No-Scale Scenario with Non-Universal Gaugino Masses

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Abstract

Phenomenological issues of no-scale structure of Kähler potential are re-examined, which arises in various approaches to supersymmetry breaking. When no-scale boundary conditions are given at the Grand Unified scale and universal gaugino masses are postulated, a bino mass is quite degenerate with right-handed slepton masses and the requirement that the lightest superparticle (LSP) be neutral supplemented with slepton searches at LEP200 severely constrains allowed mass regions of superparticles. The situation drastically changes if one moderately relaxes the assumption of the universal gaugino masses. After reviewing some interesting scenarios where non-universal gaugino masses arise, we show that the non-universality diminishes the otherwise severe constraint on the superparticle masses, and leads a variety of superparticle mass spectra: in particular the LSP can be a wino-like neutralino, a higgsino-like neutralino, or even a sneutrino, and also left-handed sleptons can be lighter than right-handed ones.

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I. INTRODUCTION

One of the most important phenomenological issues in supersymmetric (SUSY) Standard Models (SSMs) is to identify the mechanisms of supersymmetry breaking in the hidden sector and its mediation to the SSM sector (observable sector). Soft supersymmetry breaking masses which arise in effective theories after integrating over the hidden sector are in fact constrained from various requirements. For instance, they should lie in the range of \(10^2\)-\(10^3\) GeV to solve the naturalness problem in the Higgs sector which is responsible for the electroweak symmetry breaking, and satisfy mass bounds given by collider experiments. They should also satisfy flavor-changing-neutral-current (FCNC) constraints as well. Furthermore, if the lightest superparticle (LSP) is stable, which is often the case, cosmological arguments require it be electrically neutral and \(SU(3)_c\) singlet.

The structure of the soft scalar masses is characterized by the Kähler potential. In this paper, we shall focus on a special class of the Kähler structure in which the hidden sector and the observable sector are separated from each other in the Kähler potential \(K\) as follows:

\[
e^{-\frac{K}{3}} = f_{hid}(z, z^*) + f_{obs}(\phi, \phi^*),
\]

where \(z\) and \(\phi\) symbolically represent fields in the hidden and observable sector, respectively. The first example which exhibits this form of the Kähler potential is a so-called no-scale model \([1]\), and thus we call it the no-scale structure. The characteristics of the Kähler potential in the no-scale form is that the soft SUSY breaking scalar masses vanish (as the vacuum energy vanishes) and gaugino masses are a dominant source of SUSY breaking mass. Of course, this mass pattern is given at the energy scale where the soft masses are given, and the renormalization group effects due to the non-vanishing gaugino masses raise the masses of the scalar superparticles at the weak scale.

The no-scale structure of the Kähler potential is obtained in many types of models. As we will see in the next section, such models include the (tree-level) Kähler potential of simple Calabi-Yau compactification of the heterotic string theories \([2]\) both in the weak- and
strong-coupling regimes, the splitting Ansatz of the hidden and observable sectors in the superspace density in a supergravity formalism [3], and the geometrical splitting of the two sectors in a brane scenario [4,5].

In this paper we revisit some phenomenological issues of the models with the no-scale boundary conditions. This class of models has closely been investigated in the literature. A particular attention was paid to the minimal case where the boundary conditions are given at the Grand Unified Theory (GUT) scale of $2 \times 10^{16}$ GeV and the gaugino masses are assumed to be universal at this energy scale. In this case the mass spectrum of superparticles is very constrained, and the bino mass is almost degenerated with those of the right-handed sleptons. In fact it was shown that the neutralino can be the LSP only when its mass is less than about 120 GeV [3,6,7]; otherwise the stau would be the LSP which is charged, and thus not allowed if it is stable. We will revise this result, emphasizing that the present experimental bounds already exclude the large $\tan \beta$ case, leaving only $\tan \beta \lesssim 8$.

One of the main points in this paper is that slight modifications of the minimal scenario will drastically change mass spectrum of the superparticles. In particular, we shall devote ourselves to the case where the gaugino masses are non-universal at the GUT scale. We will first review several cases that the non-universality of the gaugino masses result. Then we will discuss its phenomenological implications. Most remarkably relaxing the universality condition within a factor of two or so will result in a variety of mass spectra. In particular the LSP can be not only the bino-dominant neutralino, but also a wino or higgsino-dominant neutralino, or an admixture of the gaugino and the higgsino, or even a sneutrino. Furthermore the severe upperbound on the masses of the superparticles no longer exists. Thus we expect superparticle phenomenology in this case is much richer than the minimal case.

The paper is organized as follows. In the subsequent section, we review some examples which possess the no-scale Kähler potential. In section 3, we re-examine the case where the no-scale boundary conditions are given at the GUT scale and gaugino masses are universal at the scale, and show that the superparticle mass spectrum is very restrictive and tight constraints already exclude much of the parameter space. In section 4, we argue that the
very constrained mass spectrum can be relaxed by several ways, and then we focus on one of them, namely the case with non-universal gaugino masses. After recalling some mechanisms to realize the non-universality of the gaugino masses, we consider its phenomenological implications. The final section is devoted to conclusions.

II. NO-SCALE BOUNDARY CONDITIONS

In this section, we would like to review some models which have the no-scale Kähler potential. The first model is the no-scale model \[1\] with the Kähler potential

\[
K = -3 \ln (T + T^* - \phi^* \phi)
\]

and the superpotential

\[
W = W(\phi),
\]

where \(T\) is a hidden sector field responsible for the SUSY breaking and \(\phi\) is a generic matter field. Here and in the following, we use a unit that the reduced Planck scale \(M_{pl} = 2.4 \times 10^{18}\) GeV set to unity. With the above Kähler potential and superpotential, one can compute the scalar potential in supergravity and find that

\[
V = \frac{1}{3(T + T^* - \phi^* \phi)^2} \left| \frac{\partial W}{\partial \phi} \right|^2
\]

and no supersymmetry breaking masses arise in the scalar sector. Furthermore the gravitino mass is not fixed, which can be arbitrarily heavy or light at this level. Thus the no-scale model is named after this property. Non-trivial dependence of gauge kinetic functions on the field \(T\) yields non-vanishing gaugino masses in this case.

The no-scale structure appears when one considers a Calabi-Yau compactification of weakly coupled \(E_8 \times E_8\) heterotic string theory. If one focuses on the overall modulus field whose scalar component represents the overall size of the compactified space, then one finds \[2\]
\[ K = -\ln(S + S^*) - 3\ln(T + T^* - \phi^*\phi), \]  

(5)

where \( S \) is the dilaton field and \( T \) is the overall modulus field. The superpotential in this case generally depends on the these fields \( S \) and \( T \). Now if \( T \) dominates the SUSY breaking, then one finds that the soft SUSY breaking scalar mass as well as a trilinear scalar coupling (\( A \) term) vanishes as the vacuum energy, \textit{i.e.} the vacuum expectation value of the scalar potential, vanishes. Note that the \( T \) dominant SUSY breaking occurs when the gaugino condensation triggers the SUSY breaking.

The same structure was also obtained for the heterotic M-theory [8] which corresponds to the strong coupling regime of the heterotic string theory, but this time the fields \( S \) and \( T \) have physically different meanings. In both the weak-coupling and strong-coupling cases, one has to keep in mind that quantum corrections may alter the form of the Kähler potential (5).

Severe FCNC constraints on superparticle masses may suggest that the hidden sector and the observable sector are in some way separated from each other in the Kähler potential. An assumption often taken along this line of reasoning is the separation of the two sectors in the Kähler potential itself, namely the Kähler potential is a sum of the contributions from the two sectors. This Ansatz will generate the superparticle mass spectrum of the well-known minimal supergravity model and non-zero scalar masses arise. It may be, however, more natural to consider the same separation in the superspace density in the supergravity Lagrangian [3], before making Weyl transformations to obtain the Einstein-Hilbert action for gravity part. This spirit indeed leads the form of the Kähler potential in Eq. (1). In this case and in the string cases, the gaugino masses become non-zero, provided that the hidden sector couples to the gauge multiplets via the gauge kinetic functions.

Recently it has been pointed out that the form (1) is naturally realized in a five-dimensional setting with two separated 3-branes [4,5]. Consider the five-dimensional supergravity on \( R^4 \times S^1 / Z_2 \). The geometry has two four-dimensional boundaries, \textit{i.e.} 3-branes. Suppose that the hidden sector is on one of the 3-branes and the observable sector is on the
other. Now a dimensional reduction of the theory yields, in four dimensions, the following form of the Kähler potential

\[ K = -3 \ln(T + T^* + f_{\text{hid}}(z, z^*) + f_{\text{obs}}(\phi, \phi^*)), \] (6)

where this time the real part of \( T \) stands for the length of the compactified fifth dimension.

In the brane separation scenario, the two sectors are really split geometrically and thus not only the scalar masses, but also the gaugino masses vanish. Therefore one needs to seek for another mechanism to mediate the SUSY breaking occurred in the hidden sector. One way is to invoke superconformal anomaly to obtain loop-suppressed soft masses [4,9]. This anomaly mediation is very appealing, albeit its minimal version has negative masses squared for sleptons. Many attempts to build realistic models have been made [10], and the superparticle masses obtained are in general different from those from the no-scale boundary conditions. In ref. [11], a new \( U(1) \) gauge interaction is assumed to play a role of the mediator of the SUSY breaking. The resulting mass pattern is similar to that of gauge-mediated SUSY breaking. On the other hand, if the SM gauge sector lives in the bulk, then the gauginos can play a role of the SUSY-breaking messenger [12] and the resulting mass spectrum of the superparticles exhibits the no-scale structure with non-vanishing gaugino masses, which is given at the scale of (the inverse of) the length of the fifth-dimension.

**III. MINIMAL SCENARIO**

In this section, we would like to discuss phenomenological consequences of the minimal no-scale scenario which has been mainly studied in the literature. The soft SUSY breaking masses in the minimal case are parameterized by:

- vanishing scalar masses: \( m_0 = 0 \)
- vanishing trilinear scalar couplings: \( A = 0 \)
- non-zero Higgs mixing masses: \( B \)
• non-zero universal gaugino masses: $M_{1/2}$

Note that these values are given at the GUT scale $M_{GUT} \simeq 2 \times 10^{16} \text{ GeV}$. In addition to these soft masses, we assume a non-zero supersymmetric higgsino mass, $\mu$. These masses at the weak scale are obtained by solving renormalization group equations. Given $M_{1/2}$, requiring the correct electroweak symmetry breaking relates $B$ and $\mu$ with the $Z$ boson mass $m_Z$ and the ratio of the two Higgs vacuum expectation values $\tan \beta$ as in the usual manner.

At first we roughly estimate the mass spectrum of superparticles when the Yukawa effects and the left-right mixing effects are neglected. The bino, wino and gluino masses at the weak scale are given by one parameter $M_{1/2}$ (in the following we set the renormalization point to be 500 GeV),

$$M_1^2 \simeq 0.18 M_{1/2}^2, \quad M_2^2 \simeq 0.69 M_{1/2}^2, \quad M_3^2 \simeq 7.0 M_{1/2}^2 .$$ (7)

The soft SUSY breaking masses of scalars in the first-two generations are also determined by one parameter $M_{1/2}$.

$$\tilde{m}_{u_L}^2 \simeq 5.8 M_{1/2}^2 + 0.35 m_Z^2 \cos 2\beta$$ (8)

$$\tilde{m}_{d_L}^2 \simeq 5.8 M_{1/2}^2 - 0.42 m_Z^2 \cos 2\beta$$ (9)

$$\tilde{m}_{u_R}^2 \simeq 5.4 M_{1/2}^2 + 0.15 m_Z^2 \cos 2\beta$$ (10)

$$\tilde{m}_{d_R}^2 \simeq 5.4 M_{1/2}^2 - 0.077 m_Z^2 \cos 2\beta$$ (11)

$$\tilde{m}_{\ell_L}^2 \simeq 0.51 M_{1/2}^2 - 0.27 m_Z^2 \cos 2\beta$$ (12)

$$\tilde{m}_{\ell_R}^2 \simeq 0.15 M_{1/2}^2 - 0.23 m_Z^2 \cos 2\beta$$ (13)

$$\tilde{m}_\nu^2 \simeq 0.51 M_{1/2}^2 + 0.5 m_Z^2 \cos 2\beta$$ (14)

The terms proportional to $m_Z^2 \cos 2\beta$ are $U(1)_Y$ D-term contributions. From these equations, we find that bino and right-handed slepton are light. When $M_{1/2} \gtrsim 2.8 m_Z \sim 260 \text{ GeV}$, $U(1)_Y$ D-term contribution becomes small, and then the charged right-handed slepton becomes the LSP, and this scenario contradicts to cosmological observations.
In Fig. 1, we show the numerical result. The region above the solid line is excluded cosmologically since charged stau is the LSP. For \( \tan \beta \lesssim 10 \) where left-right mixing effect is negligible, the region \( M_{1/2} \gtrsim 260 \text{ GeV} \) is excluded as we estimate above. For \( \tan \beta \gtrsim 10 \), since left-right mixing effect makes stau mass lighter, the constraint becomes stronger. In Fig. 1, we also show the value of the right-handed smuon mass. From the cosmological constraint, we find that the right-handed smuon must be lighter than about 120 GeV.

On the other hand, the LEP experiments at \( \sqrt{s} = 202 \text{ GeV} \) provide rather strong lower bound on slepton masses [13]. For smuon, except near the threshold, the cross section for smuon pair production \( \sigma(e^+e^- \rightarrow \tilde{\mu}_R^+\tilde{\mu}_R^-) \) must be smaller than 0.05 pb to survive the smuon searches at LEP. Here we impose that \( \sigma(e^+e^- \rightarrow \tilde{\mu}_R^+\tilde{\mu}_R^-) \leq 0.05 \text{ pb} \) for \( m_{\tilde{\mu}_R} \leq 98 \text{ GeV} \) and \( m_{\tilde{\chi}_1^0} \leq 0.98m_{\tilde{\mu}_R} - 4.1 \text{ GeV} \). This constraint excludes the left side of the dashed line in the Fig. 1. Combining these two constraints, we conclude that the no-scale scenario with the universal gaugino masses is allowed only for \( \tan \beta \lesssim 8 \) and \( 210 \text{ GeV} \lesssim M_{1/2} \lesssim 270 \text{ GeV} \).

IV. CASE OF NON-UNIVERSAL GAUGINO MASSES

In this section, we consider modifications of the minimal boundary conditions discussed in the previous section, and argue that slight modifications will drastically change phenomenological consequences.

The reason of the very constrained superparticle mass spectrum in the minimal case is the degeneracy of the bino mass and those of the right-handed sleptons. The degeneracy is resolved if one considers the renormalization group effects above the GUT scale [14–16]. The point is that the right-handed slepton multiplets belong to 10-plets in the minimal choice of the matter representations in the \( SU(5) \) GUT, and the large group factor in the gauge loop contributions yields large positive corrections to the slepton masses. We should note, however, that in some realistic models to attempt to explain the masses of quarks, leptons and neutrinos, matter multiplets in different generations are often taken to be in different representations of the GUT groups [17], and then the renormalization group effects
would violate the mass degeneracy among the different generations, which might cause unacceptably large FCNCs.

Secondly the stau can be the lightest superparticle in the SSM sector if it is not stable. This is indeed the case when R-parity is violated or there exists another superparticle such as a gravitino out of the SSM sector which is lighter than the stau [18].

Another possibility is to relax the universality of the gaugino masses. In the rest of this section, we will discuss this case in detail. In the next subsection, we shall review various possibilities to realize non-universal gaugino masses. In particular we will emphasize that the non-universality of the gaugino masses does not conflict with the universality of the gauge couplings. Then we will look into phenomenological implications of the non-universality.

A. Examples of Non-Universal Gaugino Masses

Once the gaugino masses are given universal at some high energy scale where the gauge groups are unified, it is shown that the gaugino mass relation $M_1 : M_2 : M_3 \simeq 1 : 2 : 6$ holds at low energy, irrespective of the breaking patterns of the GUT group [19,14]. Here we review some mechanisms in which the gaugino masses are non-universal from the beginning.

In string models with simple Calabi-Yau compactification, the gauge kinetic functions for the Standard Model gauge multiplets can be written as [20]

$$f_i = S + \epsilon_i T$$

where $i = 1, 2, 3$ represent the three Standard Model gauge groups and $\epsilon_i$ are some coefficients of one-loop order determined by the details of the compactification. If $\epsilon_i$ depends on a gauge group and the modulus field $T$ is dominantly responsible for the SUSY breaking, we will have the non-universal gaugino masses:

$$M_1 : M_2 : M_3 = \epsilon_1 : \epsilon_2 : \epsilon_3.$$  

Here we would like to emphasize that large threshold corrections are necessary for the string unification scenario in the weak coupling regime where the string scale is more than one order
of magnitude larger than the naive GUT scale, and thus appearance of the non-universal $\epsilon_i$ terms seems to be requisite. Note again that the Kähler potential may receive quantum corrections at the same order and the no-scale structure may be distorted.

The non-universality of the gaugino masses can be achieved in the conventional GUT approaches. Suppose that the gauge kinetic functions are written in the following form [21]:

$$f = c + \Sigma Z$$  \hspace{1cm} (17)

where $c$ is a universal constant, $\Sigma$ is a field which breaks the GUT group to the SM group, and $Z$ is assumed to break the SUSY. The first term respects the GUT symmetry and thus universal for all SM gauge groups, while the second term is a symmetry breaking part which depends on each SM group. As for the gauge couplings, the first term gives a dominant contribution and hence the gauge couplings are unified up to small non-universal effects from the second term. On the other hand, the gaugino masses are assumed to come from the second term in Eq. (17). They are proportional to the vacuum expectation value of $\Sigma$ and thus non-universal. The form of Eq. (17) can also be obtained through GUT threshold corrections to the gauge kinetic functions [22].

Non-universal gaugino masses can also be realized in scenarios of product GUTs [23] where the gauge group has the structure of $G_{GUT} \times G_H$ and the Standard Model gauge groups are obtained as diagonal subgroups of the two product groups. The idea of the product GUTs provides an elegant solution to the triplet-doublet splitting problem in the Higgs sector based on the missing doublet mechanism. The gauge coupling unification achieves if the gauge couplings of the $G_H$ group are sufficiently large, while contributions to the gaugino masses from the $G_H$ sector are generally sizable and destroy their universality.

The flipped $SU(5)$ is another example where the non-universality of the gaugino masses naturally arises [22]. The gauge group is $SU(5) \times U(1)$ and thus even if the $SU(5)$ part gives a universal contribution the gaugino mass from $U(1)$ in general gives a different mass, violating the universality of the $U(1)_Y$ gaugino mass with the rest two.
In summary, the non-universality of the gaugino masses is not a peculiar phenomenon even in the light of the gauge coupling unification. Motivated by this observation, we will discuss its phenomenological consequences.

B. Phenomenological Implications

In this subsection we discuss some phenomenological implications of non-universal gaugino masses. At the cutoff scale, all scalar masses are vanishing as in the minimal case, while the bino, winos and gluinos possess nonzero masses $M_{1,0}$, $M_{2,0}$, and $M_{3,0}$, respectively, and now they are no longer degenerate in general. The soft SUSY breaking mass parameters at the weak scale are obtained by solving the RGEs. In this paper we use the one-loop level RGEs. With the soft SUSY breaking masses, we evaluate the physical masses using the tree-level potential. We also obtain the value of $\mu$ from the electroweak symmetry breaking condition with tree-level Higgs potential.

Before showing numerical results, we discuss the mass spectrum of superparticles when the Yukawa effects to the RG evolutions and left-right mixings are neglected. Relations of the gaugino masses at the GUT scale $M_{GUT}$ and the electroweak scale $M_{EW}$ are

$$M_1^2 \simeq 0.18 M_{1,0}^2, \quad M_2^2 \simeq 0.69 M_{2,0}^2, \quad M_3^2 \simeq 7.0 M_{3,0}^2. \quad (18)$$

Neglecting effects of Yukawa interaction, the masses squared of sfermions at the weak scale are evaluated to be

$$\tilde{m}_{u_L}^2 \simeq 5.4 M_{3,0}^2 + 0.47 M_{2,0}^2 + 4.2 \times 10^{-3} M_{1,0}^2 + 0.35 m_Z^2 \cos 2\beta \quad (19)$$

$$\tilde{m}_{d_L}^2 \simeq 5.4 M_{3,0}^2 + 0.47 M_{2,0}^2 + 4.2 \times 10^{-3} M_{1,0}^2 - 0.42 m_Z^2 \cos 2\beta \quad (20)$$

$$\tilde{m}_{u_R}^2 \simeq 5.4 M_{3,0}^2 + 0.066 M_{1,0}^2 + 0.15 m_Z^2 \cos 2\beta \quad (21)$$

$$\tilde{m}_{d_R}^2 \simeq 5.4 M_{3,0}^2 + 0.017 M_{1,0}^2 - 0.077 m_Z^2 \cos 2\beta \quad (22)$$

$$\tilde{m}_{\ell_L}^2 \simeq 0.47 M_{2,0}^2 + 0.037 M_{1,0}^2 - 0.27 m_Z^2 \cos 2\beta \quad (23)$$

$$\tilde{m}_{\ell_R}^2 \simeq 0.15 M_{1,0}^2 - 0.23 m_Z^2 \cos 2\beta \quad (24)$$
\[ \tilde{m}_\nu^2 \simeq 0.47M_{2,0}^2 + 0.037M_{1,0}^2 + 0.5m_Z^2 \cos 2\beta . \] (25)

From the above equations, we find that if \( M_{1,0} \gtrsim 2.0M_{2,0} \), \( \tilde{m}_{i\ell}^2 \) is heavier than \( M_1^2, M_2^2 \), \( \tilde{m}_{\ell\ell}^2 \) and \( \tilde{m}_\nu^2 \). Notice that the mass of the charged left-handed slepton is heavier than the mass of the neutral sneutrino because \( \cos 2\beta \leq 0 \) for \( \tan \beta \geq 1 \). On the other hand, for \( M_{1,0}/M_{2,0} \gtrsim 2.5 \), the wino mass tends to be lighter than the sneutrino mass. Hence we expect that sneutrino can be LSP when \( 2 \lesssim M_{1,0}/M_{2,0} \lesssim 2.5 \), and wino-like neutralino can be LSP when \( M_{1,0}/M_{2,0} \gtrsim 2.5 \).

Next, we consider how \( \mu \) affects the mass spectrum of the superparticles. The value \( \mu \) is determined by minimizing the Higgs potential. At the tree-level, \( \mu \) is calculated in terms of the soft SUSY breaking masses of the Higgses and \( \tan \beta \),

\[ \mu^2 = \frac{\tilde{m}_{H_d}^2 - \tilde{m}_{H_u}^2 \tan^2 \beta}{\tan^2 \beta - 1} - \frac{1}{2}m_Z^2 . \] (26)

In order to obtain the value of \( \tilde{m}_{H_d}^2 \) and \( \tilde{m}_{H_u}^2 \), we have to include the Yukawa interaction. For the moment we consider the low \( \tan \beta \) region, i.e., we take only the top Yukawa coupling into account and neglect the bottom and tau Yukawa couplings for simplicity. In this case \( \tilde{m}_{H_d}^2 = \tilde{m}_{H_u}^2 \) and we can obtain an analytic solution for the RGE of \( \tilde{m}_{H_u}^2 \). For \( \tan \beta = 10 \), \( \mu \) is approximately

\[ \mu^2 = 2.1M_{3,0}^2 - 0.22M_{2,0}^2 - 0.0064M_{1,0}^2 + 0.0063M_{1,0}M_{2,0} \]
\[ + 0.19M_{2,0}M_{3,0} + 0.029M_{3,0}M_{1,0} - \frac{1}{2}m_Z^2 . \] (27)

From this equation, we find that the size of \( \mu \) is strongly correlated with the size of the gluino mass \( M_{3,0} \) and \( |\mu| \) becomes large as \( M_{3,0} \) increases. Hence when \( M_{3,0} \) is large enough, the left-right mixing in the slepton masses is important, which makes one of the staus, \( \tilde{\tau}_1 \), lighter than sneutrino. On the other hand if \( M_{3,0} \) is small enough, \( |\mu| \) becomes smaller than the mass of bino, wino, slepton and sneutrinos, and then a higgsino-like neutralino can be the LSP. Actually, for \( \tan \beta = 10 \), from eq. (27) we find that \( |\mu| \) is smaller than \( M_2 \) if \( M_{3,0}/M_{2,0} \lesssim 0.5 \) is satisfied.
In the non-universal case, not only the mass spectrum but also the mixing properties of the neutralinos are very different from those in the minimal case. To see this we classify the lightest neutralino $\chi_1^0$ into five cases as follows. $\tilde{\chi}_1^0$ is a linear combination of bino, wino and higgsinos and is written as

$$\tilde{\chi}_1^0 = (O_N)_{1B} \tilde{B} + (O_N)_{1W} \tilde{W} + (O_N)_{1H_d} \tilde{H}_d + (O_N)_{1H_u} \tilde{H}_u ,$$

(28)

where $O_N$ is orthogonal matrix diagonalizing the neutralino mass matrix. When $|(O_N)_{1B}|^2 > 0.8$, $|(O_N)_{1W}|^2 > 0.8$ or $|(O_N)_{1H_d}|^2 + |(O_N)_{1H_u}|^2 > 0.8$, we call these parameter region 'bino region', 'wino region' or 'higgsino region', respectively. When $|(O_N)_{1B}|^2 < 0.8$ and $|(O_N)_{1W}|^2 < 0.8$ and $|(O_N)_{1B}|^2 + |(O_N)_{1W}|^2 > 0.8$, we call the region 'bino-wino mixed region'. The other parameter region is called 'mixed region'.

In Fig. 2 we show the composition of the LSP when we relax the gaugino mass universality. Here we take $M_{2,0} = 200\text{GeV}$, $\tan \beta = 10$ and $\text{sgn}(\mu) = +1$. Recall that for the universal gaugino masses at the GUT scale, the LSP is the lighter stau and this parameter set is excluded. Once we relax the universality, however, we see that the situation drastically changes, and the composition of the LSP behaves as we have discussed with the approximate expressions eq.(18) - eq.(25). The lightest neutralino can be the LSP in a large parameter region, and furthermore unlike the universal case, it can be wino-like, higgsino-like or admixture of them as well as bino-like. When $M_{1,0}/M_{2,0} \gtrsim 2.5$ and $M_{3,0}/M_{2,0} \gtrsim 1$ the wino is the LSP. And as the ratio $M_{3,0}/M_{2,0}$ decreases, $|\mu|$ becomes comparable to $M_1$ and $M_2$ and the lightest neutralino is the admixture of bino, wino and higgsinos. Further $M_{3,0}/M_{2,0}$ becomes smaller than about 0.5, the dominant component of the lightest neutralino is higgsino. Also we find in the region $2 \gtrsim M_{1,0}/M_{2,0} \gtrsim 2.5$, the tau sneutrino is indeed the LSP. And we find that when $M_{3,0}/M_{2,0}$ is larger than 2, i.e., $|\mu|$ is large and so is the left-handed and right-handed stau mixing, sneutrino can not be the LSP, and stau is LSP even when $M_{1,0}/M_{2,0}$ is bigger than $2.5 - 3$.

In the non-universal case sfermions, as well as neutralinos and charginos, show variety of mass spectrum. From eq. (23) and eq. (24), we find that when $M_{1,0}/M_{2,0} \gtrsim 2$ left-handed
sfermions are smaller than right-handed sfermions in contrast to the universal case. For stau, the mixing angle of \( \tilde{\tau}_L \) and \( \tilde{\tau}_R \) also depends on this ratio. In Fig. 3 we show the behavior of this mixing angle \( \theta_\tau \) in the \( M_{1,0}/M_{2,0} - M_{3,0}/M_{2,0} \) plane, where \( \theta_\tau \) is defined such that the lighter stau \( \tilde{\tau}_1 \) is written as \( \tilde{\tau}_1 = \cos \theta_\tau \tilde{\tau}_L + \sin \theta_\tau \tilde{\tau}_R \). Around \( M_{1,0}/M_{2,0} \sim 2 \), the mass of the right-handed stau is as heavy as that of the left-handed stau, and they mix maximally \( (\theta_\tau = 40 - 50) \) as expected. Also masses of squarks strongly depend on \( M_3 \), and thus mass relations between squarks and sleptons drastically change. As we shall see later, some of the squarks can be lighter than the sleptons.

In Fig. 4 we show the same graph as Fig. 2 except for \( \tan \beta = 35 \). In this case the Yukawa interaction and the left-right mixing make the stau mass lighter. In fact, although wino-like, higgsino-like and mixed neutralino is LSP in large parameter region, the sneutrino can not be the LSP. To see the relation among the stau mass, the tau sneutrino mass and \( \tan \beta \), we plot in Fig. 5 the composition of the LSP in the \( M_{3,0}/M_{2,0} - \tan \beta \) plane, fixing \( M_{3,0}/M_{2,0} = 2.5 \). This figure shows that the tau sneutrino can be the LSP when \( \tan \beta \lesssim 15 \) where the left-right mixing is not so sizable. We have checked that these features are insensitive to the signs of \( \mu \) and gaugino masses.

We shall next investigate the mass spectrum of superparticles in detail, by choosing some representative parameter sets, and discuss phenomenology for each parameter set. The points we choose are listed in Table I. In table II we show contamination of \( \tilde{\chi}_1^0 \) for each point. At the points A and E, the LSP is the wino-like neutralino. At the points B, C and E the LSP is the higgsino-like neutralino. And at the point D the tau sneutrino is the LSP. In table III we list the mass spectrum of superparticles.

The wino-like neutralino is the LSP when \( M_{1,0}/M_{2,0} \gtrsim 2 \) and \( M_{3,0}/M_{2,0} \gtrsim 1 \). In the wino-like neutralino LSP case, the lighter chargino and the lightest neutralino are highly degenerate generally. This character and the resulting phenomenology has been studied in [25–30]. On top of this our scenario also predicts that the right-handed sfermions are heavier than the left-handed ones because of the inequality \( M_{1,0}/M_{2,0} \gtrsim 2 \), and colored superparticles are heavier than other superparticles because of the inequality \( M_{3,0}/M_{2,0} \gtrsim 1 \).
(see the list for the points A and E in the table [III]). The former may be an interesting feature. Anomaly-mediated SUSY breaking (AMSB) scenario also predicts the wino-like LSP. However in the minimal AMSB model where universal mass is added to all scalars to avoid negative slepton masses squared, the left-handed and right-handed sleptons in the first two generations tend to be degenerate [29]. Thus we can distinguish two scenarios with the wino LSP, the no-scale scenario with non-universal gaugino masses and the minimal AMSB, by measuring these slepton masses.

The higgsino-like neutralino is the LSP when $M_{3,0}/M_{2,0} \lesssim 0.5$ regardless of $M_{1,0}/M_{2,0}$. In the higgsino-like neutralino LSP case, the mass deference between higgsino-like neutralino LSP and chargino NLSP is generally small. The resulting phenomenology have been studied in [31–33]. Furthermore in our case, the sleptons are as heavy as the squarks due to the inequality $M_{3,0}/M_{2,0} \lesssim 0.5$. Especially the lighter stop and sbottom can be lighter than some of the sleptons. Actually, at the points B, C, and F, the lighter stop is comparable to the slepton masses, and all superparticle masses are below 400 - 450 GeV.

In the non-universal scenario, the tau sneutrino can also be the LSP when $2 \lesssim M_{1,0}/M_{2,0} \lesssim 2.5$, $1 \lesssim M_{3,0}/M_{2,0} \lesssim 5$ and $\tan \beta \lesssim 15$. From the first inequality, we find that the mass difference between left-handed and right-handed squarks in the first two generations is small, and the left-right mixing angle of the stau is big as shown in Fig. 3.

V. CONCLUSIONS

In this paper, we have revisited the no-scale scenario where the vanishing SUSY-breaking scalar masses and trilinear scalar couplings are given at the GUT scale. When the gaugino masses are given universal, the renormalization group analysis implies that the bino mass and the right-handed slepton masses are close to each other. This degeneracy leads an upperbound of the LSP mass around 120 GeV: above it the LSP would be the charged stau, which must be excluded cosmologically. Furthermore, the negative results of the slepton searches at LEP200 already excluded a large portion of the parameter space including a
large tan β region, leaving tan β ∼ 8.

We next considered various ways out to avoid the aforementioned severe constraints. Among them, we concentrated on the case of the non-universal gaugino masses. In fact the non-universality of the gaugino masses is by no means a peculiar phenomenon, rather it is realized in various scenarios, including some approaches to grand unification. We investigated some phenomenological implications of the no-scale model with the non-universal gaugino masses. We found that there is no longer severe constraint on the superparticle masses and the mass spectrum of the superparticle has much richer structure. In particular, the LSP can be the wino-like neutralino, the higgsino-like neutralino, or even the sneutrino. We also found that unlike the conventional universal gaugino mass case, the left-handed slepton masses can be lighter than the right-handed slepton masses. We expect that resulting collider signatures with these features will be quite different from the usual scenario with the universal gaugino masses. Further studies along this direction should be encouraged.

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|      | Point A | Point B | Point C | Point D | Point E | Point F |
|------|---------|---------|---------|---------|---------|---------|
| $M_{1,0}$ | 800     | 1000    | 400     | 500     | 800     | 600     |
| $M_{2,0}$ | 200     | 250     | 200     | 200     | 200     | 200     |
| $M_{3,0}$ | 400     | 125     | 100     | 300     | 300     | 100     |
| $\tan \beta$ | 10      | 10      | 10      | 10      | 35      | 35      |

**TABLE I.** Gaugino masses at the GUT scale for each point. All dimensionful parameters are given in the GeV unit.

|      | Point A | Point B | Point C | Point D | Point E | Point F |
|------|---------|---------|---------|---------|---------|---------|
| $(O_{N})_{1B}$ | -0.017 | 0.0835  | 0.241   | 0.092   | -0.022  | 0.126   |
| $(O_{N})_{1W}$ | 0.987  | -0.478  | -0.457  | -0.967  | 0.973   | -0.445  |
| $(O_{N})_{1H_d}$ | -0.149 | 0.689   | 0.710   | 0.219   | -0.213  | 0.729   |
| $(O_{N})_{1H_u}$ | 0.054  | -0.539  | -0.479  | -0.096  | 0.084   | -0.504  |

**TABLE II.** Components of the lightest neutralino $\tilde{\chi}_1^0$ which is a linear combination of bino, wino and higgsinos, $\tilde{\chi}_1^0 = (O_{N})_{1B} \tilde{B} + (O_{N})_{1W} \tilde{W} + (O_{N})_{1H_d} \tilde{H}_d + (O_{N})_{1H_u} \tilde{H}_u$. 
| particle | Point A | Point B | Point C | Point D | Point E | Point F |
|----------|---------|---------|---------|---------|---------|---------|
| $\chi^0_1$ | 160 | 106 | 70 | 156 | 159 | 72 |
| $\chi^0_2$ | 336 | 152 | 126 | 209 | 332 | 120 |
| $\chi^0_3$ | 594 | 248 | 169 | 444 | 438 | 202 |
| $\chi^0_4$ | 603 | 430 | 222 | 457 | 453 | 267 |
| $\chi^+_1$ | 160 | 113 | 81 | 157 | 159 | 81 |
| $\chi^+_2$ | 602 | 253 | 216 | 457 | 449 | 212 |
| $\tilde{u}_L$ | 929 | 336 | 263 | 701 | 702 | 265 |
| $\tilde{d}_L$ | 932 | 345 | 275 | 706 | 707 | 277 |
| $\tilde{u}_R$ | 941 | 384 | 249 | 700 | 718 | 274 |
| $\tilde{d}_R$ | 925 | 316 | 237 | 692 | 697 | 244 |
| $\tilde{\nu}$ | 196 | 250 | 144 | 155 | 196 | 167 |
| $\tilde{e}_L$ | 212 | 262 | 164 | 174 | 212 | 185 |
| $\tilde{e}_R$ | 312 | 389 | 161 | 198 | 312 | 236 |
| $\tilde{t}_1$ | 742 | 233 | 164 | 538 | 544 | 176 |
| $\tilde{t}_2$ | 925 | 409 | 339 | 721 | 709 | 338 |
| $\tilde{b}_1$ | 855 | 293 | 230 | 646 | 602 | 200 |
| $\tilde{b}_2$ | 922 | 317 | 252 | 691 | 676 | 252 |
| $\tilde{\nu}_\tau$ | 195 | 249 | 143 | 154 | 183 | 156 |
| $\tilde{\tau}_1$ | 205 | 261 | 154 | 159 | 166 | 169 |
| $\tilde{\tau}_2$ | 314 | 387 | 169 | 209 | 316 | 225 |
| $\tilde{g}$ | 1053 | 329 | 263 | 790 | 790 | 263 |

TABLE III. Mass spectrum for each point. All values are given in the GeV unit.
FIG. 1. Allowed region of the minimal no-scale scenario. The horizontal axis is the universal gaugino mass at the GUT scale $M_{1/2}$ and the vertical axis is $\tan \beta$. In the region above the solid line $\tilde{\tau}$ is the LSP and it should be cosmologically excluded. The left side of the dashed line is excluded by smuon searches by the LEP experiments at $\sqrt{s} = 202$ GeV. We also show the contour of right-handed smuon mass.

FIG. 2. The composition of the LSP in the $M_{1,0}/M_{2,0} - M_{3,0}/M_{2,0}$ plane, for $M_{2,0} = 200$GeV and $\tan \beta = 10$. The classification of the neutralino LSP is given in the text.
FIG. 3. Mixing angle $\theta_\tau$ in the $M_{1,0}/M_{2,0} - M_{3,0}/M_{2,0}$ plane for $M_{2,0} = 200\,\text{GeV}$ and $\tan\beta = 10$. We also show the region where stau and tau sneutrino are the LSP. The definition of the mixing angle is given in the text.

FIG. 4. The same as Fig. 3 but for $M_{2,0} = 200\,\text{GeV}$ and $\tan\beta = 35$. 
FIG. 5. The composition of the LSP in the $M_{3,0}/M_{2,0} - \tan \beta$ plane, for $M_{2,0} = 200\text{ GeV}$ and $M_{1,0}/M_{2,0} = 2.5$. 
REFERENCES

[1] J. Ellis, C. Kounnas and D.V. Nanopoulos, Nucl. Phys. B247, 373 (1984).

[2] E. Witten, Phys. Lett. B155, 151 (1985).

[3] K. Inoue, M. Kawasaki, M. Yamaguchi and T. Yanagida, Phys. Rev. D45, 328 (1992).

[4] L. Randall and R. Sundrum, Nucl. Phys. B557, 79 (1999).

[5] M. Luty and R. Sundrum, hep-th/9910202.

[6] S. Kelly, J.L. Lopez, D.V. Nanopoulos, H. Pois and K.-J. Yuan, Phys. Lett. B273, 423 (1991).

[7] M. Drees and M.M. Nojiri, Phys. Rev. D45, 2482 (1992).

[8] T. Banks and M. Dine, Nucl. Phys. B505, 445 (1997); H.P. Nilles, M. Olechowski and M. Yamaguchi, Phys. Lett. B415 (1997) 24; Nucl. Phys. B530, 43 (1998); Z. Lalak and S. Thomas, Nucl. Phys. B515, 55 (1998); A. Lukas, B.A. Ovrut and D. Waldram, Nucl. Phys. B532, 43 (1998); Phys. Rev. D57, 7529 (1998); K. Choi, H.B. Kim and C. Muñoz, Phys. Rev. D57, 7521 (1998).

[9] G.F. Giudice, M.A. Luty, H. Murayama and R. Rattazzi, JHEP 9812, 027 (1998).

[10] See for example, A. Pomarol and R. Rattazzi, JHEP 9905, 013 (1999); Z. Chacko, M.A. Luty, I. Maksymyk and E. Pontón, JHEP 0004, 001 (2000); E. Katz, Y. Shadmi and Y. Shirman, JHEP 9908, 015 (1999); K.-I. Izawa, Y. Nomura and T. Yanagida, Prog. Theor. Phys. 102, 1181 (1999); I. Jack and D.R.T. Jones, Phys. Lett. B482, 167 (2000).

[11] Y. Nomura and T. Yanagida, hep-ph/0005211.

[12] D.E. Kaplan, G.D. Kribs and M. Schmaltz, hep-ph/9911293; Z. Chacko, M. Luty, A.E. Nelson and E. Pontón, JHEP 0001, 003 (2000); M. Schmaltz and W. Skiba, hep-ph/0001172.
[13] G. Ganis, 'Standard SUSY at LEP', talk presented in SUSY2K, 8th Int. Conf. on Supersymmetry in Physics, CERN, Geneva, Switzerland, June 26 - July 1, 2000

[14] Y. Kawamura, H. Murayama and M. Yamaguchi, Phys. Rev. D51, 1337 (1995).

[15] N. Polonsky and A. Pomarol, Phys. Rev. D51, 6532 (1995)

[16] M. Schmaltz and W. Skiba, hep-ph/0004210.

[17] See for example, J. Sato and T. Yanagida, Phys. Lett. B430, 127 (1998); Y. Nomura and T. Yanagida, Phys. Rev. D59, 017303 (1999); M. Bando and T. Kugo, Prog. Theor. Phys. 101, 1313 (1999).

[18] T. Moroi, H. Murayama and M. Yamaguchi, Phys. Lett. B303 (1993) 289; T. Gherghetta, G.F. Giudice and A. Riotto, Phys. Lett. B446 (1999) 28; T. Asaka, K. Hamaguchi and K. Suzuki, hep-ph/0005136.

[19] Y. Kawamura, H. Murayama and M. Yamaguchi, Phys. Lett. B324, 52 (1994).

[20] A. Brignole, L.E. Ibáñez and C. Muñoz, Nucl. Phys. B422, 125 (1994); Erratum ibid. B437, 747 (1995).

[21] J. Ellis, K. Enqvist, D. Nanopoulos, and K. Tamvakis, Phys. Lett. 155B, 381 (1985); G. Anderson, C.-H. Chen, J.F. Gunion, J. Lykken, T. Moroi and Y. Yamada, in New Directions for High Energy Physics, Snowmass 96, edited by D. G. Cassel, L. Trindle Gennari, and R. H. Siemann (Stanford Linear Accelerator Center, Menlo Park, CA, 1997); K. Huitu, Y. Kawamura, T. Kobayashi, and K. Puolamäki, Phys. Rev. D61,035001 (1999); G. Anderson, H. Baer, C.-H. Chen, P. Quintana and X. Tata, Phys. Rev. D61, 095005 (2000).

[22] J. Hisano, T. Goto, and H. Murayama, Phys. Rev. D49, 1446 (1994)

[23] T. Hotta, K.-I. Izawa and T. Yanagida, Phys. Rev. D53, 3913 (1996); Prog. Theor. Phys. 95, 949 (1996); Phys. Rev. D54, 6970 (1996).
[24] N. Arkani-Hamed, H.-C. Cheng and T. Moroi, Phys. Lett. B387, 529 (1996); K. Kurosawa, Y. Nomura and K. Suzuki, Phys. Rev. D60, 117701 (1999).

[25] S. Mizuta, D. Ng, and M. Yamaguchi, Phys. Lett. B300, 96 (1993)

[26] D. Pierce and A. Papadopoulos, Nucl. Phys B430, 278 (1994), and Phys. Rev. D50, 565 (1994)

[27] C.-H. Chen, M. Dress and J. F. Gunion, Phys. Rev. Lett. 76, 2002 (1996); Phys. Rev. D55, 330 (1997)

[28] J. F. Feng, T. Moroi, L. Randall, M. Strassler, and S. Su, Phys. Rev. Lett. 83, 1731 (1999)

[29] J. F. Feng and T. Moroi, Phys. Rev. D61, 095004 (2000)

[30] J. F. Gunion and S. Mrenna, Phys. Rev. D62, 015002 (2000)

[31] S. Mizuta and M. Yamaguchi, Phys. Lett. B298, 120 (1993)

[32] G. F. Giudice and A. Pomarol, Phys. Lett. B372, 253 (1996)

[33] M. Dress, M. M. Nojiri, D. P. Roy and Y. Yamada, Phys. Rev. D56, 276 (1997)