GLUINO INDUCED EFFECTS IN STRANGENESS-CHANGING NEUTRAL CURRENTS

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A well-known feature of supersymmetry is that there are flavour-changing strong interactions mediated by gluinos. We calculate the effective $Z\bar{d}s$ coupling due to this interaction taking into account the constraints on the relevant parameters from *CP*-violation data. The consequences on the branching ratio for $K_L \to \mu\bar{\mu}$ and the *CP*-violating muon polarisation asymmetry are examined. The effects are found to be rather small and within experimentally acceptable limits.

Measurements of the bottom-quark lifetime [1] and the ratio $\bar{R} = \Gamma(b \to u e \bar{\nu})/\Gamma(b \to c e \bar{\nu})$ [2] put stringent bounds on the Kobayashi–Maskawa (KM) quark mixing angles. This, along with the possibility of a top-quark mass $(m_t) \sim 40$ GeV, may make it difficult for the KM model to accommodate the experimentally measured value of the *CP*-violation parameter ϵ [4,5]. Furthermore, the new measured values [6] of the ratio ϵ'/ϵ are smaller than the predictions of the standard $SU(3)_C \times SU(2)_L \times U(1)$ model [5]. It appears that if τ_B is larger than 1 ps with a not too large value of $\bar{R}$ (~ 0.03), then one needs to go beyond the standard model for simultaneous satisfactory explanations of the *CP*-violation parameters.

One alternative possibility that has been explored is the supersymmetric (SUSY) standard model. In this model the gluino ($\tilde{g}$) can change flavour through the gluino–squark–quark ($\tilde{g}$–$\tilde{Q}$–Q) coupling because of the mismatch between matrices that diagonalise the squark and quark mass matrices [7,8]. This new feature of supersymmetric models – strong flavour violation – is reflected in additional one-loop contributions to strangeness-changing neutral processes. In particular it has been shown [8] that with this extra SUSY contribution, τ_B as large as 2.4 ps and $\bar{R} = 0.03$ can be reconciled with the experimental values of ϵ and ϵ'/ϵ for $m_t = 40$ GeV.

It is now natural to look for additional signatures of this strong flavour violation in SUSY models vis-a-vis W-induced one-loop flavour violation in the standard model. It has been recently claimed [9], for example, that the gluino induced contribution to the rare decay of the Z-boson to heavy quarks, say, to $Z \mapsto t\bar{c}$, is considerably higher than in the standard model. Another such process is $K_L \to \mu\bar{\mu}$ which involves the effective $Z\bar{d}s$ vertex. It is highly suppressed [10], and the contribution from the standard model itself is of the same order as the experimental upper bound to the branching ratio. Thus an enhanced SUSY contribution to flavour-changing neutral currents in light quark systems would put the supersymmetric model itself in jeopardy. Our purpose in this note is to evaluate the contribution from SUSY graphs to the effective $Z\bar{d}s$ coupling arising out of flavour violation at the $\tilde{g}$–$\tilde{D}$–D vertices, where D represents the three charge-1/3 quarks and $\tilde{D}$ the six corresponding (left and right type) squarks (fig. 1). These contributions depend on, among other things, the elements of the squark mass matrix and rather sensitively on the KM mixing angles which we fix for every τ_B and $\bar{R}$ by demanding that the calculated value of ϵ agrees with experiment. We find that this yields a considerably smaller contribution from the SUSY piece than from the standard model and thus in this case SUSY is saved from being contradicted by experiments.

The SUSY interaction lagrangians that we use are

$$\mathscr{L}_{D\text{-}\tilde{D}\text{-}\tilde{g}} = i\sqrt{2}\, g_s \tilde{g}_\alpha \{\tilde{D}_i^+ (T^\alpha)[\Gamma_L^{ij}\, \tfrac{1}{2}(1-\gamma_5) + \Gamma_R^{ij}\, \tfrac{1}{2}(1+\gamma_5)] D_j\}, \quad (1)$$

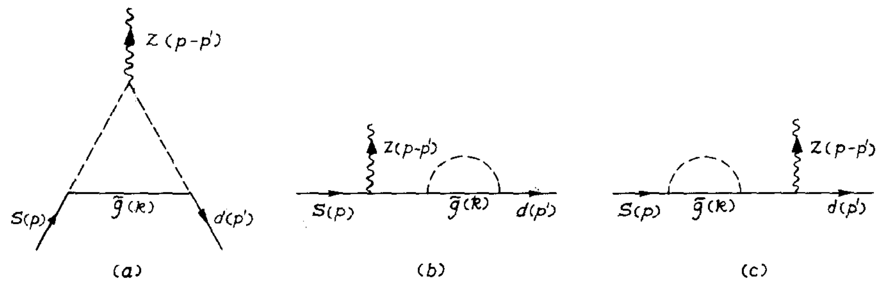

Fig. 1. The lowest order SUSY graphs that contribute to the effective Z$\bar{\rm d}$s coupling. $\tilde{\rm D}_i$ represents the ith squark (mass M_i) and $\tilde{\rm g}$ the gluino (mass $M_{\tilde{\rm g}}$).

$$\mathscr{L}_{\tilde{\rm D}\text{-}\tilde{\rm D}\text{-}{\rm Z}} = g_{\rm W}(p_1+p_2)_\mu Z^\mu \{\tilde{\rm D}_i^+ S_{ii'} \tilde{\rm D}_{i'}\}. \tag{2}$$

The flavour indices i, i' run from 1 to 6 (to take into account left and right type squarks corresponding to each flavour), and j runs from 1 to 3. T^α are the (3×3) SU$(3)_c$ generators where α runs from 1 to 8. The colour indices on the quarks and squarks have been suppressed, g_s (g_W) is the strong (weak) coupling constant and p_1 and p_2 are, respectively, the momenta of the incoming and outgoing squarks in the $\tilde{\rm D}$–$\tilde{\rm D}$–Z vertex. $\Gamma_{\rm L}$ and $\Gamma_{\rm R}$ are (6×3) matrices and S is a (6×6) matrix given by

$$\Gamma_{\rm L} = \tilde{U}_0^+ \begin{pmatrix} \mathbb{1} \\ 0 \end{pmatrix}, \quad \Gamma_{\rm R} = \tilde{U}_0^+ \begin{pmatrix} 0 \\ \mathbb{1} \end{pmatrix}, \quad S = \tilde{U}_0^+ \begin{pmatrix} a\mathbb{1} & 0 \\ 0 & b\mathbb{1} \end{pmatrix} \tilde{U}_0, \tag{3}$$

where $\mathbb{1}$ is the (3×3) unit matrix and $\tilde{U}_0$ is a (6×6) unitary matrix which diagonalises the $\tilde{\rm D}$-mass matrix in a basis in which the D-mass matrix is already diagonal. The constants a and b can be obtained from the D-quark coupling to the Z in the standard model:

$$a = (2\sin^2\theta_{\rm w}/3 - 1)/\cos\theta_{\rm w}\sin 2\theta_{\rm w}, \quad b = 2\sin^2\theta_{\rm w}/3\cos\theta_{\rm w}\sin 2\theta_{\rm w}. \tag{4}$$

Evaluation of the contributions from the three diagrams in fig. 1 is straightforward. The results we present are accurate to terms quadratic in the quark masses and the momentum transfer $q=(p-p')$ where $q\approx p$, p' and $q^2 \ll m_Z^2$ (which is valid in $K_L\to\mu\bar{\mu}$ but not in Z-decay). We write the contribution in the form

$$Z^\mu \bar{\rm d}(p')(A_\mu^{\rm a}+A_\mu^{\rm b}+A_\mu^{\rm c}){\rm s}(p),$$

with

$$A_\mu^{\rm a} = a_1\gamma_\mu + a_2\gamma_\mu\gamma_5 + a_3\bar{q}^2\gamma_\mu + a_4\bar{q}^2\gamma_\mu\gamma_5 + a_5\bar{q}_\mu + a_6\bar{q}_\mu\gamma_5 + ia_7\sigma_{\mu\nu}\bar{q}^\nu + ia_8\sigma_{\mu\nu}\bar{q}^\nu\gamma_5, \tag{5}$$

$$(A_\mu^{\rm b}+A_\mu^{\rm c}) = b_1\gamma_\mu + b_2\gamma_\mu\gamma_5, \tag{6}$$

where

$$\bar{q}=q/M_{\tilde{\rm g}}, \quad x_{\rm s}=m_{\rm s}/M_{\tilde{\rm g}}, \quad x_{\rm d}=m_{\rm d}/M_{\tilde{\rm g}}$$

and

$$a_1 = K\sum_{i,i'} \{O_1^{ii'}[(x_{\rm d}+x_{\rm s})I_1^{ii'} - I_5^{ii'}] + O_3^{ii'}(x_{\rm d}+x_{\rm s})I_7^{ii'}\},$$

$$a_2 = a_1(O_1\to -O_2, O_3\to -O_4, x_{\rm d}\to -x_{\rm d}, I_1\to I_3),$$

$$a_3 = K\sum_{i,i'} O_1^{ii'} I_6^{ii'}, \quad a_4 = a_3(O_1\to -O_2), \quad a_5 = K\sum_{i,i'} (O_1^{ii'} I_2^{ii'} + O_3^{ii'} I_8^{ii'}),$$

$$a_6 = a_5(O_1 \to O_2, O_3 \to O_4, I_2 \to I_4), \quad a_7 = a_5(I_2 \to I_1, I_8 \to I_7), \quad a_8 = a_7(O_1 \to O_2, O_3 \to O_4, I_1 \to I_3),$$

$$b_1 = \frac{K}{2} \sum_i ((a+b)\{[J_1^i - (x_s^2 + x_s x_d + x_d^2) J_2^i] O_1^i - (x_d + x_s) J_3^i O_3^i\}$$

$$+\{(b \to -b), (x_d \to -x_d), (O_1 \to O_2), (O_3 \to O_4)\}),$$

$$b_2 = b_1(a \to -a),$$

$$K = 2\alpha_s g_W / 3\pi,$$

$$O_{1,2}^{ii'} = \tfrac{1}{2}[(\Gamma_L^+)_{di'} S_{i'i} (\Gamma_L)_{is} \pm (L \to R)], \quad O_{3,4}^{ii'} = \tfrac{1}{2}[(\Gamma_L^+)_{di'} S_{i'i} (\Gamma_R)_{is} \pm (L \to R)], \tag{7,8}$$

$$O_{1,2}^{i} = \tfrac{1}{2}[(\Gamma_L^+)_{di} (\Gamma_L)_{is} \pm (L \to R)] = \tfrac{1}{2}[(\Gamma_L^+)_{di} (\Gamma_R)_{is} \pm (L \to R)]. \tag{9}$$

The integrals $I_\alpha^{ii'}$ ($\alpha = 1$–8) and J_α^i ($\alpha = 1$–3) are functions of m_s, m_d, the squark masses (M_i) and the gluino mass ($M_{\tilde{g}}$) and are presented in the appendix. Because of the unitarity of $\tilde{U}_0$ and from eq. (3) it is clear that $\sum_{ii'} O_\alpha^{ii'} = 0$ and $\sum_i O_\alpha^i = 0$. Thus the contribution from any (i, i')-independent part of the integrals I and J will vanish – this is the super-GIM mechanism. The divergences in each individual diagram are independent of i and i' and hence cancel out in this way.

As a check on our calculation it is easy to obtain the similar $\gamma\bar{d}s$ effective vertex by replacing g_W by e and choosing $a = b = 1$. In this limit the combination $(A_\mu^a + A_\mu^b + A_\mu^c)$ has a gauge invariant form with only terms proportional to $(q^2\gamma_\mu - \not{q} q_\mu)$, $(q^2\gamma_\mu - \not{q} q_\mu)\gamma_5$, $\sigma_{\mu\nu} q^\nu$ and $\sigma_{\mu\nu} q^\nu \gamma_5$. Thus current conservation is manifest.

In our numerical evaluations we choose for the $\tilde{D}$-mass matrix the standard form used, for example, in refs. [7,8]. Flavour violation in this case is essentially controlled by the parameter c which is a measure of the contribution of up-quark Yukawa coupling to the renormalisation of the $\tilde{D}$-mass. For any τ_B and $\bar{R}$, by demanding that the predicted value of ϵ be in agreement with the measured number one can set a lower bound on it, viz. c_{min}, fixing at the same time the KM matrix. For values of τ_B and $\bar{R}$ at which the non-SUSY standard model is viable, this bound is zero. A non-zero c_{min} on the other hand indicates that the standard model needs contributions from SUSY graphs to fit ϵ. We examine both light (5 GeV) and heavy (40 and 70 GeV) gluino possibilities, and for each of these, with different values of τ_B ($\bar{R} = 0.03$, $B = 0.33$), we use the KM-matrix elements and c_{min} obtained by fitting the ϵ-parameter [11].

Our results are presented in table 1. To a very good approximation the γ_μ and $\gamma_\mu\gamma_5$ couplings are of equal strength; in fact for every operator O the coupling is almost of the same magnitude as that of $O\gamma_5$. This can be understood by noting that for our choice of the $\tilde{D}$-mass matrix the left–right mixing is small and Γ_R is approximately flavour conserving. Hence $O_1 \approx O_2$ and $O_3 \approx O_4$. Moreover, the $m_d \to -m_d$ exchange which, when implemented in I_1, takes it to I_3 (see Appendix), makes only a marginal difference since terms involving m_s are dominant. The γ_μ and $\gamma_\mu\gamma_5$ terms are larger than the others by at least one order of magnitude. For comparison, the standard model (non-SUSY graph) contributions [12] are calculated. This contribution is also of (V−A) type, whose real part is $(0.31 \pm 0.03) \times 10^{-6}$ and imaginary part is $(0.44 \pm 0.15) \times 10^{-7}$ for 1 ps $\leq \tau_B \leq$ 2 ps. For this calculation one uses the KM matrix elements which yield the measured value of ϵ. For $\tau_B \geq$ 1 ps one requires in addition to the standard model a contribution from SUSY to fit ϵ, and thus supersymmetry is indirectly responsible for the standard model contribution in this range. For $\tau_B \leq$ 1 ps the non-SUSY standard model alone suffices to fit ϵ, and the $Z\bar{d}s$ contributions found therefrom are of the same order as above. In all cases, the SUSY contributions are smaller by at least two orders of magnitude than the non-SUSY ones. This can be traced to a GIM-type mechanism which is operative since the d-type squarks have nearly degenerate masses. From experimental values [10] of the upper bound to $B(K_L \to \mu\bar{\mu})$ one requires the coupling constant of the

Table 1
Four of the effective $Z\bar{d}s$ coupling constants (see eqs. (5) and (6)) as a function of τ_B.

$M_{\tilde{g}}$ (GeV)	τ_B (ps)	(a_1+b_1)		a_3		a_5		a_7	
		Re ($\times 10^{-10}$)	Im ($\times 10^{-10}$)	Re ($\times 10^{-10}$)	Im ($\times 10^{-11}$)	Re ($\times 10^{-12}$)	Im ($\times 10^{-13}$)	Re ($\times 10^{-11}$)	Im (
5	1.0	0.57	0.3	0.16	0.7	1.4	5.7	−0.09	−0.3
	1.2	1.0	0.5	0.21	0.8	1.9	7.1	−0.11	−0.4
	1.4	1.4	0.7	0.24	0.9	2.1	7.9	−0.13	−0.4
	1.6	1.9	1.0	0.27	1.0	2.4	8.6	−0.14	−0.5
	1.8	2.2	1.1	0.29	1.0	2.5	8.8	−0.15	−0.5
	2.0	2.5	1.3	0.30	1.0	2.6	9.0	−0.16	−0.5
40	1.0	1.9	0.8	0.34	1.4	−6.3	-0.29×10^2	0.8	3.8
	1.2	4.0	2.0	0.49	2.0	−8.8	-0.38×10^2	1.1	5.2
	1.4	6.2	3.4	0.60	2.4	−10.0	-0.41×10^2	1.4	5.6
	1.6	8.2	4.3	0.66	2.6	−11.0	-0.44×10^2	1.5	5.8
	1.8	12.0	6.1	0.74	2.8	−12.0	-0.47×10^2	1.7	6.3
	2.0	14.0	7.0	0.79	3.0	−13.0	-0.49×10^2	1.8	6.7
70	1.0	0.29	0.2	0.17	0.7	−0.11	−0.35	0.08	0.3
	1.2	0.81	0.5	0.24	1.0	−0.06	−0.03	0.11	0.4
	1.4	1.3	0.7	0.29	1.2	+0.04	0.49	0.11	0.6
	1.6	2.0	1.0	0.35	1.5	0.18	1.3	0.17	0.7
	1.8	2.7	1.4	0.40	1.7	0.39	2.4	0.19	0.8
	2.0	3.5	1.8	0.46	2.0	0.71	4.0	0.22	0.9

(V−A) interaction to be $\sim 0.2\times 10^{-6}$. Thus in our scheme the non-SUSY graphs alone contribute sufficiently to approach the upper bound. Hence to save SUSY from being contradicted by existing experiments it is necessary that the graphs in fig. 1 should produce relatively small effects, as opposed to what is claimed [9] for the corresponding diagrams with heavy quarks. This is precisely what we have obtained.

The actual contributions from the SUSY graphs depend sensitively on the flavour-violation parameter c. As τ_B increases, the standard model needs to be supplemented by SUSY more and more; one requires larger and larger values of c and consequently the contributions from eqs. (5) and (6) increase. For a fixed squark mass, a heavier gluino requires a larger c and the SUSY contribution increases (cf. $M_{\tilde{g}}=5$ and 40 GeV in table 1). For $M_{\tilde{g}}=70$ GeV the corresponding squark mass must also be larger [13]. This in turn reduces the SUSY contribution [‡1].

Another quantity which has recently been studied [15,16] is the muon polarisation asymmetry in $K_L\to\mu\bar{\mu}$, which is a measure of *CP*-violation in this process. This is characterised by a number $\mathscr{P}=(N_R-N_L)/(N_R+N_L)$ where N_R (N_L) is the number of muons with positive (negative) helicity emerging from the decay. It is easy to see [14] that for an effective (V−A) type interaction, $\mathscr{P}$ is vanishing. In our results the dominant term is of (V−A) form; hence we note, in passing, that this vertex does not significantly contribute to the muon polarisation asymmetry in $K_L\to\mu\bar{\mu}$.

$Z\bar{q}_1q_2$ couplings, in general, will give rise to rare decay modes for the Z which may be testable at the SLC or LEP colliders. The possibility of testing SUSY through such rare decays to *heavy* quarks has already been considered [9]. A careful calculation of this effect including *CP*-violation constraints on quark and squark mixing is in progress [17].

‡1 Recently other calculations of the process $K_L\to\mu\bar{\mu}$ have appeared in the literature [14]. The essential conclusions are similar.

In conclusion, we have studied SUSY effects on $Z\bar{d}s$ coupling taking the mixing parameters in agreement with *CP*-violation data. We find that the resultant contribution to the $K_L \to \mu\bar{\mu}$ branching ratio and the muon polarisation asymmetry are still within the stringent bounds imposed by experiment.

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Appendix. Here we list the integrals used in eqs. (5) and (6). In the notation

$$\Delta = M_i^2 - M_{i'}^2, \quad x_i = M_i/M_{\tilde{g}}, \quad H_\alpha = H_\alpha(x_i^2), \quad H'_\alpha = H_\alpha(x_{i'}^2),$$

$$I_1^{ii'} = \{(x_s - x_d)\Delta + [(x_s x_{i'}^2 - x_d x_i^2)K_1^{ii'} + (x_s - x_d)K_2^{ii'}]\}/6\Delta^2 \qquad (i \neq i'),$$

$$= (x_s + x_d)(H_1 - H_2)/2 \qquad (i = i'),$$

$$I_2^{ii'} = \{4x_s\Delta^2 + 2(x_s - x_d)(2x_i^2 + 1)\Delta + [(x_d - 3x_s)\Delta + 4(x_d - x_s)x_i^2]K_1^{ii'}$$
$$+ 2(x_d - x_s)K_2^{ii'} - [x_s\Delta^2 - (x_d - 3x_s)x_i^2\Delta + 2(x_s - x_d)x_i^4]K_3^{ii'}\}/6\Delta^3 \qquad (i \neq i'),$$

$$= -(x_s - x_d)H_2/6 \qquad (i = i'),$$

$$I_3^{ii'} = I_1^{ii'}(x_d \to -x_d), \quad I_4^{ii'} = I_2^{ii'}(x_d \to -x_d),$$

$$I_5^{ii'} = (\{6[-(1 - x_i^2)H_3 + (1 - x_{i'}^2)H'_3] - (x_s^2 - x_d^2) + 6x_d^2(H'_3 - H'_5) - 6x_s^2(H_3 - H_5)\}\Delta$$
$$- 6(x_d^2 - x_s^2)(H_4 - H'_4))/6\Delta^2 \qquad (i \neq i'),$$

$$= H_6 - H_3 - (x_d^2 + x_s^2)(H_1 - H_2)/2 \qquad (i = i'),$$

$$I_6^{ii'} = \{-\Delta^2 - (2x_{i'}^2 + 1)\Delta + 6[H_7 - \Delta(H_8 - H_4) - (i \leftrightarrow i')]\}/6\Delta^3 \qquad (i \neq i'),$$

$$= -H_2/6 \qquad (i = i'),$$

$$I_7^{ii'} = (H_3 - H'_3)/\Delta \qquad (i \neq i'),$$

$$= H_9 - H_1 \qquad (i = i'),$$

$$I_8^{ii'} = [-\Delta + \{(x_i^2 + x_{i'}^2)(H_6 - H'_6) + (2 - x_i^2 - x_{i'}^2)(H_3 - H'_3)\}]/\Delta^2 \qquad (i \neq i'),$$

$$= 0 \qquad (i = i'),$$

$$J_1^i = H_6 - H_3, \quad J_2^i = H_1 - H_2, \quad J_3^i H_9 - H_1,$$

$$K_1^{ii'} = 6(H_3 - H_5 - H'_3 + H'_5), \quad K_2^{ii'} = 6(H_5 - H'_5), \quad K_3^{ii'} = 6(H_6 - 2H_3 + H_5 - H'_6 + 2H'_3 - H'_5),$$

where

$$H_1(x) = (x^2 - 4x + 3 + 2\ln x)/2(x-1)^3, \quad H_2(x) = (2x^3 - 9x^2 + 18x - 11 - 6\ln x)/6(x-1)^4,$$

$$H_3(x) = -(3x^2 - 4x + 1 - 2x^2\ln x)/4(x-1)^2, \quad H_4(x) = -(5x^3 - 9x + 4 - 6x^3\ln x)/36(x-1)^2,$$

$$H_5(x) = -(11x^3 - 18x^2 + 9x - 2 - 6x^3\ln x)/18(x-1)^3, \quad H_6(x) = -(x - 1 - x\ln x)/(x-1),$$

$$H_7(x) = -(x^3 - 1 - 3x^3\ln x)/9(x-1), \quad H_8(x) = -(x^2 - 1 - 2x^2\ln x)/4(x-1),$$

$$H_9(x) = (x - 1 - \ln x)/(x-1)^2.$$

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