The supersymmetric flavour and CP problems can be avoided if the first two generations of sfermions are heavier than a few TeV and approximately degenerate in mass. However using flavour and CP-violating constraints on the third sfermion generation, together with the decoupling of the first two generations, can dramatically affect cosmological predictions such as the relic abundance of stable particles. In particular, we show that if the lightest supersymmetric particle is essentially bino-like then requiring that all flavour changing neutral current and CP-violating processes are adequately suppressed, imposes severe limits on the bino mass, where typically \( m_{\tilde{B}} > (200 - 300) \) GeV. This leads to difficulties for models implementing the scenario of heavy sfermion masses.

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1. Supersymmetry (SUSY) is usually invoked to solve many of the puzzles of the Standard Model such as the stability of the weak scale under radiative corrections. Furthermore, local supersymmetry provides a promising way to include gravity within the framework of unified theories of particle physics. For such reasons, supersymmetric extensions of the Standard Model have been the focus of intense theoretical activity in recent years [1].

Since experimental observations require supersymmetry to be broken, it is essential to have a knowledge of the nature and the scale of supersymmetry breaking in order to have a complete understanding of the physical implications of these theories. Unfortunately, at the moment we lack such an understanding and therefore it is important to focus on the several experimental hints which might be useful in exploring the nature of supersymmetry breaking.

The Minimal Supersymmetric Standard Model (MSSM) (or extensions of it) are characterized by the presence of new degrees of freedom, the scalar partners of the fermions (sfermions), which carry flavour number and therefore can generate potentially large contributions to the Flavour Changing Neutral Currents (FCNC’s) [2]. Moreover, new CP-violating parameters may appear in the low energy effective theory where SUSY is softly broken [3]. The requirement of consistency with the experimental data imposes strong constraints on the physics of flavour and CP violation in SUSY theories and has a profound impact on supersymmetric model building. Although the flavour changing elements in the sfermion mass matrices as well as the CP violating phases are free parameters in the MSSM, ultimately their values have to be obtained from a theory of soft supersymmetry breaking and fermion mass generation. Therefore, experimental constraints provide us with useful suggestions towards such a theory.

There are broad classes of solutions which solve the supersymmetric flavour problem and the supersymmetric CP problem. The first possibility is that for some deep theoretical reasons the pattern of the sfermion mass matrices at the weak scale is very special: they are either very close to the unity matrix in flavour space (flavour universality) [4] or they have a structure, but they are diagonal in the basis set by the quark mass matrix (alignment) [5]. Under these special conditions, the FCNC effects are tiny and the CP violating phases at the weak scale are either highly suppressed.
or efficiently screened \[3\]. Furthermore, if high degeneracy of the first two sfermion
generations occurs, their masses are bounded from below only by the present direct
searches.

The second and, \textit{a priori}, the most straightforward possibility occurs when the
masses of the first and second generation of sfermions are larger than a few TeV \[8, 9\]
and much larger than the masses of sfermions of the third generation. In principle,
this inverse hierarchy (compared to fermion masses) could be a consequence of the
supersymmetry breaking pattern at the Planck or string scale \[3\]. Other possibilities
include integrating out heavy states which give rise to extra contributions to the soft
mass terms of light particles \[7\]. Notice, though, that the contribution to $\epsilon_K$
from the first two sfermion generations is generically still too large for CP violating phases
$\sim \mathcal{O}(1)$. However, this scenario becomes tenable when further approximate degeneracy
in the mass spectrum of the first two generations of squarks is present, such as in
models with non-abelian horizontal symmetries. Explicit realizations of this possibility
are presented in \[4, 10, 11\]. In this way, the suppression of FCNC effects in the MSSM is
achieved and the SUSY contributions to CP violating observables are small even for CP
violating phases of order unity. Note also that having the first and second generation
of sfermions heavy does not necessarily lead to naturalness problems, since the first
two generations are almost decoupled from the Higgs sector and, in the absence of
universality, the naturalness upper limits on supersymmetric particle masses increase
somewhat compared to the case when universality is assumed \[12\]. Still, even without
universality, the charginos and neutralinos are likely to be accessible at LEP2.\[1\]

However, models with the first two squark generations heavy may predict in the
neutral $B$ system sizeable shifts from the Standard Model predictions of CP asymmetries
in the decays to final CP eigenstates \[4\]. In general, the supersymmetric contributions
to FCNC’s and to the CP violating observables are expected to come from the third
generation of sfermions and they are typically close to the present experimental bound.
This means that, lower bounds on the masses of the third generation sfermion may still

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\[1\] In theories where the soft SUSY breaking parameters are generated at a high scale, large masses
for the sfermions of the first and second generation may drive the scalar top mass squared to negative
values at the weak scale because of the two-loop renormalization group evolution \[13\]. This, in turn,
puts a strong lower bound on the value of the scalar top mass squared at the high scale.
be imposed from phenomenological considerations in the scenario in which the first two squark generations are decoupled.

On the other hand, it is well known that by considering the cosmological relic density of stable particles one can impose significant bounds on the parameter space of a given model. In the MSSM with R-parity conservation, the lightest supersymmetric particle (LSP), is absolutely stable and its contribution to the relic abundance $\Omega_{\text{LSP}} h^2$ in the Universe [14] may be inconsistent with the bound $\Omega_{\text{LSP}} h^2 \sim 1$ implied by a (very conservative) lower bound of at least 10 billion years on the age of the Universe. The relic abundance of the LSP is determined by its annihilation cross section, which depends sensitively upon the masses of the various particles mediating the annihilation processes. For instance, in the case when the LSP is a bino-like neutralino, which we denote by $\chi$, large sfermion masses are typically inconsistent with the cosmological bound $\Omega_{\chi} h^2 \lesssim 1$, unless the annihilation rate of the LSP into scalar and gauge bosons is efficient enough and/or near resonances. It is therefore reasonable to expect that combining the experimental bounds on FCNC and CP violating phenomena with the bounds coming from cosmological considerations will help us in significantly constraining the parameter space of the MSSM.

In the present paper we will assume that the solution to the supersymmetric flavour problem and the supersymmetric CP problem is provided by the second class discussed above, namely by the scenario where the first and second generation sfermion masses are in the $O(10)$ TeV range and approximately degenerate in mass. We will show that when parameters are chosen so that the LSP is predominantly a bino, the requirement $\Omega_{\chi} h^2 \lesssim 1$ often places a severe lower bound on the LSP mass. This result may have rich implications for the class of supersymmetric models which explain the suppression of the FCNC and CP violating effects by decoupling the first two generations of sfermions.

2. Before beginning the discussion of the cosmological bounds, let us briefly discuss what kind of limits one can infer from the FCNC and CP violating effects on the masses of the third sfermion generation. We will generically assume that the third generation sfermions are lighter than a TeV. While bounds on the stops are fairly weak, larger effects arise for the sbottom and stau. The stringest bound that one can obtain on the sbottom mass follows from the $\epsilon_K$ parameter of $K^0 - \bar{K}^0$ mixing. In the limit that
\( m_{\tilde{b}} \equiv m_{\tilde{b}_L} \simeq m_{\tilde{b}_R} \) the bound resulting from the \( \epsilon_K \) parameter is \([13, 16]\)

\[
\left( \frac{1 \text{ TeV}}{m_{\tilde{b}}} \right)^2 |V_{13}^Q V_{23}^Q V_{13}^D V_{23}^D| \sin \varphi_1 f(m_{\tilde{g}}^2/m_{\tilde{b}}^2) \lesssim 3.24 \times 10^{-5} \tag{1}
\]

where \( V^{Q,D} \) are flavour mixing matrices (that define the rotations which diagonalise the quark mass matrix in the basis where \( m_{\tilde{Q}, \tilde{D}} \) are diagonal), \( \varphi_1 = \text{Arg}(V_{13}^Q V_{23}^Q V_{13}^D V_{23}^D) \) is a CP-violating phase and \( f(x) \approx 3840 x f_6(x) - 204 \tilde{f}_6(x) \). The functions \( f_6(x) \) and \( \tilde{f}_6(x) \) are defined as \([13]\)

\[
f_6(x) = \frac{1}{6(1-x)^5}(-18x \ln x - 6 \ln x - x^3 + 9x^2 + 9x - 17) \tag{2}
\]
\[
\tilde{f}_6(x) = \frac{1}{3(1-x)^5}(-6x^2 \ln x - 6x \ln x + x^3 + 9x^2 - 9x - 1). \tag{3}
\]

Notice that the bound (1) depends on the particular details of the flavour mixing. Since we are considering models that do not have any special mechanisms for the flavour and CP-structure, we will generically assume the CP-phase to be maximal with \( \sin \varphi_1 \sim 1 \).

In order to understand how the magnitude of the off-diagonal matrix elements affects the bound we will compare our results with a CKM-like parameterisation of the mixing matrices of the form

\[
V^{Q,D} = \begin{pmatrix}
1 & \lambda & \lambda^3 \\
\lambda & 1 & \lambda^2 \\
\lambda^3 & \lambda^2 & 1
\end{pmatrix} \tag{4}
\]

where \( \lambda \sim 0.2 \) is a Cabibbo-like angle. The bound (1) is very sensitive to the amount of mixing between the first two and third generations \([4]\). For arbitrary parameterisations of the mixing matrix we will present our results by defining an average off-diagonal element \( \overline{V}_1 \equiv \left| V_{13}^Q V_{23}^Q V_{13}^D V_{23}^D \right|^{1/4} / (0.2)^{5/2} \), where \( \overline{V}_1 = 1 \) corresponds to the CKM parameterisation \([4]\). For the special limit \( m_{\tilde{b}} \simeq m_{\tilde{g}} \) the sbottom mass bound arising from (1) is

\[
m_{\tilde{b}} \gtrsim 800 \overline{V}_1^2 \text{ GeV.} \tag{5}
\]

Clearly, the bound becomes weak when the amount of flavour mixing \( \overline{V}_1 \to 0 \). This is

\[\text{(Notice also that since we are assuming CP violating phases} \sim \mathcal{O}(1) \text{ the contribution to} \epsilon_K \text{ from the first two generations is much too large (even for large squark masses). Therefore as previously mentioned in the Introduction, one requires some approximate universality to further suppress these contributions. In particular, if the first and second generation squark masses are degenerate up to} \mathcal{O}(\lambda^2), \text{ then these contributions will be sufficiently suppressed)\]
Figure 1: (a) Lower bounds on the sbottom mass for various contours of $V_1$ (solid line) and $\bar{V}_2$ (dashed line) where $A'_b = 10$ TeV and $\sin \varphi_{1,2} \sim 1$. (b) Lower bounds on the stau mass for various contours of $V_3$ (dashed line) and $\bar{V}_4$ (solid line) where $A'_\tau = 1$ TeV and $\sin \varphi_4 \sim 1$.

the case when there are special mechanisms operating such as universality or alignment. The general behaviour for arbitrary $m_{\tilde{b}}$ and $m_{\tilde{g}}$ can be seen in Fig. [1], where contours of the lower bound on the sbottom mass are shown for various values of $V_1$. In the Figure, $m_{\tilde{b}}$ is plotted as a function of the bino mass $m_{\tilde{B}}$ where at the electroweak scale $m_{\tilde{g}} \simeq 7m_{\tilde{B}}$, which follows from our assumption of gaugino mass unification.

One can see that, for values of $V_1 \gtrsim 1$, the lower bounds on either the mass of the sbottom or the gluino is quite significant, in the range of hundreds of GeV. Notice, however, that the lower mass bound from $K^0 - \bar{K}^0$ mixing disappears as the mass of the gluino or sbottom exchanged in the loop becomes very large. However, for large gluino mass a stronger lower bound can be obtained by considering the contribution of the sbottom left-right mixing to the down quark electric dipole moment (EDM) \cite{15, 16}. This contributes to the neutron EDM and gives rise to a bound

$$
\frac{e\alpha_3 m_b}{6\pi m_{\tilde{b}}^2} |A'_b| \left| V_{13}^Q V_{13}^D \right| \sin \varphi_2 m_{\tilde{g}} g(m_{\tilde{g}}^2/m_{\tilde{b}}^2) \lesssim 8.25 \times 10^{-26} \text{e cm}
$$

(6)

where $\varphi_2 = \text{Arg}(V_{13}^Q V_{13}^D A'_b)$, $A'_b = (A_b + \mu \tan \beta)$ and

$$
g(x) = \frac{1}{(1-x)^4} (2x^2 \ln x + 4x \ln x - 5x^2 + 4x + 1).
$$

(7)
The bound (6) is again sensitive to the flavour structure, which we will parameterise by $V_2 \equiv \left| V_{13}^{Q} V_{13}^{D} \right|^{1/2} /(0.2)^3$. However, unlike the bound arising from $K^0 - \bar{K}^0$ mixing, the bound (6) is also sensitive to the amount of left-right mixing in the sbottom mass matrix. Thus besides the phases from the flavour mixing there is also a CP-phase from the left-right mixing. (The effect of a CP-phase in $A'_b$ was previously considered in Ref. [17] where it was shown that the limits on the LSP mass can be relaxed by a factor of 2-3.) Note that even in the absence of a CP-phase for $A'_b$, the bound (6) still applies provided there remain nontrivial phases in the matrix elements $V_{Q,D}^{13}$. Again, without considering any special mechanism for the CP-phases, we will assume that the overall CP-phase to be maximal ($\sin \varphi_2 \sim 1$). In the special limit $m_{\tilde{b}} > \sim 410 (0.2)^3 \mathrm{GeV}$.\footnote{\label{eq:8}}

\begin{equation}
m_{\tilde{b}} \gtrsim 410 V_2^{2/3} \left( \frac{|A'_b|}{1 \ \mathrm{TeV}} \right)^{1/3} \ \mathrm{GeV}.
\end{equation}

It is clear that strong constraints on the sbottom mass can only be obtained for large $A'_b$. This can be seen in Fig. 1 where $V_2$ is plotted for $A'_b = 10 \ \mathrm{TeV}$.

Similar bounds can also be obtained for the stau and these follow from the flavour-violating process $\mu \rightarrow e\gamma$ and the electron EDM \cite{15, 16}. Again we will assume that $m_{\tilde{\tau}} \equiv m_{\tilde{\tau}_L} \simeq m_{\tilde{\tau}_R}$. The bound following from $\mu \rightarrow e\gamma$ is

\begin{equation}
\left( \frac{100 \ \mathrm{GeV}}{m_{\tilde{\tau}}} \right)^2 V_{13}^{L} V_{23}^{E} \frac{m_{\tilde{\tau}}}{m_{\tilde{\tau}}} \frac{|A'_\tau|}{|A'_b|} g(m^2_{\tilde{\tau}}/m^2_{\tilde{\tau}}) \lesssim 5 \times 10^{-4}
\end{equation}

where $A'_\tau = (A_\tau + \mu \tan \beta)$ and $V^{L,E}_{13}$ are the slepton mixing matrices. Again for comparison we will assume a CKM-like parameterisation of the matrices $V^{L,E}$ which define the rotations that diagonalise the lepton mass matrix in the basis where the slepton mass matrices are diagonal. Thus $V_3 \equiv (V_{13}^{L} V_{23}^{E})^{1/2} /(0.2)^{5/2}$ defines an average off-diagonal matrix element normalised to a CKM-like parameterisation. Notice also that there is no CP-phase since the process is CP-invariant. The typical size of the bound on the stau mass can be obtained from considering the limit $m_{\tilde{\tau}} \simeq m_{\tilde{B}}$, where

\begin{equation}
m_{\tilde{\tau}} \gtrsim 260 V_3^{2/3} \left( \frac{|A'_\tau|}{1 \ \mathrm{TeV}} \right)^{1/3} \ \mathrm{GeV}.
\end{equation}

Despite the insensitivity to CP-phases, the bounds arising from (6) only become strong for $|A'_\tau| \gg 1 \ \mathrm{TeV}$ or $V_3 \gg 1$, as can be ascertained from (10) and Fig. 1.
Much stronger constraints on the stau mass can be obtained from the electron EDM. The bound resulting from the electron EDM is [13, 16]

$$\frac{e\alpha_1 m_\tau}{2\pi m_\tau^4} |A'_\tau| \left| V_{13}^L V_{13}^E \right| \sin \varphi_4 m_\tau g(m_B^2/m_\tau^2) \lesssim 7 \times 10^{-27} \text{e cm} \quad (11)$$

where \( \varphi_4 = \text{Arg}(V_{13}^L V_{13}^E A'_\tau) \). Again assuming the CP-phase to be maximal (\( \sin \varphi_4 \sim 1 \)) and defining \( \overline{V}_4 = \left| V_{13}^L V_{13}^E \right|^{1/2}/(0.2)^3 \), lower bounds on the stau mass can be obtained for large left-right mixing (\( A'_\tau \)) in the stau mass matrix. In particular for the limit \( m_\tau \sim m_B \) we see that

$$m_\tau \gtrsim 750 \overline{V}_4^{2/3} \left( \frac{|A'_\tau|}{1 \text{ TeV}} \right)^{1/3} \text{ GeV.} \quad (12)$$

In Fig. 1, the stau mass bound is shown for contours of \( \overline{V}_4 \) and \( A'_\tau = 1 \text{ TeV} \). Unlike \( \overline{V}_3 \), the bounds arising from \( \overline{V}_4 \) are always much stronger for the same value of \( |A'_\tau| \) and \( \overline{V}_3 \approx \overline{V}_4 \), using the current experimental bounds.

In all the above bounds we have made the degenerate squark mass assumption of \( m_\tilde{b} \equiv m_{\tilde{b}_L} \simeq m_{\tilde{b}_R} \) at the electroweak scale, and similarly for the stau. If this assumption is relaxed then the bounds shown in the Figures are for the geometric mean \( \sqrt{m_{\tilde{b}_L} m_{\tilde{b}_R}} \) up to factors of \( \mathcal{O}(1) \) which follow from generalising the functions \( f(x) \) and \( g(x) \). It is also clear that if there are any special mechanisms operating in the flavour structure, such that \( \overline{V}_i \to 0 \), then all FCNC and CP-violating bounds disappear. However, we will be specifically interested in the cosmological consequences of the case where the first two sfermion generations are heavy and generically \( \overline{V}_i \sim \mathcal{O}(1) \).

3. Let us now consider the cosmological implications on the bino mass from the stringent lower bounds on the mass of the third generation sfermions resulting from the FCNC and CP-violating processes. We will be particularly interested in the cosmological relic abundance of the LSP when it is a neutralino which is predominantly bino-like, with only a small admixture of the wino and the higgsino in its composition. While in principle any superpartner could be the LSP, in the MSSM the neutralino is usually assumed to be the LSP for astrophysical reasons: it is a weakly-interacting stable massive particle for which astrophysical bounds are very weak and it can serve as an excellent dark matter candidate [18] when it is mostly a bino [19]. (Note that a higgsino-like
neutralino with a sufficiently large $\Omega_\chi h^2$ and a reasonably small mass has now been basically excluded by LEP-II, except for a small remaining region [20].)

A predominantly bino-like LSP corresponds to the case $|\mu| \gtrsim M_1$ where $M_1$ is the soft-mass of the bino. It is worth noting that a bino-like neutralino naturally arises as the only neutral LSP as a result of requiring radiative electroweak symmetry breaking (EWSB). While this has been shown to be true mainly in the case of universal soft masses at the unification scale [21], there are good reasons to believe that this will also remain valid in the case studied here. This is because $M_1$ depends on sfermion masses only at two loops, while the parameter $\mu^2$ is determined via the condition for EWSB where the sfermion masses of the first two generations enter only as a small correction [12], and are not expected to significantly alter the resulting value of $\mu$ compared to the universal case.

In order for a bino-like neutralino to give $\Omega_\chi h^2 \sim 1$, at least some sfermion masses should normally not exceed a few hundred GeV [19]. In our numerical analysis we will include all relevant final states of the neutralino annihilation and all exchange channels for the general case of any neutralino composition. However, in the nearly pure bino limit the dominant annihilation channel is into final state (ordinary) charged fermions via the (lightest) sfermion exchange and the relic abundance is approximately given by $\Omega_\chi h^2 \propto m_\tilde{f}^4/m_\chi^2$ where $m_\tilde{f}$ is the sfermion mass. Thus it is clear that for sufficiently large sfermion masses imposing the bound $\Omega_\chi h^2 < 1$ will imply a lower bound on $m_\chi$, unless other final-state channels can reduce the LSP relic abundance below one. In the pure bino limit, the annihilation cross-section into final states involving one or both gauge bosons vanishes. The final states involving the pseudoscalar $A$ and either $h$ or $H$, may be able to reduce $\Omega_\chi h^2$ below one, but they are not kinematically allowed until $m_\chi \gtrsim (m_A + m_h)/2$. This implies a rather large $m_\chi$ if $A$ is heavy.

The neutralino relic density is also reduced in the vicinity of resonances due to the exchange of the $Z$ and the Higgs bosons. Again, while the pure bino does not couple to the gauge or Higgs bosons, the small higgsino component in the nearly pure bino case allows the resonances to play some rôle in decreasing $\Omega_\chi h^2$. Of special importance is the exchange of the pseudoscalar $A$ whose coupling to the neutralino is proportional to $\tan \beta$ and therefore can become significantly enhanced, especially for larger values of
Let us now combine the stringent limits on the masses of the third sfermion generation arising from the suppression of the FCNC and CP-violating processes with the cosmological constraint \( \Omega_\chi h^2 \lesssim 1 \) for a predominantly bino-like LSP. We will consider three representative cases: \( m_\tilde{b} = m_\tilde{t} \) with \( m_\tilde{\tau} \) heavy in Fig. 2, \( m_\tilde{\tau} = m_\tilde{\tau} \) with \( m_\tilde{b} \) heavy in Fig. 3 and \( m_\tilde{b} = m_\tilde{\tau} = m_\tilde{\tau} \) in Figs. 4 and 5. In each case we have used the best possible constraint arising from FCNC and CP-violating processes. For the sbottom mass this corresponds to the \( \epsilon_K \) parameter, parameterised by contours of \( V_1 \), while for the stau mass the electron EDM parameterised by \( V_4 \) provides the most stringent constraint. The cosmological contour \( \Omega_\chi h^2 = 1 \) is shown for several choices of \( \mu \). Thus regions above and to the left of the cosmological contour are excluded.

In each Figure we see that as \( |\mu| \) decreases, the higgsino component of the neutralino increases, and consequently the two-boson (both gauge and Higgs) final states become important. This is especially true for the \( ZZ \) and \( WW \) final states which open up for relatively low \( m_\chi \) but decouple in the pure bino limit. Since we focus on a nearly pure bino as the LSP, we do not consider values of \( |\mu| \) smaller than 500 GeV in order for the bino purity (defined as the square of the bino component in the neutralino mass eigenvector) to remain above 97%. For large \( m_A \) and \( |\mu| \) of the order of 1 TeV and for small bino masses, below \( m_t \), only the tau and bottom final states are effectively open and \( \Omega_\chi h^2 \) quickly increases with the mass of their scalar partners, thus either leaving no room for \( m_\chi < m_t \) in Fig. 2 or allowing only for a relatively narrow strip below \( m_t \) in Fig. 3. When \( m_\chi > m_t \), the \( tt \) channel opens up and is enhanced by the factor \((m_t/m_W)^2\) via the higgsino component of the LSP. As the third generation sfermion masses increase further, this channel also becomes less and less effective. Finally, for \( m_\chi \) approaching \( m_A/2 \) the wide pseudoscalar resonance starts dominating quickly reducing \( \Omega_\chi h^2 \) well below one.

The combination of the exclusion curves from flavour and CP violating processes and from \( \Omega_\chi h^2 < 1 \) gives therefore strong lower limits on \( m_\chi \). The limits are particularly strong for large values of \( |\mu| \) and \( m_A \). For example, in Fig. 2 we see that for \( |\mu| \gtrsim 1000 \text{ GeV} \) the bino has to be heavier than roughly \( m_t \) even for \( V_1 = 1 \). This should be compared with the indicative upper bounds \( m_\chi \lesssim 65 \text{ GeV} \), obtained by requiring
Figure 2: Bounds on the sbottom mass as a function of the bino mass. The $\mathbf{V}_1$ contours arise from the $\epsilon_K$ parameter of $K^0 - \bar{K}^0$ mixing (regions below them are excluded). The cosmological contours $\Omega_\chi h^2 = 1$ are labelled by various values of the $\mu$ parameter. (Regions to the left and above them are excluded.) In the Figure we have assumed $m_{\tilde{t}} = m_{\tilde{b}}$, $\tan \beta = 2$ and $m_A = m_{\tilde{\tau}} = 1$ TeV.

no significant fine-tuning in the parameters of the MSSM \cite{12}. Actually, since the motivation for this scenario is to allow for basically unconstrained entries in the mixing matrices, one would expect $\mathbf{V}_1$ significantly larger than one, in which case the lower limit on $m_\chi$ would be further significantly increased.

A similar picture emerges when one considers the bounds on the stau mass arising from the electron EDM. Since the bounds on $m_{\tilde{\tau}}$ from $\mathbf{V}_4$ are more stringent than $\mathbf{V}_1$ we obtain a stronger lower limit on the bino mass. For the case plotted in Fig. 4 we find $m_\chi \gtrsim 300$ GeV for $\mathbf{V}_4 = 1$ and $|\mu| \gtrsim 1000$ GeV. Finally in Figs. 4 and 5 the sbottom and stau are now both assumed to be light and we need to simultaneously satisfy the
Figure 3: Bounds on the stau mass as a function of the bino mass. The $\mathcal{V}_4$ contours arise from the electron EDM (regions below them are excluded). The cosmological contours $\Omega_\chi h^2 = 1$ are labelled by various values of the $\mu$ parameter. (Regions to the left and above them are excluded.) In the Figure we have assumed $m_{\tilde{t}} = m_{\tilde{\tau}}$, $\tan \beta = 2$ and $m_A = m_{\tilde{b}} = A'_\tau = 1$ TeV.

There are a number of ways one can relax the bounds on $m_\chi$. This can be done by either decreasing $|\mu|$ (thus increasing the higgsino component of the LSP) or by lowering $m_A$ as can be seen in Fig. 4. In this case the lower limit on $m_\chi$ reduces to $\sim 200$ GeV for $|\mu| \gtrsim 1000$ GeV and is fairly independent of the value of $\mathcal{V}_1$ and $\mathcal{V}_4$. Another possibility is to increase $\tan \beta$ in which case the resonance effect around $m_A/2$ widens considerably. Finally, while the sbottom and stau are constrained by flavour and
Figure 4: Bounds on the stau and sbottom mass as a function of the bino mass. The contours $V_1$ are for the sbottom mass, while $V_A$ constrains the stau mass. The cosmological contours $\Omega_\chi h^2 = 1$ are labelled by various values of the $\mu$ parameter. In the Figure we have assumed $m_{\tilde{t}} = m_{\tilde{\tau}} = m_{\tilde{b}}$, $\tan\beta = 2$ and $m_A = A'_\tau = 1$ TeV.

CP constraints, there are no constraints on the stop. One can therefore choose to make the stop lighter than the sbottom and the stau. This can still only allow for $m_\chi$ above $m_t$ which is already a very strong lower bound. On the other hand, we have found that for $\mu < 0$ the bounds are even more stringent.

4. We have shown that by combining the constraints arising from the suppression of FCNC and CP-violating processes with bounds on the cosmological relic abundance, the bino mass can be severely restricted. This places severe limitations on models in which the first two sfermion generations are heavy and almost degenerate in mass and the supersymmetric contributions to the FCNC's and CP violating observables mainly
come from the third squark generation.

Such a mass spectrum has been argued to be the best from the phenomenological point of view [22] and may be obtained if the three families belong to a $2 + 1$ representation of a horizontal symmetry group $G_H$. For example, a class of models based on the group $U(2)$ predicts very heavy first and second family scalars, the CKM parameterisation [12] of the mixing matrices, i.e. $\overline{V}_i = 1$, and CP-violating phases of order unity [9]. It has also been recently pointed out that $D$-term contributions from the anomalous $U(1)$ gauge group in string theory may naturally lead to such a mass spectrum for the sfermions. On the other hand, a generic problem of this class of models is the generation of sizeable gaugino masses. In this paper we have pointed out that having the first two generations of sfermions heavy and approximately degenerate requires driving the mass of the bino-like LSP to quite large values when considerations about the present cosmo-
logical abundance of the LSP are taken into account. This leads to serious difficulties for models implementing the scenario of heavy sfermion masses.

Our constraints can be avoided in a number of ways. First, if there is a small amount of R-parity violation then the LSP can simply decay and therefore be eliminated. In this case other solutions to the dark matter problem need to be considered. Secondly, one can envisage models where the mixing between the third generation and the first two generations of sfermions is small. For example, in certain three generation string solutions, the anomalous \( U(1) \) couples universally to all three families \[23\], yielding squark degeneracy. It is also possible to increase the higgsino or wino content of the neutralino LSP, but as we have mentioned earlier this scenario may not be very natural from the point of view of mass unification.

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