Phenomenological implications of moduli-dominant SUSY breaking

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ABSTRACT

We study moduli-dominated SUSY breaking within the framework of string models. This type of SUSY breaking in general leads to non-universal soft masses, \textit{i.e.} soft scalar masses and gaugino masses. Further gauginos are lighter than sfermions. This non-universality has phenomenologically important implications. We investigate radiative electroweak symmetry breaking in the mass spectrum derived from moduli-dominated SUSY breaking, where the lightest chargino and neutralino are almost gauginos. Moreover, constraints from the branching ratio of $b \rightarrow s\gamma$ and the relic abundance of the LSP are also considered. The mass spectrum of moduli-dominated SUSY breaking is favorable to the experimental bound of the $b \rightarrow s\gamma$ decay decreasing its branching ratio. We obtain an upper bound for the gravitino mass from the cosmological constraint.
1 Introduction

Supersymmetry (SUSY) is one of the most important keywords beyond the standard model [1]. It is expected that superpartners as well as Higgs particles would be detected in the near future. Thus it is very important to study which type of mass spectra for superpartners and Higgs particles and which type of phenomenological aspects are predicted from fundamental theory such as grand unified theories (GUTs) or superstring theory.

Superstring theory is a promising candidate for unified theory including gravity. Hence it is interesting to predict the mass spectrum and other SUSY phenomenological aspects like flavor changing processes, electric dipole moment and rare decay as well as cosmological aspects within the framework of superstring theory. We have not understood SUSY breaking mechanism completely yet. However we can parameterize unknown $F$-terms and write soft SUSY breaking terms when we assume the fields that contribute to SUSY breaking [2, 3]. Generic string models include a dilaton field $S$ and moduli fields $T_m$ in their massless spectra. Vacuum expectation values (VEVs) of these fields correspond to the coupling constant and geometrical feature of a six-dimensional compactified space. When these fields contribute to SUSY breaking, soft SUSY breaking terms can be written in terms of the gravitino mass $m_{3/2}$ and some goldstino angles under the assumption of the vanishing cosmological constant [4, 5].

Recently there have been various works devoted to study phenomenological implications of these soft SUSY breaking terms which are written in terms of the gravitino mass and goldstino angles only. The dilaton-dominated SUSY breaking leads to universal soft SUSY breaking terms. Phenomenological implications of this universality were studied in Refs. [6, 7], e.g. successful electroweak symmetry breaking and the $b \to s\gamma$ process. We have non-universal soft SUSY breaking terms when both moduli fields and the dilaton field contribute to SUSY breaking. However, such non-universality is not strong. Thus such a mixed case leads to almost similar phenomenological aspects of the dilaton-dominated case. Actually some phenomenological aspects of the mixed case were studied in Refs. [8, 9], which include the case of the overall modulus $T$ with the modular weights chosen such that one can have appropriate large string threshold corrections to fit the joining of gauge coupling constants at GUT scale [4].

On the other hand, the moduli-dominated case can lead to strong non-universality among soft SUSY breaking terms [1, 3]. The moduli-dominated case with the overall moduli field leads to non-universal gaugino masses, but a universal soft scalar mass [1]. Further

\footnote{See also Ref. [1].}
multi-moduli cases can lead to non-universal soft scalar masses as well as non-universal gaugino masses \[5\]. It is important to study phenomenological aspects of this string inspired non-universality in order to understand the whole allowed parameter space of soft SUSY breaking parameters derived from superstring theory.\[5\] In Ref.\[11\] non-universal gaugino masses were discussed within the overall moduli-dominated case, showing interesting results.

In this paper we study the multi-moduli dominated SUSY breaking, especially phenomenological implications of strong non-universality among soft scalar masses as well as gaugino masses. Orbifold models are simple and interesting 4-dimensional string models \[12\] and these models include three diagonal moduli fields $T_m$ ($m = 1, 2, 3$). Here we discuss phenomenological implications of strong non-universality led in the case that these three diagonal moduli fields contribute dominantly to SUSY breaking. We consider mass spectrum, successful electroweak symmetry breaking, the $b \to s\gamma$ process and cosmological constraints, e.g., the relic abundance of the lightest neutralino. We find that there appears a significant splitting between the stop masses and this leads to interesting phenomenological implications. Moreover we show that the lightest stop is the lightest sfermion of this model.

Study on the $b \to s\gamma$ process gives a strong constraint on generic supersymmetric models \[13\], because, in general, there are additive contributions to the standard model (SM) one. On the other hand, by including the next-to-leading order in QCD correction for $b \to s\gamma$ decay, the SM prediction is above the CLEO measurement at the $1\sigma$ level \[14\]. Hence, it is likely that SUSY models provide us with a substantial destructive interference. We show that in moduli-dominated SUSY breaking the chargino contribution which gives a destructive interference is dominant and the values of the total branching ratio for all the parameter space are less than the SM branching ratio. This result corresponds to the case of $\mu > 0$ which gives strong mass splitting between the stop masses according to our sign convention.

We also show that the lightest neutralino is the lightest superpartner (LSP) and it is found to be dominantly bino. In contrast with the case of the overall modulus analyzed in Ref.\[11\] the LSP mass is not degenerate with the chargino mass. The relic abundance of the LSP imposes important constraints on the gravitino mass $m_{3/2}$ and requires a large value of $\theta_1$ (one of goldstino angles) \ie requires strong non-universality between the down

\[2\] Phenomenological implications of non-universal soft SUSY breaking terms were discussed in Ref.\[10\] within the framework of generic SUSY GUT and supergravity theory, but not superstring theory.
and up sector Higgs soft masses.

This paper is organized as follows. In section 2, we review formulae of soft SUSY breaking terms within the framework of string models. In addition, we specify our model with typical non-universal soft masses, which are derived from string models. In section 3, we study the radiative electroweak breaking and we determine the particle spectrum of the model for both cases of $\mu > 0$ and $\mu < 0$. Section 4 is devoted to the constraints on the parameter space due to $b \to s\gamma$ decay. We show that in the case of $\mu > 0$ the CLEO measurement does not impose any constraint on the parameter space. In section 5, we study the relic abundance of the lightest neutralino and we show that it imposes an upper bound of order 1.5 TeV on the gravitino mass and requires that $\theta_1$ be large. In section 6, we analyze the effect of changing the value of the parameter $T$, namely we study the implication of the small values of $T$ as well as the large value of $T$. Finally we give our conclusions in section 7.

2 Moduli-dominated SUSY breaking

In orbifold models with three diagonal moduli fields $T_m$ ($m = 1, 2, 3$) the Kähler potential is obtained as

$$ K = -\log(S + S^*) - \sum_{m=1}^{3} \log(T_m + T_m^*) + \sum_i \phi_i \phi_i^* \prod_{m=1}^{3} (T_m + T_m^*)^{n_i^m}, $$

where $n_i^m$ is the modular weight of the field $\phi_i$ for the $m$-th moduli field $T_m$. In addition, the gauge kinetic function $f_a$ is obtained at tree level as $f_a = k_a S$, where $a$ is an index for the gauge group and $k_a$ is the corresponding Kac-Moody level.

Here we assume that $S$ and $T_m$ contribute to SUSY breaking with a nonperturbative superpotential $W(S, T_m)$, although we do not specify its form. Then we parameterize unknown $F$-terms as follows [4, 5],

$$ (K^S_S)^{1/2} F^S = \sqrt{3} m_{3/2} \sin \theta, \quad (K^{T_m}_{T_m})^{1/2} F^{T_m} = \sqrt{3} m_{3/2} \cos \theta \Theta_m, $$

where $\sum_{m=1}^{3} \Theta_m^2 = 1$. Here gravitino mass $m_{3/2}$ is defined as $m_{3/2} = \langle e^{K/2}/W \rangle$. Using the above parametrization with

$$ \Theta_1 = \sin \theta_1 \sin \theta_2, \quad \Theta_2 = \sin \theta_1 \cos \theta_2, \quad \Theta_3 = \cos \theta_1 $$

we can write the soft scalar mass $m_i$ and the gaugino mass $M_a$ as

$$ m_i^2 = m_{3/2}^2 (1 + 3 \cos^2 \theta \sum_m n_i^m \Theta_m^2), $$

$$ M_a^2 = k_a m_{3/2} \cos \theta_1 \Theta_m. $$
\[ M_a = \frac{\sqrt{3} m_{3/2}^3}{Re f_a} \left[ k_a R e S \sin \theta + \cos \theta \sum_m (b^m_a - k_a \delta^m_{GS}) D(T_m, T^*_m) \Theta_m \right], \]  

where the second term on the right-hand side of eq.(4) is due to moduli-dependent threshold corrections at one-loop level \[ 16, 17 \]. Here the function \( D(T) \) is given by the use of the Eisenstein function \( \hat{G}(T) \) as \[ 3 \]

\[ D(T) = \frac{(T + T^*)}{32 \pi^3} \hat{G}(T). \]  

For example, the values of \( D(T) \) are \( 1.5 \times 10^{-3}, 2.7 \times 10^{-2}, 6.0 \times 10^{-2} \) and \( 6.6 \times 10^{-1} \) for \( T = 1.2, 5.0, 10 \) and 100, respectively. In addition \( \delta^m_{GS} \) is the Green-Schwarz coefficient \[ 19 \], which is gauge group-independent, and \( b^m_a \) denotes a duality anomaly coefficient given by \[ 17 \]

\[ b^m_a = -C(G_a) + \sum_R T(R) (1 + 2 n^m_a), \]  

where \( C(G_a) \) is the casimir of the adjoint representation and \( T(R) \) is the index of the \( R \) representation. Further we obtain the \( A \)-term associated to the \( T \)-independent Yukawa coupling as

\[ A_{ijk} = -\sqrt{3} m_{3/2} \left[ \sin \theta + \cos \theta \sum_m (1 + n^i_m + n^j_m + n^k_m) \Theta_m \right], \]  

where \( n^i_m, n^j_m \) and \( n^k_m \) are modular weights of the fields to couple. One needs a correction term in eq.(6) when the corresponding Yukawa coupling depends on moduli fields. However, \( T \)-dependent Yukawa coupling in general includes suppression factors \[ 20 \]. Thus strong Yukawa coupling such as the top Yukawa coupling is expected to be independent of \( T \). The above formula (7) seems to be reasonable.

Finally, we have to consider the scalar bilinear soft breaking term \( B_\mu H_1 H_2 \), i.e., the \( B \)-term, where \( H_1 \) and \( H_2 \) are the down and up sectors of Higgs fields, respectively. The form of the \( B \)-term strictly depends on the origin of the \( \mu \)-term in the superpotential and/or the Kähler potential. In Ref. \[ 4 \] three sources for the \( B \) parameter were considered, labeled by \( B_Z \), \( B_\mu \) and \( B_\lambda \). The source of \( B_Z \) is the presence of certain bilinear terms in the Kähler potential which can naturally induce a \( \mu \)-term of order \( m_{3/2} \) after SUSY breaking \[ 21 \]. An alternative mechanism to generate a \( B \)-term in the scalar potential is to assume that the superpotential \( W \) includes a SUSY mass term \( \mu(S, T) H_1 H_2 \) induced by a non-perturbative effect, then a \( B \)-term is automatically generated and it is called \( B_\mu \). Also it was pointed out \[ 22 \] that the presence of a non-renormalizable term

\[ \text{Several kinds of modular functions are shown in Ref.} \[ 18 \]. \]
in the superpotential $\lambda WH_1 H_2$ yields dynamically a $\mu$ parameter when $W$ acquires VEV $\mu(S, T_i) = \lambda(T_i) W(S, T_i)$ and the corresponding $B$-term is denoted by $B_\lambda$. In general, there could be an admixture of the above three possibilities. Thus we will treat $B$ as a free parameter whose value can be determined from the electroweak breaking conditions.

Here we discuss the minimal supersymmetric standard model within the framework of string models, where formulae for soft SUSY breaking parameters are given in eqs. (3-7)\(^4\). In addition we take $k_3 = 1$, $k_2 = 1$ and $k_1 = 5/3$ for $SU(3)$, $SU(2)$ and $U(1)$ among the standard model gauge group. In the dilaton-dominated SUSY breaking case, i.e. $\sin \theta \to 1$ universal soft SUSY breaking parameters are obtained as

$$m_i^2 = m_{3/2}^2, \quad M_a = \sqrt{3} m_{3/2}, \quad A_{ijk} = -\sqrt{3} m_{3/2}. \quad (8)$$

Their phenomenological implications have been studied. On the other hand, soft SUSY breaking parameters are obtained in the moduli-dominated SUSY breaking case, i.e. $\cos \theta \to 1$ as

$$m_i^2 = m_{3/2}^2(1 + 3 \sum_m n_i^m \Theta_m^2), \quad (9)$$

$$M_a = \sqrt{3} m_{3/2} \sum_m \left( \frac{b_{\mu}^m}{k_a} - \delta_{GS}^m \right) D(T_m, T_m^*) \Theta_m, \quad (10)$$

$$A_{ijk} = -\sqrt{3} m_{3/2} \sum_m \left( 1 + n_i^m + n_j^m + n_k^m \right) \Theta_m. \quad (11)$$

These soft mass formulae include $n_i^m$ and $b_{\mu}^m / k_a$, which depend on the fields and the gauge group. Thus these soft masses are in general non-universal. Therefore this strong non-universality of soft masses is an important feature of the moduli-dominated SUSY breaking. Further, $D(T)$ works as a suppressed factor for a small value of $T$, while soft scalar masses seem to be of $O(m_{3/2})$ naturally. Hence soft scalar masses are larger than gaugino masses for such value of $T$.

Here we have to take into account the $S - T$ mixing, that is, at one-loop level the dilaton field and moduli fields are mixed in the Kähler potential \[^{17}\]. In such a case, the formulae of soft masses are obtained by replacing $\cos^2 \theta \Theta_m^2$ in eq.(3) as

$$\cos^2 \theta \Theta_m^2 \to \frac{\cos^2 \theta \Theta_m^2}{1 - a_m}, \quad (12)$$

where $a_m = \delta_{GS}^m / 24\pi^2 Y$ and $Y$ is approximately written as $Y = S + S^*$. This parameter $a_m$ is estimated as $a_m \approx 10^{-2}$ when $\delta_{GS}^m = O(1)$ and $S = 2$, which corresponds to the unified gauge coupling $\alpha_X \approx 1/25$. Thus this parameter $a_m$ is negligible in most of the

\[^{4}\]Here the word “minimal” means only the minimal matter content.
parameter space of goldstino angles. However, its effect is not negligible in the purely overall moduli-dominated SUSY breaking case, which means $\cos^2 \theta = 1$ and $\Theta_m^2 = 1/3$ ($m = 1, 2, 3$) exactly. In such case the fields with the overall modular weight $\sum m n_i = -1$ have the suppressed soft scalar mass as $m_i^2 = a m_{3/2}^2$, where we take $a_m = a$ ($m = 1, 2, 3$). From the viewpoint of multi-moduli dominated SUSY breaking the parameter region of the goldstino angles leading to the suppressed soft scalar mass $m_i^2 = a m_{3/2}^2$ is rather narrow. However, in the case of these suppressed soft masses, successful electroweak breaking without color and/or charge breaking (CCB), in general requires the suppressed $A$-term with the same order of magnitude [23, 24]. Such a situation, i.e. the suppressed soft scalar masses and the suppressed $A$-term, is effectively similar to the case with large gaugino masses and non-suppressed soft scalar masses. Thus we do not consider here the case of the suppressed soft scalar masses. Instead we will discuss the effect of a large $T$ after considering the case of $T \sim O(1)$. That also means effectively the case with the suppressed soft scalar masses.

Our purpose is to study implications of moduli dominated SUSY breaking, i.e. phenomenological aspects of non-universal soft SUSY breaking terms. Thus we consider here the case leading to typically strong non-universality of soft masses. In particular non-universality between the down and up sector Higgs soft masses, $m_{H_1}$ and $m_{H_2}$, is interesting. Obviously we can obtain strong non-universality between soft scalar masses in the case that two fields have modular weights corresponding to different moduli fields. Hence we assume

$$n_{H_1} = (-1, 0, 0), \quad n_{H_2} = (0, -1, 0). \quad (13)$$

Even in the case with more than three moduli fields, similar assignments can lead to maximum non-universality, e.g. $m_{H_1}^2 - m_{H_2}^2 \approx m_{3/2}^2$. Hence this case is a good example to see what happens in generic non-universal cases derived from string models with several moduli fields. If non-universality is not strong enough, its phenomenology is similar to the universal case. In this case we find that

$$m_{H_1}^2 = m_{3/2}^2 (1 - 3 \sin^2 \theta_1 \sin^2 \theta_2), \quad (14)$$

$$m_{H_2}^2 = m_{3/2}^2 (1 - 3 \sin^2 \theta_1 \cos^2 \theta_2). \quad (15)$$

The inequality $m_{H_1}^2 \geq m_{H_2}^2$ is favorable to realize successful electroweak symmetry breaking. Thus we take here $\theta_2 = 0$ i.e. $\Theta_1 = 0$ fixing

$$m_{H_1}^2 = m_{3/2}^2. \quad (16)$$
In this case we obtain the soft scalar mass of $H_2$ as

$$m_{H_2}^2 = m_{3/2}^2 \left( 1 - 3 \sin^2 \theta_1 \right). \quad (17)$$

In the case of $\sin \theta_1 = 0$ the universal soft mass for $m_{H_1}$ and $m_{H_2}$ is obtained, while there appears strong non-universality around $\sin^2 \theta_1 \sim 1/3$. The soft mass of $H_2$ could, in principle, have a negative mass squared i.e. $m_{H_2}^2 < 0$ with a small magnitude at high energy scale, i.e. $\sin^2 \theta_1 \geq 1/3$ in a small region. However, in such a case one needs a fine-tuning for other parameters. Thus we restrict ourselves to the case of $\sin^2 \theta_1 \leq 1/3$. As will be seen, we obtain similar results around $\sin^2 \theta \approx 1/3$. Hence we can expect similar results in the case that $\sin^2 \theta_1$ exceeds a little bit $1/3$.

We need modular weights of the other fields to obtain the $A$-term and $b'^m_a$ in the gaugino mass. For simplicity we assume that the three families of quark and lepton fields have $n = (-1, 0, 0)$. For such a case, we are able to calculate $b'^m_a$ as

$$b'^3_3 = (-9, 3, 3), \quad b'^2_2 = (-8, 4, 5), \quad b'^1_i = (-10, 10, 11). \quad (18)$$

Then gaugino masses are obtained as

$$M_1 = \frac{\sqrt{3m_{3/2}}}{\text{Re} S} [(6 - \delta_{GS}) \sin \theta_1 + (33/5 - \delta_{GS}) \cos \theta_1] D(T),$$

$$M_2 = \frac{\sqrt{3m_{3/2}}}{\text{Re} S} [(4 - \delta_{GS}) \sin \theta_1 + (5 - \delta_{GS}) \cos \theta_1] D(T), \quad (19)$$

$$M_3 = \frac{\sqrt{3m_{3/2}}}{\text{Re} S} [(3 - \delta_{GS}) \sin \theta_1 + (3 - \delta_{GS}) \cos \theta_1] D(T).$$

In these equations we have assumed $T_m = T$ and $\delta_{GS}^m = \delta_{GS}$ for simplicity. Further the $A$-term is written as

$$A_1 = -\sqrt{3m_{3/2}} \cos \theta_1. \quad (20)$$

In this case soft scalar masses and $A$-terms are of $O(m_{3/2})$. On the other hand, gaugino masses include $D(T)$, which give a suppression factor except a large value of $T$, i.e., $T > O(100)$. Thus gauginos are lighter than squarks and sleptons and the lightest neutralino and chargino are dominated by gauginos. Which gaugino is lightest depends on $\delta_{GS}$. It seems to be natural that the magnitude of $\delta_{GS}$ is of $O(b'^m_a)$, i.e. $O(1)$. We have the following ratios of gaugino masses at the string scale,

$$M_3 : M_2 : M_1 = (3 - \delta_{GS}) t_1 : (1 + (4 - \delta_{GS}) t_1) : (\frac{3}{5} + (6 - \delta_{GS}) t_1) \quad (21)$$

where $t_1 = \tan \theta_1 + 1$. When $\sin^2 \theta_1 < 1/3$, $t_1$ takes a value from 1.0 to 1.7. At the weak scale we approximately have

$$M_3 : M_2 : M_1 = 7(3 - \delta_{GS}) t_1 : 2(1 + (4 - \delta_{GS}) t_1) : (\frac{3}{5} + (6 - \delta_{GS}) t_1), \quad (22)$$
because $\alpha_1 : \alpha_2 : \alpha_3 \simeq 1 : 2 : 7$ at $M_Z$. Thus $|M_3(M_Z)|$ is larger than $|M_1(M_Z)|$ except in the case with

$$\frac{15}{6} - \frac{1}{10t_1} < \delta_{GS} < \frac{27}{8} + \frac{3}{40t_1}. \quad (23)$$

Since $t_1$ is around 1.4, the above region corresponds to $2.4 < \delta_{GS} < 3.4$. This region is narrow. Similarly $|M_3(M_Z)|$ is larger than $|M_2(M_Z)|$ except in the case of $2.8 < \delta_{GS} < 3.0$. This is also very narrow. Therefore the gluino is heavier than the other gaugino masses except in the very narrow region. From a phenomenological viewpoint this region is ruled out when $D(T)$ is not large because in this region the gluino is the LSP. Further $|M_2(M_Z)|$ is larger than $|M_1(M_Z)|$ unless $2.9 < \delta_{GS} < 5.3$. Thus we have $M_3(M_Z) > M_2(M_Z) > M_1(M_Z)$ in most parameter space of $\delta_{GS}$ of $O(1)$, but we can obtain $M_3(M_Z) > M_1(M_Z) > M_2(M_Z)$ in a small region. This type of mass spectrum is derived from generic models with other values of $b'_m$, that is, generally we have $M_3(M_Z) > M_2(M_Z) > M_1(M_Z)$ in most parameter space and $M_3(M_Z) > M_1(M_Z) > M_2(M_Z)$ in a small region, but in a very narrow region $M_3(M_Z)$ is the smallest. If the wino is the lightest, both of the lightest chargino and neutralino are almost wino and these masses are degenerate [11].

![Figure 1](image.png)

Figure 1. The running values for $M_i$ with $\delta_{GS} = -5$, $m_{3/2} = 500$ GeV and $\sin \theta_1 = \sqrt{\frac{1}{3}}$.

In the following sections we discuss the case with a typical value of $\delta_{GS}$ leading to $M_3(M_Z) > M_2(M_Z) > M_1(M_Z)$. Fig.1 shows the running values for $M_i$ with $m_{3/2} = 500$ GeV and $\delta_{GS} = -5$, which is the preferred value of $\delta_{GS}$ in the orbifold models.

We impose our initial conditions (16-20) at the string scale, which differs from the gauge coupling unification scale $M_X$. However, radiative corrections between these scales
induce only small changes in the following discussions. Further the difference between these two scales could be explained by a moduli-dependent threshold correction with a certain value of $T$.

3 Radiative electroweak breaking and the particle spectrum

Given the boundary conditions in eqs. (16-20) at the compactification scale, we assume that $T \sim O(1)$. Later we will comment for the cases of large $T$ of order 100 and small $T$ as well. Here we consider $D(T) = 2.7 \times 10^{-2}$. We determine the evolution of the couplings and mass parameters according to their one loop renormalization group equations (RGE) in order to estimate the mass spectrum of the SUSY particles at the weak scale. The radiative electroweak symmetry breaking scenario imposes the following conditions on the renormalized quantities:

$$m^2_{H_1} + m^2_{H_2} + 2\mu^2 > 2B\mu, \quad (24)$$

$$\left(m^2_{H_1} + \mu^2\right)\left(m^2_{H_2} + \mu^2\right) < (B\mu)^2, \quad (25)$$

and

$$\mu^2 = \frac{m^2_{H_1} - m^2_{H_2} \tan^2 \beta}{\tan^2 \beta - 1} - \frac{M_Z^2}{2}, \quad (26)$$

$$\sin 2\beta = \frac{-2B\mu}{m^2_{H_1} + m^2_{H_2} + 2\mu^2}, \quad (27)$$

where $\tan \beta = \langle H^0_2 \rangle / \langle H^0_1 \rangle$ is the ratio of the two Higgs VEVs that give masses to the up and down type quarks and $m_{H_1}, m_{H_2}$ are the two soft Higgs masses at the electroweak scale. It was pointed out by Gamberini et al. [25] that the tree-level effective potential $V_0$ and the corresponding tree VEVs are strongly $Q$-dependent and the one-loop radiative correction to $V_0$, namely

$$\Delta V_1 = \frac{1}{64\pi^2} \sum_{\alpha} (-1)^{2s_{\alpha}} (2s_{\alpha} + 1) C_i M^4_{\alpha} \ln \left( \frac{M^2_{\alpha}}{Q^2} \right) - \frac{3}{2}, \quad (28)$$

are crucial to make the potential stable against variations of the $Q$-scale. $M^2_{\alpha}(Q)$ are the tree level mass eigenvalues, $s_{\alpha}$ is the spin of the corresponding particle and $C_i$ is the color degree of freedom.

Using eqs. (26) and (27) we can determine $\mu$ and $B$ in terms of $m_{3/2} \theta_1$ and $\tan \beta$. We take here $\tan \beta = 2$, retaining only the top Yukawa coupling. The value of $|\mu|$ as a function of the gravitino mass $m_{3/2}$ and the goldstino angle $\theta_1$ is given in Fig.2. For the
fixed value of $m_{3/2}$ the variation of $|\mu|$ in this figure corresponds to different values of $\theta_1$. In the same way, all the figures are plotted corresponding to different values of $\theta_1$.

Figure 2: The values of $|\mu|$ versus $m_{3/2}$ with $\tan \beta = 2$.

Fig. 3 shows the ratio of the coefficient of the bilinear term $B$ at the compactification scale to $m_{3/2}$ versus the gravitino mass. We note that the sign of $B$ in general is opposite to that of $\mu$ for the realization of electroweak symmetry breaking. It is also remarkable that the value of $B/m_{3/2}$ is very stable against $m_{3/2}$, i.e. $B/m_{3/2} \sim 0.35(-1.6)$ for $\mu > 0$($\mu < 0$). That could suggest the effectiveness of a certain type of $\mu$-term generation mechanism.
Figure 3: The values of $B/m_{3/2}$ versus $m_{3/2}$ with $\tan \beta = 2$.

An important constraint on the parameter space arises from the experimental lower bound on the chargino mass from LEP II, $m_{\chi^\pm} > 84$ GeV. This bound is applied as long as $m_{\chi^\pm} - m_\chi > 3$ GeV which is always satisfied in this model. Here $m_\chi$ is a neutralino mass. As we can see from Figs. 4 and 5 this bound implies that the gravitino mass has to be $m_{3/2} > 280$ GeV for $\mu > 0$ and $m_{3/2} > 420$ GeV for $\mu < 0$.

![Figure 4](image1.png)

Figure 4: The lightest chargino mass versus $m_{3/2}$ with $\mu < 0$.

![Figure 5](image2.png)

Figure 5: The lightest chargino mass of versus $m_{3/2}$ with $\mu > 0$. 
It is clear that at these values of the gravitino mass most of the scalar particles become very heavy. For instance, the right selectron mass is of order 300 GeV at this lower bound of the gravitino mass in this model, while the right selectron was found to be the lightest sfermion in the case of dilaton contribution to SUSY breaking [8]. Moreover, the off diagonal element of the stop mass matrix \( m_t (A_t + \mu \cot \beta) \) is comparable to the diagonal parts of this matrix. This gives a chance to have one of the stops to be light and Fig.6 shows the values of this stop mass versus the gravitino mass for \( \mu > 0 \) since for this case we have maximum mixing.

![Figure 6: The mass of the lightest stop quark versus \( m_{3/2} \) with \( \mu > 0 \).](image)

Actually this light stop is predicted to be the lightest sfermion in this class of models and as we will show below it has important effects on the phenomenological implication such as the prediction of the branching ratio of the \( b \to s \gamma \) decay and the relic abundance of the LSP. We would like to stress that this feature, having a significant splitting between the masses of the stop quarks, is absent in the overall modulus SUSY breaking scenario and dilaton contribution to SUSY breaking scenario.

Now we would like to investigate the composition and the mass of the lightest neutralino. The lightest neutralino \( \chi \) is a linear combination of two neutral gauginos \( \tilde{B}^0 \) (bino) and \( \tilde{W}^0_3 \) (wino) and the Higgsinos \( \tilde{H}^0_1, \tilde{H}^0_2 \),

\[
\chi = N_{11} \tilde{B}^0 + N_{12} \tilde{W}^0_3 + N_{13} \tilde{H}^0_1 + N_{14} \tilde{H}^0_2,
\]

where \( N_{ij} \) are the entries of the unitary matrix which diagonalizes the neutralino mass matrix and they are functions of \( \tan \beta, M_2 \) and \( \mu \). The gaugino 'purity' function
\( f_g = |N_{11}|^2 + |N_{12}|^2 \) describes the neutralino composition. As explained, in most of the parameter space the wino is heavier than the bino. Thus we find that the lightest neutralino is mostly bino \( e.g., f_g = 0.99 \) for \( \delta_{GS} = -5 \).

The lightest neutralino mass corresponding to the gravitino mass is plotted in Figs.7 and 8 for \( \mu > 0 \) and \( \mu < 0 \) respectively. From these figures we can easily realize that the lightest neutralino is indeed the LSP and this has an important cosmological implication as we will see in section 5.

Figure 7: The lightest neutralino as a function of the gravitino mass for \( \mu > 0 \).
Figure 8: The lightest neutralino as a function of the gravitino mass for $\mu < 0$.

Also we are interested in the SUSY Higgs spectrum, and in particular the lightest Higgs scalar $h$ and charged Higgs scalar $H^+$ whose mass is very relevant for $b \to s\gamma$ branching ratio. The lightest Higgs mass is given by

$$m_h^2 = m_{h,0}^2 + (\Delta m_h^2)_{1LL} + (\Delta m_h^2)_{\text{mix}},$$

$$m_{h,0}^2 = \frac{1}{2} \left( m_A^2 + m_Z^2 - \sqrt{(m_A^2 + m_Z^2)^2 - 4m_Z^2m_A^2\cos^2 2\beta} \right),$$

where $m_A^2 = m_{H_1}^2 + m_{H_2}^2 + 2\mu^2$,

$$\frac{3m_t^4}{4\pi^2v^2} \ln\left(\frac{m_t}{m_t}\right) \left[ 1 + O\left(\frac{m_W^2}{m_t^2}\right) \right],$$

$$\frac{3m_t^4}{8\pi^2v^2} \left[ 2h(m_{t_1}^2, m_{t_2}^2) + \tilde{A}_t^2 f(m_{t_1}^2, m_{t_2}^2) \right] \left[ 1 + O\left(\frac{m_W^2}{m_t^2}\right) \right],$$

with $\tilde{A}_t = A_t + \mu \cot \beta$. The functions $h$ and $f$ are given by

$$h(a, b) = \frac{1}{a - b} \ln\left(\frac{a}{b}\right) \quad \text{and} \quad f(a, b) = \frac{1}{(a - b)^2} \left[ 2 - \frac{a + b}{a - b} \ln\left(\frac{a}{b}\right) \right].$$

The two loop leading logarithmic contributions to $m_h^2$ are incorporated by replacing $m_t$ in eq.(29) by the running top quark mass evaluated at the scale $\mu_t$ which is given by $\mu_t = \sqrt{M_t M_s}$ where $M_t$ is the pole mass of the top quark, $M_t = 174$ GeV, and $M_s = \sqrt{M_t^2 + M_s^2}$.

Fig.9 shows the lightest Higgs mass as a function of the gravitino mass with $\mu > 0$ where we have a maximum mixing.
Figure 9: The lightest Higgs mass as a function of the gravitino mass.

It shows that the lower bound of this mass is of the order 70 GeV. The charged Higgs mass is given by

\[ m_{H^\pm}^2 = m_W^2 + m_{H_1}^2 + m_{H_2}^2 + 2\mu^2. \]

Fig.10 gives the charged Higgs mass versus the gravitino mass.

![Figure 10: The charged Higgs mass versus the gravitino mass.](image)

It is interesting to note that in this class of models the charged Higgs field is very heavy. This is due to the fact that \( m_{H_1}^2 \) is much larger (and positive) than \( m_{H_2}^2 \) (which is negative at the weak scale), and \( \mu \) is quite large. As we will see, the evidence of having heavy charged Higgs field gives an interesting implication in studying the constraints on the parameter space due to the \( b \to s\gamma \). Also we have to stress that if \( m_{H_1}^2 \) is of the order \( m_{H_2}^2 \) as in the case of the overall modulus and dilaton SUSY breaking scenarios, we get a cancellation between them and the charged Higgs field becomes much lighter. So we can conclude that the strong non-universality between the Higgs soft masses at GUT scale is preferable.

### 4 Constraints from \( b \to s\gamma \)

In this section we focus on the constraints on the parameter space \((m_{3/2}, \theta_1)\) which arise from the \( b \to s\gamma \) decay since the CLEO observation \cite{27} confirmed that \( 1 \times 10^{-4} < BR(b \to s\gamma) < 4 \times 10^{-4} \). It is well known that in supersymmetric models there are
three significant contributions of total amplitude from the $W$-loop, charged Higgs loop and chargino loop. The inclusive branching ratio for $b \to s\gamma$ is given by

$$R = \frac{BR(b \to s\gamma)}{BR(b \to c\ell\nu)}.$$  

The computation of $R$ yields \[13\]

$$R = \frac{|V_{ts}^* V_{tb}|^2 2\alpha_{em} \pi}{I(x_{cb})[1 - \frac{2}{3\pi}\alpha_s(m_b)f(x_{cb})]} \left[ \eta^{16/23} A_\gamma + \frac{8}{3}(\eta^{14/23} - \eta^{16/23}) A_g + C \right]^2.$$  

Here, $\eta$ is the ratio of the running strong-coupling constants at two energy scales $m_W$ and $m_b$, i.e. $\eta = \frac{\alpha_s(m_W)}{\alpha_s(m_b)}$. $C$ represents the leading-order QCD corrections to $b \to s\gamma$ amplitude at the $Q = m_b$ scale \[28\]. The function $I(x)$ is given by

$$I(x) = 1 - 8x^2 + 8x^6 - x^8 - 24x^4 \ln x,$$

and $x_{cb} = \frac{m_c}{m_b}$, while $f(x)$ is a QCD correction factor $f(x_{cb}) = 2.41$. The amplitude $A_\gamma$ is from the photon penguin vertex, the amplitude $A_g$ is from the gluon penguin vertex and they are given in Ref. \[13\]. It was shown that in MSSM \[13\], pure dilaton SUSY breaking \[4\] and minimal string unification \[8\] with $\tan \beta$ of order 2 the chargino contribution gives rise to a destructive interference (in case of $\mu < 0$) with SM contribution and charged Higgs contribution but it is generally smaller than the charged Higgs contribution. This leads to a severe constraint on the parameter space of these models. Also it was realized that the constraint is less severe in the case of the non-universality between the soft terms than that for the universal case. However, in minimal string unification the non-universality is very tiny since the allowed values of goldstino angle $\theta$ is very close to $\pi/2$ which corresponds to purely dilaton-dominated SUSY breaking, namely $\theta \in [0.98\text{rad.}, 2\text{rad.}]$. This constraint on $\theta$ arose from the avoidance of the tachionic mass at string scale and the conservation of the electric charge.

Before we present the result of $b \to s\gamma$ in the three moduli dominated SUSY breaking, we find that it is worthwhile to study first this constraint on the case of the overall modulus since it was omitted in Ref. \[11\]. To be able to compare the result of the overall modulus with that of $b \to s\gamma$ in our case where $\delta_{GS} = -5$ we have to use the corresponding value of $\delta_{GS}$ in the overall modulus scenario which is $\delta_{GS} = -15$. We find that in this case the chargino mass reaches 84 GeV at very large values of the gravitino mass $m_{3/2}$: about 400 GeV for $\mu > 0$ and 700 GeV for $\mu < 0$. Figs.11 and 12 show the branching ratio of $b \to s\gamma$ versus the gravitino mass in the case of the overall modulus dominated SUSY breaking for $\mu > 0$ and $\mu < 0$ respectively.
Figure 11: The branching ratio of $b \to s\gamma$ as a function of $m_{3/2}$ for $\mu > 0$.

Figure 12: The branching ratio of $b \to s\gamma$ as a function of $m_{3/2}$ for $\mu < 0$.

From these figures it is evident that the CLEO upper bound imposes severe constraints on the allowed parameter space of the overall modulus dominated SUSY breaking. This behaviour is independent of the choice of the value of $\delta_{GS}$. In the case of $\mu > 0$ the value of the branching ratio of $b \to s\gamma$ falls outside the experimentally allowed region for all the values of $m_{3/2}$ up to 1 TeV. On the other hand, for $\mu < 0$, where we have $m_{3/2} > 700$ GeV, we find that the branching ratio of $b \to s\gamma$ becomes less than the CLEO upper bound at $m_{3/2} > 1.4$ TeV. It is clear that at this value of $m_{3/2}$ the charged Higgs field
becomes very heavy and its contribution becomes small.

Now we turn to our model. We find that the chargino contribution gives rise to substantial destructive interference with SM and $H^+$ amplitude. At $\tan \beta = 2$ Figs.13 and 14 show the $b \to s\gamma$ branching ratio for $\mu > 0$ and $\mu < 0$ respectively.

![Figure 13](image13.png)

**Figure 13:** The branching ratios of $b \to s\gamma$ versus $m_{3/2}$ and $\mu > 0$.

![Figure 14](image14.png)

**Figure 14:** The branching ratios of $b \to s\gamma$ versus $m_{3/2}$ and $\mu < 0$.

From these figures we find that the branching ratios of $b \to s\gamma$ in this model for $\mu > 0$ are less than the SM value and we conclude that there is no essential constraint from
\( b \rightarrow s\gamma \) imposed on the parameter space. We would like to emphasize the reasons of the winning of the chargino contribution. First, as we have mentioned and as Fig.10 shows, the mass of the charged Higgs field is quite heavy because of the strong non-universality between Higgs masses \( m^2_{H_1} \) and \( m^2_{H_2} \). For instance, if the chargino mass is equal to 84 GeV, the charged Higgs mass is around 400 GeV, so that the charged Higgs contribution which is inversely proportional to its mass square becomes quite small. Second, in this model we have a significant splitting between the values of the stop masses, as we mentioned due to the large mixing in the stop mass matrix. In fact, the chargino amplitude crucially depends on this splitting.

This result is quite interesting since, as it was pointed out in \[14\], the SM prediction is above the CLEO measurement at the 1\( \sigma \) level. The physics beyond the SM should provide a destructive interference with the SM amplitude and our model has this feature with \( \mu > 0 \).

5 Relic abundance of the lightest neutralino

We have shown that the lightest neutralino turns out to be the LSP and it is mostly pure bino. So it could provide a natural source for the dark matter required by galactic rotational data. In this section we would like to study the relic abundance of the LSP and investigate the constraints on the parameter space by requiring the neutralino relic density to be \( 0.1 \leq \Omega_{LSP} \leq 0.9 \) with \( 0.4 \leq h \leq 0.8 \). It was shown in Ref. \[8\] that these values of \( \Omega_{LSP} \) impose a stringent upper bound on the parameter space in the minimal string unification namely it leads to an upper bound on the gravitino mass of about 600 GeV.

Since the LSP is mostly pure bino, its coupling to the lightest Higgs field and Z boson is weak and the sfermions are very heavy, many channels of the neutralino annihilation are closed or suppressed. We find that the annihilation process is dominated by the exchange of the lightest stop into up type quarks \( i.e. \) only \( u \) and \( c \). The \( t \)-channel can not be opened because the neutralino mass is smaller than the mass of the top quark. As we explained the stop is the lightest sfermion for this model. The Z-boson contribution is suppressed, except for \( m_\chi \sim m_Z/2 \) due to the small \( Z\chi\chi \) coupling \( (ig/2\cos\theta_W)(N_{13}^2 - N_{14}^2)^{\mu\gamma_5} \).

In the overall modulus dominated SUSY breaking, both stop masses are quite large. The \( \chi\chi \) annihilation is very small and this leads to a very large relic density of order \( 10^2 \) \[11\], which is, of course, an unacceptable value. In Ref.\[11\] the co-annihilation between
the LSP and the chargino was considered to reduce these values of the relic density. However, Fig.17 in this reference shows that for \( \tan \beta = 2 \) a small part of the parameter space can lead to \( \Omega_{LSP} h^2 \leq 1 \), moreover this part corresponds to chargino mass less than 88 GeV.

For the computation of the lightest neutralino relic abundance we have to determine the thermally averaged cross section \( \langle \sigma_A v \rangle \sim a + bv \) \[29\]. By neglecting the fermion masses with respect to the LSP mass we find that \( a = 0 \), and \( b \) is given in Ref.\[29\]. Given \( a \) and \( b \) we can determine the relic LSP density \( \Omega_\chi h^2 \):

\[
\Omega_\chi h^2 = \frac{\rho_\chi}{\rho_c/h^2} = 2.82 \times 10^8 Y_\infty (m_\chi/\text{GeV}), \tag{36}
\]

where

\[
Y^{-1}_\infty = 0.264 \ g_s^{1/2} \ M_P \ m_\chi \left( \frac{a}{x_F} + \frac{3b}{x_F^2} \right), \tag{37}
\]

and \( h \) is the Hubble parameter. We take \( 0.4 \leq h \leq 0.8 \), and \( \rho_c \sim 2 \times 10^{-29} h^2 \) is the critical density of the universe. In addition the freeze-out epoch \( x_F \) is written as \( x_F = m_\chi/T_F \) where \( T_F \) is the freeze-out temperature and the \( \chi \chi \) annihilation rate is smaller than the expansion rate of the universe below \( T_F \). We can iteratively compute the freeze-out temperature from

\[
x_F = \ln \left( \frac{0.0764 M_P (a + 6b/x_F) c (2 + c) m_\chi}{\sqrt{g_* x_F}} \right). \tag{38}
\]

Here \( M_P = 1.22 \times 10^{19} \text{ GeV} \) is the Planck mass and \( g_* \ (8 \leq \sqrt{g_*} \leq 10) \) is the effective number of relativistic degrees of freedom at \( T_F \).

![Figure 15](image.png)

Figure 15: The neutralino relic abundance versus \( m_{3/2} \) where \( \tan \beta = 2, \sin \theta_1 \sim \sqrt{\frac{1}{3}} \) and \( \mu > 0 \). The solid line corresponds to the maximum value of \( \Omega_\chi h^2 \) we assumed and the dotted line corresponds to the minimum value.
Fig. 15 shows the relic abundance of the lightest neutralino $\Omega LSP h^2$ as a function of the gravitino mass where the goldstino angle $\theta_1$ is equal to $\sin^{-1} \sqrt{\frac{1}{3}} \sim 0.6 \text{ rad}$. We find that the requirement of the neutralino relic density to be $0.1 \leq \Omega_{LSP} \leq 0.9$ and $0.4 \leq h \leq 0.8$ imposes an upper bound 1.5 GeV on the gravitino mass. Also we find that the neutralino relic density imposes a severe constraint on $\theta_1$, namely $\theta_1 \geq 0.4 \text{ rad}$. For smaller values of $\theta_1$ we obtain $\Omega LSP h^2 > 1$. This means that the neutralino relic abundance of the moduli SUSY breaking does not prefer the universality between $m_{H_1}$ and $m_{H_2}$.

6 Effects of other values of $T$ and $n_{i_k}^j$

In this section we analyze effects of other values of $T$ and $n_{i_k}^j$, because we have left them aside in the previous sections. They could lead to an interesting implication or there could be some constraints on these values. Obviously we can obtain similar phenomenological results in the case of $T$ of $O(1)$ except the very small $T$ case like $T = 1.2$ and $D(T) = 1.5 \times 10^{-3}$. Similar results are obtained even in the case with $T = 10$, 20 or some small value of $O(10)$. Furthermore, the results in the previous sections do not depend on a value of $\delta_{GS} \sim O(1)$ sensitively. We obtain almost the same results for other values of $\delta_{GS}$ leading to $M_3(M_Z) > M_2(M_Z) > M_1(M_Z)$.

The case of small $T$: Gaugino masses become much smaller compared with sfermion masses and $A$-term as well as the gravitino mass. The lightest chargino and neutralino become almost purely gauginos. So the experimental lower bound of the chargino mass pushes the gravitino mass to a higher value. For example, the case of $T = 1.2$ requires the gravitino mass above 4 TeV. Then the mass of the scalar particles are very heavy and this is an undesirable feature for the hierarchy problem and naturalness. Moreover in this case most of the lightest neutralino annihilation channels are suppressed or closed so that the relic abundance $\Omega LSP h^2$ is very large.

The case of large $T$: Gaugino masses become comparable to sfermion masses and $A$-term as well as the gravitino mass, e.g. in the case of $T \sim O(100)$. Note that this value of $T$ has the effective meaning including the case of $T \sim O(1)$ and the suppressed $m_i$ due to the $S-T$ mixing as said in section 2. The phenomenological prediction of this model becomes rather similar to that of the universal model. The gaugino purity function becomes smaller. In this case the experimental lower bound of the chargino mass does not always require a larger gravitino mass, that is, $O(100)$ GeV is enough. RGE effects due to gaugino masses become important. For example, non-universality of scalar
masses at the string scale is diluted at the weak scale because of RGE effects of large gaugino masses. If we get very small non-universality between two Higgs masses, fine-tuning should be required to realize successful electroweak symmetry breaking. For this type of mass spectrum, the charged Higgs field also contributes to the branching ratio of $b \to s\gamma$. In such a case we can not expect the same results as those in the case with $T = 5$.

**Other values of $n_k^i$:** We have taken $n = (-1, 0, 0)$ for all the matter fields in the analyses of the previous sections. In a similar way we can analyze the case where matter fields have other modular weights. In such case, the parameter space of the gravitino mass and the goldstino angles can have further constraints, e.g. experimental lower bounds of sfermion masses, degeneracy of sfermion masses for flavor changing processes and the constraint to avoid $m_i^2 < 0$ or CCB.

In addition, the anomaly coefficients $b'_a$ are changed when we alter the assignment of modular weights. However, the structure of the gaugino mass spectrum, i.e. $M_3(M_Z) > M_2(M_Z) > M_1(M_Z)$, is not sensitive to $b'_a$ of $O(1)$. Thus we can obtain similar results for the gaugino-Higgs sector in the case with other assignments of modular weights.

### 7 Conclusions

We have studied phenomenological implications of moduli-dominated SUSY breaking. In general, moduli-dominated SUSY breaking leads to non-universal soft scalar masses as well as non-universal gaugino masses. In addition, gauginos are lighter than most sfermions. This type of mass spectrum is very different from the one derived from dilaton-dominated SUSY breaking and the mixed dilaton/moduli breaking case as well as ordinary “minimal” supergravity with universal soft breaking terms. Thus moduli-dominated SUSY breaking leads to phenomenological aspects different from other types of SUSY breaking.

Non-universality between two Higgs masses is favorable to realize successful radiative electroweak symmetry breaking. In moduli-dominated SUSY breaking the lightest chargino and neutralino are almost gauginos and the latter is usually regarded as the LSP. Further there appears a mass splitting between the lightest Higgs field and the other Higgs fields, e.g., the charged Higgs field. Also one of the stop fields is very light compared with the other stop as well as other squarks and sleptons. This type of mass spectrum makes the branching ratio of $b \to s\gamma$ decrease, while the overall moduli case is ruled out in a wide parameter space. Furthermore, strong non-universality between $m_{H_1}$ and $m_{H_2}$ is favorable for the constraint from the relic abundance of the LSP. Hence moduli-dominated SUSY
breaking is the very interesting case in the whole parameter space of goldstino angles of string models.

When we take a very small value of $T$, a mass splitting between gauginos and others becomes large. Thus the lower bound of the chargino mass requires large sfermion masses. In this case the cosmological constraint from the relic abundance of the LSP plays a role to rule out this parameter space. On the other hand, gaugino masses are comparable to sfermion masses in the case of $T \sim O(100)$. In this case large gaugino masses dilute non-universality among soft scalar masses because of RGE effects. Thus we obtain phenomenological aspects similar to the universal case. In such case the charged Higgs field also contributes to the $b \rightarrow s\gamma$ decay increasing its branching ratio.

If gauge symmetries break reducing their ranks, there appears another type of contribution to soft scalar masses, i.e. $D$-term contributions [30, 31]. These $D$-contributions can also become sources of non-universality among soft scalar masses. Thus analyses including these $D$-term contributions would be interesting [32, 33].

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