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Probing Flavor Changing Neutral Currents in Supersymmetric Models

Thesis submitted for the degree of

"Magister Philosophiæ"

CANDIDATE	SUPERVISOR
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À minha gente.

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ABSTRACT

We analyse the constraints given by the low energy tests related to Flavor Changing Neutral Currents on the Minimal Supersymmetric Standard Model and its extensions. We explicitly show that the Minimal Model "passes" these tests intact and its non-minimal extensions are severely constrained.

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MOTIVATIONS

Supersymmetry$^{(1)}$ is one of the most appealing approaches to go beyond the Standard $SU(3)_C \times SU(2)_L \times U(1)_Y$ Model. Immediately one could ask why go beyond such a model, once it is so fantastically consistent with all experimental available data. Or, why specifically Supersymmetry is such an attractive candidate.

Answering the first question is a little embarassing since it is related with some aesthetic (or philosophical) reasons which suggest that the Standard Model is not complete. We could quote some (already very famous) features of the model which lead us to believe in its incompleteness. There is a large number of rather arbitrary parameters and assumptions. There is no manner in the Standard Model to justify the quantization of the electric charge ($q_d = \frac{1}{3} q_e$). Why are there three generations of fermions? Why three colors? Why do fermions have an asymmetric $SU(2)_L$ representation between left-handed and right-handed fermions? Surely these questions do not have an easy answer and it is also not easy to say when (or *if)* a satisfactory solution to these problems will turn up. Anyway, theorists claim they are good reasons to go beyond the Standard Model.

Why Supersymmetry is so attractive is a little easier question to answer. It is a refined tentative of describing gravity in an unified framework together the others fundamental interactions. Furthermore, it was noticed latter, low-energy Supersymmetry is able to solve the so-called Naturalness Problem in the gauge hierarchy(2), as we will see now.

Let us suppose that the Standard Model is the effective low energy approximation of a more fundamental theory (as is evidenced , e.g., by the absence of gravity in such "standard" approach) which becomes relevant at a scale μ_1. If we call the scale where $SU(2)_L \times U(1)_Y$ breaking takes place μ_2, and also consider that there is no "new physics" near the Fermi scale, then $\mu_1 >> \mu_2$, since examples of the former scale would be the Grand Unified Scale, $M_{GUT} \sim 10^{15}$ GeV, or even the Planck mass, $M_P \sim 10^{19}$ GeV, and the latter one is of order the W mass, $m_W \sim 10^2$ GeV.

Such a enormous difference between the two quoted scales creates problems. We may imagine calculating the mass of the Standard Model Higgs boson. The relevant quantity for the low energy theory is the running mass at the scale μ_2 which is related to the mass at the scale μ_1 by the following schematic form(3):

$$m_H^2(\mu_2) = m_H^2(\mu_1) + Cg^2 \int_{\mu_2}^{\mu_1} dk^2 + g^2 R + O(g^4) \tag{1}$$

where g is a coupling constant, C is a numeric adimensional factor and R is a quantity which has a logarithmic behavior as $\mu_1 \to \infty$.

The naturalness problem of the gauge hierarchy arises due to our expectation that the value of the $m_H(\mu_2)$ should not be much larger than μ_2. This results from the fact that large values for $m_H(\mu_2)$ imply large value for the Higgs self-coupling. For a Higgs mass of the order of a few TeV, its self-coupling becomes strong enough to hinder the pertubative approach in dealing with the Higgs interactions and new phenomenolgy related to this Higgs strong coupling would be expected. Here, we consider the alternative and more orthodox view that the Higgs presents perturbative coupling and so its mass cannot exceed roughly 1 TeV.

Equation (1) clearly indicates that in order that $m_H^2(\mu_2) \sim \mu_2^2 \ll \mu_1^2$, one must fine-tune the parameter $m_H^2(\mu_1)$ in an extremelly accurated way in order that it cancels the quadratically divergent second term in (1) up to at least 10^{-24} times its value. This is extremely "unnatural" and even if we were able to do it in this first order approximation, higher orders in perturbation theory could destroy this fine adjustment.

This is the famous Naturalness Problem. Composite Models[4] and Technicolor[5] are clever tentatives to solve it. But they present their own difficulties [6].

There is a third proposed approach to render such a theory "natural". It is known that chiral symmetries are able to preserve fermion masses from phenomenologically unacceptable too large values. They are, nevertheless, not sufficient to do the same with scalar boson masses since terms like $\phi\phi^*$ always respect the symmetry. One could think of a symmetry that would tie together bosons and fermions. In this way, the chiral symmetry that forbids certain fermion masses would also forbid the bosonic partner masses. This is exactly what SUPERSYMMETRY does(7). In the limit where Supersymmetry is exact and fermions are chirally "protected", masses in the scalar sector vanish. The actual small (comparing with GUT scale) value of the fermionic and scalar masses would be due to suitable breakdown of the chiral symmetry and supersymmetry.

In Supersymmetry, it is required that each known particle has a supersymmetric partner of the same mass and couplings and opposite statistics. In such framework, one finds that there is always a loop of superpartners accompanying the loop of normal particles. The extra minus sign that goes with fermion loops leads to suitable cancellations in certain Feynman diagrams and the theory presents no quadratic divergence. In terms of the equation (1), Supersymmetry guarantees that $C \equiv 0$. Therefore, in a supersymmetric case, the running mass $m_H(\mu_2)$ can be stable at a scale no much larger than μ_2. The "naturalness" problem is solved in supersymmetric theories. With Supersymmetry, however, the fundamental question as to *why* there are two hugely different scales in the theory, $\mu_1 \gg \mu_2$, is not solved yet.

It is interesting enough to be emphasized, that any solution of the "naturalness" problem *necessarily* indicates the existence of some new physics at a scale around $O(1TeV)$. This can be seen again from equation (1). Naturalness is satisfied if all the terms in such equation are of the same order. This requires that the integration which appears there must be of the same order of the scalar mass at μ_2, or, equivalently, $g^2\mu_1^2 \sim \mu_2^2 \Rightarrow \mu_1 \sim \mu_2/g \sim O(1TeV)$. This would be an expected value for a region with new physics.

In the case of Supersymmetry, this new physics consists of a spectrum of new supersymmetric particles (partners of the ordinary particles) which have masses no greater than roughly $1TeV$ and in some cases may be substantially lighter. There has been an enormous interest in searching experimentally these superpartners, and Table 1 provides the experimental limits on these supersymmetric particle masses[(8)]. Such limits are obtained in distinct ways. The most direct one consists of trying to produce the new particles. The (up to now) negative results put some lower bounds on their masses. These bounds are directly connected to the energy available in the experimental apparatus.

The observation of any one of these supersymmetric partners with production and decay through interactions dictated by Supersymmetry would be a direct confirmation of the truth of Supersymmetry in a low energy scale. A typical process of this kind would be the observation of charged scalar pair production at e^+e^- machines, which decays into electrons and photinos (which escape detection).

Normalparticles(spin)	Superparticle(spin)	Masslimits
$quark(\frac{1}{2})$	$squark(0)$	$\geq 80GeV^{\dagger}$
$lepton(\frac{1}{2})$	$slepton(0)$	$\geq 26GeV^{\ddagger}$
$photon(1)$	$photino(\frac{1}{2})$	$---$
$gluon(1)$	$gluino(\frac{1}{2})$	$\geq 90GeV^{\dagger}$
$W^{\pm}(1)$	$wino(\frac{1}{2})$	$\geq 26/40GeV^{\ddagger/*}$
$Z^0(1)$	$zino(\frac{1}{2})$	$\geq 40GeV^{*}$
$Higgs(0)$	$higgsino(\frac{1}{2})$	$---$

Table 1: Current experimental status of superpartners.

($\dagger$=CDF,$\ddagger$=$TRISTAN$,$*$=$CERN$)

There are, however, many indirect ways in which supersymmetry can be infered. They consist mainly in looking for deviations from the predictions of the Standard Model due to the eventual presence of virtual superpartners. Even if these indirect tests are *not complete* confirmations of Supersymmetry, since they have alternative explanations, they are indicative of new physics and place interesting constraints on possible extensions of the Standard Model, in particular, supersymmetric theories.

In this work we will concentrate on these indirect tests of supersymmetry. In particular, we will pursue the very interesting non-standard flavor changing interactions and their CP violation implications at the weak scale present in some models of low-energy supersymmetry.

The extremely small experimental values associated to **Flavor Changing Neutral Currents** have always been a challenge for models which intend to describe phenomenology of Particle Physics. Neutral systems, like the mixing in $\kappa^0 - \overline{\kappa}^0 (s\overline{d} \leftrightarrow d\overline{s})$, have been a good testing ground for any new theory. In the Standard Model, the success in describing such phenomenon is guaranteed by the so-called Glashow-Iliopoulos-Maiani mechanism[9], where the unitarity of the quark mixing matrix ensures a natural explanation for the suppression of the flavor changing neutral current processes. There is no $\overline{q}_i q_j Z^0$ coupling at tree level. Such an effect of flavor changing proceeds as a second order weak interaction where it is suppressed by $\Delta m_q^2 / m_q^2$, with Δm_q^2 being the difference of the squared masses of equal-charged quarks.

Many extensions of the Standard Model have been severely constrained, as Composite Models, or even ruled out, as Technicolor Theories, due to difficulties in passing the low energy flavor changing neutral current tests.

As we will explicitly verify in this work, the "Minimal"♣ low energy supersymmetric extension of the Standard Model accounts for the very small numbers associated to the suppressions in the flavor changing neutral currents. This fact can be considered as a large success of the Supersymmetric Theories.

Any extension of the Minimal Model should also respect these flavor changing

♣ see next chapter for definition and discussion of the Minimal Model.

neutral current low energy tests. Therefore, it is very useful to quantitatively determine the values which must be respected by the parameters of such models related to these phenomena. This is what will be done in this work in a model independent way. These results could be used as a guide to the choice of any realistic Supersymmetric Non-minimal Model.

We organize this work in the following way: in Chapter I the low energy supergravity models are introduced and we emphasize the new sources of flavor changing neutral currents associated to the strong gluino-quark-squark coupling. The relevant ideas about Extended Supersymmetric Models and the Mass Insertion Approximation is also introduced. Just for completeness, in Chapter II, some phenomenology related to the Mixing Phenomenon and its CP Violation effects in neutral particle-antiparticle systems is discussed. The main part of this work is Chapter III. There we explicitly calculate the supersymmetric contributions to the mixing in the kaon system and display the constraints that any supersymmetric theory must satisfy in order to be consistent with the current experimental data related to the mixing phenomenon and its CP violation effects.

CHAPTER I

Flavor Changing Neutral Currents in Supersymmetric Models

I.a. Building Models

Supersymmetry is very appealing. Besides placing bosons and fermions in aesthetically attractive common irreducible multiplets, it presents also the remarkable feature of solving the naturalness problem of the gauge hierarchy, as we have already mentioned in the "Motivations".

Local Supersymmetry is of particular interest since it naturally leads to space-time translations that differ from point to point: a general coordinate transformation. From this one might expect gravity to appear in locally supersymmetric theories and this expectation turns out to be true. Local Supersymmetry is also called Supergravity (SUGRA). The $N = 1$ supergravity models[(10)] are the most interesting ones since for $N > 1$ the fermions lie in real representation of the gauge group and thus it is not straightforward to build phenomenologically acceptable models.

Nevertheless it is phenomenologically known that supersymmetry cannot be an exact symmetry of the Nature, as is evidenced by the fact that the masses of the known particles and those of their superpartners are not the same.

In $N = 1$ supergravity models, local supersymmetry is broken spontaneously by a gauge singlet, the so-called hidden sector. It is also assumed that the hidden sector has no coupling in the superpotential with the observable sector, which contains the quarks, leptons and Higgs superfields, the gauge multiplets and, in the case of extended models (beyond the SM), the additional particles. These two sectors only communicate through a very weak interaction, which is naturally chosen to be the gravitational interaction. After the local supersymmetric breakdown, the gravitino (the spin $\frac{3}{2}$ superpartner of the spin 2 graviton) acquires a mass $m_{\frac{3}{2}}$ through the super-Higgs mechanism. The degeneracy inside the supermultiplets is destroyed and the model considered so far has all the features to be realistic, except that it is non-renormalizable!

Nevertheless, we can relieve this fact assuming the flat limit[11], where $M_P \to \infty$ (M_P is the Planck mass) maintaining $m_{\frac{3}{2}}$ fixed in the Supergravity Lagrangian. In this way we turn off gravity and regain global $N = 1$ Supersymmetry with "Soft Breaking Terms". Soft terms explicitly break the residual global supersymmetry in a very suitable way, since they preserve the theory from the unwanted quadratic divergences. A complete list of soft terms in N=1 supersymmetric theories can be found in the work of Girardello and Grisaru[12]. We can write the final Lagrangian:

$$L = L_{SUSY} + L_{SOFT} \tag{1.1}$$

where, as a first approach, L_{SUSY} is the Lagrangian of the usual "Minimal" global supersymmetric extension of the Standard Model. By Minimal we mean that the model presents three main features: The scalar kinetic terms present the canonical form, or, in other words, the Kähler metric is flat. The superpotential does not introduce new superfields apart from those which appear in the trivial supersymmetrization of the Standard Model. Finally, there is no Baryon or Lepton violating terms in the superpotential, or, equivalently, R-parity is respected in a Minimal Supersymmetric Standard Model.

L_{SOFT} is the set of supersymmetric soft breaking terms.

I.b. The Minimal Supersymmetric Standard Model

In the early attempts of constructing supersymmetric extensions of the SM, it was soon realized that the superpartners of the known particles coud not be identified with some other particle already present in the SM. 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 particle was the partner of the leptons, the V.E.V. of the sneutrinos 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 not generated; furthermore the fermionic partner of the usual Higgs doublet renders the particle content of the model anomalous. Two superfields doublets are sufficient to cancel anomalies and to provide masses to all quarks and leptons. The particle content of the minimal supersymmetric SM is given in Table 1.1.

For our purposes we need to write only the superpotential of the $N = 1$ global Supersymmetric Standard Model, which has the most general invariant supersymmetric couplings which preserves the gauge symmetries of the theory[10]:

$$W = h_U Q H_1 u^c + h_D Q H_2 d^c + h_L L H_2 e^c + \mu H_1 H_2 \tag{1.2}$$

Q, u^c, d^c,... are scalar superfields, namely they possess scalar and fermionic components with the same quantum numbers, as was previously described in Table 1.1.

The set of supersymmetric soft breaking terms takes a quite simple aspect in the case of minimal $N = 1$ supergravity theories. They are given by:

$$L_{SOFT} = m^2 \sum_{i=scalars} |\phi|^2 + [Am(h_U \tilde{Q} H_1 \tilde{u}^c + h_D \tilde{Q} H_2 \tilde{d}^c + h_L \tilde{L} H_2 \tilde{e}^c) \tag{1.3}$$

$$+ Bm\mu H_1 H_2 + \sum_a \lambda^a \lambda^a + h.c.]$$

where the sum on i extends over all the scalars and A and B are dimensionless parameters of the trilinear and bilinear contributions, respectively, λ is the gaugino, the superpartner of the gauge field, and a is a gauge group indice. Here m is the scale of the low energy supersymmetric breaking. It is very often related to $m_{\frac{3}{2}}$ and hereafter, we will use $m_{\frac{3}{2}}$ for m.

Table 1.1: Supersymmetric Standard Model particle content.

$Superfields$	$ComponentFields$	$3_C \times 2_L \times 1_Y$ $quantumnumbers$	$Name$
$MatterFields$			
Q	$\begin{pmatrix} u_L \\ d_L \end{pmatrix}$	$(3, 2, \frac{1}{6})$	$quark$
	$\begin{pmatrix} \tilde{u}_L \\ \tilde{d}_L \end{pmatrix}$		$squark$
U	u^c_L	$(3, 1, -\frac{2}{3})$	$R-up$
	$\tilde{u}^c_L$		$squark$
D	d^c_L	$(3, 1, \frac{1}{3})$	$R-down$
	$\tilde{d}^c_L$		$squark$
L	$\begin{pmatrix} \nu_L \\ e_L \end{pmatrix}$	$(1, 2 - \frac{1}{2})$	$lepton$
	$\begin{pmatrix} \tilde{\nu}_L \\ \tilde{e}_L \end{pmatrix}$		$slepton$
E	e^c_L	$(1, 1, 1)$	$R-lepton$
	$\tilde{e}^c_L$		$slepton$
$GaugeFields$			
G	$g_i(i = 1, .., 8)$	$(8, 1, 0)$	$gluon$
	$\tilde{g}_i$		$gluino$
V	$\begin{pmatrix} W^{\pm} \\ W^3 \end{pmatrix}$	$(1, 3, 0)$	W
	$\begin{pmatrix} \tilde{W}^{\pm} \\ \tilde{W}^3 \end{pmatrix}$		$W-ino$
B	B	$(1, 1, 0)$	B
	$\tilde{B}$		$B-ino$

Table 1.1: Supersymmetric Standard Model particle content (cont').

$Superfields$	$ComponentFields$	$3_C \times 2_L \times 1_Y$ $quantumnumbers$	$Name$
$HiggsFields$			
H_1	$\begin{pmatrix} \phi_1^+ \\ \phi_1^0 \end{pmatrix}$	$(1, 2, \frac{1}{2})$	$Higgs$
	$\begin{pmatrix} \tilde{\phi}_1^+ \\ \tilde{\phi}_1^0 \end{pmatrix}$		$Higgsino$
H_2	$\begin{pmatrix} \phi_2^+ \\ \phi_2^0 \end{pmatrix}$	$(1, 2, -\frac{1}{2})$	$Higgs$
	$\begin{pmatrix} \tilde{\phi}_2^+ \\ \tilde{\phi}_2^0 \end{pmatrix}$		$Higgsino$

I.c. The Squark Mass Matrix

It is well known that there is no Flavor Changing Neutral Currents (FCNC) in the Standard Model at tree level. The FCNC transitions are obtained at the one- or higher-loop order through the W exchange. The mere supersymmetrization of these one-loop Standard Model contributions shows that squarks of the same charge must be highly degenerated in order to account for the very small δm_κ in the $\kappa^0 - \overline{\kappa}^0$ system[13]. It means that the mass difference between such squarks has to be much smaller than their average mass. This is a possible justification to take equal masses for the scalars in the minimal N=1 soft-breaking terms, as in (1.3).

Nevertheless, the interest in Flavor Changing Neutral Currents in supersymmetric models grew very rapidily when was realized that another kind of FCNC[14], with no parallel in the Standard Model, was possible in the minimal version of the Supersymmetric Standard Model. Vertices involving neutral gauge fermions can, in principle, give rise to FCNC processes. The vertex quark-squark-gluino is particularly interesting since it involves the strong coupling constant α_S. Therefore, processes related to these vertices give severe constraints to the parameters of the chosen model, as we will treat later.

Let us now analyse in detail this last type of FCNC process. From the superpotential (1.2) and the soft-breaking terms (1.3), it is possible to evaluate the down-squark mass matrix at tree level, or in other words, at the superlarge scale of local supersymmtery breaking. We are able, in this way, to construct the relevant mass matrix:

$$M^2_{\tilde{d}\tilde{d}^*} = \begin{pmatrix} m^2_{\tilde{d}_L\tilde{d}^*_L} & m^2_{\tilde{d}_L\tilde{d}^{c*}_L} \\ m^2_{\tilde{d}^c_L\tilde{d}^*_L} & m^2_{\tilde{d}^c_L\tilde{d}^{c*}_L} \end{pmatrix} \tag{1.4}$$

where:

$$m^2_{\tilde{d}_L\tilde{d}^*_L} = m^2_{\tilde{d}^c_L\tilde{d}^{c*}_L} = m_d m_d^\dagger + m^2_{\frac{3}{2}}\mathbf{1} \tag{1.5}$$

and

$$m^2_{\tilde{d}_L\tilde{d}^{c*}_L} = A m_{\frac{3}{2}} m_d. \tag{1.6}$$

Up to now it is easy to realize that we don't have any flavor change since it is possible to diagonalize both $m_d m_d^\dagger$ and $m^2_{\tilde{d}_L \tilde{d}_L}$ by the same rotation matrix. This is no longer true if we renormalize $m^2_{\tilde{d}_L \tilde{d}_L}$ from its value at superlarge scale (in which equation (1.5) holds good) to much lower ones like m_W scale[15]. The coupling $h_u Q H_1 u^c$ in the superpotential (1.2) gives rise to 1-loop contributions to $m^2_{\tilde{d}_L \tilde{d}^*_L}$ proportional to $m_u m_u^\dagger$ as is evidenced by Figure 1.1.

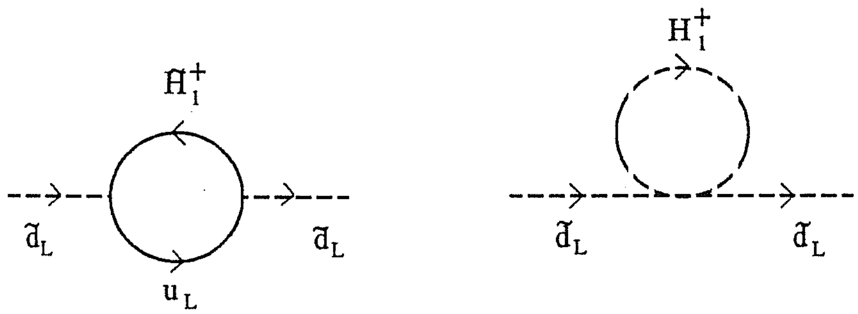

Figure 1.1: 1-loop contributions to $m^2_{\tilde{d}_L \tilde{d}^*_L}$ proportional to $m_u m_u^\dagger$.

As soon as supersymmetry is broken, no exact cancellation occurs, and we expect radiative corrections yielding a mass term for the scalar partner of the left-handed down quark, proportional to $h_u^\dagger h_u$ or, equivalently, to $c m_u^\dagger m_u$ (c depens on the SUSY breaking parameter, in this work, $m_{\frac{3}{2}}$). The proportionality coefficient

"c" is calculated evaluating the renormalization group equations. For $m_{\tilde{g}} < m_{\frac{3}{2}} \simeq 100GeV$ and $45GeV < m_t < 120GeV$, c can range in the -1 to -0.1 interval[16].

Similar terms do not arise for the "right-handed" squarks since they are $SU(2)_L$ singlets.

Nevertheless, there are radiative contributions to the off-diagonal terms in the mass matrix, as indicated by the graphs showed in Figure 2.2. The graph of Figure 2.2.a gives a contribution of the kind $g^2 h_D \langle H_2^0 \rangle \sim m_d$ and, therefore, creates no possibility of flavor changing since it can be diagonalized by the same rotation that diagonalizes m_d. That one of Figure 2.2.b gives a term proportional to (assuming $m_{\frac{3}{2}} \sim \tilde{m}$) $A^3 B \mu h_D h_U^\dagger h_U \langle H_1^0 \rangle \sim \mu m_u \frac{A^3 B m_u m_d}{\langle H_1^0 \rangle \langle H_2^0 \rangle} \sim c' \mu m_u$. The proportionality coefficient c' is usually taken to be zero in this minimal version of the supersymmetric standard model, since it is very suppressed. Therefore, the 1-loop corrected squark mass matrix is finally given by:

$$M^2_{\tilde{d}\tilde{d}^*} = \begin{pmatrix} m_d m_d^\dagger + m^2_{\frac{3}{2}} \mathbf{1} + c m_u m_u^\dagger & A m_{\frac{3}{2}} m_d \\ A m_{\frac{3}{2}} m_d & m_d m_d^\dagger + m^2_{\frac{3}{2}} \mathbf{1} \end{pmatrix} . \qquad (1.7)$$

It is useful to note that up to now everything that one has done for obtaining the down squark mass matrix would work for the up squark mass matrix. In this way, to obtain the latter from the former it is sufficient to interchange $u \leftrightarrow d$ in (1.7).

From the down squark mass matrix it is easy to infer that the only possible source of non-standard flavor changing in the minimal supersymmetric standard model is the mass term of the partner of the left-handed down quark since it is

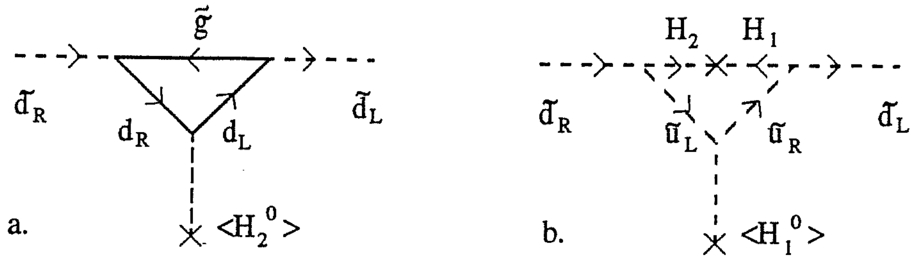

Figure 1.2: 1-loop contributions to the off-diagonal terms in the squark mass matrix.

not possible to simultaneously diagonalize m_d and m_u. Indeed, if m_d and m_u are simultaneously diagonalized, then the Cabibbo-Kobayashi-Maskawa mixing matrix vanishes. If we call $L_{d,u}$ and $R_{d,u}$ the matrices that diagonalize these quark mass matrices,

$$L^{\dagger}_{\substack{u\\d}} m_{\substack{u\\d}} R_{\substack{u\\d}} = m^{diag}_{\substack{u\\d}}, \tag{1.8}$$

we immediately conclude that:

$$m^2_{\tilde{d}_L \tilde{d}^*_L} = (m_d^{diag})^2 + m^2_{\frac{3}{2}} \mathbf{1} + cV^{\dagger}(m_u^{diag})^2 V, \tag{1.9}$$

here $V = L^{\dagger}_u L_d$ is the Cabibbo-Kobayashi-Maskawa mixing matrix.

At this point we have two ways to follow. In the first, we completely diagonalize the whole squark mass matrix:

$$\tilde{U}^{\dagger}_D M^2_{\tilde{d}\tilde{d}^*} \tilde{U}_D =$$

$$= U_D^\dagger \begin{pmatrix} (m_d^{diag})^2 + m_{\frac{3}{2}}^2 \mathbf{1} + cV^\dagger (m_u^{diag})^2 V & |A| m_{\frac{3}{2}} m_d^{diag} e^{i\delta} \\ |A| m_{\frac{3}{2}} m_d^{diag} e^{-i\delta} & (m_d^{diag})^2 + m_{\frac{3}{2}}^2 \mathbf{1} \end{pmatrix} U_D \tag{1.10}$$

where

$$\tilde{U}_D = \begin{pmatrix} e^{i\Phi_{m_{\tilde{g}}}} L_d & 0 \\ 0 & e^{-i\Phi_{m_{\tilde{g}}}} R_d \end{pmatrix} U_D \tag{1.11}$$

and

$$\delta = \Phi_A - 2\Phi_{m_{\tilde{g}}}. \tag{1.12}$$

These angles Φ are phases in the soft-breaking Lagrangian parameter A and in the gluino mass $m_{\tilde{g}}$ defined such that:

$$\begin{aligned} A &= |A|\, e^{i\Phi_A} \\ m_{\tilde{g}} &= |m_{\tilde{g}}|\, e^{-2i\Phi_{m_{\tilde{g}}}} \end{aligned} \tag{1.13}$$

The introduction of the gluino mass phase $\Phi_{m_{\tilde{g}}}$, in the way shown above, leads to the extraction of this spurious phase in the gluino mass term and in the interaction Lagrangian (1.14).

There are many interesting informations in the matrix (1.10). First, we note that the rotation matrix for the squarks contains an extra phase δ which is CP-violating. It is, nevertheless, severely constrained by the experimental upper limit on the electric dipole moment of the neutron to be roughly less than $10^{-3} \leftrightarrow 10^{-2}$ [(18)]. Second, it is useful to emphasize that down quark mass and down squark mass matrix can not be diagonalized in the same basis and so there is a possibility of flavor changing. The source of this flavor changing is $cV^\dagger (m_u^{diag})^2 V$ as is evidenced by the fact that if we take the limit $c \to 0$ the U_D matrix does not mix different families of squarks. The same consideration applies for the up mass matrix. The difference here is that the radiative term $cV^\dagger(m_d^{diag})^2V$, which appears in the up squark mass matrix, can be neglected when compared with the up-quark type masses.

This flavor changing feature of the down squark mass matrix will be directly reflected in the gluino-quark-squark vertex which Lagrangian is given by:

$$L_{q\tilde{q}\tilde{g}} = -\sqrt{2} g_S T^a \sum_{i=u,d} (\bar{\tilde{g}}^0_a P_L q^0_i \tilde{q}^{0*}_{iL} + \bar{q}^0_i P_R \tilde{g}^0_a \tilde{q}^0_{iL} - \bar{\tilde{g}}^0_a P_R q^0_i \tilde{q}^{0*}_{iR} - \bar{q}^0_i P_L \tilde{g}^0_a \tilde{q}^0_{iR}) \quad (1.14)$$

where g_S is the strong coupling constant, g^0 is gluino field, q^0 and $\tilde{q}^0$ are the three generations of up and down quark and squark fields in the current eingenstates , T^a is the $SU(3)_C$ generator and the projectors $P_{\frac{L}{R}}$ are the usual ones defined as $P_{\frac{L}{R}} = (1 \mp \gamma_5)/2$.

The transformation from current eingestates to mass eingestates is defined as follows:

$$\begin{aligned} q^0_L &= L q_L \\ q^0_R &= R q_R \\ \begin{pmatrix} \tilde{q}^0_L \\ \tilde{q}^0_R \end{pmatrix} &= \begin{pmatrix} e^{i\Phi_{m_{\tilde{g}}}} L & 0 \\ 0 & e^{-i\Phi_{m_{\tilde{g}}}} R \end{pmatrix} U_D \begin{pmatrix} \tilde{q}_L \\ \tilde{q}_R \end{pmatrix} . \end{aligned} \quad (1.15)$$

Substituting the transformations (1.15) in (1.14) we note that the matrix U_D completely determines the mixing between different families. In the general case, as U_D is non diagonal, there will exist flavor changing gluino-quark-squark interactions. As we have already mentioned, down-squark mass matrix has stronger flavor violations than those of up-squark matrix.

I.d. The Mass Insertion

There is another way to treat the flavor violations which appear in Supersymmetric Models due to the non-trivial diagonalization of the squark mass matrix. It consists of treating the term which leads to flavor changing as Mixing Mass Insertions. Mass insertion can be thought of as a perturbation in the symmetric theory. This is a very useful way to conduct calculation, specially in model independent analysis of extensions of the Minimal Supersymmetric Standard Model where we do not know exactly the squark mass matrix. We will largely apply this concept in the future chapters.

In the present case, we isolate the term proportional to c in the squark mass matrix (1.10). In this way, the interaction Lagrangian (1.14) leads to diagonal gluino-quark-squark coupling (we work in the so-called superKM basis[(18)]). The entire effect of mixing is placed in the squark propagators as a mass insertion (detailed and denoted as a cross in Figure 1.4).

In order to apply the usual mass insertion approximation we must be sure that such terms lie in the allowed region by perturbation theory. Very important transitions used in this work are those wherein the b goes to s or s goes to d. In these cases, the related mixing mass insertions can be generalized:

$$\frac{(cm_u m_u^\dagger)_{ij}}{m_{\frac{3}{2}}^2} \rightarrow \frac{c[V^\dagger (m_u^{diag})^2 V]_{ij}}{m_{\frac{3}{2}}^2} \sim cV_{ti}V_{tj}^* \frac{m_t^2}{m_{\frac{3}{2}}^2} < 1. \qquad (1.16)$$

i and j stand for b and s (or s and d). V_{ti} and V_{tj} are the entries [33] and [32] (or [32] and [31]) of the Cabibbo-Kobayashi-Maskawa matrix and we neglected terms proportional to m_c^2. We assume also $m_{\frac{3}{2}}$ around the weak scale ($m_{\frac{3}{2}} \sim 100GeV$) and $60GeV < m_t < 130GeV$ (19). The Cabibbo-Kobayashi-Maskawa matrix elements follow the Cabibbo hierarchy described in Figure 1.3. The m_t^2 contribution dominates (1.16) because we have:

$$\frac{m_t^2 V_{ti}^\dagger V_{tj}}{m_c^2 V_{ci}^\dagger V_{cj}} \geq \frac{m_t^2 V_{ts}^\dagger V_{td}}{m_c^2 V_{cs}^\dagger V_{cd}} \sim \frac{m_t^2}{m_c^2}\frac{\lambda^5}{\lambda} \sim 2.3 \times 10^{-3}\frac{m_t^2}{m_c^2} > 1,$$

since $m_t > 21 m_c$.

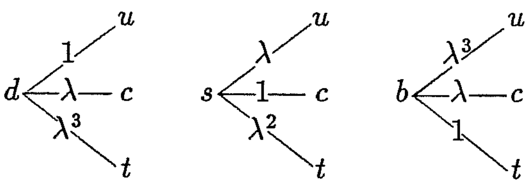

Figure 1.3: the Cabibbo hierarchy. $\lambda \sim sen\theta_{Cabibbo} \sim 0.2$.

In our calculations we will largely apply the mass insertion approximation, even in the cases where the value of the mass insertion is not known (only in the Minimal Supersymmetric Standard Model this value is known). We assume that we are in a range of validity of perturbation theory. We will conclude, *a posteriori*, that there are upper bounds for the mass insertion terms which justify this assumption.

I.e. FCNC in Non-minimal Models

In the minimal version of the supersymmetric standard model, treated hitherto, the major source of flavor changing neutral currents arises in the strong $\tilde{g} - d_{iL} - \tilde{d}_{jL}$ interactions (i and j denote flavor), since only the submatrix $m^2_{\tilde{d}_L \tilde{d}^*_L}$ receives renormalization contributions non-proportional to $m_d m_d^\dagger$. For instance, flavor transitions among squarks of the kind $\tilde{d}_{iL} \leftrightarrow \tilde{d}^c_{jL}$ are suppressed since they involve at least two mass insertions, indicated by Δ, as is schematically indicated in Figure 1.4.

The above argument is not necessarily true in extended versions of the supersymmteric standard model. There could be cases where the transitions $\tilde{d}_L \leftrightarrow \tilde{d}^c_L$ are no longer proportional to m_d and could induce flavor changing, or even, transitions like $\tilde{d}^c_L \leftrightarrow \tilde{d}^c_L$ could lead to no more negligible flavor changing terms. Thus, the above mentioned suppression may disappear.

There are some ways to arrive at Non-minimal Supersymmetric Models. In reference (18), the authors discuss the additional sources of Flavor Changing Neutral Currents in models where one gives up the requirement of putting the kinetic

$$\overset{\tilde{d}_{iL}}{---}\times\overset{\tilde{d}_{jL}}{---}\equiv[\Delta_{LL}]_{ij}=c[V^{\dagger}(m_u^{diag})^2V]_{ij}$$

$$\overset{\tilde{d}^c_{iL}}{---}\times\overset{\tilde{d}^c_{jL}}{---}\equiv[\Delta_{RR}]_{ij}\rightarrow negligible$$

$$\overset{\tilde{d}_{iL}}{---}\times\overset{\tilde{d}^c_{jL}}{---}\equiv[\Delta_{LR}]_{ij}\rightarrow\overset{\tilde{d}_{iL}}{---}\times\overset{\tilde{d}_{jL}}{---}\times\overset{\tilde{d}_{jR}}{---}$$

Figure 1.4: Squark flavor changing neutral currents in the Minimal Supersymmetric Standard Model. (i, j indicate flavor)

terms of scalar field in a canonical form. In reference (20), non-standard superfields were introduced and their couplings with the ordinary fields of the Supersymmetric Standard Model again give rise to new Flavor Changing Neutral Currents. For this reason, when one analyses the extensions of the minimal supersymmetric standard model, all possible transitions, involving the mass insertions Δ_{LL}, Δ_{LR} or Δ_{RR}, are important. That is why we are potentially interested in extensions of the supersymmetric standard model. This is a very elegant way of conducting a model independent analysis of the flavor changing phenomenon in extended versions of the supersymmetric Standard Model[(21)] and we will adopt it in this work. These extra mass insertions are shown in the Figure 1.5.

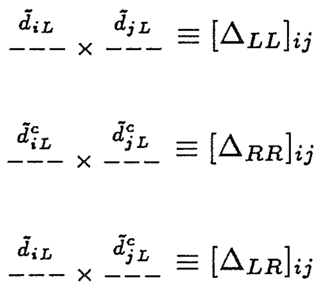

Figure 1.5: Mass insertions in non-minimal models.

CHAPTER II

PHENOMENOLOGY

II.a. Mixing phenomenon

Oscillations between particles and antiparticles were predicted for the $\kappa^0 - \overline{\kappa}^0$ system in 1955[22] and observed in 1956[23]. There are some other systems where we can expect particle-antiparticle oscillations, namely, systems containing c-, b- or t-quarks:

$$D^0(c\overline{u}) - \overline{D}^0(\overline{c}u)$$

$$B^0(\overline{b}d) - \overline{B}^0(b\overline{d})$$

$$B_s(\overline{b}s) - \overline{B}_s(b\overline{s})$$

$$T_u^0(t\overline{u}) - \overline{T}_u^0(\overline{t}u)$$

$$T_c^0(t\overline{c}) - \overline{T}_c^0(\overline{t}c)$$

$B^0 - \overline{B}^0$ and $B_s - \overline{B}_s$ systems are the only ones where oscillations present sizable effects. The original formalism introduced for the $\kappa^0 \longleftrightarrow \overline{\kappa}^0$ transition can be easily used in these other systems. The main differences are heavier masses and shorter lifetimes, and many more final states available in the decays. We will construct a very general formalism and specify it later for each case we will be interested in.

Nature has provided us with a wonderful physical situation to study the CP violation phenomenon in the κ^0 system. *Id est*, the closeness of the κ^0 mass to the three-pion masses. Due to this phase space limitation, κ_L, which mainly decays into 3π, has a much longer lifetime than κ_S and can travel a few meters before it decays away. The crucial point in this effect is that after a few 10^{-10} seconds mainly κ_L will remain, which is almost an equal mixture of κ^0 and $\overline{\kappa}^0$, and this is exactly the case of maximal mixing. The mass difference $\delta m = m_L - m_S$ and the CP violating parameters ϵ and ϵ' can be studied from the κ_L and κ_S decays separately, from $\pi^+\pi^-$ final states interference through the regeneration of κ_S in a κ_L beam, or from charge asymmetry of $l^\pm$ in the decays of $\kappa_L \to \pi^\pm l^\mp \nu$ [(24)]. It is remarkable (and a little frustrating) that so far the only CP violation observed is still the originally observed one in the system of κ_L and κ_S in 1964.

It is much harder to study CP violation effects in heavy quark systems. The lifetime of the heavy κ_L-, κ_S-like states are much shorter, since their masses are very big. The observed charmed and b-particle lifetimes, τ_c and τ_b, are of order 10^{-13} sec and 10^{-12} sec, respectively, and the lifetime τ_t, for the t-particle, is expected to be less than $O(10^{-18})$ sec[(25)]. There is no phase space limitation in the decay channels available for either heavy κ_L- or κ_S-like states, as for κ_L itself. The lifetime for these heavy quark states are expected to be comparable. Therefore, it may be interesting to study their CP violation effects via time integrated results[26].

In the construction of the general formalism for all available neutral particle-antiparticle systems, we will consider a generical and fictitious $G^0 - \overline{G}^0$ system.

Due to the possible transition $G^0 \longleftrightarrow \overline{G}^0$, induced by weak interactions, the original G^0 and $\overline{G}^0$ eigenstates are no longer physical states. These physical states are obtained after the diagonalization of the effective Hamiltonian:

$$H = \begin{pmatrix} M - \frac{i}{2}\Gamma & M_{12} - \frac{i}{2}\Gamma_{12} \\ M_{12}^* - \frac{i}{2}\Gamma_{12}^* & M - \frac{i}{2}\Gamma \end{pmatrix}$$

$$\longrightarrow \begin{pmatrix} M - \frac{i}{2}\Gamma + \frac{1}{2}(\delta m - \frac{i}{2}\delta\Gamma) & 0 \\ 0 & M - \frac{i}{2}\Gamma - \frac{1}{2}(\delta m - \frac{i}{2}\delta\Gamma) \end{pmatrix} \tag{2.1}$$

where

$$\frac{1}{2}(\delta m - \frac{i}{2}\delta\Gamma) = [(M_{12} - \frac{i}{2}\Gamma_{12})(M_{12}^* - \frac{i}{2}\Gamma_{12}^*)]^{1/2} \tag{2.2}$$

M_{ij} and Γ_{ij} are transition matrix elements and can be complex [27]:

$$M_{ij} = m_G\delta_{ij} + \langle i|H_{W,\Delta flavor=2}|j\rangle + P\sum_{\lambda} \frac{\langle i|H_{W,\Delta flavor=1}|\lambda\rangle\langle\lambda|H_{W,\Delta flavor=1}|j\rangle}{m_G - E_\lambda}, \tag{2.3}$$

$$\Gamma_{ij} = 2\pi\sum_{\lambda} \varrho\langle i|H_{W,\Delta flavor=1}|\lambda\rangle\langle\lambda|H_{W,\Delta flavor=1}|j\rangle$$

and ϱ is the density of the λ states, $\delta m \equiv m_L - m_S$, $\delta\Gamma \equiv \Gamma_L - \Gamma_S, \Gamma \equiv (\Gamma_L + \Gamma_S)/2$ and $M = (m_L + m_S)/2$. The requirement of CPT invariance and hermiticity was used in the phenomenological Hamiltonian above. Its eigenstates are found to be:

$$|G_{\substack{L\\S}}\rangle = [2(1+|\overline{\epsilon}_G|^2)]^{-1/2}[(1+\overline{\epsilon}_G)|P^0\rangle \pm (1-\overline{\epsilon}_G)|\overline{P}^0\rangle],$$

$$\overline{\epsilon}_G = \frac{\sqrt{M_{12}-\frac{i}{2}\Gamma_{12}} - \sqrt{M_{12}^*-\frac{i}{2}\Gamma_{12}^*}}{\sqrt{M_{12}-\frac{i}{2}\Gamma_{12}} + \sqrt{M_{12}^*-\frac{i}{2}\Gamma_{12}^*}} \tag{2.4}$$

Since the $|G^0\rangle$ and $|\overline{G}^0\rangle$ do not communicate through strong interactions their relative phase is not specified. $|G^0\rangle$ and $|\overline{G}^0\rangle$ are related through CP transformation up to an arbitrary phase ζ:

$$CP|G^0\rangle = e^{2i\zeta}|\overline{G}^0\rangle. \tag{2.5}$$

As a consequence of this fact, $\overline{\epsilon}_G$ is not a physical parameter. It is useful, therefore, to define an independent phase convention parameter for the measurement of CP violation in the G-system. A very interesting one is:

$$\eta_G \equiv \left|\frac{1-\overline{\epsilon}_G}{1+\overline{\epsilon}_G}\right| = \left|\frac{M_{12}^*-\frac{i}{2}\Gamma_{12}^*}{M_{12}-\frac{i}{2}\Gamma_{12}}\right|^{1/2}. \tag{2.6}$$

The amount of CP-violation is given by the deviation of η_G from unity.

The very interesting mixing phenomenon is evidenced by the time evolution of the physical states $G_{\substack{L\\S}}$:

$$|G_{\substack{L\\S}}(t)\rangle = exp(im_{\substack{L\\S}}t - \Gamma_{\substack{L\\S}}t/2)|G_{\substack{L\\S}}(0)\rangle \tag{2.7}$$

If we consider the time evolution for the states such that $|\phi(t=0)\rangle = |G^0\rangle$ and

$|\overline{\phi}(t=0)\rangle = |\overline{G}^0\rangle$, we obtain:

$$|\phi(t)\rangle = a_G(t)|G^0\rangle + a_{\overline{G}}(t)|\overline{G}^0\rangle \tag{2.8}$$

$$|\overline{\phi}(t)\rangle = \overline{a}_G(t)|G^o\rangle + \overline{a}_{\overline{G}}(t)|\overline{G}^0\rangle \tag{2.9}$$

where

$$\begin{matrix}\overline{a}_G(t)\\ a_G(t)\end{matrix} = [exp(im_Lt - \Gamma_Lt/2) \begin{matrix}-\\+\end{matrix} exp(im_St - \Gamma_St/2)]$$

$$\begin{matrix}\overline{a}_{\overline{G}}(t)\\ a_{\overline{G}}(t)\end{matrix} = \frac{1-\overline{\epsilon}_G}{1+\overline{\epsilon}_G}[-exp(im_Lt - \Gamma_Lt/2) \begin{matrix}-\\+\end{matrix} exp(im_St - \Gamma_St/2)] \tag{2.10}$$

As we have previously stated, the relevant quantity here is time integrated probability of $|\phi(t=0)\rangle = |G^0\rangle$ state goes into the $|G^0\rangle$ and $|\overline{G}^0\rangle$:

$$r \equiv \frac{(\text{“}G^0 \to \overline{G}^0\text{”})}{(\text{“}G^0 \to G^0\text{”})} = \frac{\int_0^\infty |a_{\overline{G}}(t)|^2 dt}{\int_0^\infty |a_G(t)|^2 dt} = \eta_G^2 \Delta \tag{2.11}$$

$$\overline{r} \equiv \frac{(\text{“}\overline{G}^0 \to G^0\text{”})}{(\text{“}\overline{G}^0 \to \overline{G}^0\text{”})} = \frac{\int_0^\infty |\overline{a}_G(t)|^2 dt}{\int_0^\infty |\overline{a}_{\overline{G}}(t)|^2 dt} = \eta_G^{-2} \Delta \tag{2.12}$$

$$\Delta \equiv \frac{(\delta m/\Gamma)^2 + (\frac{1}{2}\delta\Gamma/\Gamma)^2}{2 + (\delta m/\Gamma)^2 - (\frac{1}{2}\delta\Gamma/\Gamma)^2}. \tag{2.13}$$

Analysing this last equation we note that there are two situations where the mixing is maximal. Or $|\delta\Gamma/2| \approx \Gamma$ (i.e., either $\Gamma_L \gg \Gamma_S$ or $\Gamma_S \gg \Gamma_L$), or $\delta m \sim \Gamma$. An example of the first case is the κ-system, it does not matter which is the initial amount of G^0 or $\overline{G}^0$ in the initial state, after some time the system will present only G_L or G_S (respectively) which is an appoximately equal mixture of G^0 and $\overline{G}^0$, therefore maximal mixing. The second case is exampled by heavy b-quark systems, like $B_S^0 - \overline{B}_S^O$. In such case, the system oscillates very quickly between G^0 and $\overline{G}^0$ and appears as an "equal mixture" of G^0 and $\overline{G}^0$. Some simple relations are very intertesting. In both cases of maximal mixing, described above, one has:

$$r.\overline{r} \approx 1. \tag{2.14}$$

For small CP violation $\eta_G \approx 1$:

$$r \approx \overline{r}. \tag{2.15}$$

In the extreme case where both maximal mixing and small CP violation are present:

$$r \approx \overline{r} \approx 1. \tag{2.16}$$

The neutral particle-antiparticle mixing phenomenon presents the very attractive apparent effect that $|G^0\rangle$ can have decay products that belong exclusively to $|\overline{G}^0\rangle$.

II.b. The $\kappa^0 - \overline{\kappa}^0$ System

In this section we will consider the mixing phenomenon and its related CP violations in the kaon system. We, first of all, define the phase convention between the $|\kappa^0\rangle$ and $|\overline{\kappa}^0\rangle$ states as is commonly found in the literature:

$$CP|\kappa^0\rangle = |\overline{\kappa}^0\rangle. \tag{2.17}$$

Since the observed CP violation in this system is experimentally small, we are allowed to make some approximations which lead to a small value of $\overline{\epsilon}_K$. In order to obtain the mentioned small value for $\overline{\epsilon}_K$, we use the approximations $ImM_{12} \ll ReM_{12}$ and $Im\Gamma_{12} \ll Re\Gamma_{12}$. From the equation (2.2), which in the above approximations yields:

$$\delta m_K - \frac{i}{2}\delta\Gamma_K = 2ReM_{12} - iRe\Gamma_{12} \tag{2.18}$$

we can clearly find the wanted mass difference δm_K:

$$\delta m_K \simeq 2ReM_{12}. \tag{2.19}$$

The other relevant quantity is the CP violating parameter $\overline{\epsilon}_K$. From the equation (2.4) and using the quoted approximations, we get:

$$\overline{\epsilon}_K \approx \frac{(M_{12} - \frac{i}{2}\Gamma_{12}) - (M^*_{12} - \frac{i}{2}\Gamma^*_{12})}{4\sqrt{(M_{12} - \frac{i}{2}\Gamma_{12})(M^*_{12} - \frac{i}{2}\Gamma^*_{12})}} \sim \tag{2.20}$$

$$\sim \frac{ImM_{12}}{\frac{1}{2}\delta\Gamma_K - i\delta m_K} \tag{2.21}$$

here we neglected the contribution $Im\Gamma_{12}$ after examinating the possible intermediate states $\lambda = 2\pi, 3\pi, \pi e\mu, ...$ in the M_{12} and Γ_{12}, equation (2.3), for concluding that $ImM_{12} \gg Im\Gamma_{12}$.

Finally, we apply the experimentally known result in the kaon physics:

$$\delta m_K \simeq \frac{\delta \Gamma_K}{2} \tag{2.22}$$

to obtain:

$$|\bar{\epsilon}_K| \simeq \frac{1}{\sqrt{2}} \frac{Im M_{12}}{\delta m_K} \tag{2.23}$$

We end this section by presenting some interesting experimental parameters related to the decays of the kaon mesons. The basic CP violation parameters that have been measured are the amplitude ratios:

$$\eta_{\pm} \equiv \frac{\langle \pi^+ \pi^- | H | \kappa_L \rangle}{\langle \pi^+ \pi^- | H | \kappa_S \rangle} \tag{2.24}$$

$$\eta_{00} \equiv \frac{\langle \pi^0 \pi^0 | H | \kappa_L \rangle}{\langle \pi^0 \pi^0 | H | \kappa_S \rangle} \tag{2.25}$$

and the "charge asymmetry":

$$\delta \equiv \frac{\Gamma(\kappa_L \to \pi^+ l^- \overline{\nu}_l) - \Gamma(\kappa_L \to \pi^- l^+ \overline{\nu}_l)}{\Gamma(\kappa_L \to \pi^+ l^- \overline{\nu}_l) + \Gamma(\kappa_L \to \pi^- l^+ \overline{\nu}_l)}$$

From this last definition we note that it involves processes which are CP transformed one of the other. Therefore, we expect this value different from zero if CP is violated. In fact, its experimental value is $\delta = (0.33 \pm 0.012) \times 10^{-2}$.

In $\kappa_{S,L} \to 2\pi$ decays, the angular momentum vanishes. The spatial part of the wave function is therefore symmetric and, since $I(\pi) = 1$, the symmetric statistics of the pions assures that final-state pions are found in $I = 0$ or $I = 2$

isospin states. Introducing the following numbers normalized to the amplitude $\langle I = 0|H|\kappa_S\rangle$:

$$\epsilon_0 = \frac{\langle I = 0|H|\kappa_L\rangle}{\langle I = 0|H|\kappa_S\rangle} \tag{2.27}$$

$$\epsilon_2 = \frac{1}{\sqrt{2}}\frac{\langle I = 2|H|\kappa_L\rangle}{\langle I = 0|H|\kappa_S\rangle} \tag{2.28}$$

$$w = \frac{\langle I = 2|H|\kappa_S\rangle}{\langle I = 0|H|\kappa_S\rangle}. \tag{2.29}$$

One can relate the isospin states to the physical 2π states:

$$\begin{aligned} |I = 0\rangle &= \frac{1}{\sqrt{3}}|\pi^-\pi^+\rangle - \frac{1}{\sqrt{3}}|\pi^0\pi^0\rangle + \frac{1}{\sqrt{3}}|\pi^+\pi^-\rangle \\ |I = 2\rangle &= \frac{1}{\sqrt{6}}|\pi^-\pi^+\rangle + \sqrt{\frac{2}{3}}|\pi^0\pi^0\rangle + \frac{1}{\sqrt{6}}|\pi^+\pi^-\rangle \end{aligned} \tag{2.30}$$

and obtain:

$$\eta_{+-} = \frac{(\epsilon_0 + \epsilon_2)}{1 + w/\sqrt{2}} \tag{2.31}$$

$$\eta_{00} = \frac{(\epsilon_0 - 2\epsilon_2)}{1 - \sqrt{2}w} \tag{2.32}$$

Because of the validity of the $\Delta I = \frac{1}{2}$ rule for CP-conserving decays, we find $w \ll 1$ and it can be neglected. Furthermore, we can parametrize $\kappa^0 \to 2\pi$ amplitude as:

$$\langle I = n|H|\kappa^0\rangle = A_n e^{i\delta_n} \tag{2.33}$$

$$\langle I = n|H|\overline{\kappa}^0\rangle = \overline{A}_n e^{i\delta_n} \tag{2.34}$$

where δ_n is the $\pi\pi$ phase shift in the $I = n$ channel coming from the final state interactions. A commonly adopted phase convention is to choose A_0 to be real:

$$ImA_0 = 0. \tag{2.35}$$

And, by CPT Theorem, we find that:

$$\overline{A}_n = A_n^* \tag{2.36}$$

Therefore, it is straigthforward that:

$$\epsilon_0 = \frac{\langle I=0|H|\Big[(1+\overline{\epsilon}_K)|\kappa_0\rangle - (1-\overline{\epsilon}_K)|\overline{\kappa}^0\rangle\Big]}{\langle I=0|H|\Big[(1+\overline{\epsilon}_K)|\kappa_0\rangle + (1-\overline{\epsilon}_K)|\overline{\kappa}^0\rangle\Big]} = \overline{\epsilon}_K \equiv \epsilon. \tag{2.37}$$

In a similar way, we find:

$$\epsilon_2 = \frac{1}{\sqrt{2}} \frac{\langle I=2|H|\left[(1+\epsilon_K|)\kappa_0\rangle - (1-\epsilon_K)|\overline{\kappa}^0\right]}{\langle I=0|H|\left[(1+\epsilon_K|)\kappa_0\rangle + (1-\epsilon_K)|\overline{\kappa}^0\right]} = \frac{i}{\sqrt{2}} \frac{ImA_2}{A_0} e^{i(\delta_2-\delta_0)} \equiv \epsilon'. \tag{2.38}$$

With these relations for ϵ and ϵ' we can write:

$$\eta_{+-} = \epsilon + \epsilon' \tag{2.39}$$

and:

$$\eta_{00} = \epsilon - 2\epsilon' \tag{2.40}$$

ϵ' is the CP-violation parameter in decay in an analogous way as ϵ is the CP-violation parameter in states. ϵ' is said to be a direct CP-violation effect. It is a measurement of the difference in the CP violation amount in $\kappa_L \rightarrow \pi^0\pi^0$ and in $\kappa_L \rightarrow \pi^+\pi^-$ decays.

CHAPTER III

SUPERSYMMETRY AND THE KAON SYSTEM

III.a. Kaon System in the Standard Model

The $\kappa^0 - \overline{\kappa}^0$ has always been a good test for new theories. In the Standard Model, the absence of Flavor Changing Neutral Currents at tree level is insured by the GIM mechanism. Quantum corrections are suppressed by terms like $\delta m_q^2/m_q^2$, where δm_q^2 is the difference of the squared mass of the same charge quarks. Any extensions of the Standard Model must produce such a fine cancellation in order to pass the Flavor Changing Neutral Current tests.

For completeness, we will briefly present the computation of the $\kappa_L - \kappa_S$ mass difference in the context of the Standard Model and then give this calculation in the supersymmetric framework. Such mass-difference receives short distance contributions through the equation (2.19) and the effective Hamiltonian:

$$M_{12} = M_{\kappa^0\overline{\kappa}^0} = \langle \kappa^0 | H^{eff}_{\Delta S=2} | \overline{\kappa}^0 \rangle \tag{3.1}$$

From the graphs showed in the Figure 3.1, we can calculate the Standard Model one-loop effective Hamiltonian to obtain[28]:

$$\frac{G_F}{\sqrt{2}} \frac{\alpha}{4\pi sin^2\theta_W} \tilde{\Theta} (\overline{s}_L \gamma_\mu d_L)^2 . \tag{3.2}$$

where:

$$\tilde{\Theta} = \sum_{i,j=2}^{N} V_{is}^{*} V_{id} V_{js}^{*} V_{jd} \Theta(x_i, x_j) \tag{3.3}$$

V is the Kobayashi-Maskawa matrix, N is the number of quark generations, $x_i = m_{u_i}^2/m_W^2$ and, finally:

$$\Theta(x_i, x_j) = -x_i x_j \left[\frac{1}{x_i - x_j} \left[\frac{1}{4} - \frac{3}{2}\frac{1}{x_i - 1} - \frac{3}{4}\frac{1}{(x_i - 1)^2} \right] ln x_i + \right. \\ \left. +(x_i \leftrightarrow x_j) - \frac{3}{4}\frac{1}{(x_i - 1)(x_j - 1)} \right] \tag{3.4}$$

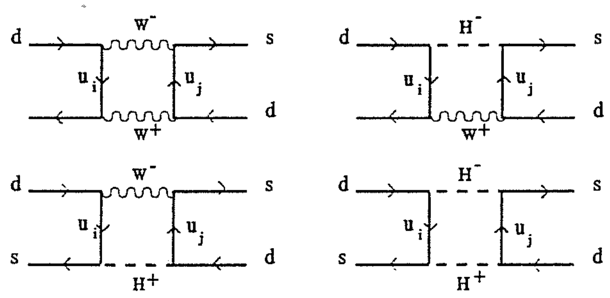

Figure 3.1: the Standard Model contributions to the effective $\overline{s}d\overline{s}d$ operator.

(i and j are family indices)

Substituting this result in the equation (3.1) we obtain:

$$\delta m_K = -\frac{G_F}{\sqrt{2}} \frac{\alpha}{6\pi sin^2\theta_W} m_K f_K^2 B \eta Re\tilde{\Theta} \tag{3.5}$$

In the equation above we find the parameter f_K is the kaon decay constant, η indicates all possible QCD corrections and B is a parameter which measures the ratio between the matrix element $\langle \kappa^0 | (\overline{s}_L \gamma_\mu d_L)^2 | \overline{\kappa}^0 \rangle$ and the same element with the vacuum insertion in it. There is some uncertainty in the calculation of this last parameter. Several authors have been working on this subject[29−35]. The found values vary from −0.4 in the MIT bag-model[33] approximation and 2.9 in the harmonic oscillator approximation[29]. We assume throughout this work the validity of the vaccum insertion approximation where $B = 1$.

The final result of this calculation is very dependent on the charm-quark mass and is written as

$$(\delta m_K)_{charm} = 2\, B \times 10^{-15} GeV. \tag{3.6}$$

We do not expect a significant deviation from this value even if the top-quark mass is very large since this contribution is highly suppressed by Cabibbo-Kobayaski--Maskawa factors. The value given above is to be compared with the experimental one:

$$(\delta m_K)_{exp} = (3.521 \pm .014) \times 10^{-15} GeV. \tag{3.7}$$

Comparing (3.7) and (3.6) we see that, if the vacuum approximation works ($B = 1$), the charm contribution to the mass difference in the kaon system doesn't saturate the experimental value (3.7). This could require an adjustment on the top-quark mass or on its Kobayashi-Maskawa couplings [36]. A fourth generation[37] or a left-right symmetry [38] were proposed as a tentative to explain the experimental enhancement showed above.

Some kind of less-conservative "new-physics" could be invoked to account for the extra contribution. Supersymmetry will be explored in the next section.

We will see that supersymmetric models, besides presenting a possible explanation to the enhancement detailed above, are also very constrained by this same experimental result.

III.b. Supersymmetric Contributions to $K^0 - \overline{K}^0$ Mixing and CP Violation

In this section we will consider the Non-minimal Supersymmetric Standard Model contributions to the mixing phenomenon in the $K^0 - \overline{K}^0$ system. The present experimental value related to this phenomenon constrains the choice of the spontaneuously broken $N = 1$ Supergravity Theories putting very powerful upper bounds in the possible values of the model independent parameters $[\Delta_{LL}]_{sd}$, $[\Delta_{LR}]_{sd}$ and $[\Delta_{RR}]_{sd}$, which are the mass insertions related to the transitions $\tilde{s}_L \leftrightarrow \tilde{d}_L$, $\tilde{s}^c_L \leftrightarrow \tilde{d}_L$ and $\tilde{s}^c_L \leftrightarrow \tilde{d}^c_L$, respectively. These parameters were introduced in Chapter I.

Particularly interesting is the gluino-quark-squark vertex which induces flavor changing, since it severely constrains the non-standard contributions to the mixing phenomena due to the presence of the strong coupling constant in it. Here, we will carefully consider the contributions of such vertices to the mixing in the $K^0 - \overline{K}^0$ system. The relevant Feynman diagrams with the gluino exchange are displayed in Figure 3.2. We emphasize that there are diagrams (e.g., 3.2.b diagram) where the Majorana nature of the gluino allows further contributions to the mixing.

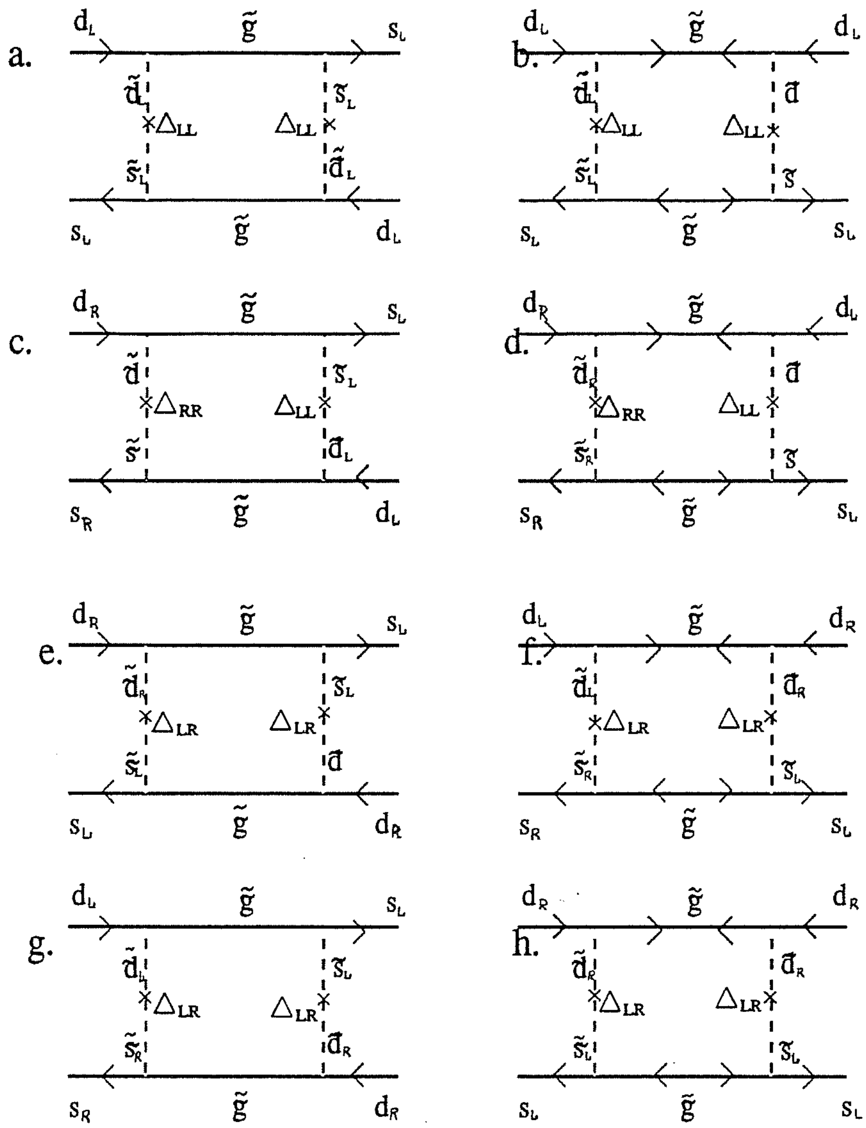

Figure 3.2: The supersymmetric gluino box diagrams contributing to the $\kappa^0 - \overline{\kappa}^0$ matrix element.

The matrix element induced by the gluino exchanging box diagrams:

$$M_{\kappa^0\overline{\kappa}^0}(\tilde{g}) = \langle \kappa^0 | H^{eff}_{\Delta S=2}(\tilde{g}) | \overline{\kappa}^0 \rangle, \tag{3.8}$$

will be explicitly evaluated, since they involve many subtleties. The final result of this calculation can be read in the equation (3.19). The relevant Feynman rules for the present calculation are obtained from the Lagrangian $L_{q\tilde{q}\tilde{g}}$, equation (1.14), and are given in the Figure 3.3.

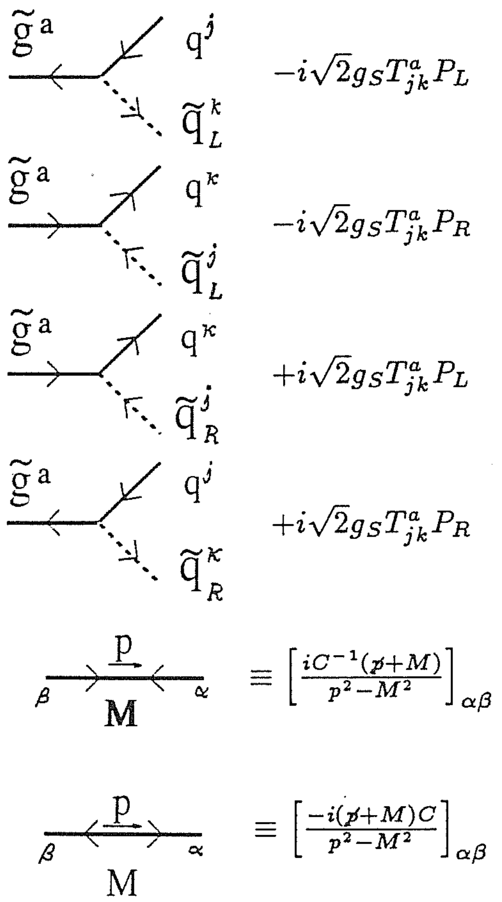

Figure 3.3: the Feynman rules for the calculation of the gluino exchanging box diagrams.

In order to calculate the above mentioned supersymmetric contributions, we will use quantities such as V's, S's, Σ's, A's and B's[17]. They are defined in the Appendices A, B and C, wherein are displayed the evaluation of the matrix elements between κ^0 and $\overline{\kappa}^0$ states in the vacuum insertion method, some useful Fierz identities, and some functional integrations, respectively. We also define M as being the gluino mass and $\tilde{m}$ as an average squark mass.

In the approximation of vanishing external momenta, the graph 2.1.a yields the following contribution:

$$\begin{aligned} a =& -2ig_S^4[\Delta_{LL}]_{sd}^2 \\ &\times \langle\overline{\kappa}^0|(\overline{s}_\gamma\gamma^\mu P_L d^\beta)(\overline{s}_\alpha\gamma_\mu P_L d^\delta)|\kappa^0\rangle(T^bT^a)_\delta^\gamma(T^bT^a)_\beta^\alpha \\ &\times \frac{1}{4}\int\frac{d^4k}{(2\pi)^4}\frac{k^2}{(k^2-M^2)^2(k^2-\tilde{m}^2)^4} \\ =& 2\alpha_S^2[\Delta_{LL}]_{sd}^2\frac{B(y)}{4\tilde{m}^6}V_+^D. \end{aligned} \tag{3.9}$$

The diagram 2.1.b also gives contributions proportional to $[\Delta_{LL}]_{sd}^2$:

$$\begin{aligned} b =& -2ig_S^4[\Delta_{LL}]_{sd}^2 M^2 \\ &\times \langle\overline{\kappa}^0|(\overline{s}_\gamma P_R C\overline{s}_\beta^T)(d^{T\delta}C^{-1}P_L d^\alpha)|\kappa^0\rangle(T^bT^a)_\alpha^\beta(T^bT^a)_\delta^\gamma \\ &\times \int\frac{d^4k}{(2\pi)^4}\frac{1}{(k^2-M^2)^2(k^2-\tilde{m}^2)^4} \end{aligned} \tag{3.10}$$

With the Fierz identity (B.1) of the Appendix B we can write:

$$\begin{aligned} b =& -2ig_S^4[\Delta_{LL}]_{sd}^2 M^2 \\ &\times \langle\overline{\kappa}^0|\frac{1}{2}(\overline{s}_\gamma\gamma^\mu P_L d^\alpha)(\overline{s}_\beta\gamma_\mu P_L d^\delta)|\kappa^0\rangle(T^bT^a)_\delta^\gamma(T^bT^a)_\beta^\alpha \\ &\times \int\frac{d^4k}{(2\pi)^4}\frac{1}{(k^2-M^2)^2(k^2-\tilde{m}^2)^4}. \end{aligned} \tag{3.11}$$

Using the formulas in Appendices A and C, we find:

$$b = 2\alpha_S^2[\Delta_{LL}]_{sd}^2\frac{A(y)M^2}{\tilde{m}^8}(\frac{1}{2}V_+^M). \tag{3.12}$$

The contributions shown above are exactly the same given by the graphs 2.1.a and 2.1.b when one interchanges the external left-handed quarks by right-handed ones. This implies that the contributions proportional to $[\Delta_{RR}]^2_{sd}$ would present the same form as those above with the change $[\Delta_{LL}]^2_{sd} \leftrightarrow [\Delta_{RR}]^2_{sd}$.

It is thus noted that the contributions of the graphs 2.1.c and 2.1.d are proportional to $\left([\Delta_{LL}]_{sd}[\Delta_{RR}]_{sd}\right)$, and are given by:

$$\begin{aligned} c = & -2ig_S^4[\Delta_{LL}]_{sd}[\Delta_{RR}]_{sd}M^2 \\ & \times \langle \overline{\kappa}^0|(\overline{s}_\alpha P_R d^\delta)(\overline{s}_\gamma P_R d^\beta)|\kappa^0\rangle (T^aT^b)^\alpha_\beta (T^bT^a)^\gamma_\delta \\ & \times \int \frac{d^4k}{(2\pi)^4} \frac{1}{(k^2-M^2)^2(k^2-\tilde{m}^2)^4} \\ = & 2\alpha_S^2[\Delta_{LL}]_{sd}[\Delta_{RR}]_{sd}\frac{A(y)M^2}{\tilde{m}^8}S_-^D. \end{aligned} \tag{3.13}$$

and

$$\begin{aligned} d = & -2ig_S^4[\Delta_{LL}]_{sd}[\Delta_{RR}]_{sd} \\ & \times \langle \overline{\kappa}^0|(\overline{s}_\alpha \gamma^\mu P_L C\overline{s}_\gamma^T)(d^{T\beta}C^{-1}\gamma_\mu P_R d^\delta)|\kappa^0\rangle (T^bT^a)^\alpha_\beta (T^aT^ba)^\gamma_\delta \\ & \times \frac{1}{4}\int \frac{d^4k}{(2\pi)^4} \frac{k^2}{(k^2-M^2)^2(k^2-\tilde{m}^2)^4} \\ = & -2ig_S^4[\Delta_{LL}]_{sd}[\Delta_{RR}]_{sd} \\ & \times \langle \overline{\kappa}^0|2(\overline{s}_\alpha P_R d^\delta)(\overline{s}_\gamma P_L(d^\beta)|\kappa^0\rangle (T^bT^a)^\alpha_\beta (T^bT^a)^\gamma_\delta \\ & \times \frac{1}{4}\int \frac{d^4k}{(2\pi)^4} \frac{k^2}{(k^2-M^2)^2(k^2-\tilde{m}^2)^4} \\ = & 2\alpha_S^2[\Delta_{LL}]_{sd}[\Delta_{RR}]_{sd}\frac{B(y)}{4\tilde{m}^6}(2S_-^M). \end{aligned} \tag{3.14}$$

In the above result we used the Fierz identity (B.2). All the remaining graphs give contributions proportional to $[\Delta_{LR}]^2_{sd}$. From the 2.1.e we obtain:

$$\begin{aligned} e = & -2ig_S^4[\Delta_{LR}]^2_{sd}M^2 \\ & \times \langle\overline{\kappa}^0|(\overline{s}_\alpha P_R d^\delta)(\overline{s}_\gamma P_R d^\beta)|\kappa^0\rangle(T^aT^b)^\alpha_\beta(T^bT^a)^\gamma_\delta \\ & \times \int\frac{d^4k}{(2\pi)^4}\frac{1}{(k^2-M^2)^2(k^2-\tilde{m}^2)^4} \\ = & 2\alpha_S^2[\Delta_{LR}]^2_{sd}\frac{A(y)M^2}{\tilde{m}^8}S_+^D. \end{aligned} \tag{3.15}$$

For the calculation of the graph 2.1.f, we will use the Fierz identity (B.3):

$$\begin{aligned} f = & -2ig_S^4[\Delta_{LR}]^2_{sd} \\ & \times \langle\overline{\kappa}^0|(\overline{s}_\gamma C\gamma^\mu P_R\overline{s}_\alpha^T)(d^{T\delta}C^{-1}\gamma_\mu P_R d^\beta)|\kappa^0\rangle(T^bT^a)^\alpha_\beta(T^bT^a)^\gamma_\delta \\ & \times\frac{1}{4}\int\frac{d^4k}{(2\pi)^4}\frac{k^2}{(k^2-M^2)^2(k^2-\tilde{m}^2)^4} \\ = & -2ig_S^4[\Delta_{LR}]^2_{sd} \\ & \times \langle\overline{\kappa}^0|(\overline{s}_\gamma\gamma^\nu P_R d^\beta)(\overline{s}_\alpha\gamma_\nu P_L d^\delta)|\kappa^0\rangle(T^bT^a)^\alpha_\beta(T^bT^a)^\gamma_\delta \\ & \times\frac{1}{4}\int\frac{d^4k}{(2\pi)^4}\frac{k^2}{(k^2-M^2)^2(k^2-\tilde{m}^2)^4} \\ = & 2\alpha_S^2[\Delta_{LR}]^2_{sd}\frac{B(y)}{4\tilde{m}^6}V_-^M. \end{aligned} \tag{3.16}$$

The latter graphs to be calculated are the 2.1.g contribution:

$$\begin{aligned} g = & -2ig_S^4[\Delta_{LR}]^2_{sd} \\ & \times \langle\overline{\kappa}^0|(\overline{s}_\gamma\gamma^\mu P_R d^\beta)(\overline{s}_\alpha\gamma_\mu P_L d^\delta)|\kappa^0\rangle(T^aT^b)^\alpha_\beta(T^bT^a)^\gamma_\delta \\ & \times\frac{1}{4}\int\frac{d^4k}{(2\pi)^4}\frac{k^2}{(k^2-M^2)^2(k^2-\tilde{m}^2)^4} \\ = & 2\alpha_S^2[\Delta_{LR}]^2_{sd}\frac{B(y)}{4\tilde{m}^6}V_-^D. \end{aligned} \tag{3.17}$$

and the 2.1.h:

$$\begin{aligned}
h =& - 2ig_S^4[\Delta_{LR}]_{sd}M^2 \\
&\times \langle \overline{\kappa}^0|(\overline{s}_\alpha C P_R \overline{s}_\gamma^T)(d^{T\beta} C^{-1} P_R d^\delta)|\kappa^0\rangle (T^b T^a)^\alpha_\beta (T^b T^a)^\gamma_\delta \\
&\times \frac{1}{4}\int \frac{d^4k}{(2\pi)^4}\frac{1}{(k^2-M^2)^2(k^2-\tilde{m}^2)^4} \\
=&2ig_S^4[\Delta_{LR}]_{sd}M^2 \\
&\times \langle \overline{\kappa}^0|\frac{1}{2}(\overline{s}_\alpha P_R d^\beta)(\overline{s}_\gamma P_R d^\delta) - \frac{1}{8}(\overline{s}_\alpha \sigma^{\mu\nu} P_R d^\beta)(\overline{s}_\gamma \sigma_{\mu\nu} d^\delta|\kappa^0\rangle (T^b T^a)^\alpha_\beta (T^b T^a)^\gamma_\delta \\
&\times \int \frac{d^4k}{(2\pi)^4}\frac{1}{(k^2-M^2)^2(k^2-\tilde{m}^2)^4} \\
=&2\alpha_S^2[\Delta_{RL}]^2_{sd}\frac{A(y)M^2}{\tilde{m}^8}(\frac{1}{2}S_+^M + \frac{1}{8}\Sigma^M).
\end{aligned} \tag{3.18}$$

Putting the equations (3.9) - (3.18) together, we write the final result for the supersymmetric contribution involving the strong gluino-quark-squark coupling to the mixing phenomenon in the $\kappa^0 - \overline{\kappa}^0$ system♠:

$$\begin{aligned}
M_{\kappa^0\overline{\kappa}^0} =& M_{12} \\
=& m_K f_K^2 \alpha_S^2 \\
&\times\Bigg[\Big([\Delta_{LL}]^2_{sd} + [\Delta_{RR}]^2_{sd}\Big)\Big(\frac{11}{108}\frac{B(y)}{\tilde{m}^6} + \frac{1}{27}\frac{A(y)M^2}{\tilde{m}^8}\Big) + \\
&+ [\Delta_{LL}]_{sd}[\Delta_{RR}]_{sd}\Big[\Big(\frac{1}{18} - \frac{1}{27}(\frac{m_K}{m_s+m_d})^2\Big)\frac{B(y)}{\tilde{m}^6} \\
&\qquad + \Big(\frac{1}{9} + \frac{16}{27}(\frac{m_K}{m_s+m_d})^2\Big)\frac{A(y)M^2}{\tilde{m}^8}\Big] + \\
&+ [\Delta_{LR}]^2_{sd}\Big[\Big(-\frac{7}{54} - \frac{2}{9}(\frac{m_K}{m_s+m_d})^2\Big)\frac{B(y)}{\tilde{m}^2} - \frac{11}{27}(\frac{m_K}{m_s+m_d})^2\frac{A(y)M^2}{\tilde{m}^8}\Big]\Bigg]
\end{aligned} \tag{3.19}$$

♠ see Appendix A for definitions.

III.c. Bounds on the Supersymmetric Parameters

The next step to be done consists in taking the current values of the quantities which enter in the equation (3.19), and are given in the Appendix D, and of comparing the result with the experimental value of δm_K through the equation (2.19). We will be able, in this way, to constrain the possible values for Δ's, putting an upper bound for each of them. In order to obtain such results, we assume that no great cancellation occurs between the $[\Delta_{LL}]^2_{sd}$, $[\Delta_{RR}]^2_{sd}$, $[\Delta_{LR}]^2_{sd}$ and $[\Delta_{LL}]_{sd}[\Delta_{RR}]_{sd}$ contributions. In doing so, we are not considering the most general case,but, significative cancellations between these terms occur only if they are very close, which is a rare situation. Barring such a case, we can safely estimate an independent bound on the Δ's. We will look for the value of each Δ which alone saturates the experimental value for δm_K.

In the limit where the gluino mass is equal to the average squark mass, $y \equiv M^2/\tilde{m}^2 \to 1$, we get:

$$\lim_{y\to 1} A(y) = \frac{1}{20} \tag{3.20}$$

$$\lim_{y\to 1} B(y) = -\frac{1}{30}. \tag{3.21}$$

In this limit we obtain the following upper bounds for the model independent parameters:

$$Re\left(\frac{[\Delta_{LR}]^2_{sd}}{\tilde{m}^4}\right) \leq 1.4 \times 10^{-6} \frac{\tilde{m}^2}{M_W^2}, \tag{3.22}$$

$$Re\left(\frac{[\Delta_{LL}]_{sd}[\Delta_{RR}]_{sd}}{\tilde{m}^4}\right) \leq 5.5 \times 10^{-7}\frac{\tilde{m}^2}{M_W^2} \tag{3.23}$$

and

$$Re\left(\frac{[\Delta_{LL}]_{sd}^2}{\tilde{m}^4}\right) \leq 1.1 \times 10^{-4}\frac{\tilde{m}^2}{M_W^2}. \tag{3.24}$$

We opted for showing these limits through the adimensional quantities $[\Delta]_{sd}^2/\tilde{m}^4$. The last result can be transformed into a limit for $[\Delta_{RR}]_{sd}^2$ simply doing $[\Delta_{LL}]_{sd}^2 \leftrightarrow [\Delta_{RR}]_{sd}^2$.

The experimental value$^\diamond$ of the CP-violation parameter $\bar{\epsilon}_K$ can be used to constrain the imaginary part of the parameters Δ's through the equation (2.23):

$$Im\left(\frac{[\Delta_{LR}]_{sd}^2}{\tilde{m}^4}\right) \leq 3.9 \times 10^{-9}\frac{\tilde{m}^2}{M_W^2}, \tag{3.25}$$

$$Im\left(\frac{[\Delta_{LL}]_{sd}[\Delta_{RR}]_{sd}}{\tilde{m}^4}\right) \leq 1.6 \times 10^{-9}\frac{\tilde{m}^2}{M_W^2} \tag{3.26}$$

and

$$Im\left(\frac{[\Delta_{LL}]_{sd}^2}{\tilde{m}^4}\right) \leq 3.2 \times 10^{-7}\frac{\tilde{m}^2}{M_W^2}. \tag{3.27}$$

We can again recover the limit on the $[\Delta_{RR}]_{sd}^2$ proceeding with the substitution $[\Delta_{LL}]_{sd}^2 \leftrightarrow [\Delta_{RR}]_{sd}^2$ in the last relation. It will be useful for our next discussion, to display the results of the equations (3.24) and (3.27) in a slightly

$^\diamond$ see Appendix D for the current experimental values used in this work.

different way. Remembering that

$$
\begin{aligned}
Re\Delta^2 =& (Re\Delta)^2 - (Im\Delta)^2 \\
Im\Delta^2 =& 2Re\Delta Im\Delta
\end{aligned}
\quad , \tag{3.28}
$$

we can write the following results:

$$
Re\frac{[\Delta_{LL}]_{sd}}{\tilde{m}^2} \leq 1.0 \times 10^{-2}\frac{\tilde{m}}{M_W} \tag{3.29}
$$

and

$$
Im\frac{[\Delta_{LL}]_{sd}}{\tilde{m}^2} \leq 2.1 \times 10^{-5}\frac{\tilde{m}}{M_W}. \tag{3.30}
$$

The results shown in the equations (3.22) - (3.30) are very general and model independent, *i.e.*, all supersymmetric models must respect these limits in order to pass the low energy tests related to flavor changing neutral currents in the kaon system. One could immediately ask: does the Minimal Supersymmetric Standard Model pass this test? To answer this question, we turn back to Chapter I to find that the relevant (non-negligible) quantity related to flavor changing neutral currents in this model is the $[\Delta]_{LL}$ given by the equation (1.9) or (1.16):

$$
[\Delta_{LL}^{min}]_{sd} = c[V^\dagger(m_u^{diag})^2V]_{sd} \sim cV_{ts}^\dagger V_{td}m_t^2. \tag{3.31}
$$

In order to write a numerical value for $Re[\Delta_{LL}^{min}]_{sd}$ and $Im[\Delta_{LL}^{min}]_{sd}$ and then to compare these values with those previously found in our model independent analysis, we must analyse the current status of the Cabibbo-Kobayashi-Maskawa angular factors which enter in the equation (3.31). We adopt here the Wolfenstein parametrization[39] of the Cabibbo-Kobayashi-Maskawa matrix:

$$\mathbf{V} = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} = \begin{pmatrix} 1-\frac{\lambda^2}{2} & \lambda & A\lambda^3\rho e^{i\phi} \\ -\lambda & 1-\frac{\lambda^2}{2} & A\lambda^2 \\ A\lambda^3(1-\rho e^{-i\phi}) & -A\lambda^2 & 1 \end{pmatrix} \tag{3.32}$$

where λ is the experimentally well known sine of the Cabibbo angle:

$$\lambda = sin\theta_{Cabibbo} = 0.221 \pm 0.002 \tag{3.33}$$

and A, ρ and ϕ are parameters to be determined.

In this parametrization, the real and the imaginary part of $[\Delta_{LL}^{min}]_{sd}$ can be determined from the equations (3.31) and (3.32). We obtain:

$$Re[\Delta_{LL}^{min}]_{sd} \simeq cm_t^2 Re[V_{ts}^\dagger V_{td}] = cm_t^2[-A^2\lambda^5(1-\rho cos\phi)] \tag{3.34}$$

and

$$Im[\Delta_{LL}^{min}]_{sd} \simeq cm_t^2 Im[V_{ts}^\dagger V_{td}] = cm_t^2[-A^2\lambda^5\rho sin\phi] \tag{3.35}$$

We now consider the available experimental information on A, ρ and ϕ. In order to infer the value of A, we must analyse some suitable experiment where the $b \to c$ transition is present, because $V_{cb} = A\lambda^2$ and λ is known. In fact, A can be fixed by the b-lifetime τ_B and the semileptonic branching ratio $B_{SL} = B(b \to e\nu X)$. We will not enter into details of these calculations. For our purposes, it is sufficient to say that the semileptonic width $\Gamma_{SL} = B_{SL}/\tau_B$ can be computed by the parton model improved by QCD corrections to obtain:

$$|V_{bc}|^2 = \xi\Gamma_{SL}[sec] \tag{3.36}$$

where the coefficient ξ is explored in reference (40). Using the current values for B_{SL} and τ_B (see Appendix D), one finds:

$$A = 1.05 \pm 0.17 \tag{3.37}$$

The parameter ρ is fixed by the ratio of the widths:

$$R = \frac{\Gamma(b \to ue\nu)}{\Gamma(b \to ce\nu)} \tag{3.38}$$

once that:

$$(0.47 \pm 0.02)R = |\frac{V_{ub}}{V_{cb}}|^2 = (\lambda\rho)^2. \tag{3.39}$$

The numerical factor is obtained from the parton model plus QCD corrections[40] and the last equality is read straightforward from equation (3.32). At 90% C.L., one has:

$$\begin{aligned} R &< 0.13^{(CRYSTALBALL)} \\ &< 0.12^{(ARGUS)} \\ &< 0.06^{(CUSB)} \\ &< 0.04^{(CLEO)} \end{aligned} \tag{3.40}$$

Taking the CLEO result (the strongest upper bound), we find:

$$R < 0.04 \Rightarrow |\frac{V_{bu}}{V_{bc}}|^2 < 0.14 \Rightarrow \rho < 0.6 \tag{3.41}$$

Finally, on $cos\phi$, one observes that it is strongly dependent on the top quark mass m_t. In fact, the Figure 3.4, taken from reference (41), shows that the exper-

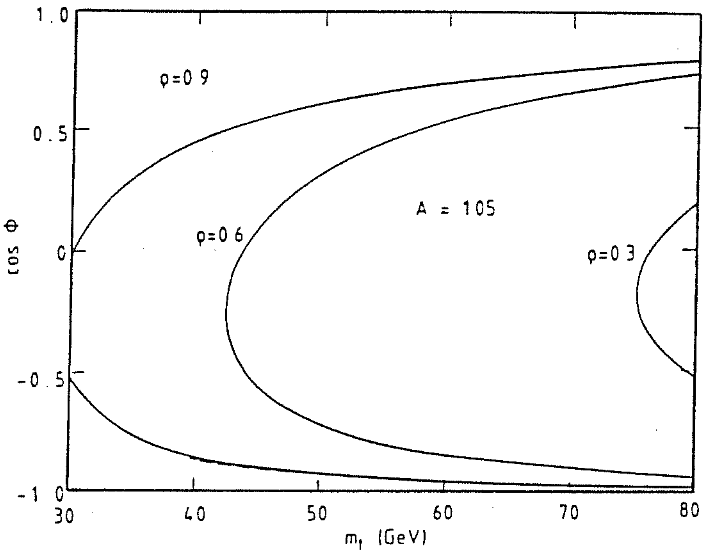

Figure 3.4: Limits on $cos\phi$ obtained from the experimental value of the CP violating parameter ϵ for the kaon system as functions of the top quark mass m_t, for various values of ρ. Here one takes $A = 1.05$ [the central value in equation (3.37)]. The indicated values of $\rho = 0.9, 0.6$ and 0.3 correspond to $R = \Gamma(b \to u)/\Gamma(b \to c) = 0.08, 0.04$ and 0.009, respectively.

imental value of the CP violating parameter ϵ for the kaon system is consistent with the currently large values of m_t ($m_t > 77GeV$) if $|cos\phi| \sim 1$.

Furthermore, analysing the equation (3.34), we observe that the most disfavorable situation concerning the value of the parameter ϕ, i.e., the value of this parameter which most pushes the absolute value of $Re[\Delta_{LL}^{min}]_{sd}$ up, towards the dangerous upper bound given in the equation (3.29), is, precisely, $cos\phi = -1$. It is interesting to notice that, for this value, all CP violating effects vanish. So, it is not a realist value. Anyway, it is a nice value to perform our comparison.

Thus, using the values for A and ρ, given by the equations (3.37) and (3.41), and $cos\phi = -1$, we obtain from (3.34):

$$\frac{Re[\Delta_{LL}^{min}]_{sd}}{\tilde{m}^2} < 2.2 \times 10^{-3} \frac{\tilde{m}}{M_W} \qquad (3.42)$$

The values of the parameters which enter in the calculation of the $Re[\Delta_{LL}^{min}]_{sd}$ were chosen such that, respecting the eventual present experimental values, again most push its value up. We used the top quark mass $m_t = 130GeV$, the average squark mass $\tilde{m} = 80GeV$ and $c = -1$.

The equation (3.42) is to be compared with the equation (3.29) which shows the value of $Re[\Delta_{LL}]_{sd}$ which saturates the experimental value of δm_K in a Supersymmetric Model. We can conclude that, even for the largest allowed values of the Cabibbo-Kobayashi-Maskawa angular factors, the Minimal Supersymmetric Standard Model yields a contribution to δm_K less than its experimental value.

It is an enthusiastic fact verifing that the Minimal Supersymmetric Model is compatible with (3.29). Nevertheless, it is also true that such a conclusion is a little frustrating since the results shown above do not lead to any interesting constraint in any of the still free parameters of the model. For instance, it does not yield any upper bound for the top mass or any other free parameter.

A different situation arises when we analyse the imaginary part of $[\Delta_{LL}]_{sd}^2$. If the top quark mass is heavy, the experimental value of $|\bar{\epsilon}_K|$ leads to a strong bound on one of the still free parameters of the model. From (3.28), (3.32) and (3.35), we can write:

$$Im[\Delta_{LL}]^2_{sd} = 2c^2 m_t^4 A^4 \lambda^{10} \rho sin\phi(1 - \rho cos\phi) \tag{3.43}$$

This value is relevant in the Minimal Supersymmetric Standard Model contribution to $|\bar{\epsilon}_K|$. Through the equations (2.23) and (3.19) we find:

$$|\bar{\epsilon}_K^{min}| = 6.86 \times 10^{-4} c^2 \rho sin\phi(1 - \rho cos\phi) \frac{M_W^2 m_t^4}{\tilde{m}^6} \tag{3.44}$$

Comparing with the experimental result of $|\bar{\epsilon}_K|$ (given in Appendix D), we obtain:

$$\rho sin\phi(1 + \rho cos\phi) \simeq 3.22 \frac{\tilde{m}^6}{c^2 M_W^2 m_t^4}. \tag{3.45}$$

For the extreme values $c^2 = 1$, $\tilde{m}$=80 GeV and m_t=130 GeV, this value becomes:

$$\rho \sin\phi(1 + \rho \cos\phi) \simeq 0.461 \tag{3.46}$$

which can be compared with that calculated in reference (40). There, in order to reproduce the experimental value for $|\bar{\epsilon}_K|$ in the Standard Model, the authors found:

$$\rho \sin\phi(1 + \rho \cos\phi) \simeq 0.335 \tag{3.47}$$

In fact, it is interesting to compare the Minimal Supersymmetric contribution to $|\bar{\epsilon}_K|$ given in equation (3.44) with the Standard Model calculations for this quantity[42]:

$$
\begin{aligned}
|\bar{\epsilon}_K^{SM}| =& 5.9 B_K \lambda^2 (A^2\lambda^4 \rho \sin\phi)[\eta_c m_c^2 + \eta_t m_t^2 \frac{A_t(x_t)}{x_t} \times \\
& \times A^2\lambda^4(1-\rho cos\phi) + \eta_{ct} m_c^2 ln\frac{m_t}{m_c}][(GeV)^{-2}]
\end{aligned}
\tag{3.48}
$$

with:

$$
\begin{aligned}
A(x_i) &= x_i[\frac{1}{4} + \frac{9}{4}(1-x_i)^{-1} - \frac{3}{2}(1-x_i)^{-2}] + \frac{3}{2}[\frac{x_i}{x_i-1}]^3 ln x_i \\
x_i &= m_i^2/M_W^2
\end{aligned}
\tag{3.49}
$$

and η_c, η_t and η_{ct} are QCD correction factors which values are, approximately, 0.8, 0.6 and 0.4, respectively. Taking $B_K = 1$, $m_c = 1.5GeV$ and $m_t \geq 80GeV$, it is straightforward to conclude that the m_t^2 term dominates (3.48). So we can calculate the ratio:

$$
\frac{|\bar{\epsilon}_K^{min}|}{|\bar{\epsilon}_K^{SM}|} \sim \frac{1}{5} c^2 \frac{M_W^4 m_t^2}{\tilde{m}^6}
\tag{3.50}
$$

Analysing the extreme case when $c^2 = 1$, $m_t = 130GeV$ and $\tilde{m} = 80GeV$, the ratio in (3.50) can be written:

$$
\frac{|\bar{\epsilon}_K^{min}|}{|\bar{\epsilon}_K^{SM}|} \sim \frac{1}{2}
\tag{3.51}
$$

Therefore, we can conclude that for suitable values of the parameters c, m_t and $\tilde{m}$, the Minimal Supersymmetric Standard Model gives non-negligible contributions to $|\bar{\epsilon}_K|$.

APPENDIX A

In this Appendix A we give the matrix elements between κ^0 and $\overline{\kappa}^0$ evaluated in the vacuum insertion method. One defines[(17)]:

$$V_{AB}^{D(M)} \equiv \langle \overline{\kappa}^0 | \overline{s}_\alpha \gamma^\mu P_A d^\beta \overline{s}_\gamma \gamma_\mu P_B d^\delta | \kappa \rangle \times (T^a T^b)^\alpha_\delta (T^{b(a)} T^{a(b)})^\gamma_\beta, \qquad (A.1)$$

$$S_{AB}^{D(M)} \equiv \langle \overline{\kappa}^0 | \overline{s}_\alpha P_A d^\beta \overline{s}_\gamma P_B d^\delta | \kappa \rangle \times (T^a T^b)^\alpha_\delta (T^{b(a)} T^{a(b)})^\gamma_\beta, \qquad (A.2)$$

$$\Sigma_{A}^{D(M)} \equiv \langle \overline{\kappa}^0 | \overline{s}_\alpha \sigma^{\mu\nu} P_A d^\beta \overline{s}_\gamma \sigma_{\mu\nu} d^\delta | \kappa \rangle \times (T^a T^b)^\alpha_\delta (T^{b(a)} T^{a(b)})^\gamma_\beta. \qquad (A.3)$$

where $A, B = L, R$ and $\alpha, \beta, \gamma, \delta = 1, 2, 3$ are color indices. The following relations are satisfied:

$$V_+^{D(M)} \equiv V_{LL}^{D(M)} = V_{RR}^{D(M)}, \qquad (A.4)$$

$$V_-^{D(M)} \equiv V_{LR}^{D(M)} = V_{RL}^{D(M)}, \qquad (A.5)$$

$$S_+^{D(M)} \equiv S_{LL}^{D(M)} = S_{RR}^{D(M)}, \qquad (A.6)$$

$$S_-^{D(M)} \equiv S_{LR}^{D(M)} = S_{RL}^{D(M)}, \qquad (A.7)$$

$$\Sigma^{D(M)} \equiv \Sigma_L^{D(M)} = \Sigma_R^{D(M)}. \qquad (A.8)$$

The final result regarding the matrix elements, remembering the relation between κ^0 and $\overline{\kappa}^0$ states, i.e., $|\overline{\kappa}^0(p)\rangle = CP|\kappa^0(-p)\rangle$, is:

$$V_+^D = \frac{11}{27} N \qquad (A.9)$$

$$V_+^M = \frac{2}{27}N \quad (A.10)$$

$$V_-^D = -(\frac{8}{27} + \frac{2}{9}R)N \quad (A.11)$$

$$V_-^M = (\frac{1}{27} - \frac{2}{9}R)N \quad (A.12)$$

$$S_+^D = -\frac{13}{54}RN \quad (A.13)$$

$$S_+^M = \frac{5}{54}RN \quad (A.14)$$

$$S_-^D = (\frac{1}{18} + \frac{8}{27}R)N \quad (A.15)$$

$$S_-^M = (\frac{1}{18} - \frac{1}{27}R)N \quad (A.16)$$

$$\Sigma^D = \frac{2}{3}RN \quad (A.17)$$

$$\Sigma^M = \frac{2}{3}RN \quad (A.18)$$

with:

$$N \equiv -\frac{1}{2}m_K f_K^2, \quad (A.19)$$

$$R \equiv \frac{m_K^2}{(m_d + m_s)^2}. \quad (A.20)$$

APPENDIX B

In this Appendix B we give some useful Fierz Identities used in the present work:

$$(\overline{a}P_R C\overline{b}^T)(c^T C^{-1} P_L d) = \frac{1}{2}(\overline{a}\gamma^\mu P_L d)(\overline{b}\gamma_\mu P_L c) \tag{B.1}$$

$$(\overline{a}\gamma^\mu P_L C\overline{b}^T)(c^T C^{-1}\gamma_\mu P_R d) = 2(\overline{a}P_R d)(\overline{b}P_L c) \tag{B.2}$$

$$(\overline{a}\gamma^\mu P_R C\overline{b}^T)(c^T C^{-1}\gamma_\mu P_R d) = (\overline{a}\gamma^\mu P_R d)(\overline{b}\gamma_\mu P_L c) \tag{B.3}$$

$$(\overline{a}C P_R\overline{b}^T)(c^T C^{-1} P_R d) = \frac{1}{2}(\overline{a}P_R d)(\overline{b}P_R c) - \frac{1}{8}(\overline{a}\sigma^{\mu\nu} P_R d)(\overline{b}\sigma_{\mu\nu} c) \tag{B.4}$$

APPENDIX C

In this Appendix C we give some useful functional integrations used in this work, we define $y = M^2/\tilde{m}^2$, where, in our calculations, M is the gluino mass and $\tilde{m}$ is an average squark mass:

$$\begin{aligned}
&\int \frac{d^4k}{(2\pi)^4} \frac{1}{(k^2 - M^2)^2 (k^2 - \tilde{m}^2)^4} \\
&= \frac{i}{16\pi^2 \tilde{m}^8} \frac{(1-y)^{-5}}{6} \left(-y^3 + 9y^2 + 9y - 17 - 6(3y+1) lny\right) \qquad (C.1) \\
&\equiv \frac{i}{16\pi^2 \tilde{m}^8} A(y),
\end{aligned}$$

$$\begin{aligned}
&\int \frac{d^4k}{(2\pi)^4} \frac{k^2}{(k^2 - M^2)^2 (k^2 - \tilde{m}^2)^4} \\
&= -\frac{i}{16\pi^2 \tilde{m}^6} \frac{(1-y)^{-5}}{3} \left(-y^3 - 9y^2 + 9y + 1 + 6y(1+y) lny\right) \qquad (C.2) \\
&\equiv \frac{i}{16\pi^2 \tilde{m}^6} B(y)
\end{aligned}$$

APPENDIX D

In this Appendix D are given experimental or theoretical values for some quantities used in this work:

$$m_K \sim 500 MeV \tag{D.1}$$

$$m_s \sim 150 MeV \tag{D.2}$$

$$m_d \sim 10 MeV \tag{D.3}$$

$$M_W \sim 80 GeV \tag{D.4}$$

$$f_K \sim 1.23 f_\pi \sim 1.23 \times 93 MeV \sim 114 MeV \tag{D.5}$$

$$\alpha_S \sim 10^{-1} \tag{D.6}$$

$$sin\theta_{Cabibbo} = 0.221 \pm .002 \tag{D.7}$$

$$\delta m_K \sim 3.52 \times 10^{-15} GeV \tag{D.8}$$

$$|\bar{\epsilon}_K| \sim (2.24 \pm 0.02) \times 10^{-3} \tag{D.9}$$

$$B_{SL} = 0.117 \pm 0.006 \tag{D.10}$$

$$\tau_B = (1.11 \pm 0.16) \times 10^{-12} sec \tag{D.11}$$

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