11): from $(\delta^u_{LL})_{31}$ and $(\delta^u_{LL})_{12} + (\delta^u_{LL})_{21}$. Assuming no cancellations between these two terms, the constraints on $(\delta^u_{LL})_{31}$ are obtained by multiplying the bounds in Tables 1 and 2 by $(A\lambda^2)^{-1} \simeq 24$. Clearly, this leaves the real part of $\left[(\delta^u_{LL})_{31}\right]^2$ essentially unconstrained, while the bound on $\sqrt{\mathrm{Im}[(\delta^u_{LL})_{31}]^2}$ is of order 10^{-1}. We note that a similar constraint on $(\delta^u_{LR})_{13}$ can be derived from the higgsino–stop contribution, however, such a constraint is typically satisfied automatically (especially if the squarks are heavy) since $(\delta^u_{LR})_{13} \sim \varepsilon m_t/m$ with $\varepsilon \ll 1$ being the 1–3 mixing the left–right sector.

Next let us consider the chargino contribution to ε' using the same approximations. The ε' parameter is a measure of direct CP violation in $K \to \pi\pi$ decays given by

$$\frac{\varepsilon'}{\varepsilon} = -\frac{\omega}{\sqrt{2}\,|\varepsilon|\,\mathrm{Re}\,A_0}\left(\mathrm{Im}\,A_0 - \frac{1}{\omega}\,\mathrm{Im}\,A_2\right), \tag{12}$$

where $A_{0,2}$ are the amplitudes for the $\Delta I = 1/2, 3/2$ transitions and $\omega \equiv \mathrm{Re}\,A_2/\mathrm{Re}\,A_0 \simeq 1/22$. Experimentally it has been found to be $\mathrm{Re}(\varepsilon'/\varepsilon) \simeq 1.9 \times 10^{-3}$ which provides firm evidence for the existence of direct CP violation. This value can be accommodated in the Standard Model although the theoretical prediction involves large uncertainties.

The effective Hamiltonian for the $\Delta S = 1$ transition is given by

$$H^{\Delta S=1}_{\mathrm{eff}} = \sum_i \widehat{C}_i(\mu)\widehat{Q}_i. \tag{13}$$

Just as in the Standard Model, two operators, $\widehat{Q}_6$ and $\widehat{Q}_8$, play the dominant role. They originate from the gluon and electroweak penguin diagrams (Fig. 2) and are defined by

$$\widehat{Q}_6 = \left(\bar{s}^\alpha d^\beta\right)_{V-A} \sum_{q=u,d,s} \left(\bar{q}^\beta q^\alpha\right)_{V+A}, \qquad \widehat{Q}_8 = \frac{3}{2}\left(\bar{s}^\alpha d^\beta\right)_{V-A} \sum_{q=u,d,s} e_q\left(\bar{q}^\beta q^\alpha\right)_{V+A}, \tag{14}$$

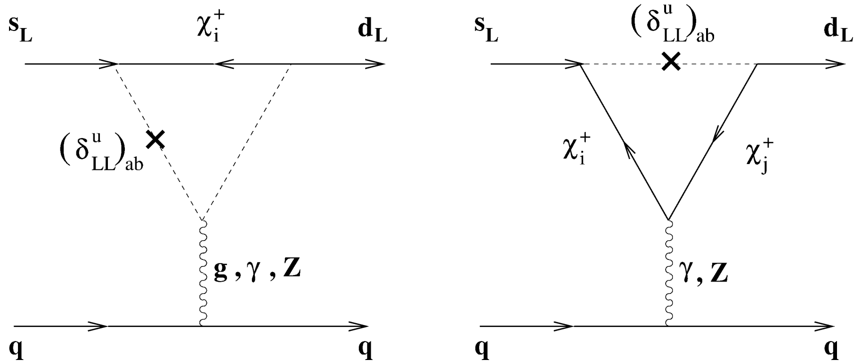

Fig. 2. Leading chargino–up-squark contributions to ε' (a "mirror" diagram is not shown).

with $(\bar{f}f)_{V-A} \equiv \bar{f}\gamma_\mu(1-\gamma_5)f$. Their matrix elements are enhanced by $(m_K/m_s)^2$ compared to those of the other operators:

$$\langle(\pi\pi)_{I=0}|Q_6|K^0\rangle = -4\sqrt{\frac{3}{2}}\left[\frac{m_K}{m_s(\mu)+m_d(\mu)}\right]^2 m_K^2(f_K-f_\pi)B_6,$$
$$\langle(\pi\pi)_{I=2}|Q_8|K^0\rangle = \sqrt{3}\left[\frac{m_K}{m_s(\mu)+m_d(\mu)}\right]^2 m_K^2 f_\pi B_8, \tag{15}$$

where $B_{6,8}$ are the bag parameters. In addition, the contributions of these operators are enhanced by the QCD corrections. Although the Wilson coefficient of $\widehat{Q}_8$ is suppressed by α/α_s compared to that of $\widehat{Q}_6$, its contribution to ε' is enhanced by $1/\omega$ and is significant. In fact, it provides the dominant contribution in our analysis.

The relevant QCD corrections in the context of the MSSM with minimal flavor violation have been studied in Ref. [5] and later, in more detail, in Ref. [6]. To account for a general flavor structure in the mass insertion approximation, only the loop functions of Ref. [6] are to be modified. In our numerical analysis we use the parameterization of Ref. [6] and express the chargino contribution to ε'/ε as

$$\left(\frac{\varepsilon'}{\varepsilon}\right)^{\tilde{\chi}^+} = \mathrm{Im}\left(\sum_{a,b} K^*_{a2}\left(\delta^u_{LL}\right)_{ab}K_{b1}\right)F_{\varepsilon'}, \tag{16}$$

where

$$F_{\varepsilon'} = (P_X+P_Y+P_Z)F_Z + \frac{1}{4}P_Z F_\gamma + P_E F_g. \tag{17}$$

Here we have omitted the box diagram contributions which are negligible [6]. The parameters P_i include the relevant matrix elements and NLO QCD corrections, and are given by $P_X = 0.58$, $P_Y = 0.48$, $P_Z = -7.67$, and $P_E = -0.82$. The quantities F_i are functions of supersymmetric parameters resulting from the gluon, photon, and Z penguin diagrams (Fig. 2) and are calculated in the mass insertion approximation. Explicitly,

$$F_g = 2\frac{m_W^2}{m^2}\sum_i |V_{i1}|^2 f_g(x_i), \qquad F_\gamma = 2\frac{m_W^2}{m^2}\sum_i |V_{i1}|^2 f_\gamma(x_i),$$
$$F_Z = \frac{1}{8} - 2\sum_i |V_{i1}|^2 f_Z^{(1)}\left(\frac{1}{x_i},\frac{1}{x_i}\right) + 2\sum_{i,j} V^*_{j1}V_{i1}\left[U_{i1}U^*_{j1}f_Z^{(2)}(x_j,x_i) - V_{j1}V^*_{i1}f_Z^{(1)}(x_j,x_i)\right], \tag{18}$$

Table 3
Bounds on $\left|\mathrm{Im}(\delta^u_{LL})_{21}\right|$ from ε'. For some parameter values the mass insertions are unconstrained due to the cancellations of different contributions to ε'. These bounds are largely insensitive to $\tan\beta$ in the range 3–40; μ is set to 200 GeV

$M_2 \setminus m$	300	500	700	900
150	0.11	0.11	0.13	0.16
250	0.17	–	0.87	0.64
350	0.12	0.29	0.74	–
450	0.12	0.23	0.42	0.79

where $x_i \equiv m^2_{\tilde\chi_i^+}/m^2$ and the loop functions are given by

$$
\begin{aligned}
f_g(x) &= \frac{1-6x+18x^2-10x^3-3x^4+12x^3\ln x}{18(x-1)^5},\\
f_\gamma(x) &= \frac{22-60x+45x^2-4x^3-3x^4+3(3-9x^2+4x^3)\ln x}{27(x-1)^5},\\
f_Z^{(1)}(x,y) &= \frac{(y-1)\left[(x-1)(x^2-x^2y+xy^2-y^2)+x^2(y-1)\ln x\right]-(x-1)^2y^2\ln y}{16(x-1)^2(y-1)^2(y-x)},\\
f_Z^{(2)}(x,y) &= \sqrt{xy}\,\frac{(y-1)\left[(x-1)(x-y)+x(y-1)\ln x\right]-(x-1)^2y\ln y}{8(x-1)^2(y-1)^2(y-x)}.
\end{aligned}
\tag{19}
$$

As noted in Ref. [6], the dominant contribution typically comes from the Z-penguin diagram, especially if the SUSY particles are heavy. This can be seen as follows. Due to the gauge invariance, the $g\bar s_L d_L$ and $\gamma\bar s_L d_L$ vertices are proportional to the second power of the momentum transfer, i.e., $(q_\mu q_\nu - g_{\mu\nu}q^2)/m^2$. This momentum dependence is cancelled by the gluon (photon) propagator which leads to the suppression factor $1/m^2$ in the final result. On the other hand, the $Z\bar s_L d_L$ vertex exists at $q^2=0$ due to the weak current nonconservation and is momentum-independent to leading order. It is given by a dimensionless function of the ratios of the SUSY particles' masses. The Z propagator then leads to the suppression factor $1/M_Z^2$ which is much milder than $1/m^2$ appearing in the gluon and photon contributions.

The resulting bounds on $\mathrm{Im}(\delta^u_{LL})_{21}$ are presented in Table 3. Note that there is no SM contribution to ε' since we assume a vanishing CKM phase. In the limit of heavy superpartners, these bounds become insensitive to the SUSY mass scale. This occurs due to the dominance of the Z-penguin contribution. Indeed, the contributions of the photon and gluon penguins fall off as $1/(\text{SUSY scale})^2$ as can be seen from Eq. (18). On the other hand, the Z-penguin contribution stays constant in the "decoupling" limit. This may seem to conflict with the intuitive expectation of the decoupling of heavy particles. However, the proper decoupling behaviour is obtained when the flavor violating parameters δm^2_{ab} are kept constant (or if they grow slower than the masses of the superpartners). To put it in a slightly different way, the decoupling should be expected when the mass splittings among the squarks grow slower than the masses themselves.

It is noteworthy that (for universal GUT scale gaugino masses of around 200 GeV) the bounds on $\mathrm{Im}(\delta^u_{LL})_{21}$ are slightly stronger than those on $\mathrm{Im}(\delta^d_{LL})_{21}$ derived from the gluino contribution to ε' [3]. The suppression due to the weaker coupling is compensated by a larger loop function mainly due to the presence of the diagram on the right in Fig. 2.

These results show that to have a chargino-induced ε' would require a relatively large LL mass insertion ($\mathcal{O}(10^{-1})$) which typically violates the constraints from ΔM_K and ε. Yet, it is possible to saturate ε and ε' with the chargino contributions in corners of the parameter space. For instance, taking $M_2 = 450$ GeV and $m = 300$ GeV, ε requires

$$2\,\mathrm{Im}\big(\delta^u_{LL}\big)_{21}\,\mathrm{Re}\big(\delta^u_{LL}\big)_{21} \simeq 2.3\times 10^{-4}. \tag{20}$$

Then, assuming $\mathrm{Im}(\delta^u_{LL})_{21} \simeq 0.12$ to produce ε', $\mathrm{Re}(\delta^u_{LL})_{21}$ has to be 9×10^{-4}. These values are in marginal agreement with the ΔM_K bound:

$$\sqrt{\Big|\big[\mathrm{Re}\big(\delta^u_{LL}\big)_{21}\big]^2 - \big[\mathrm{Im}\big(\delta^u_{LL}\big)_{21}\big]^2\Big|} \leqslant 0.12. \tag{21}$$

The main lesson, however, is that combining the chargino and the gluino contributions can provide fully supersymmetric ε and ε' in considerable regions of the parameter space. [1] For example, only a small ($\mathrm{Im}(\delta^d_{LR})_{21} \sim 10^{-5}$) mass insertion in the down-sector is required to generate the observed value of ε' [3]. Then ε can be entirely due to the mass insertions in the up-sector: $\mathrm{Im}(\delta^u_{LL})_{21} \sim 10^{-3}$ and $\mathrm{Re}(\delta^u_{LL})_{21} \sim 10^{-2}$. Generally, this does not require large SUSY CP-phases and may be accommodated in the framework of approximate CP symmetry [13], which is motivated by the strong EDM bounds (for a review see [14]). Alternatively, the CP-phases can be order one but enter only flavor-off-diagonal quantities which occurs in models with Hermitian flavor structures [15]. Clearly, the regions of the parameter space where CP violation can be fully supersymmetric are significantly enlarged if both the gluino and the chargino contributions are included. Of course, it remains a challenge to build a realistic well-motivated model with all of the required features.

References

[1] L.J. Hall, V.A. Kostelecky, S. Raby, Nucl. Phys. B 267 (1986) 415.
[2] M.J. Duncan, Nucl. Phys. B 221 (1983) 285;
J.F. Donoghue, H.P. Nilles, D. Wyler, Phys. Lett. B 128 (1983) 55;
J.S. Hagelin, S. Kelley, T. Tanaka, Nucl. Phys. B 415 (1994) 293.
[3] F. Gabbiani, E. Gabrielli, A. Masiero, L. Silvestrini, Nucl. Phys. B 477 (1996) 321.
[4] M. Ciuchini et al., JHEP 9810 (1998) 008;
J.A. Bagger, K.T. Matchev, R. Zhang, Phys. Lett. B 412 (1997) 77.
[5] E. Gabrielli, G.F. Giudice, Nucl. Phys. B 433 (1995) 3;
E. Gabrielli, G.F. Giudice, Nucl. Phys. B 507 (1995) 549, Erratum.
[6] A.J. Buras, P. Gambino, M. Gorbahn, S. Jäger, L. Silvestrini, Nucl. Phys. B 592 (2001) 55.
[7] J. Ellis, D.V. Nanopoulos, Phys. Lett. B 110 (1982) 44;
G.C. Branco, G.C. Cho, Y. Kizukuri, N. Oshimo, Phys. Lett. B 337 (1994) 316;
M. Misiak, S. Pokorski, J. Rosiek, hep-ph/9703442;
T. Goto, Y. Okada, Y. Shimizu, Phys. Rev. D 58 (1998) 094006.
[8] A.J. Buras, hep-ph/0101336;
G. Buchalla, hep-ph/0103166.
[9] A. Masiero, H. Murayama, Phys. Rev. Lett. 83 (1999) 907;
S. Khalil, T. Kobayashi, Phys. Lett. B 460 (1999) 341;
M. Brhlik, L. Everett, G.L. Kane, S.F. King, O. Lebedev, Phys. Rev. Lett. 84 (2000) 3041;
S. Baek, J.H. Jang, P. Ko, J.H. Park, Phys. Rev. D 62 (2000) 117701.
[10] S.A. Abel, J.M. Frere, Phys. Rev. D 55 (1997) 1623;
S. Khalil, T. Kobayashi, A. Masiero, Phys. Rev. D 60 (1999) 075003;
A.L. Kagan, M. Neubert, Phys. Rev. Lett. 83 (1999) 4929;
S. Khalil, T. Kobayashi, O. Vives, Nucl. Phys. B 580 (2000) 275;
A.J. Buras, G. Colangelo, G. Isidori, A. Romanino, L. Silvestrini, Nucl. Phys. B 566 (2000) 3.
[11] A. Masiero, O. Vives, Phys. Rev. Lett. 86 (2001) 26.
[12] J.A. Casas, S. Dimopoulos, Phys. Lett. B 387 (1996) 107.

[1] In principle it is possible to saturate both ε' and ε with gluino contributions [9], see also [10]. It was noted in Ref. [11] that the LR mass insertions of the required size may lead to the charge and color breaking minima. However, this is not true in general, i.e., when the Higgs, squark, and slepton mass parameters are unrelated, see Ref. [12].

[13] M. Dine, E. Kramer, Y. Nir, Y. Shadmi, Phys. Rev. D 63 (2001) 116005;
G. Eyal, Y. Nir, Nucl. Phys. B 528 (1998) 21.
[14] S. Abel, S. Khalil, O. Lebedev, hep-ph/0103320, to appear in Nucl. Phys. B.
[15] S. Abel, D. Bailin, S. Khalil, O. Lebedev, Phys. Lett. B 504 (2001) 241.