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Direct CP violation in $K \to 3\pi$ decays induced by SUSY chromomagnetic penguins*

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Abstract

An analysis of the CP violating asymmetry in $K^{\pm} \to (3\pi)^{\pm}$ decays in the Standard Model and, by means of the mass insertion approximation, in a wide class of possible supersymmetric extensions, is presented. We find that the natural order of magnitude for this asymmetry is $\mathcal{O}(10^{-5})$ in both cases. Within supersymmetric models effects as large as $\mathcal{O}(10^{-4})$ are possible, but only in a restricted range of the relevant parameters.

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1. The origin of CP violation is one of the fundamental questions of particle physics and cosmology which remains an open problem to date. The recent measurements of ε'/ε [1] (see also [2]) represent an important step forward in our understanding of this phenomenon, since they have ruled out superweak scenarios. Nonetheless we are still far from a quantitative description of the dynamics which generate the amount of CP violation observed in hadronic processes. Indeed, even within the Standard Model (SM) it is very hard to predict the value of ε'/ε in terms of the parameters of the Cabibbo-Kobayashi-Maskawa (CKM) matrix (see e.g. Refs. [3] and references therein). Given the large theoretical uncertainties affecting the calculation of this quantity, it is very useful to collect additional experimental information about CP violation in $|\Delta S| = 1$ transitions. In this respect charge asymmetries in non-leptonic decays, as the difference in $K^+ \to (3\pi)^+$ and $K^- \to (3\pi)^-$ Dalitz plot distributions [4], represent an interesting class of observables, since they are straight direct CP-violating effects free from $|\Delta S| = 2$ contaminations. Moreover, contrary to ε'/ε, these asymmetries stem from the interference of two $\Delta I = 1/2$ amplitudes and do not necessarily suffer the suppression of $\Delta I = 3/2$ transitions. In spite of these advantages, however, within the SM such observables are expected to be very small, of $\mathcal{O}(10^{-5})$, due to the constraints from ε'/ε and the smallness of final-state interactions [5,6]. A natural question is whether extensions of the SM could enhance these CP-violating asymmetries at such a level that they could be recognized as a clear signal of new physics.

Good candidates to provide new large CP violating effects are the supersymmetric extensions of the SM with generic flavour couplings and minimal particle content. In this framework, among the possible contributions which may be envisaged, it has been recently recognized the importance of the chromomagnetic operator (CMO). Its CP-odd contribution can become large in the presence of misalignment between quark and squark mass matrices, and, without conflict with the experimental determination of the K^0–$\bar{K}^0$ mixing amplitude, it can account for the largest part of the measured ε'/ε [7]–[10]. Actually a non-standard CMO is the only possibility, within this framework, to considerably affect the CP-violating part of $\Delta I = 1/2$ amplitudes without serious fine-tuning problems in $|\Delta S| = 2$ processes [10].

In this paper we investigate the possibility of using the CMO to enhance CP violating effects in $K \to 3\pi$ decays. We work under the assumption that the Wilson coefficient of this operator is mostly due to left-right mixing among down-like squarks. This implies that a large chromomagnetic term is necessarily accompanied by sizeable corrections to ε and $K_L \to \pi^0 e^+ e^-$. We thus perform a combined analysis of these processes together with ε'/ε and $K \to 3\pi$ decays. We have not considered other possible supersymmetric sources of CP violation in $K \to 3\pi$ decays, such as left-left or right-right squark mixing, since only the left-right ones trigger the enhancement of the CMO which is the most promising candidate to give an observable effect.

Our main conclusions are that, even within supersymmetry, effects at the level of $\mathcal{O}(10^{-4})$ may be observed only under special circumstances, among which large cancellations of different contributions in ε'/ε. Otherwise, similar to the SM, the natural order of magnitude of the charge asymmetries will remain of $\mathcal{O}(10^{-5})$. As we shall discuss, this conclusion is rather general and applies also to other direct CP-violating observables in non-leptonic processes.

The paper is organized as follows. We first derive the main formulae needed to evaluate the CMO contribution to CP-odd asymmetries in $K^\pm \to (3\pi)^\pm$ decays. Then we discuss the role of down-type left-right mass insertions, including contributions from the CMO, in the K^0–$\bar{K}^0$ mixing amplitude. Finally a combined analysis of CP-violating effects in $K^\pm \to (3\pi)^\pm$ decays, taking into account the constraints imposed by ε, ε'/ε and $K_L \to \pi^0 e^+ e^-$, is presented. The results are summarized in the conclusions.

2. We start by analyzing the charge asymmetries in $K^\pm \to \pi^\pm\pi^\pm\pi^\mp$ decays. As discussed in Ref. [5,6], the most interesting CP-violating observable is the asymmetry in the Dalitz plot slopes $g_\pm$. Neglecting the suppressed $\Delta I = 3/2$ contributions, this can be written as

$$\frac{g_+ - g_-}{g_+ + g_-} = \left[\frac{\mathrm{Im}b}{\mathrm{Re}b} - \frac{\mathrm{Im}a}{\mathrm{Re}a}\right] \sin(\alpha_0 - \beta_0)\,, \tag{1}$$

where the weak amplitudes a and b are defined by the momentum expansion of $A(K^+ \to \pi^+\pi^+\pi^-)$ around the center of the Dalitz Plot,

$$A(K^+ \to \pi^+\pi^+\pi^-) = a e^{i\alpha_0} + b e^{i\beta_0} Y + \mathcal{O}(Y^2)\,, \tag{2}$$

$$Y = \frac{3(p_K - p_{\pi^-})^2 - M_K^2 - 3M_\pi^2}{M_\pi^2}\,, \tag{3}$$

and $\alpha_0,\ \beta_0$ are the small rescattering phases, known from chiral perturbation theory (ChPT) [11], evaluated at $Y = 0$. In the limit where we neglect $\Delta I = 3/2$ contributions, the slope asymmetries of $K^\pm \to \pi^\pm\pi^\pm\pi^\mp$ and $K^\pm \to \pi^0\pi^0\pi^\pm$ modes are identical.

Since a and b are $\Delta I = 1/2$ amplitudes, on general grounds one expects $\mathrm{Im}a/\mathrm{Re}a$ and $\mathrm{Im}b/\mathrm{Re}b$ to be both of the same order as the weak phase of $A_0 = A(K \to (2\pi)_{I=0})$, namely

$$\frac{\mathrm{Im}a}{\mathrm{Re}a} \sim \frac{\mathrm{Im}b}{\mathrm{Re}b} \sim \frac{\mathrm{Im}A_0}{\mathrm{Re}A_0} \sim \mathrm{Re}\left(\frac{\varepsilon'}{\varepsilon}\right)|\varepsilon|\frac{\mathrm{Re}A_0}{\mathrm{Re}A_2} \sim 10^{-4}\,. \tag{4}$$

Given that $\sin(\alpha_0 - \beta_0) \lesssim 0.1$ [11], this sets the "natural" order of magnitude for the asymmetry to 10^{-5} [5]. Actually, within the SM, the situation is even worse: neglecting the CMO, which in this case has a very small coefficient, the asymmetry vanishes at the lowest order in the chiral expansion. This happens because at this order there is only one octet operator which generate the same weak phase to all the $\Delta I = 1/2$ amplitudes [6]. Clearly the situation may improve if the contribution of the CMO is enhanced by supersymmetric effects, which we now discuss.

Let us start from the relevant piece of the effective Hamiltonian. This can be written as [10]

$$\mathcal{H}_{\rm mag} = C_g^+ Q_g^+ + C_g^- Q_g^- + \text{h.c.}\,, \tag{5}$$

where

$$Q_g^{\pm} = \frac{g}{16\pi^2}\left(\bar{s}_L \sigma^{\mu\nu} t^a G^a_{\mu\nu} d_R \pm \bar{s}_R \sigma^{\mu\nu} t^a G^a_{\mu\nu} d_L\right) \tag{6}$$

and the dominant contribution to Wilson coefficients, generated by gluino exchange diagrams, is given by [10,12]

$$C_g^{\pm}(m_{\tilde{g}}) = \frac{\pi\alpha_s(m_{\tilde{g}})}{m_{\tilde{g}}}\left[\left(\delta^D_{LR}\right)_{21} \pm \left(\delta^D_{LR}\right)^*_{12}\right] G_0(x_{gq})\,. \tag{7}$$

Here $(\delta^D_{LR})_{ij} = (M^2_D)_{i_L j_R}/m^2_{\tilde{q}}$ denote the off-diagonal entries of the (down-type) squark mass matrix in the super-CKM basis [13] and $x_{gq} = m^2_{\tilde{g}}/m^2_{\tilde{q}}$ the ratio of gluino and (average) squark mass squared. The explicit expression of $G_0(x)$ can be found in [10].

The realization of $Q_g^{\pm}$ in terms of meson fields, to the lowest order in $1/N_c$ and in the derivative expansion, can be written as

$$Q_g^{\pm} = \frac{11}{256\pi^2}\frac{f_\pi^2 M_K^2}{m_s + m_d}\left[U D_\mu U^\dagger D^\mu U \pm D_\mu U^\dagger D^\mu U U^\dagger\right]_{23}\,, \tag{8}$$

where $U = \exp(i2\phi/f_\pi)$, $f_\pi = 132$ MeV, ϕ is the octet field of pseudoscalar mesons, and the overall coupling has been fixed by the chiral quark model estimate of Ref. [14]. Using Eq. (8) we can derive the following matrix elements

$$\langle \pi^0(p)|Q_g^+|K^0(p)\rangle = -\frac{11 B_{g1}}{32\sqrt{2}\pi^2}\frac{M_K^2 p^2}{m_s + m_d}\,, \tag{9}$$

$$i\langle \pi^+\pi^-|Q_g^-|K^0\rangle = -\frac{11 B_{g2}}{32\pi^2}\frac{M_K^2 M_\pi^2}{f_\pi(m_s + m_d)}\,, \tag{10}$$

$$\langle \pi^+\pi^+\pi^-|Q_g^+|K^+\rangle = -\frac{11 B_{g3}}{16\pi^2}\frac{M_K^2 M_\pi^2}{f_\pi^2(m_s + m_d)}\,. \tag{11}$$

The B-factors, B_{gi}, have been introduced to parametrize our ignorance of the precise overall coefficient in Eq. (8) and of possible higher-order terms. For practical purposes, in the following we shall also set $m_s + m_d = 110$ MeV in Eqs. (9)–(11), encoding in the B_{gi} the remaining uncertainty on the true value of the quark masses.[1]

For the $K \to 3\pi$ amplitudes defined in (2) we obtain

$$\left|\frac{\mathrm{Im}a}{\mathrm{Re}a}\right| = \frac{3}{2M_K^2 G_8} \times \frac{11}{16\pi^2}\frac{M_\pi^2 M_K^2}{f_\pi^2(m_s+m_d)} \left|B_{g3}\mathrm{Im}C_g^+\right| , \qquad \frac{\mathrm{Im}b}{\mathrm{Re}b} = 0 , \tag{12}$$

where we have used the lowest-order chiral relation between Rea and ReA_0, expressing the latter in terms of the standard coupling $G_8 = 9.1 \times 10^{-6}$ GeV^{-2} (i.e. Rea = $2M_K^2 G_8/3$). In view of the numerical analysis, it is convenient to introduce the following simple expression, which can be readily derived from Eq. (12) and $\mathcal{H}_{\mathrm{mag}}$ in (5)

$$\left|\frac{g_+ - g_-}{g_+ + g_-}\right| \simeq 1.97 \times \left[\frac{\eta\ \alpha_s(m_{\tilde g})}{\alpha_s(500\mathrm{GeV})}\frac{500\mathrm{GeV}}{m_{\tilde g}}\frac{G_0(x_{gq})}{G_0(1)}\right] \left|B_{g3}\mathrm{Im}\delta_+\right| , \tag{13}$$

where we have defined $\delta_\pm = (\delta^D_{LR})_{21} \pm (\delta^D_{LR})^*_{12} = (\delta^D_{LR})_{21} \pm (\delta^D_{RL})_{21}$. We found very useful to introduce $\delta_\pm$ since these are the natural couplings appearing at first order in any parity conserving $(+)$ or parity violating $(-)$ observable. In the evaluation of the numerical coefficient above we have used $\alpha_s(500\mathrm{GeV}) = 0.096$ and, as anticipated, $m_s + m_d = 110$ MeV. The parameter $\eta \simeq 0.9$ [10] is the correcting factor due to the running of the Wilson coefficient from $m_{\tilde g}$ to the operator renormalization scale.

3. An important constraint on the couplings $\delta_\pm$ comes from K^0–$\bar K^0$ mixing. Besides the usual gluino-box amplitudes widely discussed in the literature (see e.g. Refs. [12,16,17] and references therein) further contributions arise from the single and double insertion of the CMO. Schematically we can write

$$\mathcal{A}^{\mathrm{susy}}(K^0 \to \bar K^0) = \mathcal{A}_{\mathrm{boxes}} + \mathcal{A}_{1-\mathrm{mag}} + \mathcal{A}_{2-\mathrm{mag}} . \tag{14}$$

$\mathcal{A}_{\mathrm{boxes}}$, which is dominated by short-distance contributions, can conveniently be written in the form

$$\begin{aligned}\mathcal{A}_{\mathrm{boxes}} \;=\; & \frac{\alpha_S^2}{m_{\tilde g}^2}\frac{1}{432}\left(\frac{M_K}{m_s+m_d}\right)^2 M_K^2 f_K^2 \Big[x_{gq}^2 f_6(x_{gq})\,(85\eta_2 B_2 + 3\eta_3 B_3)\,(\delta_+^2 + \delta_-^2) \\ & + x_{gq}\tilde f_6(x_{gq})\,(33\eta_4 B_4 + 15\eta_5 B_5)\,(\delta_+^2 - \delta_-^2)\Big] ,\end{aligned} \tag{15}$$

where the definitions of $f_6(x)$, $\tilde f_6(x)$ and of the B_i parameters have been taken from Ref. [16]. In the numerical analysis we will use the simplified expression

$$\mathcal{A}_{\mathrm{boxes}} = 2.9 \times 10^{-11}\mathrm{GeV}^2 \times \left[\frac{\alpha_s(m_{\tilde g})}{\alpha_s(500\mathrm{GeV})}\frac{500\mathrm{GeV}}{m_{\tilde g}}\right]^2 \left(B_+\,\delta_+^2\ +\ B_-\,\delta_-^2\right) , \tag{16}$$

[1] For an extensive discussion about the possible chiral realizations of the CMO see Ref. [15]. The factor B_{g3} could in principle be a function of Y, but for simplicity in the following we will ignore this possibility; B_{g2} coincides with the B_G of [10] for $m_s + m_d = 110$ MeV.

where $B_\pm$ are coefficients of $\mathcal{O}(1)$, which may vary by a factor of $2 \div 3$ depending on the values of the B parameters of the $\Delta S = 2$ operators, the precise value of x_{gq}, the perturbative QCD corrections, etc. [16].

The single and double insertions of the CMO are expected to be dominated by long distance contributions. For illustration we give here the expression of the single π^0 contribution, already considered in Ref. [19]

$$\begin{aligned} \mathcal{A}^{\pi^0}_{1-\mathrm{mag}} &= 2\,\langle \bar{K}^0 | \mathcal{H}^{\Delta S=1}_{\mathrm{SM}} | \pi^0 \rangle \frac{1}{M_K^2 - M_\pi^2} \langle \pi^0 | \mathcal{H}_{\mathrm{mag}} | K^0 \rangle \\ &= \frac{1}{M_K^2 - M_\pi^2} \left(\frac{11 B_{g1}}{32\pi} \frac{M_K^6}{m_s + m_d} G_8 f_\pi^2 \right) \frac{\alpha_S}{m_{\tilde{g}}} \eta G_0(x_{gq}) \delta_+ \,. \end{aligned} \tag{17}$$

One can then generalize the above expression to include the contribution from other one-meson states, such as the η and η'. In the following we will use the simplified expression

$$\mathcal{A}_{1-\mathrm{mag}} = 4.8 \times 10^{-13} \mathrm{GeV}^2 \times \left[\frac{\eta\, \alpha_s(m_{\tilde{g}})}{\alpha_s(500\mathrm{GeV})} \frac{500\mathrm{GeV}}{m_{\tilde{g}}} \frac{G_0(x_{gq})}{G_0(1)} \right] \kappa_1\, \delta_+ \,, \tag{18}$$

where the numerical coefficient has been computed from Eq. (17) and we have absorbed the hadronic uncertainties, namely B_{g1} and contributions from intermediate states other than the π^0, in the factor κ_1. Using the Gell-Mann-Okubo mass formula the π^0 and η contributions would cancel, thus one may argue that κ_1 should be substantially smaller than one. We know, however, that a similar argument fails for $K_L \to \gamma\gamma$, where the effective coupling, corresponding to our κ_1, is of $\mathcal{O}(1)$. Note that intermediate states with parity opposite to the one pseudoscalar-meson state may give contributions proportional to δ_-. We have neglected these effects in our analysis.

Similarly, in the case of the double insertion one gets

$$\mathcal{A}^{\pi^0}_{2-\mathrm{mag}} = \langle \bar{K}^0 | \mathcal{H}_{\mathrm{mag}} | \pi^0 \rangle \frac{1}{M_K^2 - M_\pi^2} \langle \pi^0 | \mathcal{H}_{\mathrm{mag}} | K^0 \rangle \,, \tag{19}$$

that proceeding as before leads to

$$\mathcal{A}_{2-\mathrm{mag}} = 1.9 \times 10^{-11} \mathrm{GeV}^2 \times \left[\frac{\eta\, \alpha_s(m_{\tilde{g}})}{\alpha_s(500\mathrm{GeV})} \frac{500\mathrm{GeV}}{m_{\tilde{g}}} \frac{G_0(x_{gq})}{G_0(1)} \right]^2 \kappa_2\, \delta_+^2 \,. \tag{20}$$

4. For the other two quantities which are used in our analysis, namely ε'/ε and the $\mathrm{BR}(K_L \to \pi^0 e^+ e^-)$, rather than giving the explicit analytic expressions, for which we refer the reader to Ref. [10], we only list here two convenient expressions:

$$\mathrm{Re} \left(\frac{\varepsilon'}{\varepsilon} \right)_{\mathrm{mag}} = 92.6 \times \left[\frac{\eta\, \alpha_s(m_{\tilde{g}})}{\alpha_s(500\mathrm{GeV})} \frac{500\mathrm{GeV}}{m_{\tilde{g}}} \frac{G_0(x_{gq})}{G_0(1)} \right] B_{g2}\, \mathrm{Im}\delta_- \,, \tag{21}$$

$$\mathrm{BR}(K_L \to \pi^0 e^+ e^-)_{\mathrm{mag}} = 6.1 \times 10^{-4} \left(\frac{\tilde{y}_\gamma(m_{\tilde{g}}, x_{gq}) G_0(x_{gq})}{\tilde{y}_\gamma(500\mathrm{GeV}, 1) G_0(1)} \right)^2 B_T^2 (\mathrm{Im}\delta_+)^2 \,, \tag{22}$$

where the definitions of B_T and $\tilde{y}_\gamma$ can be found in Ref. [10], and in $K_L \to \pi^0 e^+ e^-$ we have neglected the interference with the SM contribution.

Besides the numerical expressions given above, for our study we also use the following experimental inputs:

$$|\varepsilon| \;=\; (2.28 \pm 0.02) \times 10^{-3} \,, \qquad \mathrm{Re}\left(\frac{\varepsilon'}{\varepsilon}\right) \;=\; (21.2 \pm 4.6) \times 10^{-4} \,, \tag{23}$$

$$\mathrm{BR}(K_L \to \pi^0 e^+ e^-) \;<\; 5.6 \times 10^{-10} \;[18] \,. \tag{24}$$

We are now ready to discuss the bounds in the Imδ_+–Imδ_- plane imposed by the experimental measurements and the theoretical expressions given in Eqs. (16), (18), (20)–(22). As usual, we impose the constraints by requiring that all CP violating observables are saturated by the supersymmetric effects considered above, i.e. neglecting the SM contributions. By setting the values of all hadronic parameters (namely the B_i's and the κ_i's) to one, at our reference values of gluino and squark masses ($m_{\tilde{g}} = 500$ GeV and $x_{gq} = 1$), we find Im$\delta_- \sim$ Im$\delta_+ \sim 10^{-5}$. According to Eq. (13), this implies that the $K \to 3\pi$ asymmetry is of the same order, as found within the SM. It is interesting to note that values of Im$\delta_\pm \sim 10^{-5}$ are consistent with the approximate flavor-symmetry scenario of Ref. [7], where $|\delta_\pm| \lesssim \sin\theta_c m_s/m_{\tilde{q}}$. In this framework one could therefore find a "natural" supersymmetric explanation for both ε and ε'.

To obtain larger values of the charge asymmetry in $K \to 3\pi$ decays one has to relax the bound on Imδ_+. To this purpose we note that ε'/ε put an explicit constraint only on Imδ_- but not on Imδ_+, whereas $\mathrm{BR}(K_L \to \pi^0 e^+ e^-)$ and $\mathcal{A}_{\mathrm{boxes}}$ put upper bounds on Imδ_+ only at the level of 10^{-4}. The strongest limit is set by the contribution of $\mathcal{A}_{1-\mathrm{mag}}$ to ε, since this term is linear in Imδ_+ (the quadratic terms may become competitive only for Reδ/Im$\delta \gg 1$ or if $\kappa_1 \ll \kappa_2$). This contribution is subject to a large uncertainty which is parametrized by κ_1. Therefore one can relax the bound by taking for κ_1 a value sensibly smaller than one, as for example done in Ref. [19]. In Fig. 1, we display the bounds obtained for $\kappa_1 = 1.0$, 0.3 and 0.1. Only in the latter (optimistic) case one may obtain $K \to 3\pi$ asymmetries in the 10^{-4} range. This is very similar to what has been found in Ref. [19] for the CP asymmetries of hyperon decays.

Even accepting values of Im$\delta_+ \sim 10^{-4}$, it remains to be explained the large cancellation between Im$(\delta^D_{LR})_{21}$ and Im$(\delta^D_{RL})_{21}$ necessary to satisfy the 10^{-5} bound on Imδ_- imposed by ε'/ε. An underlying mechanism forcing this cancellation exists, however, in the $U(2)$ models considered in Ref. [9].

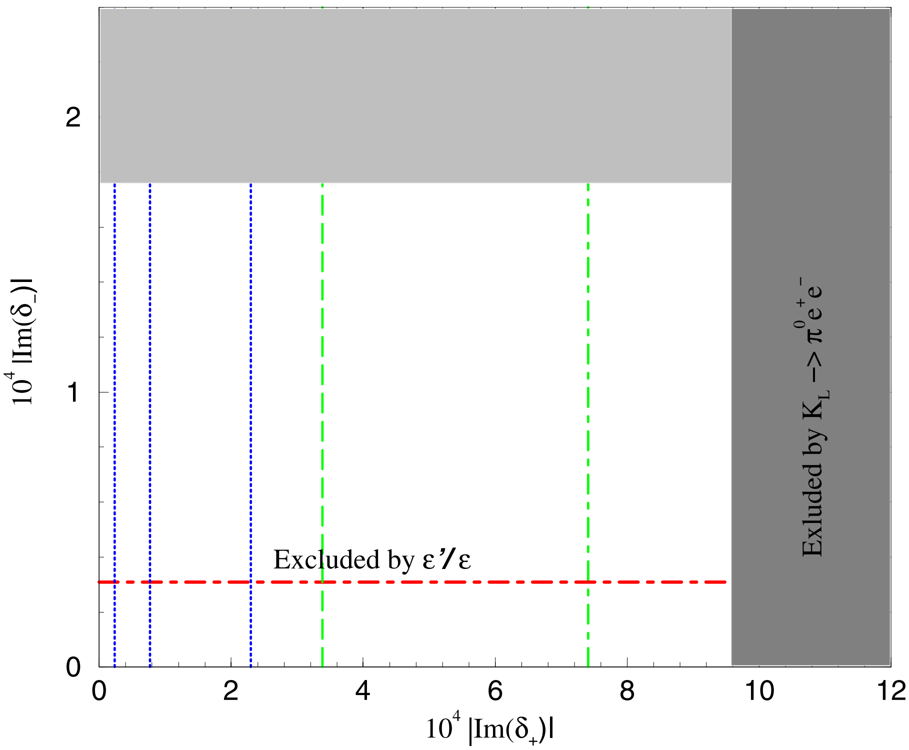

Figure 1: Constraints in the $\mathrm{Im}\delta_+$–$\mathrm{Im}\delta_-$ plane imposed by ε'/ε, ε and $\mathrm{BR}(K_L \to \pi^0 e^+ e^-)$. All bounds are obtained for $x_{gq} = 1$ and scale linearly with $(500\mathrm{GeV}/m_{\tilde{g}})$. The vertical dotted lines correspond to the ε constraint on $\mathcal{A}_{1-\mathrm{mag}}$ for $\kappa_1 = 1.0$, 0.3 and 0.1 (from left to right); the other vertical lines are obtained from the ε bound on $(\mathcal{A}_{\mathrm{boxes}} + \mathcal{A}_{2-\mathrm{mag}})$ for $\delta_- = 0$, $\mathrm{Re}\delta_+ = \mathrm{Im}\delta_+$ and $B_+ = \kappa_2 = 1$ (dashed) or $B_+ = -\kappa_2 = 1$ (dot-dashed). The constraint from $\mathrm{BR}(K_L \to \pi^0 e^+ e^-)$ is obtained for $B_T = 1$. The limits from ε'/ε are obtained for $B_{g2} = 1$ and $(\varepsilon'/\varepsilon)_{\mathrm{SM}} = 0$ (horizontal dash-dotted line) or the extreme case $(\varepsilon'/\varepsilon)_{\mathrm{SM}} \leq -120 \times 10^{-4}$ [20] (horizontal shadowed region).

In principle one could relax the ε'/ε constraint on Imδ_- by allowing other contributions to ε'/ε to be large. For example if one accept the striking result of Ref. [20],[2] $(\varepsilon'/\varepsilon)_{\rm SM} \sim -120 \times 10^{-4}$, one would need a large supersymmetric contribution, corresponding to Im$\delta_- \sim 10^{-4}$ to reconcile the theory with the experimental number. This, however, leads to a new fine tuning problem, because then the natural order of magnitude of ε'/ε would be 10^{-2} rather than 10^{-3}.

5. Large CP violating effects triggered by misalignments between quark and squark mass matrices are among the most promising phenomena to uncover Supersymmetry at low energy. Among the possible effects those driven by the chromomagnetic operator are particularly interesting since they could completely account for the already measured CP violating parameters, ε and ε'/ε. In this paper we have studied the possibility that effects of the CMO are detectable from the enhanced asymmetry in $K \to 3\pi$ decays. We find that this is possible only if several conditions, on which we have a poor theoretical control, conspire in the same direction. Moreover, even if this were the case, one should then face a fine-tuning problem in ε'/ε.

Our analysis of supersymmetric CP-violating effects in $K \to 3\pi$ decays is parallel to those recently performed in $K \to \pi\pi\gamma$ [21] and hyperon decays [19]. In all these cases the conclusions are hampered by poor knowledge of some hadronic parameters, which in the future will hopefully be computed on the lattice. We stress that this problem is absent (or at least much simpler) in rare K decays like $K_L \to \pi^0 \nu\bar{\nu}(e^+e^-)$ [10], whose experimental investigation will definitely provide useful and unambiguous information about the nature of CP violation.

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