Rare kaon decays in SUSY with non-universal $A$ terms

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Abstract

We study the rare kaon decays in the framework of general SUSY models. Unlike the results in the literature, we find the contributions from the gluino exchange to the branching ratio of $K^+ \rightarrow \pi^+ \nu \bar{\nu}$ can reach the central value ($\sim 1.5 \times 10^{-10}$) of the new E787 data while the predicted value of standard model is less than $10^{-10}$. We also find that the same effects also enhance the decays of $K_L \rightarrow \pi^0 \nu \bar{\nu}$, $K_L \rightarrow \pi^0 e^+ e^-$ and $K_L \rightarrow \mu^+ \mu^-$. 
One of the essential reasons for the success of the standard model (SM) is that it naturally satisfies all measured phenomena of flavor changing neutral current (FCNC) in $K$ and $B$ meson systems. The FCNC processes are forbidden at the tree level and suppressed at the loop level by the Glashow-Iliopoulos-Maiani (GIM) mechanism and by the small quark mixing matrix elements which involve the transitions between the third and the first two generations in the Cabibbo-Kobayashi-Maskawa (CKM) matrix. Therefore, those models which have no GIM mechanism and flavor mixing suppression could largely enhance the FCNC decays. Through these decays, one can search for the existence of physics beyond the SM.

Recently, the processes associated with $|\Delta B| = 1$ and $|\Delta K| = 1$ have further progresses in experiments. The representative of the former is the decay modes of $B \to K\ell^+\ell^-$ ($\ell = e, \mu$), observed at the Belle detector in the KEKB $e^+e^-$ storage ring with the branching ratio (BR) of $Br(B \to K\ell^+\ell^-) = (0.75^{+0.25}_{-0.21} \pm 0.09) \times 10^{-6}$, while the SM expectation is around $0.5 \times 10^{-6}$. The latter is $K^+ \to \pi^+\nu\bar{\nu}$. The new data on BR from E787 at the Alternating Gradient Synchrotron (AGS) of Brookhaven National Laboratory is given by

$$Br(K^+ \to \pi^+\nu\bar{\nu}) = (1.57^{+1.75}_{-1.02}) \times 10^{-10},$$

but the predicted BR in the SM is $(0.72 \pm 0.21) \times 10^{-10}$. By comparing the theoretical uncertainties in both processes, one can find that the BRs of $B \to K\ell^+\ell^-$ involve three independent form factors and their errors can not be reduced by the exiting measurements. Therefore, before studying the new physics effects on such processes, it is necessary to know how large the true theoretical uncertainties are. On the contrary, $K^+ \to \pi^+\nu\bar{\nu}$ has less errors from the hadronic matrix element and is free of the long-distance uncertainty. For eliminating the effect of $K^+ \to \pi^+\nu\bar{\nu}$ matrix element, the BR of $K^+ \to \pi^+\nu\bar{\nu}$ can be related to that of $K^+ \to \pi^0e^+\nu \bar{\nu}$ with measured BR of 0.0482 by using the isospin symmetry. Even the corrections to the isospin limit have also been included by Ref. Besides $K^+ \to \pi^+\nu\bar{\nu}$, the relevant decays, such as $K_L \to \mu^+\mu^-$, $K_L \to \pi^0\nu\bar{\nu}$, and $K_L \to \pi^0e^+\nu\bar{\nu}$, have the similar characters; especially the BRs of the last two are related to CP violation (CPV). By means of loop induced effects and less hadronic uncertainties, such kind of rare kaon decays provides good candidates to test the SM. On the other hand, it was pointed out by Ref. that $\beta$ or $\phi_1$, one of the three angles in the CKM matrix, can be described by combining the BRs of $K^+ \to \pi^+\nu\bar{\nu}$ and $K_L \to \pi^0\nu\bar{\nu}$, denoted by $[\sin 2\phi_1]_{\pi\nu\bar{\nu}}$. As well known, this angle is also related to the time-dependent CP asymmetry of the $B \to J/\Psi K_s$ decay, expressed by $[\sin 2\phi_1]_{J/\psi K_s}$. It is obvious $[\sin 2\phi_1]_{\pi\nu\bar{\nu}} = [\sin 2\phi_1]_{J/\psi K_s}$ in the SM. However, once introducing new CP phases, the identity will be changed so that $[\sin 2(\phi_1 + \theta_K)]_{\pi\nu\bar{\nu}} \neq [\sin 2(\phi_1 + \theta_B)]_{J/\psi K_s}$, where $\theta_K$ and $\theta_B$ are the effects of new physics on $K$ and $B$ decays, respectively. Hence, by comparing $\sin 2\phi_1$ measured from rare kaon decays and the asymmetry in the $B$ meson system, we can also tell whether there exist new physics.

Supersymmetric (SUSY) theory not only supplies an elegant mechanism for the breaking of the electroweak symmetry and a solution to the hierarchy problem, but possesses abundant flavor and CP structure. Besides the original CKM matrix, the SUSY models introduce new flavor mixing effects, such as the upper and down type squark mixing matrices. The new CP violating phases in SUSY models can arise from the trilinear and bilinear SUSY soft breaking $A$ and $B$ terms, the $\mu$ parameter for the scalar mixing as well as gaugino masses. Unfortunately, it has been shown that with the universal assumption on the soft breaking parameters, these phases are severely bounded by electric dipole moments (EDMs) so that the contributions to $\epsilon$ and $\epsilon'$ are far below the experimental values. In the literature,
some strategies to escape the constraints of EDMs have been suggested. They are mainly
(a) by setting the squark masses of the first two generations to be as heavy as few TeV
but allowing the third one to be light; (b) by including all possible contributions to
EDMs such that somewhat cancellations occur in some allowed parameter space
and (c) with the non-universal soft $A$ terms instead of universal ones. In particular, those
models with non-universal parameters have been demonstrated that they can be realized in
some string-inspired models. Moreover, without the universal assumption, the corresponding
off-diagonal terms for the left-right mixing of the squark mass matrix are unnecessary to be proportional to
the light quark mass directly. In this paper, we will show the implication of the generalized
$A$ terms on the BRs of $K^+ \rightarrow \pi^+ \nu \bar{\nu}$ and other relevant rare kaon decays.

For simplicity, we adopt the mass-insertion approximation in which the masses of squark
are taken as degenerate approximately. We will only concentrate on the $Z^\mu \tilde{d}_R \gamma_\mu \tilde{P}_L$ effective interactions introduced by Z-penguin diagrams with $P_{L(R)} = (1 \mp \gamma_5)/2$. Although box diagrams also contribute to $s \rightarrow d \nu \bar{\nu}$, one can easily check that the effects compared to those of the SM are suppressed by $M_W^2/M_X^2$ with $M_X$ being the typical SUSY mass scale. Here, we will not discuss them. As to the dipole vertices such as $Z^\mu \tilde{d}_L \gamma_\mu q_s$, because they are suppressed by $q/M_Z$ and by the light quark masses $m_s, m_d$, their contributions are also negligible. Hence, there are two mechanisms to generate effective couplings for $Z^\mu \tilde{d}_R \gamma_\mu \tilde{P}_L$. One is from the penguin of chargino exchange and another one is from the same diagram but with gluino exchange. Comparing to the contributions of chargino and gluino, due to smaller couplings, the effects of neutralino are always negligible.

In the literature it is claimed that due to the magnitude of the relevant mass-insertion
parameters being proportional to the light quark masses or proportional to the
suppressed quark mixing matrix elements, the effects from gluino are much smaller than those of the SM. As a consequence, the dominant one comes from the chargino sector that the mass-insertion terms are associated with stop-quark effects and the corresponding
constraints on the squark mixing matrix elements are loose. However, it is shown by
the new data of E787 that the central value has been close to the predicted value of the
SM. In spite of the still large errors, if the central value hints the existence of new effects,
actually it implies that the models with very large contributions to $K^+ \rightarrow \pi^+ \nu \bar{\nu}$ should be
further constrained. On the contrary, it increases the possibility for the new physics which
is compatible with the SM. Due to the enormous progress in the SUSY flavor physics, what
we want to emphasize is that the contributions from the Z-penguin with gluino exchange are
not negligible and can have sizable effects if the conventional soft SUSY breaking $A$ terms
possess more general flavor structures or are non-universal (nondegenerate). To accomplish
the purpose, in the following analysis, we consider the case that the contributions from
chargino are small.

We start by writing the effective interactions for $Z$ coupling to quarks and squarks as

$$
\mathcal{L} = -i \frac{g}{c_W} \sum_f \left[ \tilde{f}_L \partial_\mu \left( T_3 f - s_W^2 Q_f \right) f_L \right. \\
- \left. \tilde{f}_R \partial_\mu \left( s_W^2 Q_f \right) f_R \right] Z^\mu - \frac{g}{c_W} \sum_f \tilde{f}_\gamma \gamma_\mu \\
\left[ \left( T_3 f - s_W^2 Q_f \right) P_L - s_W^2 Q_f P_R \right] f Z^\mu,
$$

(2)

where $f$ denotes the fermion and $\tilde{f}_{L(R)}$ corresponds to the superpartner of $f$ with the chirality
soft breaking couplings and are written as

\[ \mathcal{M}_D^2 = \begin{pmatrix} (m_D^{2})_{LL} & (m_D^{2})_{LR} \\ (m_D^{2})_{RL} & (m_D^{2})_{RR} \end{pmatrix}, \tag{3} \]

where we have adopted the so-called super-CKM basis that the quarks have been the mass eigenstates so that \( m_D \) is the diagonalized down quark mass matrix, \( \hat{1} \) denotes the \( 3 \times 3 \) unit matrix, the definition of the 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 and down type quarks, respectively, \( \mu \) is the mixing effect of \( \Phi^u \) and \( \Phi^d \), (\( m_D^{2} \))_{LL(RR)} stand for the soft breaking masses for down type squarks and (\( m_D^{2} \))_{LR} describe the trilinear soft breaking couplings and are written as

\[ (m_D^{2})_{LR} = \frac{v_d}{\sqrt{2}} V_{DL} A_{D}^{\alpha} V_{DR}^{\dagger}, \tag{5} \]

where \( V_{DL(R)} \) transform the left(right)-handed quarks from weak eigenstates to mass eigenstates and \( A_{\alpha j}^d = Y_{ij}^d A_{ij}^{d\alpha} \) with \( Y_{ij}^d \) and \( A_{ij}^{d\alpha} \) being Yukawa and soft SUSY breaking matrix, respectively.

Compared to the \( M_W \) scale, due to the smallness of the involved momentum or momentum transfer and the masses of external legs, we drop them in our considerations so that one can easily get the conclusion by the similar situation to the gauge invariant requirement on \( \gamma - s - d \) vertex that the one-loop contributions from the left-left (LL) mixing of the down squark mass matrix are vanished. Therefore, in order to get the effective interaction \( Z - s - d \) that the incoming and outgoing particles carry the same chirality, it needs a double mass-insertion in the squark propagator, \( \nu \), the appearance of \( \sum_j (M_D^{2})_{d2j} A_B (M_D^{2})_{2j1} B_A \) with \( A(B) = L(R) \) or \( R(L) \) and \( j = 1, 2, 3 \) is necessary [20]. In general, we also consider \( \tilde{g} - \tilde{q}_L - q_L \) and \( \tilde{g} - \tilde{q}_R - q_R \) vertices simultaneously. The situation is different from the case of chargino exchange in which only the left-handed couplings give the main contributions.

According to the interactions in Eqs. (2) and (4), by including the self-energy diagrams and all possible emissions of \( Z \)-boson from the propagators of internal squarks, the effective interactions for \( s \to dZ \) can be derived as

\[ \mathcal{L}_{\tilde{g}} = C_Z [\tilde{Z}_L(x) \tilde{d}_{\gamma \mu} P_L s - \tilde{Z}_R(x) \tilde{d}_{\gamma \mu} P_R s] Z^\mu \tag{6} \]

with

\[ C_Z = \frac{G_F e}{\sqrt{2} \pi^2} M_Z^2 \frac{\cos \theta_W}{\sin \theta_W}, \]

\[ \tilde{Z}_A(x) = \frac{\alpha_s \sin^2 \theta_W}{12 \alpha_{em}} C_F G(x) \sum_{j=1,2,3} (\delta_{2j}^d)_{AB} (\delta_{j1}^d)_{BA}, \tag{7} \]
where $C_F = 4/3$, $(\delta^d_{ij})_{AB} \equiv (M^2_{Dij})_{AB}/m_q^2$, $m_q$ is the average mass of squark in the super-CKM basis and

$$G(x) = \frac{2x^2 + 5x - 1}{2(x-1)^3} - \frac{3x^2 \ln x}{(x-1)^2}$$

with $x = m^2_{\tilde{g}}/m_q^2$ and $m_{\tilde{g}}$ being the gluino mass. We note that according to Eq. (3) the effect of the different chirality is opposite in sign each other. As known, the associated matrix elements for relevant kaon decays is $\langle \pi | \bar{d} \gamma^\mu s | K \rangle$ so that the opposite sign actually reflects somewhat cancellation between different chiral couplings. We will see later that in some SUSY models, the cancellation is almost complete.

Altogether, the effective interactions combined with those of the SM can be written as

$$L = C_Z \left[ (X_{SM}(x_t) + \tilde{Z}_L(x)) \bar{d} \gamma^\mu P_L s - \tilde{Z}_R(x) \bar{d} \gamma^\mu P_R s \right] Z^\mu$$

with

$$X_{SM}(x_t) = \lambda_c P_0 + \lambda t X_0(x_t).$$

The explicit expressions of functions $X_0(x_t)$ and $Y_0(x_t)$ can be found in Ref. [23]. $P_0$ is given in Ref. [24]. From the formulas in Ref. [23], the BR is expressed as

$$Br(K^+ \rightarrow \pi^+ \nu \bar{\nu}) = \frac{1}{3} \kappa_+ \left[ (\text{Im}F)^2 + (\text{Re}F)^2 \right]$$

with

$$F = \frac{\lambda_c}{\lambda} P_0 + \frac{\lambda t}{\lambda^5} X_0(x_t) + \frac{1}{\lambda^5} X_{\tilde{g}}(x),$$

$$\kappa_+ = \frac{3\alpha^2 Br(K^+ \rightarrow \pi^0 e^+ \nu)}{2\pi^2 \sin^4 \theta_W} \lambda^8,$$

$$X_{\tilde{g}}(x) = \tilde{Z}_L(x) - \tilde{Z}_R(x),$$

where $x_t = m^2_t/m_W^2$, $\lambda_i = V_{is}^* V_{id}$ with $\lambda_c$ being real to a very high accuracy, and $r_{K^+} = 0.901$ summarizes isospin breaking corrections in relating $K^+ \rightarrow \pi^+ \nu \bar{\nu}$ to $K^+ \rightarrow \pi^0 e^+ \nu$. Here, the isospin breaking corrections from the quark mass effects and electroweak radiative corrections have been calculated in Ref. [8]. By means of Eq. (10), we clearly see that if the $SU(2)_L$ breaking effects for left- and right-handed coupling are the same, the influence on $Br(K^+ \rightarrow \pi^+ \nu \bar{\nu})$ also vanishes. With the same effective interactions, the contributions from the gluino exchange to BRs of $K_L \rightarrow \pi^0 e^+ e^-$, $K_L \rightarrow \pi^0 \nu \bar{\nu}$ and $K_L \rightarrow \mu^+ \mu^-$ can also be described by

$$Br(K_L \rightarrow \pi^0 e^+ e^-)_{\text{dir}} = \kappa \left[ |\text{Im}Y_A|^2 + |\text{Im}Y_V|^2 \right]$$

$$Br(K_L \rightarrow \pi^0 \nu \bar{\nu}) = \frac{1}{3} \kappa_L (\text{Im}F)^2$$

$$Br(K_L \rightarrow \mu^+ \mu^-)_{SD} = \kappa_\mu (\text{Re}D)^2$$

(11)
\[ y_{TA} = \frac{\lambda_t}{\lambda_5} Y_0(x_t) + \frac{1}{\lambda_5} X_\tilde{g}(x), \]
\[ y_{TV} = \frac{\lambda_t}{\lambda_5} [1 + (1 - 4 \sin^2 \theta_W) C_0(x_t)] + \frac{1}{\lambda_5} (1 - 4 \sin^2 \theta_W) X_\tilde{g}(x), \]
\[ D = \frac{\lambda_t}{\lambda_5} Y_0(x_t) + \frac{\tilde{\Delta}}{\lambda_5} + \frac{X_\tilde{g}}{\lambda_5} \]

and

\[ \kappa_L = r_{KL} \frac{\tau_{KL}}{\tau_{K^+}} \frac{3 \alpha^2 B r(K^+ \to \pi^0 e^+ \nu)}{2 \pi^2 \sin^4 \theta_W} \lambda^8 = 1.89 \cdot 10^{-10}, \]

where \( \kappa = \kappa_L / 6, \kappa_\mu = 1.68 \times 10^{-9} \), and \( C_0(x_t) \) and \( \tilde{\Delta}_c \) from the charmed loop can be found in Ref. [23].

The essential question is whether the involving gluino contributions can yield the BR of \( K^+ \to \pi^+ \nu \bar{\nu} \) as large as that given by E787. To explore the possibility, we have to analyze the constraints on the relevant parameters. For simplicity, we just consider the case for \( j=3 \) in Eq. (10) so that the involving mass-insertion parameters are only \( (\delta^d_{23})_{AB} \) and \( (\delta^d_{31})_{AB} \) although this assumption is unnecessary. To estimate the CP violating effects, we take that the real and imaginary parts of relevant parameters are approximately the same in order of magnitude, such as \( |\text{Im}(\delta^d_{ij})_{AB}| \approx |\text{Re}(\delta^d_{ij})_{AB}| / 2 \). Under our assumption, it is known immediately that the bounds on the parameters are from \( Br(B \to X_s \gamma) \) and the \( B_d - \bar{B}_d \) mixing. The former constrains \( (\delta^d_{23})_{AB} \) while the latter is \( (\delta^d_{31})_{AB} \). In addition, it is worth mentioning that charge and color breaking (CCB) minima and the potential unbounded from below (UFB) may also give strict bounds [25]. To relax the constraints from the vacuum instability, we adopt the following two strategies: (a) According to the result in Ref. [23], the model independent upper bounds on \( (\delta^d_{31})_{LR} \) can be described by \( m_\tilde{q} \sqrt{2 + m_T^2 / m_\tilde{q}^2} \). By taking the slepton mass to be few TeV, the constraints are compatible with those directly obtained by FCNC decays. (b) Our universe is resting a false vacuum. From the analysis in Ref. [24], the constraints from the conditions of CCB minima and UFB can be relaxed if the life time of the metastable vacuum is as long as that of the present age of the universe. Although the requirements of CCB and UFB in [25] are necessary, after all, they are not sufficient. Hence, in our numerical calculations, we use the constraints gotten from FCNC processes.

Since the relevant constraints have been considered in Ref. [27], we display the bounds in Table 1 for fixing the specific chirality. The values in the entries of Table 1 are for \( m_\tilde{q} = 500 \) GeV. For the different choices, the values for the second column need to be multiplied by \( m_\tilde{q} / 500 \) but it is \( (m_\tilde{q} / 500)^2 \) for the third column. The parameters that all R(L) are replaced by L(R) depend on the details of SUSY models. As illustrations, we show three possible situations in SUSY models:

- Scenario I: \( \tilde{Z}_L \simeq \tilde{Z}_R \). This approximation is equivalent to \( (\delta^d_{ij})_{RL} \simeq (\delta^d_{ij})_{RL} \). It has been shown recently that if the Yukawa and soft-breaking \( A \) matrices are hermitian, the identity could be realized [28, 29, 30]. As a result of Eq. (10), the contributions to rare kaon decays all vanish. Nevertheless, this scenario implies that the hyperon CP asymmetry in SUSY is one order of magnitude larger than that of the SM [31].
In order to deal with the small $\epsilon$ in scenario I, another shortcoming is that the predicted $\epsilon$ is also far smaller than that of the experimental measurement \cite{29}. In order to deal with the small $\epsilon$ problem and escape the constraint from the EDM, the asymmetric soft-breaking $A$ matrix is proposed by Ref. \cite{32} and it is also found that such kind of the asymmetric property could be realized in some string-inspired supergravity models. Hence, the situation of scenario II can be reached if the $A$ matrix is asymmetric. According to the results of Eqs. (9) and (11) and the constraints in Table 1, the BRs of relevant rare kaon decays are found as follows:

\begin{align}
Br(K^+ \rightarrow \pi^+\nu\bar{\nu}) & = 1.55 \times 10^{-10}, \\
Br(K_L \rightarrow \pi^0\nu\bar{\nu}) & = 2.50 \times 10^{-10}, \\
Br(K_L \rightarrow \pi^0e^+e^-)_{dir} & = 3.47 \times 10^{-11}, \\
Br(K_L \rightarrow \mu^+\mu^-)_{SD} & = 1.85 \times 10^{-9}.
\end{align}

(12)

Because our purpose is to demonstrate that the gluino contributions from penguin diagrams to rare kaon decays could reach the current experimental ranges, the values in Eq. (12) are obtained by setting $m_{\tilde{q}} = 640$ GeV with $x = 0.3$ and the sign of $\tilde{Z}_{L(R)}$ is the same as that of SM. And also we choose the value of $\lambda_i$ such that the results of the SM for $K^+ \rightarrow \pi^+\nu\bar{\nu}$, $K_L \rightarrow \pi^0\nu\bar{\nu}$, $K_L \rightarrow \pi^0e^+e^-$, and $K_L \rightarrow \mu^+\mu^-$ are $0.58 \times 10^{-10}$, $0.46 \times 10^{-10}$, $0.7 \times 10^{-11}$ and $5.26 \times 10^{-10}$, respectively.

- Scenario II: $|\tilde{Z}_R - \tilde{Z}_L| \sim |\tilde{Z}_{L(R)}|$. Besides the negligible contributions to rare K decays in scenario I, another shortcoming is that the predicted $\epsilon$ is also far smaller than that of the experimental measurement \cite{29}. In order to deal with the small $\epsilon$ problem and escape the constraint from the EDM, the asymmetric soft-breaking $A$ matrix is proposed by Ref. \cite{32} and it is also found that such kind of the asymmetric property could be realized in some string-inspired supergravity models. Hence, the situation of scenario II can be reached if the SUSY soft-breaking $A$ matrix is asymmetric. According to the results of Eqs. (9) and (11) and the constraints in Table 1, the BRs of relevant rare kaon decays are found as follows:

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Br(K^+ \rightarrow \pi^+\nu\bar{\nu}) & = 1.55 \times 10^{-10}, \\
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- Scenario III: $|\tilde{Z}_R| \ll |\tilde{Z}_L|$ or vice versa. In this situation, as expected, the numerical predictions should be similar to the case of scenario II.

Finally, we give a brief discussion on chargino contributions. According to the analysis of Ref. \cite{29}, one can understand that unlike gluion case, the main effects of chargino on $Zd$s effective interaction come from the wino component where only left-handed couplings are involved. That is, only $(\delta_{2j}^u)_{LR}(\delta_{1j}^u)_{RL}$ have the significant contributions. However, if the trilinear soft breaking $A_{ij}^u$ and $A_{ij}^d$ parameters arise from the same origin and have the same order of magnitude, due to the smaller weak couplings, the effects of chargino could be much smaller than those of gluino if both superparticles have the same masses. On the other hand, if we allow that $A^u$ is different from $A^d$ entirely, in order to guarantee that the effects of chargino are negligible, the mass of lightest chargino is taken as heavy as $O(\text{TeV})$ so that although the constraints of relevant mass insertion parameters are not strict, all contributions from chargino will become insignificant.

In summary, we have studied the rare kaon decays in the framework of general SUSY models. In terms of the rich flavor structure, we find that the effects of the gluino exchange can make $Br(K^+ \rightarrow \pi^+\nu\bar{\nu})$ up to the central value in the new BNL-E787 result. With the same effects, the BRs of the short distance contributions to $K_L \rightarrow \pi^0\nu\bar{\nu}$, $K_L \rightarrow \pi^0e^+e^-$ and $K_L \rightarrow \mu^+\mu^-$ are larger than those in the SM.
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