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Flavour changes in locally supersymmetric theories

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FLAVOUR CHANGES IN LOCALLY SUPERSYMMETRIC THEORIES

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Locally supersymmetric models lead in general to flavour-changing gluino exchanges. We derive the necessary expressions for describing them in detail. We show that the $K^0-\bar{K}^0$ mixing puts a lower bound on the gravitino mass and gives new restrictions on the Kobayashi–Maskawa angles.

1. Supersymmetric [1–3] theories imply a large number of new particles. Most of them are relatively heavy and are not produced directly. However, they can contribute as intermediate states to low energy processes which can be measured with high precision, and thus might give stringent restriction on the new particles. Such a restriction comes from the $K^0-\bar{K}^0$ $\Delta S = 2$ mass mixing which we treat in this paper.

In the standard model of the electroweak interactions only the W-bosons can contribute to the box diagram which generates $K^0-\bar{K}^0$ mixing. In supersymmetric (Susy) models, both charged and neutral gauge fermions (gluinos $\tilde{g}$, W-inos, $\tilde{w}$, z-inos, ...) can in principle give rise to flavour changing processes. In particular, gluino exchange might give large contributions.

Most recent models do not have flavour changing gluinos at tree level (that is, for the tree level mass matrix of the particles). But in many models radiative corrections are important, in particular to the masses of the scalar quarks (squarks). In these models, the radiative corrections can induce flavour changing gluino couplings.

In this paper we study the gluino contributions to

$K^0-\bar{K}^0$ mixing. Several authors [4] have given estimates of possible flavour violations due to supersymmetric intermediate particles, in particular the $\tilde{W}$. They follow the work of Gaillard and Lee [5], and base their results on the approximation that scalar quarks are much lighter than gauge fermions. This is, however, not applicable to the recently proposed models [6–8] with local supersymmetry where one expects the masses of the gluinos to be less or equal to the gravitino mass, and hence to the masses of squarks and scalar leptons. We derive therefore the exact expressions and evaluate also the required hadronic matrix elements.

2. The important coupling for us is that between gauge fermions λ_A , quarks ψ and squarks ϕ , given as

$$g_A \lambda_A \phi_i^* \tau_A \Gamma_{ij} \psi_j + \text{h.c.} \quad (1)$$

where g_A is the gauge coupling of λ_A and τ_A the respective coupling matrix in the representation of ψ and ϕ . (Group indices are suppressed.) The indices i, j denote mass eigenstates, and the couplings Γ^{ij} have the form

$$\Gamma^{ij} = \Gamma_L^{ij} \frac{1}{2}(1 + \gamma_5) + \Gamma_R^{ij} \frac{1}{2}(1 - \gamma_5). \quad (2)$$

The $\Gamma_{L,R}$ are obtained from the unitary matrices

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Fig. 1. Box diagram giving the effective $\Delta S = 2$ hamiltonian sq_i, sq_j are scalar quarks, λ_a, λ_b are gauge fermions.

U_ϕ, U_ψ needed to diagonalize the ϕ, ψ mass matrices. They have typically the form $U_\phi^\dagger U_\psi$. (The U_ψ is, for example, related to the KM matrix [9].)

It is obvious that if $U_\phi = U_\psi$, gluinos conserve flavour. On the other hand, the W will cause flavour changes, just as the W with $\Gamma = \text{KM}$ matrix. In any case, eq. (1) gives rise to the $\Delta S = 2$ effective operator of fig. 1. To evaluate it, we note that in supersymmetric theories there are two scalars with the same $SU(3)_{\text{colour}} \times U(1)_{\text{QED}}$ quantum numbers, ϕ_R and ϕ_L , associated with the right- and left-handed quarks, respectively. The possible mass terms are $\phi_L \phi_L^*, \phi_R^* \phi_R$, and $\phi_R^* \phi_L + \text{h.c.}$ The first two do not break any quantum numbers while the last breaks $SU(2) \times U(1)$ and is always proportional to a quark mass. Mixing of ϕ_L and ϕ_R will then be of order m_q/m_W or so and thus small, and we neglect it. (We will see below that this is also justified for another reason.)

The dominant $\Delta S = 2$ hamiltonian has then only even powers of Γ_L and Γ_R . We define

$$G_L^{ij} = \Gamma_L^{si} \Gamma_L^{*dj} L + \Gamma_R^{si} \Gamma_R^{*dj} R \\ + \Gamma_R^{si} \Gamma_L^{*dj} R + \Gamma_L^{si} \Gamma_R^{*dj} L$$

$$L, R = \frac{1}{2}(1 \pm \gamma_5), \quad (3)$$

then, the effective $\Delta S = 2$ hamiltonian is

$$\mathcal{H} = (\alpha_A^2/8m_g^2) [(\bar{d}G_2^{ij}\tau_A\tau_B s)(\bar{d}G_2^{ji}\tau_A\tau_B s) I^{ij} \\ + \frac{1}{4}(\bar{d}\gamma_\mu G_1^{ij}\tau_A\tau_B s)(\bar{d}G_1^{ji}\tau_B\tau_A s) K^{ij}] + \text{h.c.} \quad (4)$$

where

$$I^{ij} = \frac{m_g^2}{m_j^2 - m_i^2} \left(\frac{m_j^2 \ln(m_j^2/m_g^2)}{(m_j^2 - m_g^2)^2} - \frac{m_i^2 \ln(m_i^2/m_g^2)}{(m_i^2 - m_g^2)^2} \right. \\ \left. + \frac{m_j^2 - m_i^2}{(m_j^2 - m_g^2)(m_i^2 - m_g^2)} \right), \quad (5)$$

$$K^{ij} = I^{ij} + \frac{m_g^2}{(m_i^2 - m_j^2)} \left(\frac{m_g^2 \ln m_g^2 - m_i^2 \ln m_i^2}{(m_g^2 - m_i^2)} \right. \\ \left. - \frac{m_g^2 \ln m_g^2 - m_j^2 \ln m_j^2}{(m_g^2 - m_j^2)} \right), \quad (6)$$

where m_g is the common mass of the gauginos in the loop.

Since one is often interested in squarks which are nearly degenerate in mass, it is useful to give the expansion of these formulae for $m_j^2 \simeq m_i^2$. We have

$$I^{ii} = \frac{m_g^2}{(m_i^2 - m_g^2)^2} \left[2 - \frac{m_i^2 + m_g^2}{m_i^2 - m_g^2} \ln \left(\frac{m_i^2}{m_g^2} \right) \right], \quad (7)$$

$$I^{ij} \simeq I^{ii} + \frac{1}{2}(m_j^2 - m_i^2) J + \frac{1}{6}(m_j^2 - m_i^2)^2 \tilde{J}, \quad (8)$$

$$I^{ji} \simeq I^{ii} + (m_j^2 - m_i^2) J + \frac{1}{2}(m_j^2 - m_i^2)^2 \tilde{J}, \quad (9)$$

$$K^{ii} = I^{ii} + \frac{m_g^2}{m_g^2 - m_i^2} \left[\frac{m_g^2}{m_g^2 - m_i^2} \ln \left(\frac{m_g^2}{m_i^2} \right) - 1 \right], \quad (10)$$

$$K^{ij} \simeq K^{ii} + \frac{1}{2}(m_j^2 - m_i^2) K + \frac{1}{6}(m_j^2 - m_i^2)^2 \tilde{K}, \quad (11)$$

$$K^{ji} \simeq K^{ii} + (m_j^2 - m_i^2) K + \frac{1}{2}(m_j^2 - m_i^2)^2 \tilde{K}, \quad (12)$$

where

$$J = [m_g^2/(m_i^2 - m_g^2)^3] \\ \times \left[\frac{6m_i^2}{(m_i^2 - m_g^2)} \ln \left(\frac{m_i^2}{m_g^2} \right) - 5 - \frac{m_g^2}{m_i^2} \right. \\ \left. - 4 \ln \left(\frac{m_i^2}{m_g^2} \right) \right], \quad (13)$$

$$\tilde{J} = [m_g^2/(m_i^2 - m_g^2)^4] \\ \times \left[-\frac{24m_i^2}{(m_i^2 - m_g^2)} \ln \left(\frac{m_i^2}{m_g^2} \right) + 17 + 8 \frac{m_g^2}{m_i^2} \right. \\ \left. + 18 \frac{m_i^2}{m_g^2} - \frac{m_g^4}{m_i^4} \right], \quad (14)$$

$$K = J + [m_g^2/(m_g^2 - m_i^2)^2] \\ \times \left[\frac{2m_g^2}{(m_g^2 - m_i^2)} \ln \left(\frac{m_g^2}{m_i^2} \right) - 1 - \frac{m_g^2}{m_i^2} \right], \quad (15)$$

$$\tilde{K} = \tilde{J} + [m_g^2/(m_g^2 - m_i^2)^3] \times \left[\frac{6m_g^2}{(m_g^2 - m_i^2)} \ln\left(\frac{m_g^2}{m_i^2}\right) - 2 - \frac{5m_g^2}{m_i^2} + \frac{m_g^4}{m_i^4} \right]. \quad (16)$$

The apparent poles in eqs. (13) to (16) disappear, when we expand in powers of $(m_i^2 - m_g^2)$. For later use we record

$$\tilde{J} \simeq 3/10 m_g^4, \quad (14.1)$$

in leading order in $m_g^2 - m_i^2$.

3. We must next discuss the values of the Γ 's in eq. (1) for gluons. For box diagrams with $\tilde{W}$ exchange we must take expressions (3) to (16) and substitute essentially the KM matrix for Γ .

In global supersymmetry it has proved difficult to construct realistic models. The only acceptable ones are those where radiative corrections are important for the masses of the "light" particles. A typical example is a model of Ovrut and Nappi [10]. There, the leading contribution to the masses of the squarks are not tree graphs (fig. 2), but the loop graphs (figs. 3 and 4). The contributions of the latter are generally large. However they are the same for all flavours and thus uninteresting for the mixing between mass eigenstates and flavour violations. In fig. 2, only electrically neutral Higgs fields H can couple. This means that the $(\text{mass})^2$ matrix of the squarks for a given charge is of the form $MM^\dagger$ where M is the mass matrix of the corresponding quarks. Clearly $MM^\dagger$ is diagonalized by the same unitary matrix as $M^{\dagger+1}$ and no flavour changing gluino exchange arises. But in fig. 3 charged Higgs fields can also occur. This means that the $(\text{mass})^2$ matrix for one charge of squarks will obtain also a contribution proportional to the square of the mass matrix of the quarks of the other charge. [Note, that with the standard particle assignment (left-handed doublets and right-handed singlets), this only applies

^{*1} In general, if $UMV^\dagger = \text{diagonal}$ (U, V , left- and right-handed unitary mixing matrices), $UMM^\dagger U = \text{diag}(VM^\dagger MV^\dagger) = \text{diagonal}$.

Fig. 2. Tree mass for scalar quarks. The crosses are the vacuum expectation values of the Higgs fields H .

Fig. 3. Radiative mass for scalar quarks due to Yukawa couplings.

to "left-handed" squarks ϕ_L]. Thus, the $(\text{mass})^2$ matrix for the down squarks is

$$M^2 = M_d M_d^\dagger + c M_u M_u^\dagger, \quad (17)$$

where M_d, M_u are the mass matrices for the down and up quarks and c a number, calculated from diagrams such as in fig. 3. In eq. (17), the flavour-unity piece from fig. 4 is omitted. Obviously, M^2 is not diagonalized by the same unitary matrix as M_d (from the left), and flavour changing gluino coupling can arise.

We consider this example further since it will be important in the following. In the standard model M_u can be chosen diagonal and M_d such that $(\text{KM matrix}) \cdot M_d = \text{diagonal}$ ^{*2}. If $c \geq 1$, then, since $m_c \gg m_s$ and $m_t \gg m_b$, M^2 in eq. (17) is essentially diagonal, and the coupling between down quarks and squarks is given by the KM matrix. As mentioned, for "right-handed" squarks and quarks the coupling is the unit matrix, and for up particles the effect is tiny, unless c is very large. We will therefore only treat the left-handed down field in the following.

In order to estimate the effects due to this mixing, we need to know c . We will now discuss its size in models based on local supersymmetry [6–8].

In these models, supersymmetry is broken at a very high energy scale of $\sim (M_W \cdot M_P)^{1/2} \simeq 10^{11}$ GeV. Even though at tree level (appropriate for this scale)

^{*2} For example, $M_u = \text{diag}(m_u, m_c, m_t)$, $\text{KM} \cdot M_d = \text{diag}(m_d, m_s, m_b)$.

Fig. 4. Radiative mass for scalar quarks due to gauge fermions and gauge bosons.

quarks and squarks have the same mixing matrices (and consequently no flavour changing gluino couplings) the masses of the particles at the relevant scale of 100 GeV are highly affected by renormalization effects. The (mass)² matrix of the squarks satisfies [8,11] the renormalization group equation

$$dM^2/dt = (1/8\pi^2)(C_d C_d^\dagger M_1^2 + C_u C_u^\dagger M_2^2) - (1/2\pi^2)(\frac{4}{3}\tilde{m}_3^2 e_3^2 + \frac{3}{4}\tilde{m}_2^2 e_2^2 + \frac{1}{36}\tilde{m}_1^2 e_1^2), \quad (18)$$

where $\tilde{m}_i$, e_i are the masses of the gauge fermions and the coupling constants of the three gauge groups [$i=1$: U(1), $i=2$: SU(2), $i=3$: SU(3)] and M_1^2, M_2^2 are the sums of the (mass)² of the particles in the appropriate loop diagrams of fig. 3. Eq. (18) implies a mass matrix of the form

$$M^2 = a1 + bM_d M_d^\dagger + cM_u M_u^\dagger, \quad (19)$$

for the down squarks, where b is understood to include also a tree mass proportional to $M_d M_d^\dagger$. Eq. (19) is of the form (17).

To find a reasonable value of c we first consider the model of Ellis, Nanopoulos and Tamvakis [8]. The Higgs mass is given by

$$M_H^2 = M_a^2 - M_b^2, \\ M_a^2 = \alpha_2^2 (9/20\pi\alpha_G) M^2, \\ M_b^2 = (18/61\pi)(h_t^2/\alpha_G^2) \alpha_3 M^2, \quad (20)$$

where h_t is the Yukawa coupling of the top quark, α_G is the common coupling at the unification point of SU(3), SU(2), and U(1), $\alpha_i = e_i^2/4\pi$ at M_W and M an input parameter. When $M_H^2 < 0$ spontaneous symmetry breaking occurs, and the Higgs field develops a vacuum expectation value v , with

$$v^2 = -m_H^2/4\pi\alpha_2. \quad (21)$$

The (mass)² matrix of the down quarks has, beside a term proportional to the unit matrix, the contributions (figs. 2 and 3)

$$\delta M^2 = M_d M_d^\dagger, \quad (22)$$

$$\delta M^2 = (C_u C_u^\dagger + C_d C_d^\dagger)(36/1225\pi)(\alpha_3/\alpha_G^2) M^2, \quad (23)$$

where, again, M_d is the down quark mass matrix and C_u, C_d the Yukawa coupling matrices for up and down quarks. We now take $M_H^2 = -M_a^2$ ^{†3}. Using eq. (21),

the radiative term (23) is comparable to the tree level contribution (22) if $\alpha_3 \geq 4\alpha_2$. If we take $\alpha_3 \simeq 0.1$ and $\alpha_2 \sim \frac{1}{30}$, we see that the two terms are indeed comparable, and $c \sim 1$ in eq. (17).

As apparent from eq. (20), a relatively large top mass is required in this model, But as stressed above, we expect c to be large in all supergravity models. We have therefore solved eq. (18) (and similar equations [11,8] for the other parameters) for one generation. We find that for reasonable values of them

$$c \approx O(1). \quad (24)$$

If this is taken for all flavours, we are exactly in the situation described after eq. (17). The (mass)² shifts of the down quarks are essentially the (mass)² of the corresponding up quarks. We will use this in our numerical estimates.

To calculate the effect of the $\Delta S = 2$ hamiltonian, we need the hadronic matrix elements of the various four quark operators which appear in eq. (4). One of these is the same operator which appears in the standard $\Delta S = 2$ transition found by Gaillard and Lee [5]. This has recently [12] been extracted from $K \rightarrow 2\pi$ data using PCAC and SU(3). The other operators, however, must be calculated directly. We evaluate these in the MIT bag model [13]. The result is

$$\langle K^0 | \mathcal{H} | \bar{K}^0 \rangle = \frac{4}{3} \frac{\alpha_s^2}{m_g^2} \left[\frac{2}{3} (\Gamma_L^{si} \Gamma_R^{*di} \Gamma_R^{sj} \Gamma_L^{*dj}) I^{ij} (7A - B) \right. \\ \left. + (\Gamma_L^{si} \Gamma_L^{*di} \Gamma_L^{sj} \Gamma_L^{*dj}) + \Gamma_R^{si} \Gamma_R^{*di} \Gamma_R^{sj} \Gamma_R^{*dj} \right] I^{ij} (A + B) \\ \left. + \frac{11}{192} (\Gamma_R^{si} \Gamma_L^{*di} \Gamma_R^{sj} \Gamma_L^{*dj}) + \Gamma_L^{si} \Gamma_R^{*di} \Gamma_L^{sj} \Gamma_R^{*dj} \right] K^{ij} M \\ - \frac{1}{6} (\Gamma_R^{si} \Gamma_R^{*di} \Gamma_L^{sj} \Gamma_L^{*dj}) K^{ij} (7A + B) \Big], \quad (25)$$

where A and B are bag model overlap integrals [13]

$$A = \frac{m_K}{2\pi R^3} \int dx x^2 (u^2 - l^2)^2 = 0.055 m_K, \quad (26)$$

$$B = \frac{m_K}{2\pi R^3} \int dx x^2 (2ul)^2 = 0.063 m_K, \quad (27)$$

with $u(l)$ being the upper (lower) component of the quark's Dirac wave function. M is taken from ref. [12] and has the value

$$M = 0.1 m_K. \quad (28)$$

^{†3} This is certainly reasonable, since eq. (20) only makes sense if $M_b^2 \sim M_a^2$.

4. We can now give the size of the flavour change. Using eqs. (7) to (16) and eqs. (25) to (28) we obtain

$$\langle K^0 | \mathcal{H} | \bar{K}^0 \rangle = (4\alpha_s^2/3m_{\tilde{g}}^2) \tilde{J}(A+B) [\frac{1}{6}(u_{13}u_{23})^2 \Delta_{31}^2 + \frac{1}{6}(u_{12}u_{22})^2 \Delta_{21}^2 - u_{12}u_{22}u_{13}u_{23}\Delta_{31}\Delta_{21}]. \quad (29)$$

Here, $m_{\tilde{g}}$ is the mass of the gluino, the u_{ij} are the elements of the KM matrix between the i th down quark and the j th up quark, and $\Delta_{21} = m_c^2$, $\Delta_{31} = m_t^2$, as discussed below eq. (24). In deriving eq. (29) we use the fact that only the "left" squarks occur in the loop of fig. 1; this implies the absence of the leading terms $\sim \Delta_{ij}$ of eqs. (8) and (9).

Taking $\alpha_s \approx 0.1$, and replacing m_K in eq. (28) by $M_W/160$ (corresponding to $M_W = 80$ GeV) [14] we obtain

$$(M_W^2/m_{\tilde{g}}^2) \Delta_{i1} \Delta_{j1} \tilde{J}(u_{1i}u_{2i}u_{1j}u_{2j}) \lesssim 6 \times 10^{-7}, \quad (30)$$

where the factor 6 on the right hand side only enters if $i = j$. As mentioned, the mass of the gluino is expected to be not far from the mass of the squarks or of the gravitino, $m_{3/2}$. Using therefore the approximate expression (14.1), we get

$$(M_W^2/m_{\tilde{g}}^2)(\Delta_{i1})^2(u_{1i}u_{2i})^2 \lesssim 2 \times 10^{-6}. \quad (31)$$

(Note, that if $m_{3/2} > m_{\tilde{g}}$, $m_{\tilde{g}}$ in eq. (31) should be replaced by $m_{3/2}$, as is apparent from eqs. (8) and (14) and the connection of $m_{3/2}$ to the masses of the squarks [6-8].)

Consider first eq. (31) for $i = 2$. Then

$$(u_{12}u_{22})^2 m_c^4 M_W^2 / m_{\tilde{g}}^6 \lesssim 2 \times 10^{-6}. \quad (32)$$

Taking $u_{12}^2 \approx \sin^2 \theta_c = \frac{1}{20}$, $u_{22}^2 \approx \cos^2 \theta_c \approx 1$, $M_c = 2$ GeV and $M_W = 80$ GeV, eq. (32) yields

$$m_{\tilde{g}} \gtrsim 37 \text{ GeV}. \quad (33)$$

In view of the above comment, this can be viewed as a lower bound on the gravitino mass.

Secondly, use eq. (31) with $i = 3$. Then

$$(u_{13}u_{23})^2 m_t^4 M_W^2 / m_{\tilde{g}}^6 \lesssim 2 \times 10^{-6}. \quad (34)$$

As mentioned, m_t needs to be fairly large in order to trigger symmetry breaking. Our numerical evaluation of eq. (22) suggests then $m_t^2/m_{\tilde{g}}^2 \approx \frac{1}{3} - \frac{1}{2}$. Using the conservative number $\frac{1}{3}$, eq. (34) implies

$$(u_{13}u_{23})^2 (M_W^2/m_{\tilde{g}}^2) \lesssim 2 \times 10^{-5}. \quad (35)$$

For $m_{\tilde{g}} \sim M_W$, this puts a mild restriction on the mix-

ing angles u_{13} and u_{23} . If one uses the well known super CKM assumption (Ellis and Nanopoulos [14]) where $u_{13} \sim \sin^2 \theta_c$, $u_{23} \sim \sin \theta_c$, then eq. (35) implies

$$m_{\tilde{g}} \gtrsim 200 \text{ GeV},$$

a rather dramatic number. The angles are, however, not yet determined and we would rather view eq. (35) as a condition for them.

5. To summarize our results: Locally supersymmetric theories have flavour changing gluino couplings. They lead to a $K-\bar{K}^0$ mixing analogous to that of the W -inos, if α_{weak} is replaced by α_{strong} and M_W by the gluino (or gravitino, if heavier) mass. We obtain a lower bound on this mass of 37 GeV, as well as new restrictions on the KM angles. The effect of these gluino exchanges is absent in other highly suppressed observables, such as $K \rightarrow \mu\mu$ or $\mu \rightarrow e\gamma$, but it might contribute to quantities such as the electric dipole moment of the neutron.

Beside the renormalization effect discussed here, there is potentially another effect which leads to flavour-changing gluino couplings, namely the renormalization of the so-called A -term in the scalar potential [8], which contributes to the mixing of left- and right-handed squarks. For a reasonable range of the parameters, the numerical evaluation of the renormalization group equations indicates that this effect is not large.

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