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FLAVOR CHANGING NEUTRAL CURRENTS IN N=1 SUPERGRAVITY MODELS

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CONTENTS

1. INTRODUCTION

It is expected that in the next few years the experimental results, coming from the new generation of accelerators, could shed new light on the fundamental theory of elementary particles. This is the reason why it is particularly interesting to study specific physical processes in order to compare the Standard Model (SM) prediction with the predictions of its possible extensions. From a theoretical point of view, the introduction of Supersymmetry (SUSY) predicts important novelties with respect to the usual SM. It has been noticed recently that SUSY models have new sources of flavor changing neutral currents (FCNC). In particular, as I will show later in more detail, the gluino can couple to quarks and squarks of different families, giving strong interaction contributions to FCNC. These couplings arise from quantum effects, since they become effective after the masses of the squarks are evaluated at low energy solving the renormalization group equations. In view of this pecularity of the SUSY models, it becomes very interesting to revisit the whole phenomenology of FCNC. This is the purpose of my work, in which I want to point out the limits on the masses of SUSY partners that can be recovered from the present experimental results and the discrepancies with the SM that might be interpreted as signatures of SUSY in future experiments.

The work is organized as follows. In section 2 I summarize the SUSY versions of the SM, which are the framework in which I work. In section 3 I discuss the structure of the squark mass matrix and the FCNC interactions coming from its diagonalization. Section 4 is devoted to the present experimental bounds on the masses of SUSY particles, which will be useful later in setting new constraints on SUSY models. In section 5 the K^o-$\bar{K}^o$ system is studied, analysing the bounds coming from the K_L-K_S mass difference and the ε parameter and pointing out the differences between minimal and non-minimal models. In section 6 I study if the SUSY contributions can significantly enhance the FCNC mediated process $K^+ \to \pi^+ \nu \bar{\nu}$. In section 7 the SUSY contributions to the correlated decay $K_L \to \mu^+ \mu^-$ are considered. In section 8 I analyse if the branching ratio of the process $K^+ \to \pi^+$ + "missing energy" can be enhanced assuming that the particles produced in the final state were some SUSY partners instead of neutrinos. Finally in section 9 FCNC decays of heavier particles are discussed: I consider the case of B and D mesons and of Z^o gauge boson. Some integral formulas, used in the evaluation of the relevant Feynman graphs, are contained in the Appendix.

2. SUPERGRAVITY MODELS

The SM seems to be a satisfying description of elementary particles and funtamental forces. In spite of all the successes of the SM, from a theoretical point of view, there are still open problems (many free parameters, different coupling constants, family repetition, quantum gravity, Higgs sector,...). Then, even if there is no experimental evidence against the SM, it seems reasonable to search for solutions to those theoretical and aesthetic problems in the context of new physics. SUSY looks very promising in providing new answers. There are good reasons for incorporating SUSY as a fundamental symmetry of Nature. SUSY provides a unified description of particles with different spin. It helps in taming the infinities of the perturbative expansion and, for the first time, it gives an interesting hint how to connect gravity with the other forces. Furthermore SUSY solves the so called hierarchy problem [1]. Let us consider this issue in more detail.

GUTs imply an incredibly tiny ratio between the Fermi and the grand unification mass scales, $M_W^2/M_X^2 \lesssim 10^{-24}$. If masses are generated by the Higgs mechanism, one would normally expect all the masses to be of the same order of magnitude. Even if one starts from a potential with different Higgs sectors, characterized by dreadfully different parameters, the quantum corrections will mix the Higgs sectors and spoil the hierarchy. In order to obtain a mass ratio of 10^{-24} it appears to be necessary to fine tune the parameters in the Higgs potential to 24 decimal places. The hierarchy problem does not concern only GUTs, but also the SM, because it is due to the quadratic divergences always present in a simple scalar particle theory. To understand why this happens, consider a theory with dimensionless bare couplings g_o and bare masses m_o . I assume that the theory has a fundamental mass scale K, which behaves as a cutoff. Because of our ignorance about quantum gravitational effects, I assume K to be of order of the Planck mass (10^{19} GeV). Now, if the self-energy corrections are quadratic in K, the physical mass of scalar particles very schematically becomes:

$$m^2 \equiv m_o^2 + \Delta m^2 = m_o^2 + K^2 g_o^2 \qquad (2.1)$$

If one wants a physical mass of the particle of order M_W^2 , the dimensionless bare mass must be:

$$\mu_o^2 \equiv \frac{m_o^2}{K^2} = -g_o^2\left(1 - \frac{m^2}{K^2}\right) = -g_o^2\left(1 - 10^{-34}\right) \qquad (2.2)$$

This means that μ_o^2 must be adjusted to the 34th decimal place. If not, the mass of the scalar particle will come out to be order K. Furthermore a readjustment must be performed at each order in the perturbative expansion. This violates the intuitive concept of naturalness which requires the observable properties of the theory to be stable against minute variations of the fundamental parameters. Here one has two problems, one of which is more technical, the other more philosophical: first of all how to keep the hierarchy despite the radiative corrections, then why the huge hierarchy is present in Nature. A possible solution to the technical problem is to ban the scalars from the theory, assuming them to be fermion-antifermion condensates. This leads to introduce a new QCD-like force, called technicolor. Another strategy is to keep the scalars alive, but to cancel their dangerous quadratic divergences. The only symmetry able to do so is

SUSY. In addition to the scalars, one has fermionic degrees of freedom, which contribute to the loops with a minus sign. In presence of exact SUSY, the quadratic divergences cancel. Another way to show this result is the following: bosons and fermions are related by SUSY; since the chiral symmetry forbids fermion masses at any perturbative order, the quantum corrections can not generate masses for the scalar particles. Once SUSY is broken, scalar masses can be produced, while fermion masses are still protected by chiral symmetry. The cancellation is no longer exact, and one gets a correction to boson mass $\delta m^2 \sim \alpha \Delta m^2$, where Δm^2 is the typical splitting inside the supermultiplet. To solve the hierarchy problem, I impose that the mass correction is not larger than the Higgs mass itself, which implies:

$$\Delta m \lesssim \frac{100\ GeV}{\sqrt{\alpha}} \simeq O(TeV) \qquad (2.3)$$

Therefore, if the splitting between superpartners is less than few TeV, the technical part of the hierarchy problem is solved. Still one could wonder why the huge hierarchy is there. It is important to stress that it is not necessarily unnatural to have a large hierarchy present in the theory. I use an example for illustration. Consider the asymptotically free QCD, where the coupling constant g_s satisfies the equation:

$$\frac{\partial g_s}{\partial \ln q} = \beta(g_s) \simeq -c\, g_s^3 + O(g_s^5) \qquad c > 0 \qquad (2.4)$$

Solving the equation, I find a relation between the QCD scale Λ, where $g_s(\Lambda) \simeq 1$, and the grand unification scale M_x, where $g_s(M_x) \ll 1$:

$$\frac{\Lambda}{M_x} \simeq \exp\left(-\frac{1}{2c\, g_s^2(M_x)}\right) \ll 1 \qquad (2.5)$$

A large hierarchy is obtained naturally. Note that the ratio Λ/M_x is stable for small variations of $g_s(M_x)$. SUSY can be useful in solving completely the hierarchy problem. Supergravity models where the SUSY breakdown scale is dynamically generated have been proposed [2]. In these models the electroweak scale can be extracted from the SUSY scale, and all the mass scales in the theory might be determined starting from the only fundamental one, the Planck scale.

I will restrict myself only to N=1 supergravity models, because for $N>1$ the fermions lie in real representation of the gauge group and it is not straightforward to build models compatible with the low energy phenomenology.

The N=1 broken supergravity models must always be understood as an effective limit of a more fundamental theory, after integrating out the extra degrees of freedom. The proliferation of models is due to the poor knowledge of the mechanism of SUSY breakdown. Different models predict different values for the sparticle masses. However, since I treat the SUSY masses as free parameters, most of my conclusions will be model independent. The model dependence comes in when I make particular assumptions on the form of the mixings between current and mass eigenstates.

Although there are no experimental hints about the choice of the right model, it is possible to make detailed calculations and predictions for SUSY theories. This is due to the fact that the form of the interaction is fixed by the symmetry and the coupling constants are related to those of the SM, therefore allowing a perturbative expansion.

A common feature of all SUSY models is that the number of particles is at least doubled with respect to the SM. Each known particle has a SUSY partner, which belongs to the same N=1 supermultiplet. In the early attempts of constructing realistic SUSY models, it was soon realized that it is not possible to put known particles in the same supermultiplet. The fermionic partners of the gauge bosons are in a real gauge representation and this is not the case for quarks and leptons. If the Higgs particles were the partners of the leptons, the v.e.v. of the sneutrino would break the lepton number; moreover the sneutrino is not able to give mass to both up and down quarks. If one introduces only one Higgs doublet superfield, masses for both up and down quarks are generated. Furthermore the fermionic partners of the usual Higgs doublet render the particle content of the model anomalous. Two superfield doublets are sufficient to cancel anomalies and to provide masses to quarks and leptons. The minimal particle content of a SUSY model is summarized in table 2.I.

Gauge group content SU(3)	SU(2)	U(1)	Particle	SUSY partner		Mass eigenstate
3	2	1/3	$q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix}$	$\tilde{q}_L = \begin{pmatrix} \tilde{u}_L \\ \tilde{d}_L \end{pmatrix}$	squarks	
3	1	-4/3	u_R	$\tilde{u}_R$		
3	1	2/3	d_R	$\tilde{d}_R$		
1	2	-1	$\ell_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix}$	$\tilde{\ell}_L = \begin{pmatrix} \tilde{\nu}_L \\ \tilde{e}_L \end{pmatrix}$	sleptons	
1	1	2	e_R	$\tilde{e}_R$		
8	1	0	g	$\tilde{g}$	gluino	
1	3	0	γ	$\tilde{\gamma}$	photino	neutralino $\tilde{\chi}^0$
	1		Z^0	$\tilde{Z}^0$	z-ino	
			$W^\pm$	$\tilde{W}^\pm$	w-ino	
1	2	1	$\begin{pmatrix} H_1^+ \\ H_1^0 \end{pmatrix}$	$\begin{pmatrix} \tilde{H}_1^+ \\ \tilde{H}_1^0 \end{pmatrix}$	higgsinos	chargino $\tilde{\chi}^\pm$
1	2	-1	$\begin{pmatrix} H_2^0 \\ H_2^- \end{pmatrix}$	$\begin{pmatrix} \tilde{H}_2^0 \\ \tilde{H}_2^- \end{pmatrix}$		

Table 2.I.

The model is invariant under the usual gauge group SU(3) X

SU(2) x U(1) and under local SUSY. Given the particle content and the invariance group, the Lagrangian of the theory is not uniquely determined (remember that I did not require renormalizability). All the arbitrarity of the model is contained into two functions, the Kähler potential $G(z,z^*)$ and the coupling function of the vector superfield kinetic term, $f_{\alpha\beta}(z)$. The model is specified by a choice of these two functions. I will not make any assumptions on the form of $f_{\alpha\beta}(z)$. A non-vanishing v.e.v. of $\frac{\partial f_{\alpha\beta}(z)}{\partial z}$ gives rise to tree level gaugino masses. In this way we may think to relate the function $f_{\alpha\beta}(z)$ to the gaugino masses. Since I treat the masses of the SUSY partners as free parameters, it is consistent to keep $f_{\alpha\beta}(z)$ completely general.

The choice of the Kähler potential is rather critical. I will focus my interest on the so called "minimal models", characterized by a flat Kähler metric [3]:

$$\frac{\partial^2 G(z,z^*)}{\partial z_i \, \partial z^{*j}} \equiv G^i_j = -\delta^i_j \tag{2.6}$$

(Sometimes the minimal models are defined with the further constraint $f_{\alpha\beta}(z) = \delta_{\alpha\beta}$). The constraint (2.6) fixes the Kähler potential up to a function of z alone, the superpotential. Note that there is no symmetry reason to impose (2.6), i.e. the gravitational radiative corrections do not preserve (2.6). The choice (2.6) is dictated only by the practical reason of simplifying the calculations and giving determinate predictions.

Since there is no evidence in Nature for a degeneracy in supermultiplets, the next step in building a realistic model is to find a way to break SUSY. I will assume that local SUSY is spontaneously broken in a hidden sector of the theory. The hidden sector contains additional particles that, for semplicity, I assume to be neutral with respect to the standard gauge group and with no couplings in the superpotential with the observable sector. No specific assumption on the hidden sector is made, except that it is responsible for local SUSY breakdown and it communicates with the observable sector only through gravitational interactions.

After the local SUSY breakdown, the gravitino acquires a mass through the super-Higgs mechanism. The SUSY breakdown destroys the degeneracy inside the supermultiplets, but the masses of the theory are constrained by the formula (in the minimal case) [4] :

$$\sum_J (-)^{2J} (2J+1)\, m_J^2 = 2(N-1)\, m_{3/2}^2 \tag{2.7}$$

where N is the number of chiral fermions in the theory and $m_{3/2}$ is the gravitino mass. The sum is over all the particles of the theory with spin J and mass m_J .

The model considered so far seems to have all the features to be realistic, but unfortunately one can not do much with it, since it is not renormalizable. To have a predictive theory, one considers the so called flat limit, where $M_{P\ell} \to \infty$ with $m_{3/2}$ fixed. This means that gravity is turned off after SUSY breakdown and so low energy gravitational effects can survive. Assuming that the hidden sector gets v.e.v.s order $M_{P\ell}$, the scalar potential of the minimal models becomes simply, in the flat limit [5]:

$$V = \left| \frac{\partial f}{\partial y_i} \right|^2 + m_{3/2}^2 |y_i|^2 + m_{3/2} \left[y_i \frac{\partial f}{\partial y_i} + (A-3) f(y) + h.c. \right] \tag{2.8}$$

where f is the rescaled superpotential, A a parameter determined by the hidden sector and y_i are the observable scalar fields. We have obtained the scalar potential of a global SUSY theory plus soft breaking terms. In the case of a tri-linear potential (2.8) reduces to:

$$V = \left|\frac{\partial f}{\partial y_i}\right|^2 + m_{3/2}^2 \, |y_i|^2 + A\, m_{3/2} \left(f(y) + h.c. \right) \qquad (2.9)$$

It is important to stress that the potential (2.8) is a consequence of the assumption that the superpotential is simply given by the sum of two superpotentials involving fields from each of the two sectors. Different combinations give different flat limit potentials.

Starting from non minimal Kähler metrics, the flat limit gives the following scalar potential [6]:

$$V = \left|\frac{\partial f}{\partial y_i}\right|^2 + m_{ij}^2 \, y_i y_j^* + m \left(h(y) + h.c. \right) \qquad (2.10)$$

m_{ij}^2 is a general matrix and h(y) is a general function of the scalar fields; the mass parameters are of the same order of the gravitino mass. Therefore the non minimal models predict a flat limit potential of the same form of (2.9), but now the parameters are unconstrained. In particular the scalar mass terms are completely general. As I will show later, the K^o-$\overline{K}^o$ system implies that the squarks must be highly degenerate in masses. An interesting feature of the minimal models is that they automatically give an equal contribution to sfermion masses, thus accomplishing a good degeneracy. Note that these masses are defined at the scale of SUSY breakdown; the low energy values can be evaluated using the renormalization group equations.

The low energy limit is then well defined. I choose the most general superpotential invariant under SU(3) x SU(2) x U(1), under the global symmetries B and L and satifing the requirement of renormalizability:

$$f = \Gamma_e^{ij} L_{L_i} H_1 E_{R_j}^{+} + \Gamma_d^{ij} Q_{L_i} H_1 D_{R_j}^{+} + \Gamma_u^{ij} Q_{L_i} H_2 U_{R_j}^{+} + \hat{m} H_1 H_2 \quad (2.11)$$

where the superfields are denoted by capital letters. The scalar potential of the model is

$$V = V_{susy} + V_{soft} \quad (2.12)$$

where V_{susy} is the global SUSY invariant potential and

$$\begin{aligned} V_{soft} = {} & M^2_{Q_{ij}} (\tilde{u}^{+}_{L_i} \tilde{u}_{L_j} + \tilde{d}^{+}_{L_i} \tilde{d}_{L_j}) + M^2_{U_{ij}} (\tilde{u}^{+}_{R_i} \tilde{u}_{R_j}) + M^2_{D_{ij}} \tilde{d}^{+}_{R_i} \tilde{d}_{R_j} + \\ & M^2_{L_{ij}} \tilde{\ell}^{+}_{L_i} \tilde{\ell}_{L_j} + M^2_{E_{ij}} \tilde{e}^{+}_{R_i} \tilde{e}_{R_j} + M^2_{H_1} H_1^{+} H_1 + M^2_{H_2} H_2^{+} H_2 + \rho H_1 H_2 + \\ & + \eta_{U_{ij}} H_2 \tilde{q}_{L_i} \tilde{u}^{+}_{R_j} + \eta_{D_{ij}} H_1 \tilde{q}_{L_i} \tilde{d}^{+}_{R_j} + \eta_{E_{ij}} H_1 \tilde{\ell}_{L_i} \tilde{e}^{+}_{R_j} + \text{h.c.} \end{aligned} \quad (2.13)$$

In the minimal case the parameters defined at the scale of SUSY breakdown are constrained by the relations:

$$\begin{aligned} & M^2_Q = M^2_U = M^2_D = m^2_{3/2} \mathbb{1} \qquad && M^2_{H_1} = M^2_{H_2} = m^2_{3/2} \\ & \eta_U = m_{3/2} A \Gamma_U \quad \eta_D = m_{3/2} A \Gamma_D \qquad && \rho = m_{3/2} (A-1) \hat{m} \end{aligned} \quad (2.14)$$

To show that the previous SUSY model is compatible with the standard phenomenology we have to prove that it predicts the right electroweak breakdown. A tree level SU(2) x U(1) breakdown needs the introduction of an extra gauge singlet superfield [7], which is harmful for the solution of the hierarchy problem [8]. Anyway the renormalization group equations can drive a negative square mass for the Higgs scalar, breaking spontaneously SU(2) x U(1).

This model represents the framework in which I work. It has all the general characteristics to be considered as a realistic model. A more detailed analysis is necessary to point out the possible drawbacks or the new predictions.

3. SQUARK MASS MATRIX

In this section I want to focus my attention on the FCNC interactions arising in SUSY theories. In the SM the GIM mechanism insures that no FCNC occurs at tree level, and the one loop contributions are suppressed by a factor $\frac{\Delta m_q^2}{m_q^2}$. In the SUSY case, as I will show later, the quark and squark mass matrices are not diagonalized by the same transformation and this leads to FCNC at tree level. Let us consider in more detail the mass matrix of the squarks.

From the globally SUSY invariant potential one recovers, as soon as SU(2) x U(1) is spontaneously broken, squark mass terms proportional to $m_q^+ m_q$, where m_q is the corresponding quark mass matrix. These mass terms are of the kind $\tilde{q}_L^+ \tilde{q}_L + \tilde{q}_R^+ \tilde{q}_R$, where $\tilde{q}_{L,R}$ denote the scalar partners of left and right-handed quarks. Since in this sector SUSY is still exact, quarks and squarks are obviously degenerated.

When we turn on the soft breaking terms, we get masses for the scalar quarks. These terms will be called $\mu_L^2 \tilde{q}_L^+ q_L + \mu_R^2 \tilde{q}_R^+ \tilde{q}_R$; in the minimal models $\mu_{L,R}^2$ are proportional to the unity matrix in generation space. Moreover, after the electroweak spontaneous breakdown, the trilinear interactions give rise to mass terms like $\tilde{q}_L^+ \tilde{q}_R + h.c.$. In the case of minimal models, the coefficient is $A m\, m_q$, where m_q is the quark mass matrix. As these terms mix scalar partners of quarks with different chirality, from now on they will be called left-right squark mixings. Then the tree level mass term of the squarks is:

$$\left(\tilde{q}_L^{(0)+} \; \tilde{q}_R^{(0)+}\right) M_q^2 \begin{pmatrix} \tilde{q}_L^{(0)} \\ \tilde{q}_R^{(0)} \end{pmatrix} \tag{3.1}$$

with

$$M_q^2 = \begin{pmatrix} 1\!\!1 \mu_L^2 + m_q^+ m_q & A^* m\, m_q^+ \\ A m\, m_q & 1\!\!1 \mu_R^2 + m_q^+ m_q \end{pmatrix} \tag{3.2}$$

The matrix (3.2) has been obtained under the assumption that $<H_1> = <H_2>$. If this is not the case, there are other contributions coming from the D-terms. Thus I have to add to the left and right sectors of the up squark mass matrix respectively

$$\left(\frac{1}{2} g_2^2 - \frac{1}{3} g_1^2\right)\left(\frac{<H_1>^2 - <H_2>^2}{2}\right) \quad \text{and} \quad \frac{4}{3} g_1^2 \left(\frac{<H_1>^2 - <H_2>^2}{2}\right) \tag{3.3}$$

and anologously for down squarks. However note that these terms are equal for all the generations. I can take them into account simply rescaling the parameters μ_L^2 and μ_R^2 . Furthermore, from the F-term I can get a mass proportional to $\hat{m}$, in the left-right mixing sector. Also this term can be reabsorbed in the definition of the parameters. Actually the complete left-right mixing term for up squarks in the minimal version is:

$$A m_{3/2} \Gamma_U <H_2> + \hat{m} \Gamma_U <H_1> \equiv A m m_U \tag{3.4}$$

and analogously for down squarks. These redefinitions of the parameters are consistent with the way I treat the model. Actually μ_L , μ_R, m and A , unknown quantities of the matrix (3.2), are treated as free parameters to be determined by experiments. As I have already mentioned, I study the low energy limit as a phenomenological model, independently of the theory from which it comes from. In conclusion, (3.2) is the tree level mass matrix of the squarks; for definitness, I take μ_L, μ_R and m to be of the same order, and I denote this scale by μ .

The radiative corrections can give significant contributions [9]. The quantum corrections to squark masses diagonal in generation space can be absorbed in the unknown parameters. Let us focus our attention on the graphs with exchange of charged higgs and higgsino of fig.3.1 (I consider first the case of down squarks).

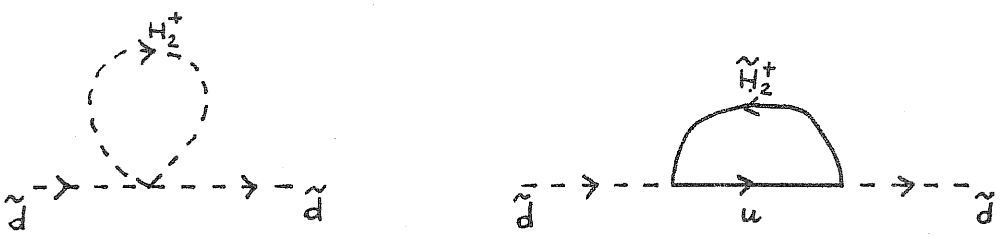


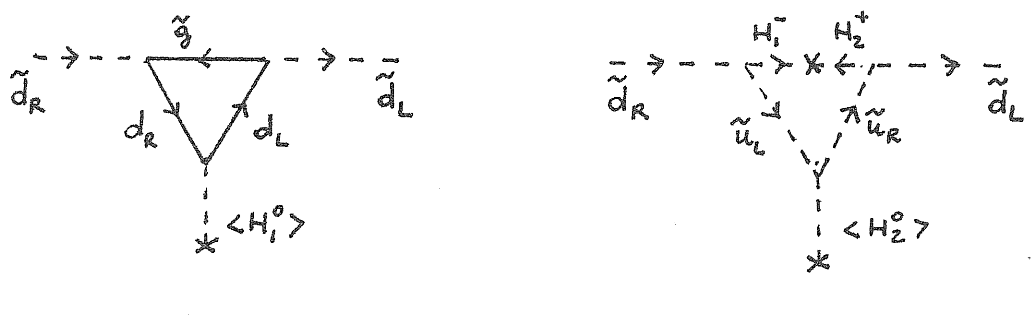
Fig.3.1

The relevant term in the superpotential (2.11) is:

$$\Gamma_U^{ij} D_{L_i} H_2^{(+)} U_{R_j}^+ \tag{3.5}$$

which gives rise to a four scalar interaction $\tilde{d}_L H_2^{(+)} H_2^{(+)+} \tilde{d}_L^+$ proportional to $\Gamma_U^+ \Gamma_U$ and to a Yukawa interaction $\tilde{d}_L \bar{u}_R \tilde{H}_2^{(+)}$ with coupling constant Γ_U. Therefore, as soon as SUSY is broken, so that no exact cancellation occurs, the radiative corrections yield a mass term for the scalar partners of the left-handed down quarks of the kind $c m_U^+ m_U$. The coefficient c can be computed evaluating the renormalization group equations for squark masses. For $\mu \simeq 40$ GeV, c is in the range between 0.5 and 0.7. Anologously, a term proportional to $m_d^+ m_d$ is found for the partners of the left-handed up quarks; however the radiatively induced elements of M_U^2 are numerically less important. Note that these terms arise only for partners of the left-handed quarks, since the right-handed quarks are SU(2) singlets.

Radiative corrections to the left-right mixings can arise through the graphs of fig. 3.2. The graph of fig. 3.2(a) gives a contribution of the kind $g_s^2 \Gamma_D \langle H_1^0 \rangle \sim m_D$, which can be absorbed into $A m m_D$. However the graph of fig. 3.2(b) gives



(a) (b)

Fig. 3.2

a term of the kind $\rho\, \eta_D \eta_U^+ \eta_U \langle H_2^0 \rangle \sim \Gamma_D \langle H_2^0 \rangle \Gamma_U^+ \Gamma_U \sim m_D \dfrac{m_U^+ m_U}{m_W^2}$.
The left-right mixing term finally becomes:

$$A\, m_{3/2}\, m_D \left(1 + \frac{c'}{m_W^2}\, m_U^+ m_U \right) \tag{3.6}$$

In my discussion I will neglect the radiatively induced term (i.e. I take c'=0), since it is suppressed by the factor $1/m_W^2$. This term becomes interesting in the study of CP violation, because it introduces new observable phases [10].

After taking into account the quantum corrections, the mass matrix for the down squarks becomes

$$M_D^2 = \begin{pmatrix} \mathbb{1}\mu_L^2 + m_D^+ m_D + c\, m_U^+ m_U & A^* m\, m_D^+ \\ A\, m\, m_D & \mathbb{1}\mu_R^2 + m_D^+ m_D \end{pmatrix} \tag{3.7}$$

For the up squarks the mass matrix is identical after the interchange between U and D. Let us call $L_{D,U}$ and $R_{D,U}$ the matrices that diagonalize the quark mass matrices in the following way:

$$R_{D,U}^+\, m_{D,U}\, L_{D,U} = \hat{m}_{D,U} \tag{3.8}$$

with $\hat{m}_{D,U}$ diagonal and $L_U^+ L_D = V$ is the Kobayashi-Maskawa (KM) matrix. The squark mass matrix is diagonal after the transformation (everything works analogously for m_U^2):

$$\tilde{U}_D^+ m_D^2 \tilde{U}_D = U_D^+ \begin{pmatrix} \mu_L^2 1\!\!1 + \hat{m}_D^2 + cV^+ \hat{m}_U^2 V & |A| m_{3/2} \hat{m}_D e^{i\delta} \\ |A| m_{3/2} \hat{m}_D e^{-i\delta} & \mu_R^2 1\!\!1 + \hat{m}_D^2 \end{pmatrix} U_D \qquad (3.9)$$

where $$\tilde{U}_D \equiv \begin{pmatrix} e^{i\varphi} L_D & 0 \\ 0 & e^{-i\varphi} R_D \end{pmatrix} U_D \qquad (3.10)$$

with $\delta = \phi - 2\varphi$, combination of the phases of the A parameter ($A \equiv |A| e^{-i\phi}$) and of the gluino Majorana mass ($M_{gl} \equiv |M_{gl}| e^{-2i\varphi}$). Neglecting the left-right mixing, assuming c to be O(1) and noting that the eigenvalues of $\hat{m}_U^2$ are larger than those of $\hat{m}_D^2$, a good approximation for U_D can be

$$U_D = \begin{pmatrix} V^+ & 0 \\ 0 & 1\!\!1 \end{pmatrix} \qquad (3.11)$$

(and analogously $U_U = 1\!\!1$).

As can be seen from (3.9), the rotation matrix for the squarks contains an extra phase. As I will show in section 5, the limit on the dipole electric moment of the neutron severely constraints the new phase.

The trasformation from current eigenstates ($\psi^{(0)}$) to mass eigenstates (ψ) is defined as follows:

$$d_L^{(0)} = L_D d_L \qquad \begin{pmatrix} \tilde{d}_L^{(0)} \\ \tilde{d}_R^{(0)} \end{pmatrix} = \begin{pmatrix} e^{i\varphi} L_D & 0 \\ 0 & e^{-i\varphi} R_D \end{pmatrix} U_D \tilde{d} \qquad (3.12)$$

$$d_R^{(0)} = R_D d_R$$

In terms of mass eigenstates, one finds various interactions which mix different generations and the two different scalar sections. It is particularly interesting for my purposes to note that the gluino can mediate FCNC at tree level. Remember however that quantum effects are already contained in the squark mass matrix. The gluino-quark-squark vertex is:

$$-\sqrt{2}\, g_s \left(\bar{d}_L^{(0)} \quad \bar{d}_R^{(0)} \right) \tilde{g}_\alpha^{(0)}\, T^\alpha \begin{pmatrix} \tilde{d}_L^{(0)} \\ \tilde{d}_R^{(0)} \end{pmatrix} + h.c. =$$

$$= -\sqrt{2}\, g_s \left(\bar{d}_L \quad \bar{d}_R \right) \tilde{g}_\alpha\, T^\alpha U_D\, \tilde{d} + h.c. \tag{3.13}$$

where $\tilde{g}_\alpha$ is the gluino fields and T^α is the color generator in the down quark representation. Since U_D is, in general, non diagonal, the interaction in (3.13) mixes different families. An identical vertex arises for neutralino interaction. The chargino-quark-squark vertex is:

$$-g\left[\bar{u}_L^{(0)}\, \tilde{\chi}^{(+)}\, \tilde{d}_L^{(0)} + \bar{d}_L^{(0)}\, \tilde{\chi}^{(-)}\, \tilde{u}_L^{(0)} + h.c. \right] =$$

$$= -g\, e^{i\varphi} \left[(\bar{u}_L \quad 0)\, \tilde{\chi}^{(+)} \begin{pmatrix} V & 0 \\ 0 & 1\!\!1 \end{pmatrix} U_D\, \tilde{d} + (\bar{d}_L \quad 0)\, \tilde{\chi}^{(-)} \begin{pmatrix} V^+ & 0 \\ 0 & 1\!\!1 \end{pmatrix} U_U\, \tilde{u} \right] + h.c. \tag{3.14}$$

SUSY theories display new kinds of FCNC interactions with respect to the SM. In the next sections, after a digression of the limits on the masses of SUSY particles, I will study the consequences of this novelty.

4. LIMITS ON SPARTICLE MASSES

In this section I want to discuss the present limits on the masses of the new particles introduced by SUSY. These bounds will be used later, when I study the consequences of SUSY in low energy FCNC phenomenology. The limits can be extracted from negative results in experiments where one tries to produce the new particles. Other experimental bounds come from the study of the contribution of the sparticles to known processes. Furthermore cosmology can be useful in setting new constraints on masses and finally, from the SUSY models themselves, it is possible to derive theoretical bounds. I will discuss in detail all the limits and compare the results.

An important issue is to know which sparticle is the lightest. Actually most of the SUSY models so far proposed have an exact R-parity which distinguishes the usual particles from their partners. This symmetry, which survives after the SU(2) x U(1) breakdown, avoids the possibilty of mass mixings between ordinary and SUSY particles. The R-parity implies that the lightest SUSY particle is stable and then it can disappear only through pair-annihilation. If a charged, uncolored sparticle (winos, charged higgsinos, charged sleptons) were the lightest, its present number density would imply that its mass is larger than 350 GeV. [11]. Suppose the gluino is the lightest: since gluinos are confined, the lightest of these new hadrons would be stable. If this state is charged, it can be excluded by the same argument as before. A stable neutral new hadron is acceptable. However in all the SUSY models, the use of the renormalization group equations predicts a photino lighter than the gluino. Therefore the idea of the gluino as the lightest partner can be disregarded. With a similar argument I can exclude the squark as a good candidate: again the renormalization group equation predicts that they are heavier than some sleptons. I am left only with the possibility of the sneutrino or of a linear combination of neutral gauginos and higgsinos. Both the possibilities give rise to acceptable models. Assuming the presence of the R-parity, I can infer that the lightest SUSY particle behaves essentially like a neutrino. Actually in a laboratory experiment, where one starts from non-SUSY particles, only an even number of partners can be produced. This implies that the cross section for the lightest SUSY particle is of weak interaction size, because it can only interact with non-SUSY matter via the exchange of a slepton or a squark, which are heavy, as I am going to show. This characterizes the signature of SUSY. Note that this would not be true if the gluino were the lightest SUSY particle. Because of the $\tilde{g}\tilde{g}g$ vertex, gluinos can interact without the exchange of heavy SUSY particles.

I will describe now the existing limits on SUSY partner masses, discussing each sparticle separately. I will consider the cases of photino, gluino, charged slepton, squark, sneutrino, chargino and neutralino.

<u>Photino</u>

The photino is the spin 1/2 partner of the photon. I assume that the mixing with the z-ino and neutral higgsino is small and that the photino is an approximate mass eigenstate. Since it escapes direct experimental detection, the most stringent bounds on its mass come from cosmological arguments. The limits on photino mass depend on the model under consideration. I discuss the three following classes of models:

a) The photino is the lightest sparticle, it is stable and its mass is $\lesssim$ 1 MeV. In this case the photinos would have survived without annihilation during the expansion of the Universe, since 1 MeV corresponds to the decoupling temperature of weakly interacting particles. The requirement that the photino contribution to the mass density of the Universe is smaller than the critical mass density, which is believed to be an upper limit on the current mass density of the Universe, leads to the bound

$$\tilde{m}_\gamma \lesssim 100 \text{ eV} \tag{4.1}$$

It seems however unnatural, in the models proposed so far, to have such light photinos, after the quantum corrections are taken into account. The authors of ref.[12] find that the radiative corrections for the photino mass are in the range of 100 MeV. Then, even if photino masses of few eV are not experimentally excluded, the scenario presented here is considered to be disfavoured.

b) I suppose again the photino to be the lightest sparticle, but its mass is now larger than ~ 1 MeV. In this case the photino is heavy enough to annihilate into a pair of fermions through the exchange of a sfermion. With an analysis à la Lee-Weinberg, the following results are derived in ref.[13]:

i) If the photino is lighter than the τ lepton, then

$$\tilde{m}_\gamma > 28 \left(\frac{\tilde{m}_f}{100}\right)^2 \tag{4.2}$$

where all the masses are expressed in GeV. $\tilde{m}_f$ are the masses of the exchanged sfermions, which are all set equal. The bound (4.2) is consistent only for $\tilde{m}_f < 25$ GeV.

ii) If the photino is heavier than the τ lepton and $\tilde{m}_f \lesssim 45$ GeV, then one simply recovers the spectral bound:

$$\tilde{m}_{\gamma} > m_{\tau} \qquad (4.3)$$

iii) If the photino is heavier than the τ lepton and $\tilde{m}_f >$ 45 GeV, the bound grows approximately linearly until $\tilde{m}_f \simeq 100$ GeV:

$$\tilde{m}_{\gamma} > \frac{\tilde{m}_f}{2.7} - 16 \qquad \text{for } 50\ \text{GeV} < \tilde{m}_f < 100\ \text{GeV} \qquad (4.4)$$

where the masses are expressed in GeV.

In conclusion, in this scenario, for light sfermions (~ 40 GeV) ref.[13] sets the lower bound:

$$\tilde{m}_{\gamma} \gtrsim 1.8\ \text{GeV} \qquad (4.5)$$

A less stringent bound is found in ref.[14]. The discrepancy is due to a different factor of 2 in the photino annihilation cross section ($\sigma_{\tilde{\gamma}\tilde{\gamma}}$) and to the inclusion of masses for light quarks, without which $\sigma_{\tilde{\gamma}\tilde{\gamma}}$ would be purely P-wave below the charm threshold. Finally the result of ref.[14] is

$$\begin{array}{lll} \tilde{m}_{\gamma} \gtrsim 0.5\ \text{GeV} & \text{for} & \tilde{m}_f \sim 20\ \text{GeV} \\ \tilde{m}_{\gamma} \gtrsim 1.8\ \text{GeV} & \text{for} & \tilde{m}_f \sim 40\ \text{GeV} \\ \tilde{m}_{\gamma} \gtrsim 5.0\ \text{GeV} & \text{for} & \tilde{m}_f \sim 100\ \text{GeV} \end{array} \qquad (4.6)$$

Note however that the two analyses coincide for $\tilde{m}_f \sim 40$ GeV.

c) A possible alternative is to consider an unstable photino [15]. If the R-parity is not present, the photino can decay into ordinary particles, evading the cosmological bounds. Still we have to deal with the theoretical limits inferred from the computation of the radiative corrections, which suggest a photino heavier than 100 MeV. Furthermore a photino decaying $\tilde{\gamma} \to \gamma \tilde{G}$ ($\tilde{G}$ is the Goldstino) has been studied at PETRA. Assuming a scale of SUSY breaking of 100 GeV, they find [16] that the photino must be heavier than 20 GeV, for $\tilde{m}_{e_L} = \tilde{m}_{e_R} =$

= 50 GeV. Lighter photinos are allowed by heavier selectrons, and the photino is unconstrained for $\tilde{m}_e > 100$ GeV.

Gluino

The gluinos are the spin 1/2 partners of the gluon. They are Majorana spinors and, since they carry color, they are assumed to be confined in new SUSY hadrons. These states can be fermionic composites of quark, antiquark and gluino which are neutral in color with electric charge 0, $\pm$ 1.Other possible new neutral composites are the gluon-gluino fermions or the gluino-gluino bosons. If the gluinos are very light, then they could be present as partons in the low energy hadrons. MIT-bag model [17] and Bethe-Saltpeter model calculations [18] suggest that these new bound states should have masses in the range of 1 GeV, if the gluinos are massless or very light. If the gluinos are heavy, the composite will have approximately the same mass and the gluinos can be treated as free particles. Most of these new "R-hadrons" (they transform non-trivially under R-parity) are unstable under strong interactions, but, for light gluinos, the analyses of ref.[17-18] predict the existance of low energy gluino hadrons which decay only weakly. Experimentally charged hadrons with lifetimes greater than 10^{-8} sec. in the mass range between 1.5 and 9 GeV are excluded [19].

A way to put bounds on gluino masses is to use the beam dump experiments. In these experiments the gluinos, if light, are copiously produced in $pN \to \tilde{g}\tilde{g}X$, where X denotes any other particle. Subsequently the gluino is assumed to decay before interacting. Either it decays $\tilde{g} \to q\bar{q}\tilde{\gamma}$ or $\tilde{g} \to g\tilde{\gamma}$, and among the products we always find a photino, which is supposed to be stable. The photino escapes the apparatus and can occasionally interact in a neutrino detector. The BEBC experiment [20] sets the limits

$$M_{gl} > 3 \text{ GeV} \quad \text{for} \quad \tilde{m}_q \lesssim 150 \text{ GeV}$$
$$M_{gl} > 4 \text{ GeV} \quad \text{for} \quad \tilde{m}_q \lesssim 65 \text{ GeV} \tag{4.7}$$

It should be noted that these experiments can not rule out extremely small M_{gl} (approximately 0.5 GeV or less) or large $\tilde{m}_q$, because the lifetime of the gluino would increase until it can interact with the dump before it decays ($\tau \gtrsim 5\cdot 10^{-11}$ sec).

One can also try to use the $p\bar{p}$ experiments to find new limits on gluino mass. As it is well known, few monojets and dijets with missing energy are reported in 1983 [21] and 1984 [22] UA1 data. These events have large missing transverse energy (15-50 GeV) and a hadronic jet of 15-25 GeV. If no other hadronic jet has energy larger than 12 GeV, the event is labelled as a monojet. The SM may provide some explanations for such events: QCD dijet with one jet missed by the detector; $p\bar{p} \to W + X$, $W \to \tau\nu$ and $\tau \to \nu$ + jet; $p\bar{p} \to Z^0 + g\,(\text{or } q) + X$, $Z \to \nu\bar{\nu}$; $p\bar{p} \to W + g\,(\text{or } q) + X$, $W \to e\nu$ where the electron is slow or inside the jet; $p\bar{p} \to q\bar{q}$ followed by semileptonic decays with slow leptons. However there is a possible excess of such events. SUSY is a candidate for an alternative solution, since it provides new sources of missing energy. First of all I consider the monojet production coming from gluino decays. The cross section of the gluino pair production $p\bar{p} \to \tilde{g}\tilde{g} + X$ is large, for light gluinos, as can be seen from the analysis of ref.[19]. The gluino most likely decays $\tilde{g} \to q\bar{q}\tilde{\gamma}$ through exchange of a squark or $\tilde{g} \to g\tilde{\gamma}$ with a one loop graph mediated by a quark and a squark. In order to give significant missing energy, the two gluinos should carry sizable momenta. Therefore they are expected to emerge back-to-back, and so are the photinos coming from their decay. In this way the two photino momenta cancel and we are left with no sources of large missing energy. On the contrary, if one photino carries a large fraction of the gluino momentum and the other photino carries very little of its gluino momentum, one finds two jets and missing transverse energy (from photinos). One of the two jets is energetic but the other one will have energy less than 12 GeV. The result is the appearance of a single jet with a large missing transverse energy, i.e a monojet. It has been observed [23] that an efficient source of monojet events in the collider is given by $gg \to \tilde{g}\tilde{g}g$, where the gluon is energetic and the two gluinos are back-to-back with the gluon. Since the gluinos are emitted in the same direction, the two photinos will go parallel and the momenta add. We are in the situation of an energetic jet coming from the harder gluon recoiling against missing energy. Thus there are various sources of monojet events coming from gluino production. The numerical study of these processes confirms this expectation. The authors of ref.[24] find mainly monojet configurations for $M_{gl} < 40$ GeV. The observed relatively absence of such events suggests the bound [24]:

$$M_{gl} \gtrsim O(40)\ \mathrm{GeV} \qquad (4.8)$$

However, because of all the uncertainties in the computation of $\tilde{g}\tilde{g} \to$ monojet event, other authors [25] conclude that gluino masses O(5) GeV are still allowed by 1983 UA1 data. In this way it is still open the possibility of a scenario [26] with light gluinos ($\sim$ few GeV) and heavy squarks ($\sim$ O(100) GeV). Anyway, apart from the collider results, this scenario has another difficulty. Actually one expects the photino mass to be almost one order of magnitude less than the gluino mass, in these models. Such a light photino would contradict the cosmological bounds previously discussed.

A more complete analysis of the missing energy events is performed in ref.[27], keeping into account the 1984 UA1 data and all monojets, dijets and multijets. As a result they plot the missing energy event in a gluino mass - squark mass plane. They impose the limit on the number of events according to the UA1 collaboration results and thus the following bounds are set [27]:

$$\begin{aligned} \tilde{m}_q &> 65 \text{ GeV} \quad &\text{for} \quad M_{gl} &\simeq 150 \text{ GeV} \\ \tilde{m}_q &> 75 \text{ GeV} \quad &\text{for} \quad M_{gl} &\simeq 80 \text{ GeV} \\ M_{gl} &> 60 \text{ GeV} \quad &\text{for} \quad \tilde{m}_q &\simeq 100 \text{ GeV} \\ M_{gl} &> 70 \text{ GeV} \quad &\text{for} \quad \tilde{m}_q &\simeq 80 \text{ GeV} \end{aligned} \tag{4.9}$$

Furthermore the 1984 UA1 run fairly reduces the possibility of having light gluinos ($\sim$ 5 GeV) and heavy squarks ($\sim$ 100 GeV).

The next issue is whether SUSY could explain any excess of monojets (if such excess actually exists). For a gluino lighter than 40 GeV, there are far to many monojets compared with what is observed. For a gluino heavier than 80 GeV (and heavy squarks) no monojets are predicted. With a gluino mass in the range 60-70 GeV, it is possible to get the correct monojet rate. Unfortunately in this case the dijet to monojet rate is 2-1 to 3-1, while 1984 UA1 run has reported 23 monojet events (15-17 of which estimated from SM background) and 2 dijet events (2 of which estimated from SM background). In conclusion the $p\bar{p}$ results set stringent limits on gluino (and squark) masses. In the present situation the SUSY models can not provide a solution to the possible excess of the missing transverse energy events.

Charged sleptons

The charged sleptons are the spin 0 partners of the electron, muon, τ ,... For each lepton there are two corresponding scalar particles, partners of the left and right-handed fermion. The mass eigenstates are linear combinations of these two states, since mixings between the partners of the right and left-handed leptons are present in the SUSY models under consideration.

The dominant decay is $\tilde{\ell}^{\pm} \to \ell^{\pm} \tilde{\gamma}$, if the photino is light enough. Since the photino is undetected, the signal will be given by a lepton and missing energy of order 50%. More exotic decays are possible (like the β -decay $\tilde{\ell}^{\pm} \to \ell^{\pm} \nu \tilde{\nu}$), but the decay into photino is certainly the most popular.

The slepton could be produced in $p\bar{p}$ experiments via $q\bar{q} \to Z^{\circ} \to \tilde{\ell}^{+} \tilde{\ell}^{-}$. The signal would be $\ell^{+} \ell^{-}$ + missing energy: no such events have been reported, but no limits have been extracted by their absence. Another possibility is $W^{\pm} \to \tilde{\ell}_L^{\pm} \tilde{\nu}$ assuming $\tilde{\ell}_L$ and $\tilde{\nu}$ light enough to allow the decay. The signal is $\ell^{\pm}$ + missing energy, to be distinguished from $W^{\pm} \to \ell^{\pm} \bar{\nu}$. The negative results imply [22]:

$$\tilde{m}_{e_L} \gtrsim 26 \text{ GeV} \quad \text{for} \quad \tilde{m}_{e_L} = \tilde{m}_{\nu} \tag{4.10}$$

The best place to look for sleptons is in $e^{+}e^{-}$ experiments. The unsuccessful searches of JADE [28] and Mark J [29] collaborations for sleptons in the process $e^{+}e^{-} \to \tilde{\ell}^{+} \tilde{\ell}^{-}$, the slepton decaying $\tilde{\ell}^{\pm} \to \ell^{\pm} \tilde{\gamma}$, set the bounds:

$$\begin{array}{llll} \tilde{m}_{e} > 22\text{-}23 \text{ GeV} & \text{for} & \tilde{m}_{\gamma} \lesssim 19 \text{ GeV} & \tilde{m}_{e_L} = \tilde{m}_{e_R} \\ \tilde{m}_{\mu} > 20.9 \text{ GeV} & \text{for} & \tilde{m}_{\gamma} \lesssim 15\text{-}16 \text{ GeV} & \tilde{m}_{\mu_L} = \tilde{m}_{\mu_R} \\ \tilde{m}_{\tau} > 18 \text{ GeV} & \text{for} & \tilde{m}_{\gamma} \lesssim 13 \text{ GeV} & \tilde{m}_{\tau_L} = \tilde{m}_{\tau_R} \end{array} \tag{4.11}$$

Another way to look for sleptons in $e^{+}e^{-}$ experiments is to consider the process $e^{+}e^{-} \to e^{\pm} \tilde{e}^{\mp} \tilde{\gamma}$. JADE [28] and CELLO [16-30] collaborations find:

$$\tilde{m}_e > 25 \text{ GeV} \quad \text{for} \quad \tilde{m}_\gamma = 0 \quad \tilde{m}_{e_L} = \tilde{m}_{e_R}$$
$$\tilde{m}_e > 20 \text{ GeV} \quad \text{for} \quad \tilde{m}_\gamma = 10 \text{ GeV} \quad \tilde{m}_{e_L} = \tilde{m}_{e_R} \tag{4.12}$$

The best limits on slepton masses come from $e^+e^- \to \gamma\tilde{\gamma}\tilde{\gamma}$, where the photon is used to tag the process, otherwise invisible. However these limits are very sensitive to $\tilde{m}_\gamma$. The results from ASP experiments [31] are:

$$\tilde{m}_{e_L} \gtrsim 42 \text{ GeV} \quad \text{for} \quad \tilde{m}_\gamma = 0 \quad \tilde{m}_{e_R} >> \tilde{m}_{e_L}$$
$$\tilde{m}_e \gtrsim 51 \text{ GeV} \quad \text{for} \quad \tilde{m}_\gamma = 0 \quad \tilde{m}_{e_R} = \tilde{m}_{e_L}$$
$$\tilde{m}_e \gtrsim 48 \text{ GeV} \quad \text{for} \quad \tilde{m}_\gamma = 5 \text{ GeV} \quad \tilde{m}_{e_R} = \tilde{m}_{e_L}$$
$$\tilde{m}_e \gtrsim 33 \text{ GeV} \quad \text{for} \quad \tilde{m}_\gamma = 10 \text{ GeV} \quad \tilde{m}_{e_R} = \tilde{m}_{e_L}$$

Sneutrino

The sneutrino is the spin 0 partner of the neutrino. Since only left-handed neutrinos are observed in nature, I assume one scalar particle for each generation of neutrinos. The sneutrino is a possible candidate to be the lightest SUSY particle. Differently from the case of the photino and higgsino, the cosmological bounds on a stable neutrino are very mild. One finds [32] no restriction on the sneutrino mass, as long as the Majorana z-ino mass is larger than 3.5 GeV. If the z-ino is lighter than 3.5 GeV (which is very unlikely) any stable sneutrino is constrained in the range $m_\tau \lesssim \tilde{m}_\nu \lesssim 3.5$ GeV, where m_τ is the τ lepton mass. The absence of stringent cosmological bounds on sneutrino masses is due to the fact that sneutrinos can pair annihilate into neutrinos through neutralino exchange, which is a process non P-wave suppressed differently from the photino case (which is a Majorana particle).

The only experimental limit on sneutrino mass is recovered from the analysis of τ decays [33]. If $\tilde{m}_{\nu_\tau} < m_\tau$ and if there exists another $\tilde{\nu}_\ell$ with $\tilde{m}_{\nu_\ell} < m_\tau$, the process $\tau^\pm \to \ell^\pm$ + "nothing" has a new SUSY channel, as shown in fig.4.1.

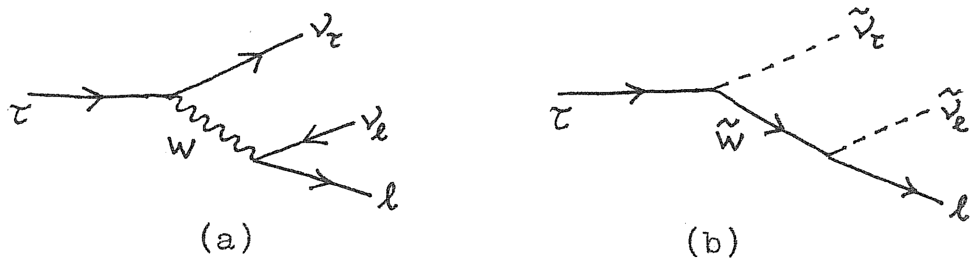


Fig. 4.1. SM (a) and SUSY (b) contribution to $\tau^\pm \to \ell^\pm$ + "nothing".

The contribution of the graph of fig. 4.1(b) is constrained by the experimental results on the τ lifetime, the lepton energy spectrum and the branching ratio of the decay. The numerical analysis sets a combined limit on the masses of $\tilde{\nu}_\tau$, $\tilde{\nu}_\ell$ and $\tilde{W}$ [33], which can be approximately summarized by the requirement $\tilde{m}_{\nu_\tau} + \tilde{m}_{\nu_\ell} > m_\tau$, if $\tilde{m}_W \lesssim m_W$. If $\tilde{m}_W > m_W$, lighter $\tilde{m}_\nu$ are allowed.

If the sneutrino is the lightest sparticle, the photino is allowed to decay $\tilde{\gamma} \to \nu\tilde{\nu}$. It is interesting to remark that, even in this case, the photino could remain unobservable. If the sneutrino is not the lightest sparticle, then it can decay. The dominant mode is $\tilde{\nu} \to \nu\tilde{\gamma}$, but also 3-body and 4-body decays with leptons and/or quarks in the final state are allowed.

The sneutrino could be produced in e^+e^- experiment either directly or through the decay of Z^0 or $W^\pm$ bosons, but no experimental limits are given.

<u>Squarks</u>

The squarks are the spin 0 partners of the quarks. Analogously to what happens with leptons, two scalars correspond to each fermion. For instance the partners of the down quark are two spin 0 particles, color triplets, one of which is a SU(2) singlet while the other one belongs to a SU(2) doublet.

The dominant decay of the squark is $\tilde{q} \to \tilde{g} q$, since it is a strong interaction mediated process. Subsequently the gluino decays with a photino in the final products, leaving missing energy as a signature of the process. If the gluino is heavier than the scalar squark, the dominant decay becomes $\tilde{q} \to \tilde{\gamma} q$. The signal is, in this case, much clearer since the missing energy is O(50 %). If both the decay modes are kinematically forbidden, and a sneutrino is the lightest sparticle, the squark can still decay $\tilde{q} \to \tilde{\nu} \ell q'$ (via w-ino exchange) or $\tilde{q} \to \tilde{\nu} \nu q$ (via z-ino exchange). Finally, as I showed previously, the squark could be the lightest SUSY particle, assuming the lightest hadron state to be neutral. JADE collaboration [16] has studied all these possibilities in e^+e^- experiments. For 2/3 charge squark mass, the following limits are set [16]:

If $\tilde{q} \to q\tilde{g}$, $\tilde{m}_q > 19.2$ GeV for $\tilde{m}_{q_L} = \tilde{m}_{q_R}$, $M_{g\ell} = 3$ GeV, $\tilde{m}_\gamma = \frac{M_{g\ell}}{6} = 0.5$ GeV

" " , $\tilde{m}_q > 20.0$ GeV for $\tilde{m}_{q_L} = \tilde{m}_{q_R}$, $M_{g\ell} = 10$ GeV, $\tilde{m}_\gamma = \frac{M_{g\ell}}{6} = 1.7$ GeV

" " , $\tilde{m}_{q_L} < 11.3$ GeV or $\tilde{m}_{q_L} > 17.8$ GeV for $\tilde{m}_{q_L} \ll \tilde{m}_{q_R}$, $M_{g\ell} = 3$ GeV, $\tilde{m}_\gamma = \frac{M_{g\ell}}{6} = 0.5$ GeV

If $\tilde{q} \to q\tilde{\gamma}$, $\tilde{m}_q > 21.4$ GeV for $\tilde{m}_{q_L} = \tilde{m}_{q_R}$, $\tilde{m}_\gamma < 10$ GeV

" " , $\tilde{m}_{q_L} < 3.2$ GeV or $\tilde{m}_{q_L} > 21.0$ GeV for $\tilde{m}_{q_L} \ll \tilde{m}_{q_R}$, $\tilde{m}_\gamma < 10$ GeV

If $\tilde{q} \to q'\tilde{\nu}\ell$ or $\tilde{q} \to q\tilde{\nu}\nu$, $\tilde{m}_q > 20.8$ GeV for $\tilde{m}_{q_L} = \tilde{m}_{q_R}$, $\tilde{m}_Z = \tilde{m}_W \simeq m_W$, $\tilde{m}_\nu = 1$ GeV

" " , $\tilde{m}_{q_L} < 5.0$ GeV or $\tilde{m}_{q_L} > 20.0$ GeV for $\tilde{m}_{q_L} \ll \tilde{m}_{q_R}$, $\tilde{m}_Z = \tilde{m}_W \simeq m_W$, $\tilde{m}_\nu = 1$ GeV

If $\tilde{q}$ stable, $\tilde{m}_q > 15.0$ GeV for $\tilde{m}_{q_L} = \tilde{m}_{q_R}$

" " , $\tilde{m}_{q_L} < 2.0$ GeV or $\tilde{m}_{q_L} > 15.0$ GeV for $\tilde{m}_{q_L} \ll \tilde{m}_{q_R}$

Note that the best limits come from $\tilde{q} \to q\tilde{\gamma}$ decay mode.

The $p\bar{p}$ experiments give even more stingent bounds than those recovered from e^+e^- experiments. The analysis of the $p\bar{p}$ results has already been discussed in the section devoted to gluinos. I showed that the result is a plot of the number of transverse missing energy events in a gluino mass - squark mass plane. Then limits on SUSY particle masses can be inferred from UA1 data. The results are those of formula (4.9).

Chargino

The mass eigenstates of the mixtures of the spin 1/2 partners of charged gauge and Higgs bosons are referred as charginos. The two Weyl spinors of the $W^{\pm}$ real superfields and the two Weyl spinors of the chiral superfields containing the charged Higgs mix together giving rise to two Dirac fermions (the charginos),as soon as the electroweak symmetry is spontaneously broken. This is a consequence of the Higgs mechanism in presence of SUSY. Since I want to discuss the limits on the masses in a model independent way, I will not specify the mixings among the fermionic states, which are sensitive to the specific cases.

Also the decay modes are very model dependent; I summarize possible chargino decays in fig. 4.2.

Charginos have been searched in e^+e^- experiments. They can be pair produced $e^+e^- \to \tilde{\chi}^+\tilde{\chi}^-$ through γ or Z^0 exchange in the s-channel or through sneutrino exchange in the t-channel. Assuming that they decay $\tilde{\chi}^{\pm} \to \ell^{\pm}\tilde{\nu}$, as shown in fig.4.2(a), and that $\tilde{\nu} \to \tilde{\gamma}\nu$, Mark J [34] and JADE [35] collaborations exclude the following region for chargino masses:

$$\tilde{m}_{\nu} < \tilde{m}_{\chi} < 22.5 \text{ GeV} \quad \text{(excluded)} \qquad (4.13)$$

If the sneutrino is heavy and the gluino is lighter than the chargino, the decay mode in fig. 4.2(b) with gluino production may become dominant. In this case JADE [35] finds the limit:

$$\tilde{m}_\chi > 22.6 \text{ GeV} \quad \text{for } M_{g\ell} = 10 \text{ GeV}, \tilde{m}_\gamma = \frac{M_{g\ell}}{6} = 1.7 \text{ GeV} \qquad (4.14)$$

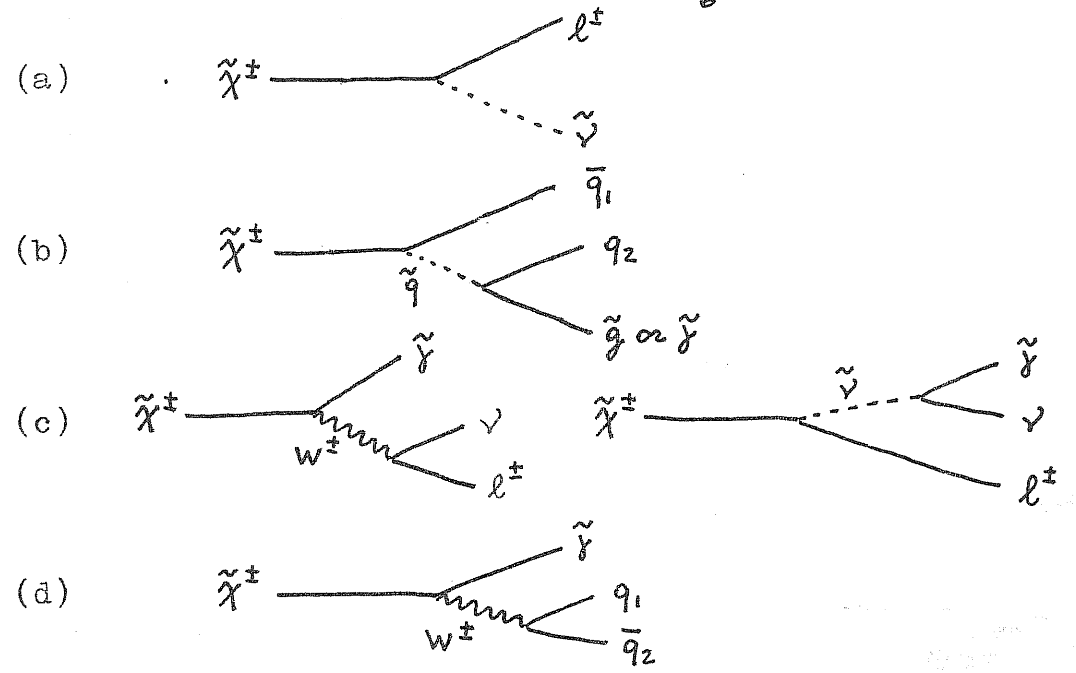


Fig. 4.2. Decay modes of the chargino ($\tilde{\chi}^\pm$).

If both these decay modes are kinematically excluded, the chargino decays as in fig. 4.2(c) or (d). Limits are derived by JADE collaboration [35], which depend on unknown branching fractions. However, independently of these unknown parameters and the photino mass, one can set the bound:

$$\tilde{m}_\chi > 22.5 \text{ GeV} \qquad (4.15)$$

In conclusion, from e^+e^- experiments a chargino mass lower bound of 22.5 GeV is found, independently of the decay modes.

Signatures of the chargino could be found at the $p\bar{p}$ collider, too. For instance the process $W^\pm \to \tilde{\gamma}\, \tilde{\chi}^\pm$, $\tilde{\chi}^\pm \to \ell^\pm \nu \tilde{\gamma}$ has identifiable characteristics. Furthermore it was found [36] that the decay $Z \to \tilde{\chi}^+\tilde{\chi}^-$ can give spectacular dileptons events with high rate, for slepton masses O(20-30) GeV. Non observation of this decay could give soon the new bound $\tilde{m}_\chi >$ O(40) GeV.

Neutralino

The neutralinos are the mass eigenstates, mixtures of the spin 1/2 partners of the two neutral Higgs and two gauge bosons. They can be arranged as four neutral Majorana spinors (one of which is generally assumed to be the photino). As for the charginos, the unknown masses and mixings imply significant uncertainties in couplings and cause difficulties in calculations. Decay modes are model dependent; the possibilities shown in fig. 4.3 have been studied.

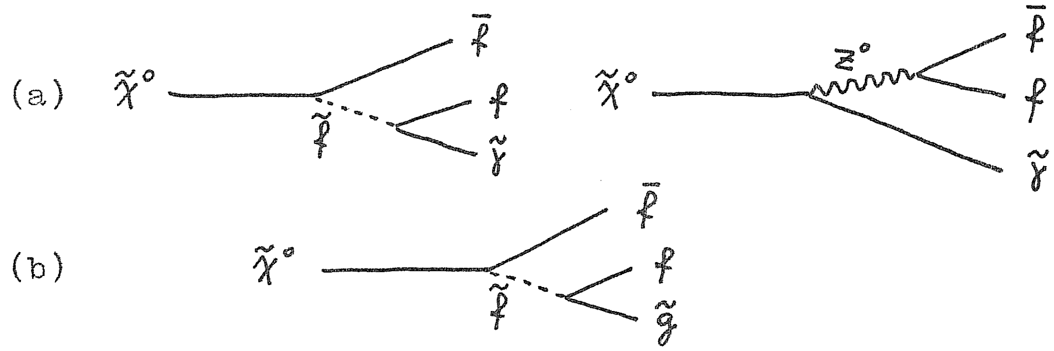


Fig. 4.3. Neutralino ($\tilde{\chi}^{o}$) decay modes. f denotes quarks and leptons.

The process $e^+e^- \to \tilde{\gamma}\tilde{\chi}^o, \tilde{\chi}^o \to \ell^+\ell^-\tilde{\gamma}$ leads to a dramatic signature with 2/3 of the total energy missing. Studying the electron production Mark J collaboration [34] excludes the region

$$6 \text{ GeV} < \tilde{m}_{\chi^o} < 33 \text{ GeV} \quad \text{(excluded)} \qquad (4.16)$$

assuming $\tilde{m}_\gamma$ = 2-5 GeV, $\tilde{m}_\ell \lesssim$ 50 GeV and BR($\tilde{\chi}^o \to e^+e^-\tilde{\gamma}$) $\simeq$ 5 %. Similar regions are excluded by JADE collaboration [37], with an analysis of the $\tilde{\chi}^o$ decay modes $e^+e^-\tilde{\gamma}$, $\mu^+\mu^-\tilde{\gamma}$, $q\bar{q}\tilde{\gamma}$ and $q\bar{q}\tilde{g}$. The event $e^+e^- \to \tilde{\chi}^o\tilde{\chi}^o$ could give a signal $e^+e^-\mu^+\mu^-$ with 1/3 missing energy. The neutralinos may be produced at hadron collider too, but no limits have been given.

As I have already mentioned, the neutralino is allowed to be the lightest sparticle. The phenomenology of such models is discussed in ref.[38]. Ref.[14] studies the cosmological consequences of these possibilities. Analogously to the cosmological bounds on photino masses discussed previously, a limit on an approximate higgsino eigenstate is found, excluding the region [14]:

$$100 \text{ eV} \lesssim \tilde{m}_{\chi^0} \lesssim 5 \text{ GeV} \qquad (4.17)$$

5. K^o-$\bar{K}^o$ SYSTEM

The K^o-$\bar{K}^o$ system has always been a good testing ground for new theories. The GIM mechanism insures a natural explanation for the suppression of FCNC processes. No $\bar{q}_1 q_2 Z^o$ (with $1 \neq 2$) couplings are present at tree level and the one loop quantum corrections are further suppressed by $\Delta m_q^2/m_q^2$, where Δm_q^2 is the difference of squared masses of quarks with the same charge. Any extensions of the SM must give as well a natural explanation of these delicate cancellations. The K^o-$\bar{K}^o$ system seems to be able to provide constraints on any new physics and, in particular, on SUSY. I will turn now to discuss the impact of SUSY in the computation of K_L-K_S mass difference.

The short distance contribution to the K_L-K_S mass difference is given by

$$\Delta m_K \equiv m_{K_L} - m_{K_S} = \frac{1}{m_K} \mathrm{Re} \langle \bar{K}^o | H_{eff} | K^o \rangle \tag{5.1}$$

The one loop effective Hamiltonian in the SM is recovered by computing the graphs of fig. 5.1.

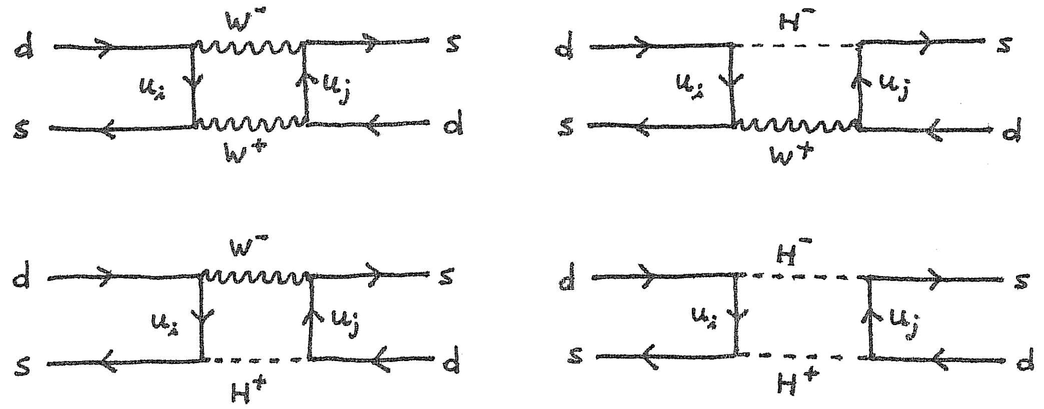


Fig. 5.1. SM contributions to $\bar{s}d\bar{s}d$ effective operator; i and j are generation indeces.

The direct computation gives the effective operator:

$$\frac{G_F}{\sqrt{2}} \frac{\alpha}{4\pi \sin^2\theta_W} \tilde{E} \left(\bar{s}_L \gamma_\mu d_L\right)^2 \tag{5.2}$$

where $\tilde{E}$ can be written, exploiting the unitarity of the KM matrix V:

$$\tilde{E} = \sum_{i,j=2}^{N} V_{is}^* V_{id} V_{js}^* V_{jd} \; E(x_i, x_j) \tag{5.3}$$

N is the number of quark generations, $x_i = \frac{m_{U_i}^2}{m_W^2}$, and the quarks of the first generation are assumed to be massless. Finally:

$$E(x_i, x_j) = -x_i x_j \left\{ \frac{1}{x_i - x_j} \left[\frac{1}{4} - \frac{3}{2} \frac{1}{x_i - 1} - \frac{3}{4} \frac{1}{(x_i - 1)^2} \right] \ln x_i + \right. \tag{5.4}$$
$$\left. + \left(x_i \to x_j \right) - \frac{3}{4} \frac{1}{(x_i - 1)(x_j - 1)} \right\}$$

Substitution in (5.1) yields:

$$\left(\frac{\Delta m_K}{m_K}\right)_{SM} = -\frac{G_F}{\sqrt{2}} \frac{\alpha}{6\pi \sin^2\theta_W} f_K \, B \, \mathrm{Re}\,\tilde{E} = -3.8 \cdot 10^{-10} B \, \mathrm{Re}\,\tilde{E} \tag{5.5}$$

B is the parameter measuring the ratio between $\langle \bar{K}^0 | (\bar{s}_L \gamma^\mu d_L)^2 | K^0 \rangle$ and the same matrix element with vacuum insertion. f_K is the usual meson decay constant. B is unknown, even if several attempts to compute its value have been performed. Different values are found, using different approximations:

Approximation	Ref.	B
vacuum insertion		1
harmonic oscillator	[39]	1.4÷2.9
vacuum + 1 particle intermediate states	[40]	1.2÷1.6
" "	[41]	0.9÷1.2
static quark model	[42]	≃ 1
leading order of 1/N expansion	[43]	0.75
MIT-bag model	[39-42-44]	-0.4÷0.5
chiral SU(3) and PCAC	[45]	0.33
sum rules	[46]	0.24÷0.42
lattice (first calculations)	[47]	≃ O(1)
lattice	work in progress	?

Experimentally

$$\left(\frac{\Delta m_K}{m_K}\right)_{exp} = (7.074 \pm 0.029)\cdot 10^{-15} \qquad (5.6)$$

Keeping into account only the charm quark contribution to Δm_K, eq. (5.5) yields:

$$\left(\frac{\Delta m_K}{m_K}\right)_{charm} \simeq 4\cdot 10^{-15}\ B \qquad (5.7)$$

Therefore, if B is less than 1, the charm quark is insufficient by itself to account for Δm_K. This could mean that the top quark KM couplings are large enough and the top quark mass small enough to explain the enhancement of the experimental result [48]. The discrepancy could also be removed by introducing a fourth generation [49] or by invoking a left-right symmetric extension of the SM [50]. The long distance contributions could play an interesting role as well, but the actual computation has large uncertainties [51]. From a less conservative point of view, one could think that the extra contribution is provided by SUSY effects. I then turn to study if SUSY can give new answers or if Δm_K sets bounds on SUSY parameters.

In a SUSY theory, in addition to the graphs of fig. 5.1, new graphs contributing to Δm_K are present. First of all one extra Higgs is exchanged in fig. 5.1, since in a SUSY theory two Higgs doublets are required to give mass to all the quarks. Then there is the supersymmetrization of the graphs of fig. 5.1, where the chargino is exchanged (fig. 5.2(a)) and the new gluino and neutralino exchange diagrams (fig. 5.2(b)). As shown in section 3, the gluino interaction with quarks and squarks of different generations comes from the non simultaneous diagonalization of quark and squark mass matrices and

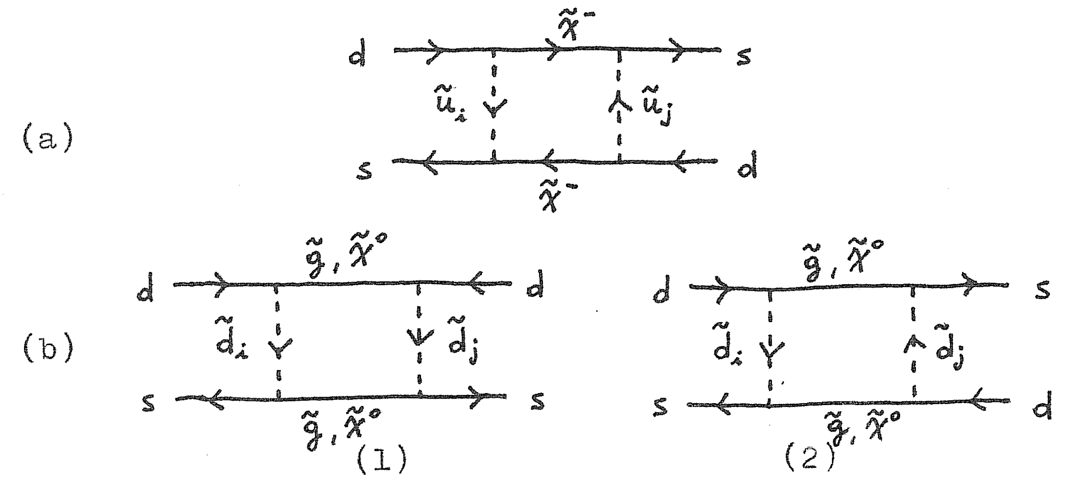


Fig. 5.2. SUSY contributions to $\bar{s}d\bar{s}d$ effective operator.

provides new sources of FCNC. The dominant contribution comes from the gluino exchange, because of the presence of the strong coupling constant. The graph of fig. 5.2(a) has been computed in ref.[52], but no major constraints for the minimal models have been found.

Let us compute the gluino exchange diagrams[53]. Because of the Majorana nature of the gluino spinor, there are two different diagrams which give rise to the $\bar{s}d\bar{s}d$ effective operator, as shown in fig. 5.2(b 1-2). If I treat the terms called left-right squark mixings in section 3 as mass insertions in the squark propagators, then the major contribution is obtained when the external quarks are all left-handed. From a perturbative point of view, this requires no mass insertions. If external quarks are all right-handed, no flavor mixing is possible unless left-right insertions are introduced. Choosing left-handed quarks, graph 5.2(b 1) yields, in the approximation of vanishing external momenta:

$$-i\,4\left(\frac{g_s}{\sqrt{2}}\right)^4 \bar{s}_\omega (1+\gamma_5) C \bar{s}_\psi^T \,(T^\alpha T^\beta)_{\psi\chi}\, d_\chi^T C^{-1}(1-\gamma_5)(T^\alpha T^\beta)_{\omega\varphi}\, d_\varphi \qquad (5.8)$$
$$\times \int \frac{d^4k}{(2\pi)^4} \frac{M_{gl}^2}{(k^2-M_{gl}^2)^2} \sum_{i,j=1}^{6} \frac{U_{2i}\, U_{1i}^*\, U_{2j}\, U_{1j}^*}{(k^2-m_i^2)\,(k^2-m_j^2)}$$

Greek indeces from the beginning of the alphabet (α , β) run from 1 to 8; Greek indeces from the end of the alphabet (ω , ψ , χ , φ) run from 1 to 3; summation over color indeces is understood. M_{gl} is the gluino mass and m_i are the eigenvalues of the down squark mass matrix. U is the down squark rotation matrix defined in (3.10). Using the following identities:

$$\bar{\psi}_1 (1+\gamma_5) \bar{\psi}_2^T \psi_3^T C^{-1} (1-\gamma_5) \psi_4 = -2 \bar{\psi}_{1L} \gamma^\mu \psi_{4L} \bar{\psi}_{2L} \gamma_\mu \psi_{3L} \tag{5.9}$$

$$\bar{\psi}_L T^\alpha T^\beta \gamma_\mu \phi_L \bar{\psi}_L T^\alpha T^\beta \gamma^\mu \phi_L = \frac{1}{9} \left(\bar{\psi}_L \gamma^\mu \phi_L \right)^2 \tag{5.10}$$

eq. (5.8) reduces to

$$\frac{2}{9} g_s^4 \, \bar{s}_L \gamma^\mu d_L \, \bar{s}_L \gamma_\mu d_L \; i \int \frac{d^4k}{(2\pi)^4} \frac{M_{gl}^2}{(k^2 - M_{gl}^2)^2} \left(\sum_{i=1}^{6} \frac{U_{2i} U_{1i}^*}{k^2 - m_i^2} \right)^2 \tag{5.11}$$

The contribution of the graph in fig. 5.2(b 2) is

$$i \, 4 \left(\frac{g_s}{\sqrt{2}} \right)^4 \bar{s}_\omega (T^\alpha T^\beta)_{\omega\varphi} \gamma^\mu (1-\gamma_5) d_\varphi \, \bar{s}_\chi (T^\beta T^\alpha)_{\chi\psi} \gamma^\nu (1-\gamma_5) d_\psi \times$$
$$\times \int \frac{d^4k}{(2\pi)^4} \frac{k_\mu k_\nu}{(k^2 - M_{gl}^2)^2} \sum_{i,j=1}^{6} \frac{U_{2i} U_{1i}^* U_{2j} U_{1j}^*}{(k^2 - m_i^2)(k^2 - m_j^2)} \tag{5.12}$$

With a Fierz rearrangement and with the help of the identity

$$\bar{\psi}_L T^\alpha T^\beta \gamma_\mu \phi_L \bar{\psi}_L T^\beta T^\alpha \gamma^\mu \phi_L = \frac{11}{18} \left(\bar{\psi}_L \gamma^\mu \phi_L \right)^2 \tag{5.13}$$

eq. (5.12) reduces to

$$-\frac{11}{18} g_s^4 \, \bar{s}_L \gamma^\mu d_L \, \bar{s}_L \gamma_\mu d_L \; i \int \frac{d^4k}{(2\pi)^4} \frac{k^2}{(k^2 - M_{gl}^2)^2} \left(\sum_{i=1}^{6} \frac{U_{2i} U_{1i}^*}{k^2 - m_i^2} \right)^2 \tag{5.14}$$

Therefore the effective operator is given by:

$$\frac{8\pi^2}{9} \alpha_s^2 \, \bar{s}_L \gamma^\mu d_L \, \bar{s}_L \gamma_\mu d_L \; i \int \frac{d^4k}{(2\pi)^4} \frac{(4 M_{gl}^2 - 11 k^2)}{(k^2 - M_{gl}^2)^2} \left(\sum_{i=1}^{6} \frac{U_{2i} U_{1i}^*}{k^2 - m_i^2} \right)^2 \tag{5.15}$$

The sum over squark eigenvalues can be evaluated at the zero order in left-right mixing insertions (i.e. partners of right-handed quarks decouple) and keeping only the dominant stop contribution (the stop is the less degenerated squark):

$$\sum_{i=1}^{6} \frac{U_{2i} U_{1i}^*}{k^2 - m_i^2} \simeq \frac{c \, V_{ts} V_{td}^* \, m_t^2}{(k^2 - \mu^2)^2} \tag{5.16}$$

where μ is a common down squark mass. Finally the computation of the integral yields the following effective operator

$$-\alpha_s^2 \;\frac{\left(c\, V_{ts} V_{td}^{*}\, m_t^2\right)^2}{M_{gl}^6}\; F(x)\;\left(\bar{s}_L \gamma^\mu d_L\right)^2 \qquad (5.17)$$

where $x = \frac{\mu^2}{M_{gl}^2}$ and

$$F(x) = \frac{\left(1 - \frac{5}{2}x + 38x^2 + \frac{11}{2}x^3\right)}{27\, x^2 (1-x)^4} + \frac{\left(x + \frac{5}{9}\right)\ln x}{(1-x)^5} \qquad (5.18)$$

A good check of the result is to verify that F(x) is finite in the limit $x \to 1$.

Given the expression (5.17) of the effective operator, I can write down the SUSY contribution to Δm_K, coming from gluino exchange:

$$\left(\frac{\Delta m_K}{m_K}\right)_{SUSY} \simeq \frac{2}{3}\, B\, f_K^2\, \alpha_s^2\; \frac{c^2 m_t^4}{M_{gl}^6}\; F(x)\; \mathrm{Re}\left[\left(V_{ts} V_{td}^{*}\right)^2\right] \qquad (5.19)$$

I impose now that the SUSY contribution does not exceed the difference between the experimental value of Δm_K and the value found in the SM:

$$\left(\frac{\Delta m_K}{m_K}\right)_{SUSY} \lesssim \left(\frac{\Delta m_K}{m_K}\right)_{exp} - \left(\frac{\Delta m_K}{m_K}\right)_{SM} \qquad (5.20)$$

Treating M_{gl} and μ as the only unknown quantities, the constraints (5.20) can be used to exclude a region in the gluino mass - squark mass plane. Few remarks about the parameters used are in order. I assume $\alpha_s \simeq 0.1$ at the typical loop momentum scale, c = 0.6 and m_t = 40 GeV. For such value of the top mass and for the KM angles given in [54], the top quark is not able to account for an enhancement of the charm quark result for Δm_K. In the numerical analysis I kept the charm contribution to $(\Delta m_K)_{SUSY}$, which a priori is not negligible since the mixing angles are much larger than the top quark case. B is taken 0.33 which is a propitious situation. The numerical computation of translating the bound (5.20) into a limit on masses yields to the graph of fig. 5.3 (°).

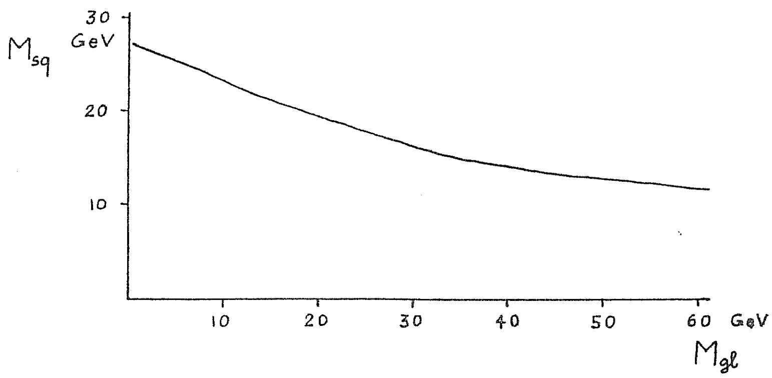


Fig. 5.3. Bound on squark and gluino masses coming from Δm_K; the region below the line is excluded.

The result is in agreement with the analysis of ref.[55], where the case of only two generations is studied. From fig. 5.3 it is apparent that Δm_K does not constraint the parameters of the minimal N=1 supergravity models. Low ($\lesssim$ 20 GeV) squark masses are excluded, but, as I showed in section 4, the experimental limits are much stronger. Even very light gluinos are allowed by the SUSY Δm_K computation. Furthermore, in view of this result, it is clear that the minimal SUSY models can not account for a sizeble contribution to Δm_K and are not good candidates to explain the possible discrepancy with the experimental result.

One can analyse these results from a different point of view. Instead of starting with a squark mass matrix relative to the minimal models, with no further constraints coming from Δm_K , one can consider a more general squark mass

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(°) I thank G. Degrassi for help in the numerical computation.

matrix coming from non minimal models and look for bounds set by Δm_K . These limits do not affect the squark masses themselves, but only the mass difference between squarks of same electric charge and different generation. From the diagrams with gluino exchange, bounds on the down squark splitting are obtained, as shown in fig. 5.4.

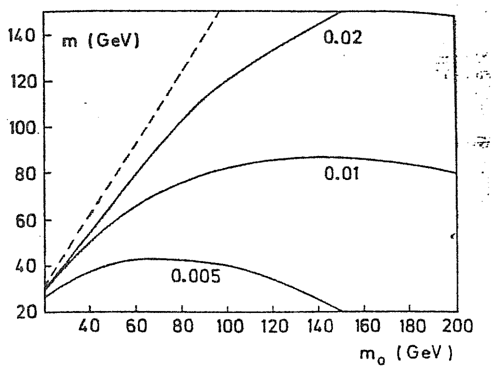


Fig. 5.4. Upper bounds on the down squark splitting $\Delta\mu/\mu$ in a two generations model [56]. m is the gluino mass, m_o is the average squark mass and B=1. No limits can be derived in the region above the dashed line.

Note that, in the minimal models previously discussed, the down squark splitting is order $\frac{c\, m_c^2}{2\mu^2}$, which is clearly below the upper bounds set by fig. 5.4.

Limits on the up squark mass splitting can be derived from the study of the diagrams with chargino exchange of fig. 5.2(a). Due to the presence of the weak coupling constant, the bounds on up squark mass splitting are approximately one order of magnitude smaller than those in fig. 5.4 [57]. The minimal models predict an up squark mass splitting order $\frac{m_c^2}{2\mu^2}$, in the two generations case.

In conclusion I found that the real part of the K^o-$\bar{K}^o$ mass matrix gives no further constraints in the case of minimal models. Bounds on squark mass splittings are found for non minimal models, where the squark mass matrix is more arbitrary.

I discuss now the impact of SUSY in the computation of the immaginary part of the K^o-$\bar{K}^o$ mass matrix. Because of new physical phases introduced by SUSY, one may expect that CP violation can be explained even choosing a vanishing δ_{KM} (the phase of the KM matrix with three generations). However the constraint on the electric dipole moment of the neutron (d_n) rules out the possibility of a large contribution to ε coming from phases appearing in the mass matrix of charginos (through the graph of fig. 5.2(a)) [52] or coming from the extra phase of the squark mass [55]. This phase is a combination of the A parameter and the phase of the (complex) Majorana mass of the gluino. The constraint $d_n < 10^{-24}$ e cm implies $|\delta| < 10^{-2}$ and the effect of the new phase δ in the computation of ε becomes negligible.

The dominant SUSY contribution to Im(K^o-$\bar{K}^o$) comes from the gluino exchange diagrams in fig. 5.2(b), and it is due to the KM phase δ_{KM}. Since the computation has already been illustrated previously, here I only quote the result for the SUSY contribution to the ε parameter [58]:

$$\varepsilon_{SUSY} = 12.5\ B\ \mathrm{Im}\, V_{ts} V_{td}^* \ c^2 \alpha_s^2 \left(\frac{M_W}{M_{gl}}\right)^6 \left(\frac{m_t}{m_c}\right)^2 \times \\ \times \left[-\mathrm{Re}\, V_{us} V_{ud}^* + \left(\frac{m_t}{m_c}\right)^2 \mathrm{Re}\, V_{ts} V_{td}^* \right] F\left(\frac{\mu^2}{M_{gl}^2}\right) e^{i\frac{\pi}{4}} \tag{5.21}$$

where F(x) is given in eq. (5.18). It can be noticed from (5.21) that ε_{SUSY} grows like m_t^4, while the SM contribution to ε goes like m_t^2. This means that the agreement with the experimental result can be achived even for a light top quark. Actually in view of the long lifetime of the b quark [54]:

$$\tau_B = (1.13 \pm 0.16)\cdot 10^{-12} \text{ sec} \qquad (5.22)$$

the SM, assuming B = 0.33, can not account for the observed value of ε :

$$Re\,\varepsilon = (1.621 \pm 0.088)\cdot 10^{-3} \qquad (5.23)$$

if the top quark is too light. SUSY can solve the possible impasse. Fig. 5.5 taken from ref.[58] shows the lower bound on top quark mass, coming from the computation of ε , as a function of the bottom quark lifetime. One can see that the

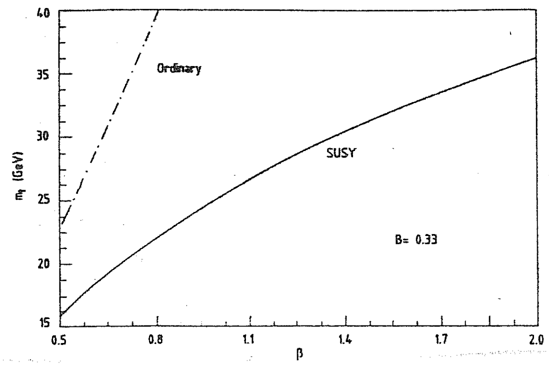


Fig. 5.5. Lower bounds on the top quark mass in the SM and in SUSY as a function of β (the bottom lifetime is $\tau_B = \beta\cdot 10^{-12}$ sec.), assuming B=0.33 [58].

SM is compatible only for $m_t \gtrsim 50$ GeV, while the SUSY models can explain the experimental value of ε , even with low top quark masses. Note that the result is very sensitive to c, since ε_{SUSY} is proportional to c^2 . Furthermore it is worthwhile mentioning that the SM can still account for the right value of ε , assuming B to be larger than 0.33.

The contribution of the gluino exchange diagrams to the electric dipole moment of the neutron, coming from δ_{KM} , does not exceed the experimental limit. However in SUSY models one expects values for the electric dipole moment of the neutron much larger than in the conventional case [59].

Finally I want to mention that SUSY can reduce the SM prediction for ε'/ε , because of the presence of the "pengluinos" (penguin graphs involving gluinos) [59].

6. $K^+ \to \pi^+ \nu \bar{\nu}$

Recently the kaon decay $K^+ \to \pi^+$ + "missing energy" has gained a lot of importance. The first reason is that the experimental situation will be soon dramatically improved: for the decay with a pair of undetected neutral particles, the 787 BNL experiment is expected to improve the branching ratio from the present limit of $1.4 \cdot 10^{-7}$ [60] to a value $\simeq 2 \cdot 10^{-10}$ [61]. From the theoretical point of view, the situation looks very interesting. The lack of significant long distance contributions to $K^+ \to \pi^+ \nu \bar{\nu}$ makes this process, unlike K^0-$\bar{K}^0$ mixings or $K_L \to \mu^+ \mu^-$ (which will be discussed later), a clear test of higher order corrections in the SM or of appearance of new physics. It has been suggested that the pion energy spectrum from the decay $K^+ \to \pi^+ \nu \bar{\nu}$ can give information about the mass of the τ neutrino if the latter lies in the range of tens of MeV [62]. It has also been pointed out that the pion energy spectrum sensitively depends on the spinorial nature (Dirac or Majorana) of the neutrinos [63]. Furthermore the value of the branching ratio, or new limits on it, can constraint the number of neutrinos (i.e. generations) or give hints in the direction of new physics.

As far as the SM is concerned, the process $K^+ \to \pi^+$ + "missing energy" is identified with $K^+ \to \pi^+ \nu \bar{\nu}$. This FCNC mediated process occurs at one loop level through the Feynman graphs depicted in fig. 6.1. Computing the graphs of fig. 6.1 in the approximation of vanishing external momenta, assuming the down and strange quarks and all the neutrinos to be massless and exploiting the SU(3) symmetry, the branching ratio for $K^+ \to \pi^+ \nu \bar{\nu}$ in the SM becomes [64]:

$$BR(K^+ \to \pi^+ \nu \bar{\nu}) = \frac{\alpha^2}{8\pi^2 \sin^4\theta_W V_{us}^2} \sum_{i=1}^{N} \left| \sum_{j=2}^{N} V_{js}^* V_{jd} D(x_j, y_i) \right|^2 \times \qquad (6.1)$$

$$\times BR(K^+ \to \pi^0 e^+ \nu_e)$$

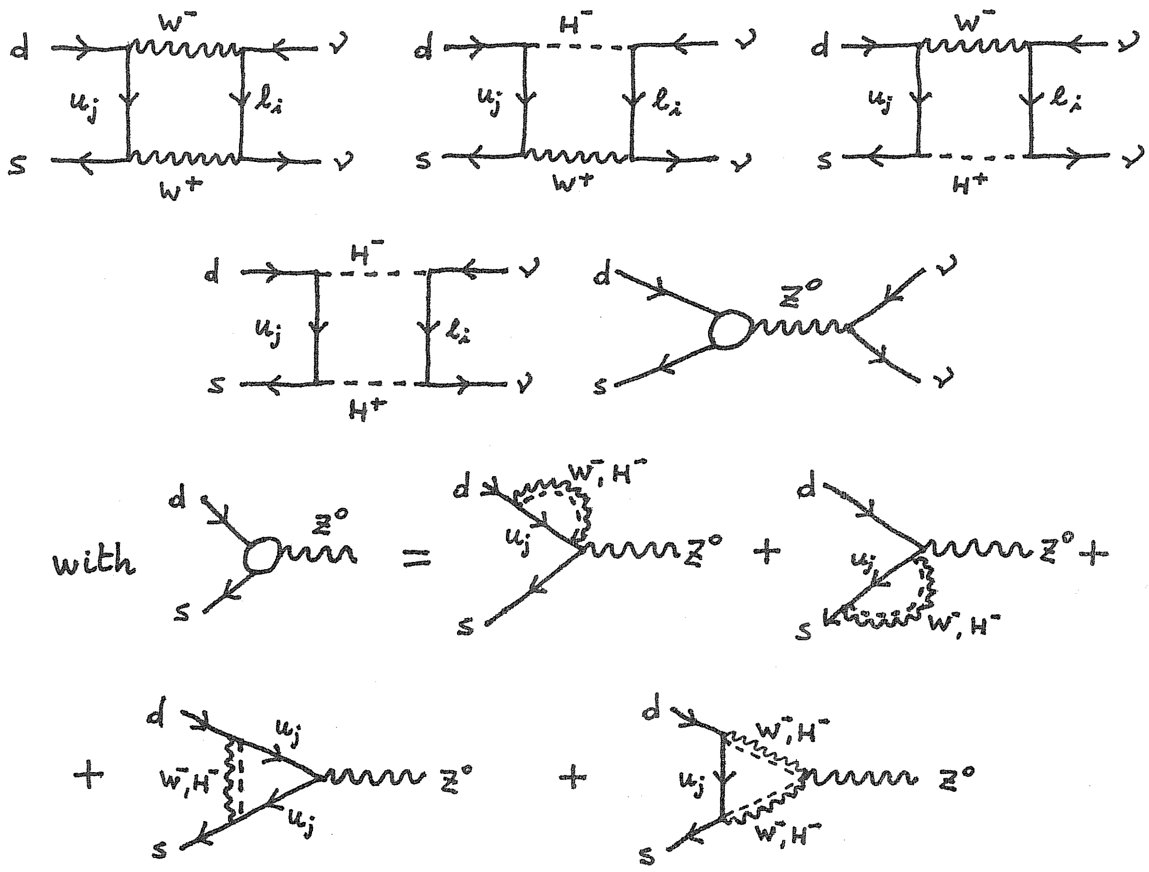


Fig. 6.1. SM graphs contributing to $K^+ \to \pi^+ \nu \bar{\nu}$

N is the number of generations of quarks and leptons, $x_j = \frac{m_{u_j}^2}{m_W^2}$, $y_i = \frac{m_{\ell_i}^2}{m_W^2}$ and

$$D(x,y) = -\frac{1}{8}\frac{yx}{y-x}\left(\frac{y-4}{y-1}\right)^2 \ln y + \frac{1}{8}\left[\frac{x}{y-x}\left(\frac{x-4}{x-1}\right)^2 + 1 + \frac{3}{(x-1)^2}\right] x \ln x + \frac{x}{4} - \frac{3}{8}\left(1 + \frac{3}{y-1}\right)\frac{x}{x-1} \qquad (6.2)$$

Note that the hadronic matrix element present in the calculation of BR($K^+ \to \pi^+ \nu \bar{\nu}$) is related to the matrix element of the decay $K^+ \to \pi^0 e^+ \nu$ which, in turn, is deduced from experiments; no large uncertainties due to hadrons are present. The QCD corrections for $K^+ \to \pi^+ \nu \bar{\nu}$, computed in [65] do not considerably modify my result. Neglecting lepton masses and replacing numbers, I obtain

$$BR(K^+ \to \pi^+ \nu\bar{\nu}) = 1.3 \cdot 10^{-5} N \left| \sum_{j=2}^{N} V_{js}^* V_{jd} D(x_j, 0) \right|^2 \qquad (6.3)$$

For N=3 and $m_t = 40$ GeV, the branching ratio is $O(10^{-10})$, much below the experimental limit.

Assuming the existance of a fourth generation, an extra decay channel in the new neutrino is open. Ref.[66] shows that it is possible to choose mixing angles of the fourth generation such that the decay is dominated by the exchange of the new heavy quark. Although there are various constraints to be satisfied by a 4x4 KM matrix (ε , ε'/ε and $BR(K_L \to \mu^+\mu^-)$) [67], a solution with large mixing angles between third and fourth generation is allowed. This situation in which the fourth generation has larger mixings with the third than with the first two generations is called "flavor flip" [68]. In this case, the presence of the new heavy quark is numerically more important than the fourth neutrino type in the final state. For a mass of the new quark between 200 GeV and 400 GeV, the following upper bound on the branching ratio is found [66]

$$BR(K^+ \to \pi^+ \nu\bar{\nu})_{4\ \text{generations}} < 8 \cdot 10^{-9} \qquad (6.4)$$

Then, with the impact of a fourth generation, it is possible to increase the $K^+ \to \pi^+ \nu\bar{\nu}$ branching ratio of more than one order of magnitude with respect to the SM with three generations.

Experiments have searched for a decay $K^+ \to \pi^+$ + "missing energy" with a pion spectrum from a two body decay, finding the limit [60](°):

$$BR(K^+ \to \pi^+ + \text{"undetected particle"}) < 3.8 \cdot 10^{-8} \qquad (6.5)$$

(°) The experiment 787 at BNL will improve the limit on the branching ratio up to an initial sensibility of $4 \cdot 10^{-11}$.

Although these experiments were set to search for the axion, these results can be used today to find limits on the presence of a possible new force. Recently a reanalysis of the Eötvös experiments has discovered a possible compatibility with the existence of a new intermediate-range force (the so called fifth force) [69]. Assuming that the new force is mediated by a photonlike light vector particle, called hyperphoton (γ'), one can compute the branching ratio of the process $K^+ \to \pi^+ + \gamma'$ [70]. The experimental result (6.5) sets the following limit:

$$\frac{f^2}{4\pi m_{\gamma'}^2} < 0.7 \cdot 10^{-25}\ eV^{-2} \qquad (6.6)$$

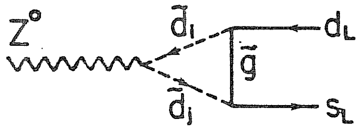
where $m_{\gamma'}$ is the hyperphoton mass and f measures its coupling. This bound is considered to be conservative since both the long-distance effects and the inclusion of the penguin diagram contributions make the limit stronger. The reanalysis of Eötvös data suggests an hyperphoton with [69] $m_{\gamma'} \simeq 10^{-9}$ eV and $\frac{f^2}{4\pi} \simeq 10^{-40}$, which gives

$$\frac{f^2}{4\pi m_{\gamma'}^2} \simeq 10^{-22}\ eV^{-2} \qquad (6.7)$$

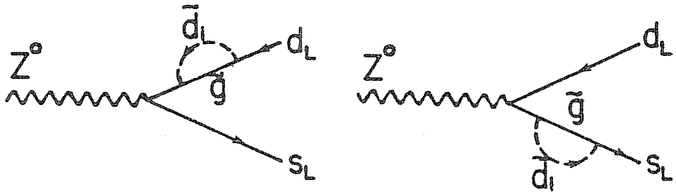
Thus the limit (6.6) rules out the presence of the hyperphoton suggested by the speculations of ref.[69].

I come back now to $K^+ \to \pi^+ \nu \bar{\nu}$, and I look for the consequences of SUSY to this process. I may have new contributions coming from SUSY virtual intermediate states, but SUSY may contribute also through new decay channels $K^+ \to \pi^+$ + "missing energy", which are experimentally undistinguishable. Here I confine myself to the first kind of contributions, leaving the second for a next section. Since there is a large number of possible diagrams with exchange of SUSY particles, to be clearer, I will subdivide all the graphs into different classes and discuss them separately. I will start with the graphs with gluino exchange which, at first sight, seem the most promising ones. Then I will turn to the box diagrams with w-ino and z-ino exchange and finally I will discuss the graphs with a $d\bar{s}Z$ effective vertex coming from chargino exchange.

As I have previously discussed, the gluinos have flavor changing interactions, which give rise to an effective $d\bar{s}Z$ vertex through the graphs of fig. 6.2



(a)



(b)

Fig. 6.2. Gluino contribution to the $d\bar{s}Z$ effective vertex.

Evaluating the graphs at vanishing external momenta, I obtain:

$$(a) = \frac{2 g_s^2 g}{\cos\theta_W} \bar{s}_L \gamma^\mu d_L \, i \int \frac{d^d k}{(2\pi)^d} \frac{k^2}{k^2 - M_{gl}^2} \sum_{i,j=1}^{6} \frac{U_{2i}\left[\sum_{k=1}^{3} U^*_{ki} U_{kj} - \frac{2}{3}\sin^2\theta_W \delta_{ij}\right] U^*_{1j}}{(k^2 - \tilde{m}_i^2)(k^2 - \tilde{m}_j^2)} \tag{6.8}$$

$$(b) = -\frac{2 g_s^2 g}{\cos\theta_W}\left(1 - \frac{2}{3}\sin^2\theta_W\right) \bar{s}_L \gamma^\mu d_L \, i \int \frac{d^d k}{(2\pi)^d} \frac{k^2}{k^2 - M_{gl}^2} \sum_{i=1}^{6} \frac{U_{2i} U^*_{1i}}{(k^2 - \tilde{m}_i^2)^2} \tag{6.9}$$

where $\tilde{m}_i^2$ is the i-th eigenvalue of M_D^2, M_{gl} is the gluino mass and U is defined in (3.10). The sum of the two contributions yields:

$$(a)+(b) = -\frac{2 g_s^2 g}{\cos\theta_W}\, \bar{s}_L \gamma^\mu d_L \, i \int \frac{d^d k}{(2\pi)^d} \frac{k^2}{k^2 - M_{g\ell}^2} \sum_{i,j=1}^{6} \frac{U_{2i}\left(\sum_{k=4}^{6} U^*_{ki} U_{kj}\right) U^*_{1j}}{(k^2 - \tilde{m}_i^2)(k^2 - \tilde{m}_j^2)} \qquad (6.10)$$

It is important to notice that in the final result the electromagnetic part of the current coupled to Z vanishes (i.e. the term proportional to $\sin^2\theta_W$ cancels in (6.10)). This holds true because the effective vertex of the photon, at zero external momentum, vanishes due to current conservation [71].

If I assume A = 0, there is no left-right mixing, and the matrix U is 3x3 block diagonal. In this case (6.10) vanishes because, if i=4,5,6, then U_{2i} =0 and, if i=1,2,3, then U_{ki} =0 (since k=4,5,6). Since two matrix elements of each sector are present, to obtain a non-vanishing result, one needs two left-right mixings in the perturbative expansion of U.

To be clearer, I will show the same result, using a different argument, which is more physical and probably more clarifing. If I have external left quarks and no left-right squark mixings, the graphs in fig. 6.2 will not feel any SU(2)xU(1) breaking and the Z^0 coupling will vanish because of the Ward identity which forces the photon coupling to be zero. Therefore, to have a vanishing result, right-handed quarks or their scalar partners must appear in fig.6.2. If I insist in having left-handed quarks, two left-right insertions are necessary [71], as shown in fig. 6.3.

The sum over squarks in (6.10) can be computed with a perturbative expansion in the left-right mixings, treating these terms as mass insertions in the squark propagators:

$$\sum_{i,j=1}^{6} \frac{U_{2i}\left(\sum_{k=4}^{6} U^*_{ki} U_{kj}\right) U^*_{1j}}{(k^2 - \tilde{m}_i^2)(k^2 - \tilde{m}_j^2)} \simeq \sum_{i=1}^{3} V^*_{id} V_{is} \frac{A^2 m^2 m_{d_i}^2}{(k^2 - \tilde{m}_i^2)^4} \qquad (6.11)$$

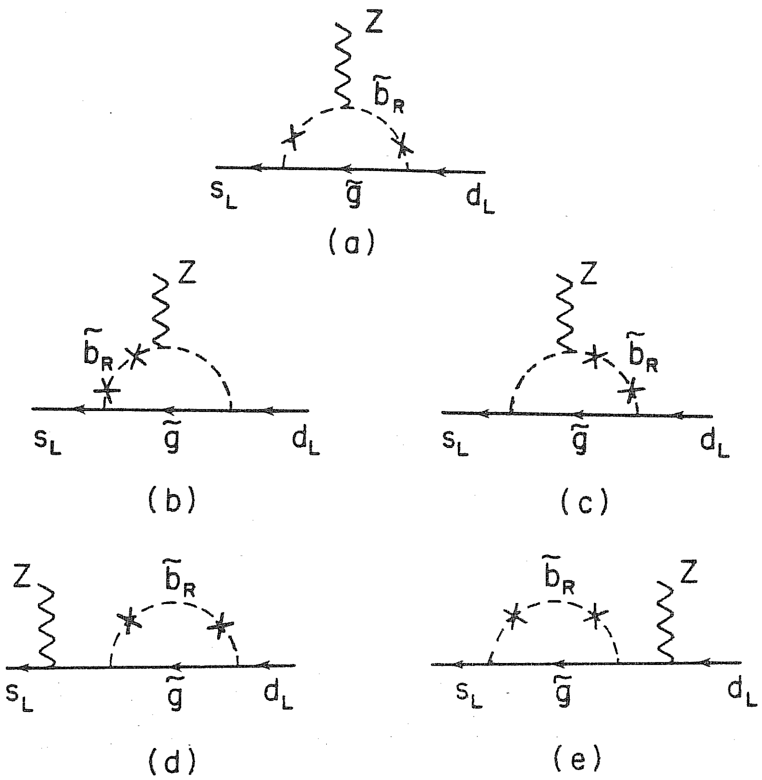


Fig. 6.3. Mass insertions in the $d\bar{s}Z$ effective vertex with gluino exchange.

Computing the integral, the effective vertex becomes:

$$-\frac{\alpha_s}{6\pi}\frac{g}{\cos\theta_W}\sum_i V_{id}V_{is}^* \frac{A^2 m^2 m_{d_i}^2}{M_{gl}^4} F(x_i)\ \bar{s}_L\gamma^\mu d_L \qquad (6.12)$$

m_{d_i} are the eigenvalues of the down quarks mass matrix, $x_i=\frac{\tilde{m}_i^2}{M_{gl}^2}$ with $\tilde{m}_i$ down squark masses and

$$F(x)=\frac{1}{(x-1)^4}\left[\frac{x^2}{2}-3x+3\ln x+\frac{1}{x}+\frac{3}{2}\right] \qquad (6.13)$$

Keeping only the gluino contributions, the branching ratio of $K^+\to\pi^+\nu\bar{\nu}$ becomes (with three generations of neutrinos):

$$BR(K^+\to\pi^+\nu\bar{\nu})_{gluinos}=\frac{\alpha_s^2}{6\pi^2}\left|\sum_i V_{id}V_{is}^*\frac{A^2m^2m_{d_i}^2}{M_{gl}^4}F(x_i)\right|^2\frac{BR(K^+\to\pi^0e^+\nu_e)}{|V_{us}|^2} \qquad (6.14)$$

Using A=3, m_b =4.7 GeV, α_s =0.1 and the angles of the KM matrix obtained in ref.[54], I get a ratio with the SM branching ratio:

$$\frac{BR(K^+\to\pi^+\nu\bar{\nu})_{gluinos}}{BR(K^+\to\pi^+\nu\bar{\nu})_{SM}} = \begin{cases} 2\cdot 10^{-2} & \text{for } m=\mu=M_{gl}=30\,GeV \\ 5\cdot 10^{-2} & \text{for } m=\mu=30\,GeV,\ M_{gl}=10\,GeV \end{cases} \qquad (6.15)$$

Then, even with low values of sparticle masses, the gluino contribution is far from detection in the BNL experiment.

There are two reasons for the suppression of the gluino graphs. First of all there is the super-GIM effect, which is more efficient than in the case of quarks, because squarks are more degenerate. The second reason is the suppression coming from the double left-right mixing. I could avoid these insertions if I had computed the graph with an external momentum different from zero. Of course the amplitude in the decay rate would be now suppressed by a factor q^2/M_Z^2 , where q is the Z momentum, which is of the order of the K meson mass.

Searching for larger contributions, I am studying now the box diagrams where one avoids left-right mixing (or momentum transfer) suppressions, and the effective vertex with higgsino exchange, where no super-GIM suppression is present.

Let us consider the box diagrams depicted in fig. 6.4, assuming that $\tilde{w}$ and $\tilde{z}$ (fermionic partners of W and Z) are essentially mass eigenstates. Neglecting the squark left-right mixing, the graphs of fig. 6.4 are easily computed:

$$(a) = -\frac{g^4}{64\pi^2}\,\frac{1}{\tilde{m}_w^2}\sum_{i=1}^{3} V_{id}V_{is}^{*}\; G_w(x_i,y)\; \bar{s}_L\gamma^\mu d_L\; \bar{\nu}_L\gamma_\mu\nu_L \qquad (6.16)$$

with $$G_w(x,y) = \frac{1}{y-x}\left[\left(\frac{x}{x-1}\right)^2\ln x - \frac{1}{x-1} - \left(\frac{y}{y-1}\right)^2\ln y + \frac{1}{y-1}\right] \qquad (6.17)$$

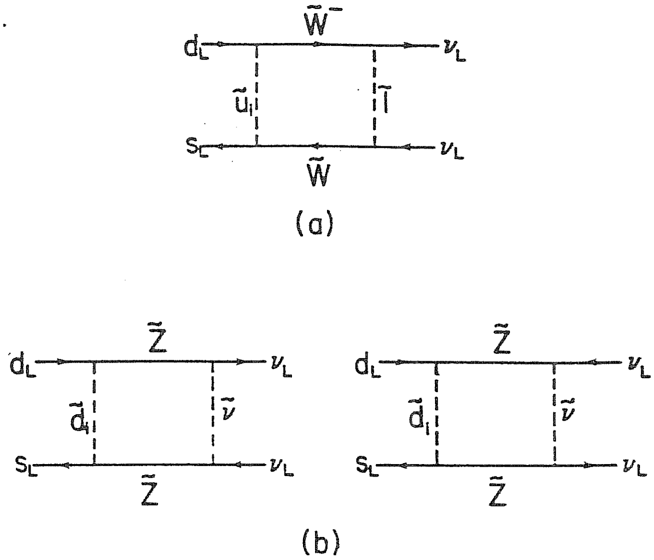


Fig. 6.4. SUSY box diagrams for $K^+ \to \pi^+ \nu \bar{\nu}$.

$x_i = \dfrac{\tilde{m}_{u_i}^2}{\tilde{m}_W^2}$, $y = \dfrac{\tilde{m}_\ell^2}{\tilde{m}_W^2}$ where $\tilde{m}_{u_i}$, $\tilde{m}_\ell$, $\tilde{m}_W$ are the masses of the i-th up squark, the slepton and the w-ino.

$$(b) = -\frac{g^4 (1 - \frac{2}{3}\sin^2\theta_W)^2}{(16\pi)^2 \cos^4\theta_W \, \tilde{m}_Z^2} \sum_{i=1}^{3} V_{id} V_{is}^* \, G_Z(x_i, y) \, \bar{s}_L \gamma^\mu d_L \, \bar{\nu}_L \gamma_\mu \nu_L \qquad (6.18)$$

with

$$G_Z(x,y) = \frac{1}{y-x}\left[\left(1 - \frac{1}{(x-1)^2}\right)\ln x + \frac{1}{x-1} - \left(1 - \frac{1}{(y-1)^2}\right)\ln y - \frac{1}{y-1}\right] \qquad (6.19)$$

$x_i = \dfrac{\tilde{m}_i^2}{\tilde{m}_Z^2}$, $y = \dfrac{\tilde{m}_\nu^2}{\tilde{m}_Z^2}$ where $\tilde{m}_i$, $\tilde{m}_\nu$, $\tilde{m}_Z$ are the masses of the i-th down squark, the sneutrino and the z-ino.

The ratio between the squared modulus of the amplitude of the graphs of fig. 6.4(a,b) and the corresponding quantity computed in the SM gives:

$$(a) = 0.5 \quad \text{for} \quad \tilde{m}_W = \tilde{m}_\ell = \mu = m = 30 \text{ GeV}$$

$$(b) = \begin{cases} 0.2 & \text{for} \quad \tilde{m}_Z = \mu = m = 30 \text{ GeV} \\ 0.5 & \text{for} \quad \mu = m = 30 \text{ GeV} \quad \tilde{m}_Z = 10 \text{ GeV} \end{cases} \qquad (6.20)$$

where the sneutrino mass has been neglected.

Therefore, even if the box diagrams can dominate over the gluino penguin graphs, they can not exceed the SM contribution. One may have expected a larger contribution from these diagrams, since they have a similar structure to those with W and Z exchange, and I am assuming that the w-ino and z-ino could be considerably lighter than their bosonic partners. The super-GIM effect has again to be blamed for that.

I consider now the $d\bar{s}Z$ effective vertex with chargino exchange. For chargino I mean a mass eigenstate, linear combination of the w-ino and the charged higgsinos. Since the w-ino contribution suffers of the super-GIM suppression, I consider the most propitious situation where the higgsino is almost an exact mass eigenstate and I focus my attention on its exchange. Note that a neutral higgsino can not give a sizable contribution, because its coupling constant will always be proportional to lepton or down quark mass. Then the $d\bar{s}Z$ effective vertex is mediated by the graph of fig. 6.5.

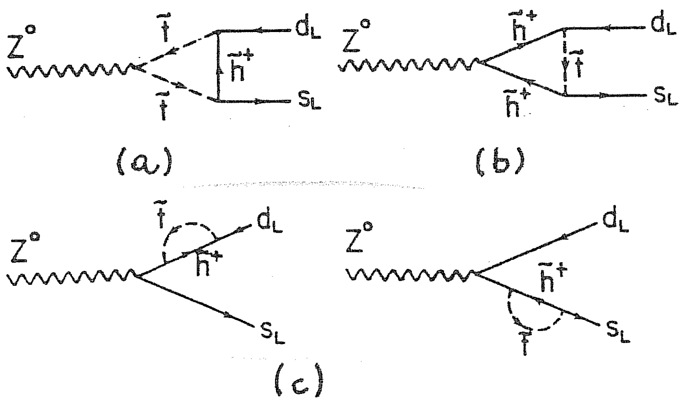


Fig. 6.5. $d\bar{s}Z$ effective vertex through higgsino exchange.

Keeping into account only the dominant elements of the up squark mass matrix (i.e. like $\mu^2, m_t^2, Am\,m_t$), the rotation matrix of the up squarks becomes:

$$U_v = \left(\begin{array}{ccc|ccc} 1 & & & 0 & & \\ & 1 & & & 0 & \\ & & -\frac{1}{\sqrt{2}} & & & \frac{1}{\sqrt{2}} \\ \hline 0 & & & 1 & & \\ & 0 & & & 1 & \\ & & \frac{1}{\sqrt{2}} & & & \frac{1}{\sqrt{2}} \end{array}\right) \tag{6.21}$$

With the approximation of U_v as in eq. (6.21), the contributions of the diagrams of fig. 6.5 are:

$$(a) = \Gamma^\mu i \int \frac{d^d k}{(2\pi)^d} \frac{k^2}{k^2 - \tilde{m}_h^2} \left[\left(-1 + \frac{8}{3}\sin^2\theta_W\right)\left(\frac{1}{(k^2 - m_+^2)^2} + \frac{1}{(k^2 - m_-^2)^2}\right) + \right.$$
$$\left. + \frac{2}{(k^2 - m_-^2)(k^2 - m_+^2)} \right]$$

$$(b) = -4\Gamma^\mu (1 - 2\sin^2\theta_W)\, i \int \frac{d^d k}{(2\pi)^d} \frac{\left(\tilde{m}_h^2 - \frac{k^2}{2}\right)}{(k^2 - \tilde{m}_h^2)^2} \left[\frac{1}{(k^2 - m_+^2)} + \frac{1}{(k^2 - m_-^2)}\right]$$

$$(c) = -2\Gamma^\mu \left(1 - \frac{2}{3}\sin^2\theta_W\right) i \int \frac{d^d k}{(2\pi)^d} \frac{k^2}{k^2 - \tilde{m}_h^2} \left[\frac{1}{(k^2 - m_+^2)^2} + \frac{1}{(k^2 - m_-^2)^2}\right] \tag{6.22}$$

$$\Gamma^\mu \equiv \frac{g^3 m_t^2}{16 M_W^2 \cos\theta_W} V_{td} V_{ts}^* \bar{s}_L \gamma^\mu d_L$$

$m_\pm$ are the masses of the two contributing squarks, i.e. $m_\pm^2 = \mu^2 + m_t^2 \pm A m m_t$. Again I find that the electromagnetic part of the current coupled to Z vanishes when I sum up all the contributions. The evaluation of the integrals yields the following finite result for the effective vertex:

$$\frac{g^3 m_t^2}{(16\pi)^2 M_W^2 \cos\theta_W} V_{td} V_{ts}^* \bar{s}_L \gamma^\mu d_L \; F(x_+, x_-) \tag{6.23}$$

with
$$F(x,y) = \frac{(x^2 - 2x)}{(x-1)^2}\ln x + \frac{1}{x-1} + \frac{(y^2 - 2y)}{(y-1)^2}\ln y +$$
$$+ \frac{1}{y-1} + \frac{2}{x-y}\left(\frac{y^2}{y-1}\ln y - \frac{x^2}{x-1}\ln x\right) \tag{6.24}$$

Plugging in numbers, I find that, even for light sparticles, the higgsino contribution is not larger than the one coming from gluino penguins.

In conclusion, I have analysed in detail all the contributions to $K^+ \to \pi^+ \nu \bar{\nu}$ coming from SUSY intermediate states, finding that, for plausible sparticle masses, no visible effects are expected.

7. $K_L \to \mu^+\mu^-$

Dealing with $K^+ \to \pi^+ \nu \bar{\nu}$, one must always consider the related process $K_L \to \mu^+\mu^-$. The two decays are actually deeply connected, because they are originated by very similar one loop Feynman graphs. Unfortunately the limits on $K_L \to \mu^+\mu^-$ are less reliable, because they are affected by long distance contributions, which have large uncertainties at present days. The experimental value of the branching ratio of $K_L \to \mu^+\mu^-$ is:

$$BR(K_L \to \mu^+\mu^-)_{exp} = (9.1 \pm 1.9) \cdot 10^{-9} \qquad (7.1)$$

Subtracting the contribution of the two-photon intermediate state, one can set an upper bound on the short distance contribution [72]:

$$BR(K_L \to \mu^+\mu^-)_{sd} \lesssim 6 \cdot 10^{-9} \qquad (7.2)$$

In the SM the decay occurs through Z^0 exchange graphs or through box diagrams, as shown in fig. 7.1.

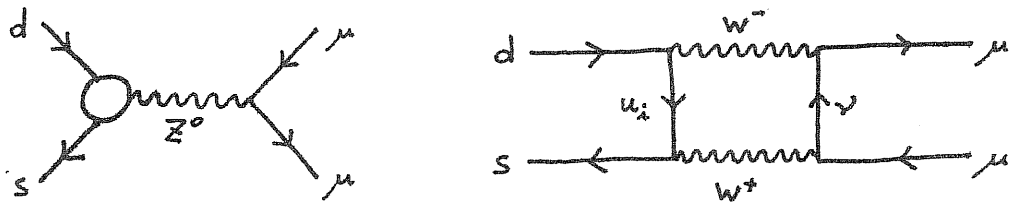


Fig. 7.1. Diagrams for $K_L \to \mu^+\mu^-$. The blob represents the effective $d\bar{s}Z$ vertex.

The diagrams with Z^0 exchange have the same structure of those contributing to $K^+ \to \pi^+ \nu \bar{\nu}$ after replacing the lepton current with the neutrino current. In the case of $K_L \to \mu^+\mu^-$ the scalar Higgs do not contribute to the box diagrams once one neglects the muon mass. Moreover, as can be seen comparing fig. 7.1 with fig. 6.1, the two box diagrams with $W^\pm$ exchange have a different structure in the fermionic indeces. This leads to a factor 1/4 in the amplitude for $K_L\to\mu^+\mu^-$ with respect to $K^+\to\pi^+\nu\bar{\nu}$ [73]. Summing up all the graphs, one obtains the following short-distance decay width [64]:

$$\Gamma(K_L\to\mu^+\mu^-)_{sd} = \frac{\alpha^2}{4\pi^2\sin^4\theta_W}\,\frac{\left[\sum_{i=2}^{N} Re\left(V_{is}^* V_{id}\right) C(x_i)\right]^2}{V_{us}^2}\,\Gamma(K^+\to\mu^+\nu_\mu) \qquad (7.3)$$

with $$C(x) = \frac{3}{4}\left(\frac{x}{x-1}\right)^2 \ln x + \frac{x}{4} - \frac{3x}{4(x-1)} \qquad (7.4)$$

Replacing numbers, I find:

$$BR(K_L\to\mu^+\mu^-)_{sd} = 1.4\cdot 10^{-3}\left[\sum_{i=2}^{N} Re\left(V_{is}^* V_{id}\right) C(x_i)\right]^2 \qquad (7.5)$$

Therefore, enhancing the branching ratio of $K^+\to\pi^+\nu\bar{\nu}$, also BR($K_L\to\mu^+\mu^-$) is automatically increased; however BR($K_L\to\mu^+\mu^-$) has to satisfy the bound (7.2) and then a limit on BR($K^+\to\pi^+\nu\bar{\nu}$) can be inferred.

I want to study now what happens to the bound (7.2) when SUSY models are considered. Let us start with the supersymmetrization of the box diagram (fig. 7.2).

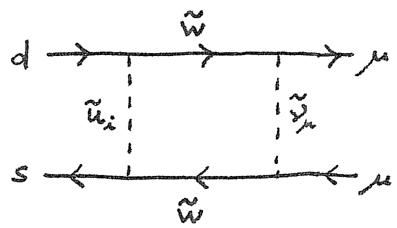


Fig. 7.2. SUSY box diagram for $K_L\to\mu^+\mu^-$.

In the case of w-ino exchange, the structure of the fermionic indeces is the same, no matter if the muon is replaced by the neutrino. Therefore the factor 1/4 is not present any more. Furthermore, since the sneutrino can be lighter than the charged slepton, we could have excluded a priori a large contribution to $K^+\to\pi^+\nu\bar{\nu}$ coming from the box diagrams, if we believe in the bound (7.2). No limits on sparticle masses can be extracted, due to the uncertainties present in (7.2).

The amplitude for the decay $K_L \to \mu^+\mu^-$ due to Z^0 exchange diagrams is easily computed, substituting the gluino mediated effective vertex given in the previous section:

$$\mathcal{A}(K_L \to \mu^+\mu^-)_{gluinos} = -\frac{G_F}{\sqrt{2}} \frac{2\alpha_s}{3\pi} (\bar{s}_L \gamma^\mu d_L)(\bar{\mu}_L \gamma_\mu \mu_L) \times \times \sum_i V_{id} V_{is}^* \frac{A^2 m^2 m_{d_i}^2}{M_{gl}^4} F(x_i) \tag{7.6}$$

I checked that, for allowed sparticle masses, the amplitude (7.6) does not exceed the SM amplitude. Therefore, even with the gluino contribution, the bound (7.2) is still satisfied.

Now I want to show how strigent is the bound (7.2) for the decay $K^+ \to \pi^+ \nu \bar{\nu}$. First I compute the ratio of the SM and gluino exchange amplitudes for $K_L \to \mu^+\mu^-$ and $K^+ \to \pi^+ \nu \bar{\nu}$:

$$R(K_L \to \mu^+\mu^-) \equiv \frac{\mathcal{A}(K_L \to \mu^+\mu^-)_{gluinos}}{\mathcal{A}(K_L \to \mu^+\mu^-)_{SM}} = \frac{2}{3} \frac{\alpha_s}{\alpha} \sin^2\theta_W \frac{\left| \sum_i V_{id} V_{is}^* A^2 \frac{m^2}{M_{gl}^4} m_{d_i}^2 F(x_i) \right|}{\left| \sum_i V_{id} V_{is}^* C(x_i) \right|} \tag{7.7}$$

and

$$R(K^+ \to \pi^+ \nu \bar{\nu}) \equiv \frac{\mathcal{A}(K^+ \to \pi^+ \nu \bar{\nu})_{gluinos}}{\mathcal{A}(K^+ \to \pi^+ \nu \bar{\nu})_{SM}} = \frac{2}{3} \frac{\alpha_s}{\alpha} \sin^2\theta_W \frac{\left| \sum_i V_{id} V_{is}^* A^2 \frac{m^2}{M_{gl}^4} m_{d_i}^2 F(x_i) \right|}{\left| \sum_i V_{id} V_{is}^* D(x_i, 0) \right|} \tag{7.8}$$

The ratio of the two quantities depends only on SM quantities, and thus I obtain:

$$R(K^+ \to \pi^+ \nu \bar{\nu}) \simeq 0.2 \; R(K_L \to \mu^+\mu^-) \tag{7.9}$$

This implies that no large contributions to $K^+ \to \pi^+ \nu \bar{\nu}$ from gluino exchange can be obtained without spoiling the bound (7.2).

In conclusion I found that $K_L \to \mu^+\mu^-$ gives a severe bound on SUSY contributions to $K^+ \to \pi^+ \nu \bar{\nu}$. A way to increase the branching ratio of the process $K^+ \to \pi^+$ + "missing energy", without spoiling the limit on $K_L \to \mu^+\mu^-$, is to imagine that the neutral particles, produced in pair, are some SUSY partners. Now I turn to analyse this possibility.

8. $K^+ \to \pi^+ +$ "neutral sparticles"

Let us suppose now that the undetected particles in the process $K^+ \to \pi^+ +$ "missing energy" are some SUSY partners [65-74]. Experimentally the situation is not distinguishable from the decay into neutrinos, and so the branching ratio of the process can be enhanced through these new channels. I will consider the case of sneutrinos, photinos and higgsino. Let us start with the sneutrinos.

Since the sneutrinos are scalars, a first remark is that the pion energy spectrum is different from the case of fermionic undetected particles [75]. In this way the sneutrinos could leave an indirect signature in favour of SUSY. The decay $K^+ \to \pi^+ \tilde{\nu} \bar{\tilde{\nu}}$ can occur through the graphs depicted in fig. 8.1. Since these diagrams have an analogous structure

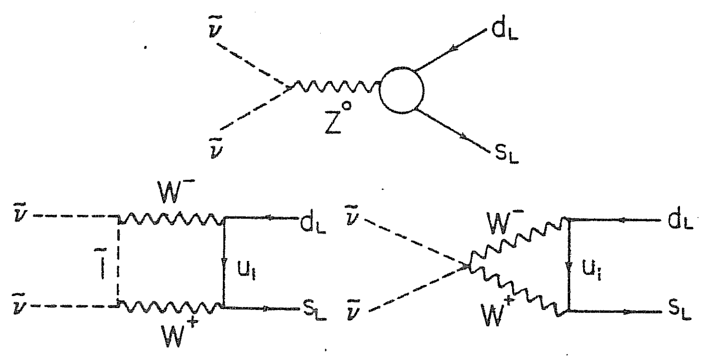


Fig. 8.1. Feynman graphs for the decay $K^+ \to \pi^+ \tilde{\nu} \bar{\tilde{\nu}}$.

to those contributing to the decay $K^+ \to \pi^+ \nu \bar{\nu}$ in the SM, I expect a branching ratio of the same order of magnitude.

I want to study now whether, assuming rather light SUSY particles, the decay into sneutrinos may become dominant over the SM decay into neutrinos. First I write down the general formula for the branching ratio and then I discuss the possible contributions from SUSY intermediate states.

The matrix element of the process can be written in the following way, using the vanishing momentum transfer approximation:

$$\mathcal{M} = \frac{G_F}{\sqrt{2}} \lambda \langle \pi^+ | \bar{s} \gamma^\mu (1-\gamma_5) d | K^+ \rangle \left[a (p_1+p_2)_\mu + b (p_1-p_2)_\mu \right] \quad (8.1)$$

a, b and λ are parameters depending on the dynamics and p_1, p_2 are the sneutrino momenta. The hadronic matrix element is written in the usual parametrization:

$$\langle \pi^+ | \bar{s} \gamma_\mu (1-\gamma_5) d | K^+ \rangle = f_+(q^2) (K+\pi)_\mu + f_-(q^2) (K-\pi)_\mu \quad (8.2)$$

where q is the momentum transfer. At $q^2=0$, $f_-(0) = 0$ and $f_+(0) = 0.98$, obtained by correcting the current algebra result. Neglecting the sneutrino mass, the decay width is:

$$\Gamma = \frac{2 G_F^2 |\lambda|^2 f_+^2(0) M_K^5}{(16\pi)^3} \left\{ a^2 \left[2x(1-x)^2 \ln x + (1-x^4) - 2x(1-x^2) \right] + b^2 \left[-2x^2 \ln x + \frac{(1-x^4)}{6} - \frac{4}{3} x (1-x^2) \right] \right\} \quad (8.3)$$

where $x = \frac{m_\pi^2}{M_K^2}$. Substituting the numbers, I obtain:

$$BR(K^+ \to \pi^+ \tilde{\nu} \bar{\tilde{\nu}}) = |\lambda|^2 (0.6 a^2 + 0.1 b^2) \tilde{N}_\nu \quad (8.4)$$

I introduced $\tilde{N}_\nu$ to specify the number of available decay channels, i.e. $\tilde{N}_\nu$ counts how many neutrino generations are light enough to make the decay kinematically possible. Even if there are no stringent bounds on sneutrino mass, as I showed in section 4, it seems difficult, in broken N=1 supergravity models, to keep sneutrino masses so low, without an unnatural fine tuning. Here I suppose the sneutrino mass matrix to allow at least one tiny eigenvalue so as to allow for the decay.

I am turning now to the analysis of the specific SUSY contributions. First of all I consider the graph with Z^0 exchange and gluino mediated $d\bar{s}Z$ effective vertex, where

$$a = 1 \qquad b = 0 \qquad \lambda = \frac{\alpha_s}{3\pi} \sum_{i=1}^{3} V_{id} V_{is}^{*} \; A^2 \frac{m^2}{M_{gl}^4} m_{d_i}^2 \, F(x_i) \tag{8.5}$$

Even for light sparticles ($\mu = M_{gl} = 30$ GeV), the branching ratio is three orders of magnitude less than that of $K^+ \to \pi^+ \nu\bar{\nu}$ in the SM.

A further possibility is to consider box diagrams with exchange of gauge fermions, as shown in fig. 8.2, assuming them to be much lighter than their boson partners. For instance, from the w-ino exchange with a vanishing lepton mass, I get

$$a = b = 1 \qquad \lambda = \frac{\alpha}{2\pi \sin^2\theta_w} \frac{M_W^2}{\tilde{m}_w^2} \sum_{i=1}^{3} V_{id} V_{is}^{*} \; G(x_i) \tag{8.6}$$

with $x_i = \dfrac{\tilde{m}_{\nu_i}^2}{\tilde{m}_w^2}$ and $G(x) = \dfrac{1}{(x-1)^2} \ln x - \dfrac{1}{x-1}$ (8.7)

Replacing the numbers in (8.6), I find that, even for

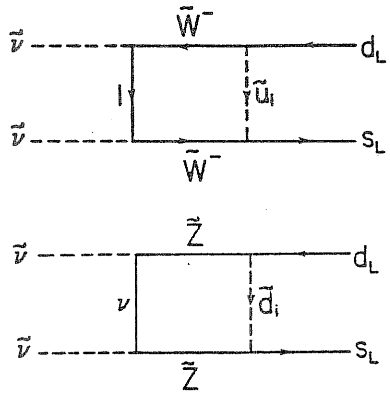


Fig. 8.2. Box diagrams with gaugino exchange contributing to $K^+ \to \pi^+ \tilde{\nu}\bar{\tilde{\nu}}$.

$\tilde{m}_w$ = 30 GeV, the ratio of the branching ratios of $K^+ \to \pi^+ \tilde{\nu}\bar{\tilde{\nu}}$ and $K^+ \to \pi^+ \nu\bar{\nu}$ is $0.3\,\tilde{N}_\nu$, and these new graphs are not dominant over the ones in fig. 8.1.

In conclusion, from graphs analogous to the SM, I expect a branching ratio for $K^+\to\pi^+\tilde{\nu}\bar{\tilde{\nu}}$ of the same order of the one for $K^+\to\pi^+\nu\bar{\nu}$ in the SM; however the energy spectrum of the pion coming from neutrino decay is, in principle distinguishable from the one for neutrino decay.

If the two light neutrals are photinos, the decay $K^+\to$ $\to\pi^+$ + "missing energy" can already occur at tree level, as shown in fig. 8.3, where the Majorana nature of the pho-

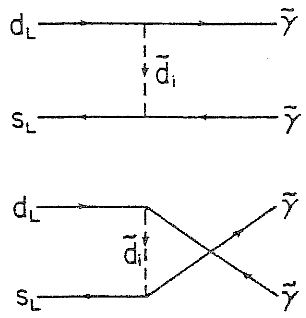


Fig. 8.3. Feynman graphs for the decay $K^+\to\pi^+\tilde{\gamma}\tilde{\gamma}$.

tino is taken into account. The two diagrams of fig. 8.3 yield, at zero momentum transfer:

$$\begin{aligned} (a) &= \left(-\frac{\sqrt{2}\,e}{3}\right)^2 \bar{s}\,\frac{(1+\gamma_5)}{2}\,\tilde{\gamma}\,\bar{\tilde{\gamma}}\,\frac{(1-\gamma_5)}{2}\,d\,\sum_{i=1}^{6}\frac{U_{2i}U_{1i}^*}{\tilde{m}_i^2} \\ (b) &= -\left(-\frac{\sqrt{2}\,e}{3}\right)^2 \bar{s}\,C\,\frac{(1+\gamma_5)}{2}\,\bar{\tilde{\gamma}}^T\,\tilde{\gamma}^T\,C^{-1}\,\frac{(1-\gamma_5)}{2}\,d\,\sum_{i=1}^{6}\frac{U_{2i}U_{1i}^*}{\tilde{m}_i^2} \end{aligned} \qquad (8.8)$$

Summing up, I obtain the effective operator

$$\frac{4\pi\alpha}{9}\,\bar{s}_L\gamma^\mu d_L\,\bar{\tilde{\gamma}}\,\gamma_\mu\gamma_5\,\tilde{\gamma}\,\sum_{i=1}^{6}\frac{U_{2i}U_{1i}^*}{\tilde{m}_i^2} \qquad (8.9)$$

Effective vertices containing right-handed quarks are necessarily suppressed by left-right insertions, because no flavor changing is present in the right sector alone. Neglecting such mixings, I can use the KM matrix as squark rotation matrix. Since in the zero mass limit Dirac and Majorana spinors behave in the same way, the branching ratio of $K^+\to\pi^+\tilde{\gamma}\tilde{\gamma}$ can be related to the one for $K^+\to\pi^0 e^+\nu$ in the same way I did for the decay into neutrinos, assuming a negligible photino mass. Then in the approximation of vanishing left-right squark mixing and zero photino mass, I get the branching ratio:

$$BR(K^+\to\pi^+\tilde{\gamma}\tilde{\gamma}) = \left(\frac{2\sqrt{2}\,\pi\alpha}{9G_F V_{us}}\right)^2 BR(K^+\to\pi^0 e^+\nu)\left|\sum_{i=1}^{3}\frac{V_{di}V_{si}^*}{\tilde{m}_i^2}\right|^2 \qquad (8.10)$$

where I included a factor 1/2! which accounts for the fact that a Majorana particle is undistinguishable from its antiparticle. The numerical results, for different values of μ, are summarized in table 8.I, assuming c = 0.6 and m_t =40 GeV.

μ	$BR(K^+\to\pi^+\tilde{\gamma}\tilde{\gamma})$
30 GeV	1.3 10^{-7}
40 GeV	1.7 10^{-8}
60 GeV	9.1 10^{-10}
80 GeV	1.0 10^{-10}

Table 8.I.

It is worth mentioning that the pion energy spectrum is different if the fermions pair are Majorana or Dirac [63]. Since the nature of the neutrino spinor is not known, I can not claim that this helps in characterizing a SUSY signature.

From table 8.I it can be seen that the present experimental limit gives already a combined bound: either the lightest down squark mass is larger than 30 GeV or the photino is heavier than 170 MeV (for m_t = 40 GeV). Moreover a negative result in an experiment with sensibility of 10^{-10} in the branching ratio could improve this bound up to 80 GeV for down squarks. As I showed in section 4, the negative results in experiments searching for $e^+e^- \to \gamma\tilde{\gamma}\tilde{\gamma}$ imply that, for negligible photino masses, the selectron must be heavier than $\sim$ 50 GeV. Since squarks have strong radiative corrections, the squarks are expected to be heavier than the selectrons. Then, even if the squarks satisfy this constraint, an experiment for the branching ratio of $K^+ \to \pi^+$ + "missing energy" up to 10^{-10} can still probe an interesting range of down squark masses.

Another possibility is that the decay $K^+ \to \pi^+ \tilde{\gamma}\tilde{\gamma}$ occurs through the cascade process $K^+ \to \pi^+\pi^0$, $\pi^0 \to \tilde{\gamma}\tilde{\gamma}$. Now the decay into photinos is no longer super-GIM suppressed, but the helicity suppression is present. The decay $\pi^0 \to \tilde{\gamma}\tilde{\gamma}$ occurs at tree level and is mediated by squark exchange. The decay width is easily computed to be [74]:

$$\Gamma(\pi^0 \to \tilde{\gamma}\tilde{\gamma}) = 4\pi\alpha^2 f_\pi^2\, \tilde{m}_\gamma^2\, \frac{m_\pi}{\mu^4} \sqrt{1 - 4\frac{\tilde{m}_\gamma^2}{m_\pi^2}} \qquad (8.11)$$

where f_π = 93 MeV and μ is the common mass for the up and down squark which mediate the interaction. Thus I find the branching ratio:

$$BR(K^+ \to \pi^+\tilde{\gamma}\tilde{\gamma})_{cascade} = 2.5 \cdot 10^{-9} \sqrt{1 - 4\frac{\tilde{m}_\gamma^2}{m_\pi^2}} \left(\frac{40\ \mathrm{GeV}}{\mu}\right)^4 \left(\frac{\tilde{m}_\gamma}{50\ \mathrm{MeV}}\right)^2 \qquad (8.12)$$

Analogously to $\pi^0 \to \tilde{\gamma}\tilde{\gamma}$, one could consider the decay $K_L \to \tilde{\gamma}\tilde{\gamma}$, which is however flavor changing and super-GIM suppressed (and, of course, still helicity suppressed). The decay width is:

$$\Gamma(K_L \to \tilde{\gamma}\tilde{\gamma}) = \frac{|\lambda|^2}{8\pi} f_K^2\, m_{K^0}\, \tilde{m}_\gamma^2 \sqrt{1 - 4\frac{\tilde{m}_\gamma^2}{m_{K^0}^2}} \qquad (8.13)$$

with $$\lambda = \frac{4\pi\alpha}{9} \sum_i \frac{V_{di} V_{si}^*}{\tilde{m}_i^2} \qquad (8.14)$$

The branching ratio for the disappearance of the kaon into two photinos, for different values of μ and $\tilde{m}_\gamma$ are:

BR($K_L \to \tilde{\gamma}\tilde{\gamma}$)	μ = 30 GeV	μ = 40 GeV	μ = 50 GeV
$\tilde{m}_\gamma$ = 100 MeV	6.9 10^{-6}	1.2 10^{-6}	2.6 10^{-7}
$\tilde{m}_\gamma$ = 200 MeV	1.8 10^{-5}	3.0 10^{-6}	6.7 10^{-7}

Until now I assumed that the photino is light enough to allow the decay $K^+ \to \pi^+ \tilde{\gamma}\tilde{\gamma}$. As I showed in section 4, this is not so easy to be achieved. There are cosmological bounds on the photino mass which exclude values between ~100 eV and ~ 1 GeV. However these constraints may be evaded in models with unstable photinos. Furthermore there are theoretical limits coming from the computation of the radiative corrections that suggest photinos heavier than 100 MeV. Anyway, since light photinos are not experimentally ruled out, I believe that it is still interesting to study the phenomenological consequences of photinos produced in intermediate energy experiments. In view of this difficulty and of the interesting bounds extracted from the analysis of $K^+ \to \pi^+ \tilde{\gamma}\tilde{\gamma}$, we are urged to study heavier systems in order to probe the mass of the photino in a wider range of values. Before that, I want to discussed briefly the possibility of a decay into neutral higgsinos.

In section 4 I discussed all the constraints on neutral higgsino mass, which may forbid the decay. Anyway the decay may occur through the same graphs of fig. 8.3, replacing the photino with the higgsino. If the external quarks are left-handed, the graphs are suppressed by left-right squark mixings, since the scalar partners of right-handed quarks have no flavor mixing terms. Conversely, if the external quarks are right-handed, the coupling constants are proportional to down and strange quark mass. The contributions are then irrelevant.

9. FCNC DECAYS OF HEAVIER PARTICLES

In this section I discuss FCNC decays of heavier particles: I start by extending the previous analysis to the case of the B meson. The relevant effective operator that gives rise to FCNC decays with photino production can be parametrized in the following way:

$$\lambda_{(i)}\ \bar{b}_L \gamma^\mu d_L^{(i)}\ \bar{\tilde{\gamma}}\, \gamma_\mu \frac{(a+b\gamma_5)}{2} \tilde{\gamma} \qquad (9.1)$$

where i is a generation index. Since in this section I consider a non vanishing photino mass, the Majorana nature of the photino plays an essential role. Then, for comparison, I show the two different cases for Majorana and Dirac spinors (equal in the limit $\tilde{m}_\gamma \to 0$). Thus I have a=0, b=2 for Majorana and a=-b=1 for Dirac. The effective operators come from squark exchange (fig. 9.1). The computation of the graphs in

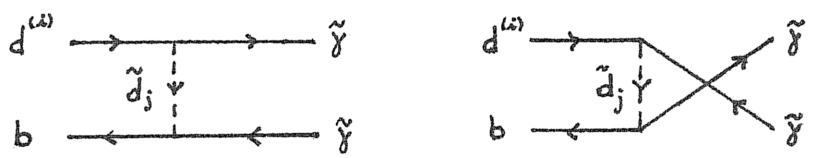


Fig. 9.1. Graphs giving rise to the effective operator $\bar{b}d^{(i)}\tilde{\gamma}\tilde{\gamma}$.

fig.9.1 is identical to the one illustrated in section 8. However, the sum over quark generations can be further simplified, using an approximation valid when the bottom quark is present. Actually the contribution of the first generation can be expressed in terms of the other two, exploiting the unitarity of the KM matrix, and then it can be noticed that the contribution of the second generation is numerically negligible. In formulae (i=1,2):

$$\sum_{j=1}^{3} \frac{V_{jb}^{*} V_{ji}}{\tilde{m}_j^{2}} = \sum_{j=2}^{3} V_{jb}^{*} V_{ji} \left(\frac{1}{\tilde{m}_j^2} - \frac{1}{m_1^2} \right) \simeq V_{ti} \frac{c\, m_t^2}{\mu^2 (\mu^2 + c\, m_t^2)} \qquad (9.2)$$

Finally the coefficient $\lambda_{(i)}$ can be written

$$\lambda_{(i)} = \frac{4\pi\alpha}{g} V_{ti} \frac{c\, m_t^2}{\mu^2(\mu^2 + c\, m_t^2)} \qquad (9.3)$$

Since the bottom quark is heavy, in first approximation, I can describe it as a free parton inside the B meson.

The B^+ meson decays into two photinos and non charmed hadrons through the spectator graph of fig. 9.2(a). The de-

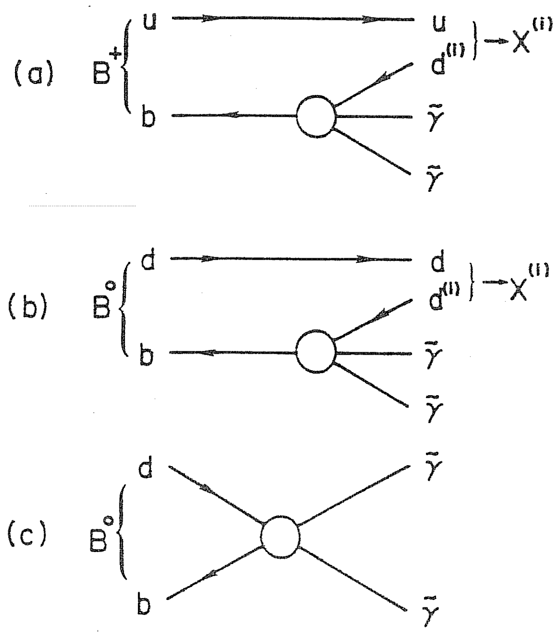


Fig. 7.2. FCNC decays of B meson with photino production.

cay width is:

$$\Gamma(B^+ \to X^{(i)} \tilde{\gamma}\tilde{\gamma}) = \frac{|\lambda_{(i)}|^2 m_b^5}{3\,(8\pi)^3} \rho\left(\frac{\tilde{m}_{\tilde{\gamma}}^2}{m_b^2} \right) \qquad (9.4)$$

$\tilde{m}_{\gamma}$ and m_b are the photino and bottom quark mass, $X^{(i)}$ denotes all possible physical states originated from $u\bar{d}^{(i)}$ and

$$\rho(x) = \left[1 - \frac{2}{3}(17-2\alpha)x + \frac{2}{3}(17+10\alpha)x^2 - 4(7+2\alpha)x^3\right]\sqrt{1-4x} + $$
$$+ 8x^2\left[1-\alpha+(4+2\alpha)x-(7+2\alpha)x^2\right] \ln \frac{1+\sqrt{1-4x}}{1-\sqrt{1-4x}} \qquad (9.5)$$

α =1 in the Majorana case, while α =-2 for Dirac.

The B^o meson can decay either with a graph similar to the B^+ one, or through direct annihilation into two photinos as shown in fig. 7.2(b-c). The graph of fig. 7.2(b) mediates the process $B^o \to X^{(i)}\tilde{\gamma}\tilde{\gamma}$, with $X^{(i)}$ decay products of $d\bar{d}^{(i)}$ and its decay width is given again by (9.4). The graph of fig. 7.2(c) exhibits helicity suppression and gives the following decay width:

$$\Gamma(B^o \to \tilde{\gamma}\tilde{\gamma}) = \beta \frac{|\lambda_{(d)}|^2}{16\pi} f_B^2 M_B \tilde{m}_{\gamma}^2 \sqrt{1 - 4\frac{\tilde{m}_{\gamma}^2}{M_B^2}} \qquad (9.6)$$

M_B is the B^o meson mass and f_B is the decay constant, computed in ref.[76] to be 53 MeV. β parametrizes the Majorana (β =1) and the Dirac (β =1/2) case.

The total width of the B meson is computed evaluating the non leptonic and semileptonic decays through W-exchange spectator graphs, s and t channel annihilation graphs, gluon emission at lowest order, keeping into account the QCD corrections. The result is [77]:

$$\Gamma(B^+) = (8.791\ |V_{ub}|^2 + 2.871\ |V_{cb}|^2)\,\Gamma_o$$
$$\Gamma(B^o) = (8.845\ |V_{ub}|^2 + 3.896\ |V_{cb}|^2)\,\Gamma_o \qquad (9.7)$$
$$\Gamma_o \equiv \frac{G_F^2\, m_b^5}{192\,\pi^3}$$

For m_t = 40 GeV, I find the branching ratios summarized in the following table:

	$\tilde{m}_{\tilde{\gamma}}$ = 0 GeV	$\tilde{m}_{\tilde{\gamma}}$ = 1 GeV	$\tilde{m}_{\tilde{\gamma}}$ = 2 GeV
$B^+ \to X^{(ns)} \tilde{\gamma}\tilde{\gamma}$			
μ = 30 GeV	$3.2\ 10^{-4}$	$1.7\ 10^{-4}$	$4.5\ 10^{-6}$
μ = 40 GeV	$5.3\ 10^{-5}$	$2.9\ 10^{-5}$	$7.5\ 10^{-7}$
$B^+ \to X^{(s)} \tilde{\gamma}\tilde{\gamma}$			
μ = 30 GeV	$1.0\ 10^{-2}$	$5.5\ 10^{-3}$	$1.4\ 10^{-4}$
μ = 40 GeV	$1.7\ 10^{-3}$	$9.2\ 10^{-4}$	$2.4\ 10^{-5}$
$B^0 \to X^{(ns)} \tilde{\gamma}\tilde{\gamma}$			
μ = 30 GeV	$2.4\ 10^{-4}$	$1.3\ 10^{-4}$	$3.3\ 10^{-6}$
μ = 40 GeV	$3.9\ 10^{-5}$	$2.2\ 10^{-5}$	$5.6\ 10^{-7}$
$B^0 \to X^{(s)} \tilde{\gamma}\tilde{\gamma}$			
μ = 30 GeV	$7.4\ 10^{-3}$	$4.1\ 10^{-3}$	$1.0\ 10^{-4}$
μ = 40 GeV	$1.3\ 10^{-3}$	$6.9\ 10^{-4}$	$1.8\ 10^{-5}$
$B^0_L \to \tilde{\gamma}\tilde{\gamma}$			
μ = 30 GeV	0.0	$1.7\ 10^{-6}$	$4.7\ 10^{-6}$
μ = 40 GeV	0.0	$2.8\ 10^{-7}$	$7.8\ 10^{-7}$

For comparison, I have computed the branching ratio for the analogous process with emission of neutrinos in the SM:

$$BR\,(B \to X \nu\bar{\nu}) \sim 10^{-8} \qquad (9.8)$$

Other FCNC B decays with photon or gluon emission in the SM are discussed in ref.[78]. The results are very sensitive to the top quark mass, but one can see that the decays into photinos may give sizable contributions.

In conclusion I found that the inclusive hadronic decays of the B meson, with no charmed particles in the final state, can easily reach branching ratios of order 10^{-4} for photino masses in the range 1÷2 GeV and light squarks. The results seem promising, but a large production of B and a way to distinguish them from the background (which can be rather large [78]) are necessary to reveal these processes in experiments.

In principle one could consider FCNC processes also for the D meson. As I have already shown, the flavor mixings in the up squark mass matrix are proportional to the down quark mass matrix. Then these models do not provide relevant contributions either to D^0-$\bar{D}^0$ mass difference or to $D^+ \to \pi^+ +$ "missing energy".

Another interesting process is the flavor changing decay of Z^0, which has been studied in the framework of the SM [79]. The decay $Z^0 \to t\bar{q}$, where q is a light quark, can reach a branching ratio at most of 10^{-10} . If LEP produces 10^8 Z^0 per year, the decay is not detected. A fourth generation could improve the situation: the presence of a new, heavy enough, quark b' avoids the strong GIM cancellation among down, strange and bottom quarks.

The decay $Z^0 \to b\bar{q}$ is less suppressed by GIM effect, because the top quark is exchanged. For m_t = 40 GeV, the branching ratio is of order 10^{-7} , but, in this case, the signal is less clear, since the bottom quark is light at the Z^0 scale.

In the SUSY case, a flavor changing Z^0 decay can occur through the graphs of fig. 9.3, which involve gluino exchange.

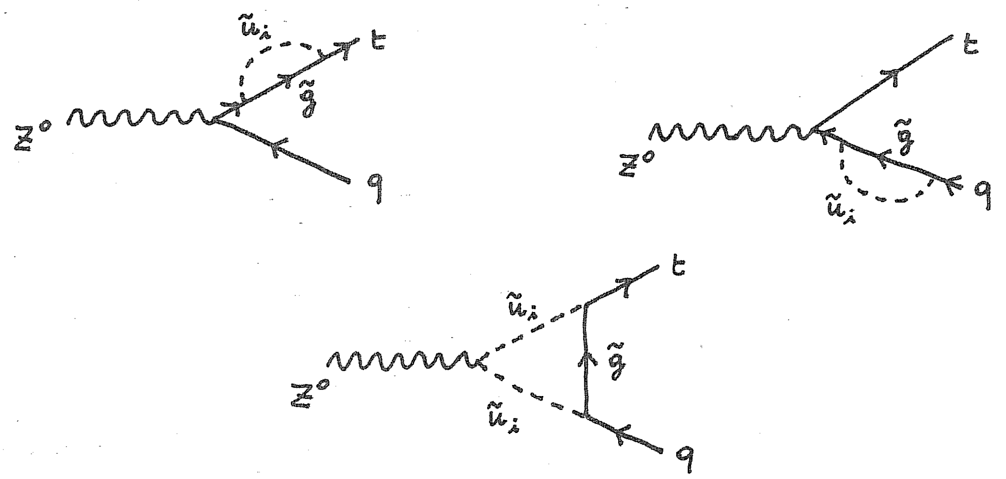


Fig. 9.3. SUSY contributions to $Z^0 \to t\bar{q}$ via gluino exchange.

The charged higgsino exchange diagrams could also be important, since the super-GIM effect is avoided. Now, differently from the case of the SUSY contributions to $K^+ \to \pi^+ \nu \bar{\nu}$, it is not necessary to use left-right mixings, since the external momenta are no longer negligible. From the analysis of ref.[80], it seems possible to have a large contribution to $Z^0 \to t\bar{q}$ coming from these SUSY graphs. However Gamberini and Ridolfi [81] have recently reanalised the decay and, in view of the present bounds on squark masses, have shown that, in the minimal models, the extra SUSY contributions can not exceed the SM result for both $Z^0 \to t\bar{q}$ and $Z^0 \to b\bar{q}$. The contribution to $Z^0 \to t\bar{q}$ is negligible because the flavor changing interaction of the gluino with the up squark sector is highly suppressed. Ref. [81] also finds that it is possible to enhance the decay $Z^0 \to b\bar{s}$ in models with a non minimal content of superfields. These models [82], where a large flavor mixing is present, can be thought as SUSY extensions of some grand unified models.

In conclusion I have studied in this section the impact of SUSY in FCNC decays of heavier systems finding that new effects might be manifest in the future for the decays of the B meson.

APPENDIX

I give here some one loop integrals used to solve the relevant Feynman graphs:

$$\int \frac{d^4k}{(2\pi)^4} \frac{1}{(k^2-M^2)(k^2-m^2)^2} = \frac{i}{16\pi^2 M^2}\left[\frac{1}{(x-1)^2}\ln x - \frac{1}{x-1}\right]$$

$$\int \frac{d^4k}{(2\pi)^4} \frac{1}{(k^2-M^2)^2(k^2-m^2)} = \frac{i}{16\pi^2 M^2}\left[-\frac{x}{(x-1)^2}\ln x + \frac{1}{x-1}\right]$$

$$\int \frac{d^dk}{(2\pi)^d} \frac{k^2}{(k^2-M^2)(k^2-m^2)^2} = \frac{i}{16\pi^2}\left[\frac{1}{\varepsilon} - \gamma + \ln\frac{4\pi}{M^2} - \frac{(x^2-2x)}{(x-1)^2}\ln x - \frac{1}{x-1}\right]$$

$$\int \frac{d^dk}{(2\pi)^d} \frac{k^2}{(k^2-M^2)^2(k^2-m^2)} = \frac{i}{16\pi^2}\left[\frac{1}{\varepsilon} - \gamma + \ln\frac{4\pi}{M^2} - \frac{x^2}{(x-1)^2}\ln x + \frac{1}{x-1} + 1\right]$$

$$\int \frac{d^4k}{(2\pi)^4} \frac{1}{(k^2-M^2)(k^2-m^2)^4} = \frac{i}{16\pi^2 M^6}\left[\frac{(-2x^2-5x+1)}{6x^2(x-1)^3} + \frac{1}{(x-1)^4}\ln x\right]$$

$$\int \frac{d^4k}{(2\pi)^4} \frac{k^2}{(k^2-M^2)(k^2-m^2)^4} = \frac{i}{16\pi^2 M^4}\left[\frac{(x^2-5x-2)}{6x(x-1)^3} + \frac{1}{(x-1)^4}\ln x\right]$$

$$\int \frac{d^4k}{(2\pi)^4} \frac{1}{(k^2-M^2)^2(k^2-m^2)^4} = \frac{i}{16\pi^2 M^8}\left[\frac{(-17x^2-8x+1)}{6x^2(x-1)^4} + \frac{(3+x)}{(x-1)^5}\ln x\right]$$

$$\int \frac{d^4k}{(2\pi)^4} \frac{k^2}{(k^2-M^2)^2(k^2-m^2)^4} = \frac{i}{16\pi^2 M^6}\left[-\frac{(x^2+10x+1)}{3x(x-1)^4} + \frac{2(x+1)}{(x-1)^5}\ln x\right]$$

$$\int \frac{d^dk}{(2\pi)^d} \frac{k^\mu}{(k^2-m^2)[(p-k)^2-M^2]} = \frac{i}{16\pi^2}\frac{p^\mu}{2}\left[\frac{1}{\varepsilon} - \gamma + \ln\frac{4\pi}{M} - \frac{x^2}{(x-1)^2}\ln x + \frac{x}{x-1} + \frac{1}{2}\right] + O(p^3)$$

with $d = 4-2\varepsilon$ and $x = \frac{m^2}{M^2}$.

$$\int \frac{d^4k}{(2\pi)^4} \frac{1}{(k^2-m_1^2)(k^2-m_2^2)(k^2-M^2)^2} = \frac{i}{16\pi^2 M^4}\frac{1}{(y-x)}\left[\frac{x}{(x-1)^2}\ln x - \frac{1}{x-1} - \frac{y}{(y-1)^2}\ln y + \frac{1}{y-1}\right]$$

$$\int \frac{d^4k}{(2\pi)^4} \frac{k^2}{(k^2-m_1^2)(k^2-m_2^2)(k^2-M^2)^2} = \frac{i}{16\pi^2 M^2}\frac{1}{(y-x)}\left[\left(\frac{x}{x-1}\right)^2\ln x - \frac{1}{x-1} - \left(\frac{y}{y-1}\right)^2\ln y + \frac{1}{y-1}\right]$$

with $x = \frac{m_1^2}{M^2}$, $y = \frac{m_2^2}{M^2}$

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