The gaugino code

Kiwoon Choi

Department of Physics, Korea Advanced Institute of Science and Technology
Daejeon 305-701, Korea
E-mail: `kchoi@hep.kaist.ac.kr`

Hans Peter Nilles

Physikalisches Institut, Universität Bonn
Nussallee 12, D-53115 Bonn, Germany
E-mail: `nilles@th.physik.uni-bonn.de`

ABSTRACT: Gauginos might play a crucial role in the search for supersymmetry at the Large Hadron Collider (LHC). Mass predictions for gauginos are rather robust and often related to the values of the gauge couplings. We analyse the ratios of gaugino masses in the LHC energy range for various schemes of supersymmetry breakdown and mediation. Three distinct mass patterns emerge.

Contents

1. Introduction

Soft breaking terms are the signals of the various schemes of supersymmetry (SUSY) breakdown and its mediation to the superpartners of the standard model (SM) particles [1]. With the upcoming experiments at LHC we might hope to identify these superpartners and get information about their spectrum and interactions. This in turn would allow us to infer the pattern of the soft supersymmetry breaking terms. The crucial question then concerns our ability to identify supersymmetry as the underlying scheme and then draw conclusions about the mechanism that is responsible for supersymmetry breakdown.

This could be a difficult task, as the relation between superpartner spectra and the underlying scheme could be quite complicated and model dependent, especially in the case of incomplete experimental knowledge of the spectra. Strategies to decode the spectra and determine the mechanism of supersymmetry breakdown have to be developed. Do there exist some model independent properties that reveal special schemes? Can we make useful statements without the knowledge of the heavy particle spectrum far beyond the

TeV-scale? In general, of course, we will not be able to answer these questions in detail, but we might hope to identify some basic characteristic patterns of soft terms.

From all the soft terms known, the gaugino masses have the simplest form and appear to be the least model dependent, as e.g. compared to the mass terms of squarks and sleptons. Therefore, an identification of gaugino-like particles would be a first step in favour of a potential supersymmetric interpretation of physics beyond the standard model at the LHC. In addition, gaugino masses M_a ($a = 1, 2, 3$) are often simply related to the gauge coupling constants g_a of the SM gauge group $SU(3) \times SU(2) \times U(1)$. These gauge coupling constants have been measured at the TeV-scale with the (approximate) result:

$$g_1^2 : g_2^2 : g_3^2 \simeq 1 : 2 : 6. \tag{1.1}$$

In the minimal supersymmetric extension of the standard model (MSSM) the renormalization group evolution of gaugino masses and gauge coupling constants is related in a simple way: M_a/g_a^2 does not run at the one-loop level. In a basic scheme like gravity mediation [2] with a universal gaugino mass $M_{1/2}$ at the grand unified (GUT) scale, i.e. mSUGRA scenario, the MSSM gaugino masses at the TeV scale would thus obey the relation:

$$\text{mSUGRA pattern:} \quad M_1 : M_2 : M_3 \simeq 1 : 2 : 6 \tag{1.2}$$

which we call the mSUGRA pattern of low energy gaugino masses in the subsequent discussion.

Would such a pattern of gaugino masses uniquely point back to gravity mediation with universal $M_{1/2}$ at the GUT-scale? Or in other words, how common is the mSUGRA pattern to various SUSY-breaking schemes? Does it depend on the universality of gauge couplings and/or gaugino masses at a large scale? What are other possible patterns of gaugino masses which would result from a reasonable theoretical scheme that could at the same time be clearly distinguished from the mSUGRA pattern?

In the present paper we would like to address these questions and try to identify various patterns of gaugino masses at the TeV-scale which might be obtained within a reasonable theoretical framework. Our main goal is to see what kind of information on SUSY breakdown can be extracted once one can determine the low energy gaugino mass ratios by future collider experiments. We shall find that only a few distinct patterns emerge, some of them characteristic for a specific scheme, others shared by quite different underlying schemes. It is encouraging to see that in many cases the values of low energy gaugino masses are independent of the particle spectrum at a large scale (like the GUT or intermediate scale) and can therefore give us precious model independent information.

Here is a summary of our results. The mSUGRA pattern (1.2) of low energy gaugino masses is shared by many different schemes of SUSY breaking. These include, of course, the mSUGRA-type SUSY-breaking scenarii realized in different higher dimensional supergravity (SUGRA) or string theories with a large string and compactification scales near $M_{\text{GUT}} \simeq 2 \times 10^{16}$ GeV or $M_{\text{Pl}} \simeq 2.4 \times 10^{18}$ GeV, e.g. the dilaton/moduli-mediated SUSY breakdown in heterotic string/M theory compactified on Calabi-Yau manifolds [3, 4], flux-induced SUSY breakdown in Type IIB string theory [5], as well as the gaugino mediation

realized in higher dimensional brane models [6]. These mSUGRA-type scenarios either predict, or assume if necessary, the unification of gauge couplings and/or gaugino masses at a large scale near $M_{\rm GUT}$, and they all give rise to the gaugino mass pattern (1.2) at the TeV scale.

The scheme known as gauge mediation [7], although quite different from mSUGRA scenarii in other aspects, also gives rise to the low energy gaugino mass pattern (1.2) under the assumption of gauge coupling unification at $M_{\rm GUT}$. In fact, even a broader class of different SUSY breaking schemes can lead to the mSUGRA pattern. A nontrivial example of this kind is the large volume compactification of Type IIB string theory discussed in [8]. Although the internal manifold of this compactification has an exponentially large volume and the string scale has an intermediate scale value $M_{\rm st} \sim 10^{11}$ GeV, this scheme gives the mSUGRA pattern (1.2) irrespective of whether or not the gauge couplings are unified at the intermediate string scale.

There are two other simple patterns of low energy gaugino masses distinct from the mSUGRA pattern that emerge under suitable theoretical assumptions that seem to be particularly appealing and well-motivated. The first one is the one appearing in the scheme of anomaly mediation [9]:

$$\text{Anomaly pattern:} \quad M_1 : M_2 : M_3 \simeq 3.3 : 1 : 9, \tag{1.3}$$

and the second one is mirage mediation [10]:

$$\text{Mirage pattern:} \quad M_1 : M_2 : M_3 \simeq (1 + 0.66\alpha) : (2 + 0.2\alpha) : (6 - 1.8\alpha), \tag{1.4}$$

where α is a parameter of order unity that will be defined later.

The anomaly pattern in (1.3) requires that all SUSY-breaking fields X^I are sequestered from the visible sector. Such sequestering might be naturally achieved in certain class of theories with extra dimensions [11] or CFT sector [12]. Then the gaugino masses are dominated by the SUGRA compensator-mediated contribution, such that

$$\frac{M_a(\mu)}{g_a^2(\mu)} = \frac{b_a}{16\pi^2} m_{3/2}, \tag{1.5}$$

where $m_{3/2}$ is the gravitino mass and b_a are the one-loop beta-function coefficients at the scale μ. Independently of the UV structure at scales above TeV, $b_a = (33/5, 1, -3)$ at TeV if the effective theory at TeV is given by the MSSM, which yields the anomaly pattern of low energy gaugino masses.

The mirage pattern in (1.4) is a kind of hybrid between the mSUGRA pattern and the anomaly pattern as it arises from SUSY-breaking schemes in which the soft terms receive comparable contributions from both moduli mediation and anomaly mediation. It has been observed that such scheme is naturally realized in KKLT-type moduli stabilization [13] and its appropriate generalizations [14, 15] which yield

$$\frac{M_a(\mu)}{g_a^2(\mu)} = \left(1 + \frac{\ln(M_{\rm Pl}/m_{3/2})}{16\pi^2} g_{\rm GUT}^2 b_a \alpha\right) \frac{M_0}{g_{\rm GUT}^2}, \tag{1.6}$$

where $M_0 \sim 1\,\mathrm{TeV}$ is a mass parameter characterizing the moduli mediation and $\alpha = \mathcal{O}(1)$ is a parameter representing the ratio of anomaly mediation to moduli mediation. In this case, gaugino masses are *unified at a mirage scale* [16]:

$$M_{\mathrm{mir}} = M_{\mathrm{GUT}} \left(\frac{m_{3/2}}{M_{\mathrm{Pl}}} \right)^{\alpha/2}, \qquad (1.7)$$

and the resulting low energy values take the mirage pattern in (1.4) for $g^2_{\mathrm{GUT}} \simeq 1/2$. In fact, the original KKLT-type moduli stabilization and some of its generalizations predict $\alpha \simeq 1$, although different values of $\alpha = \mathcal{O}(1)$ are possible in other generalizations. As a consequence, the gaugino mass pattern (1.4) with $\alpha = 1$ can be considered as a benchmark point of the mirage pattern.

Another example of SUSY-breaking scheme leading to the mirage pattern of gaugino masses is deflected anomaly mediation [17] in which M_a/g_a^2 receive contributions from anomaly-mediation and gauge-mediation of comparable size. In this case, α represents the ratio of anomaly to gauge mediation. Although it gives the same pattern of low energy gaugino masses as mirage mediation, deflected anomaly mediation can be distinguished from mirage mediation as it gives a different pattern of soft scalar masses for a given value of α [10, 16, 17].

Our scan of well motivated mediation scenarii thus leads to a remarkably small number of distinct patterns of gaugino masses that could be tested at the LHC. At LHC, the cascade decays of gluino and squarks are expected to provide information on the various combinations of the gluino, squark, slepton, and neutralino masses [18]. Using the kinematic edges and thresholds of various invariant mass distributions, the gluino mass M_3 and the two lightest neutralino masses $m_{\chi_1^0}$ and $m_{\chi_2^0}$ are expected to be determined with a reasonable accuracy. In case that $\chi_{1,2}^0$ are mostly the $\mathrm{SU}(2) \times \mathrm{U}(1)$ gauginos, this would mean that the three MSSM gaugino masses are determined. In the other case that χ_1^0 and χ_2^0 correspond mostly to the neutral Higgsinos, the dilepton invariant mass distribution of $\chi_2^0 \to \chi_1^0 l l^+$ shows a quite distinctive feature [19], and then one might be able to extend the kinematic analysis to determine the heavier neutralino masses $\chi_{3,4}^0$ which are then expected to be the $\mathrm{SU}(2) \times \mathrm{U}(1)$ gaugino masses. At any rate, under the assumption that the neutralino mixings are not sizable, the mSUGRA pattern and the anomaly pattern predict $M_3/m_{\chi_1^0} \simeq 6$ and $M_3/m_{\chi_1^0} \simeq 9$, respectively. Thus, regardless of the nature of the LSP neutralino χ_1^0, they could be excluded if the data indicates the gluino to LSP mass ratio significantly smaller than 6. On the other hand, even when one finds the gluino to LSP mass ratio close to 6 or an even larger value, one still needs to check the possibility of Higgsino-like LSP in order to exclude the mirage pattern predicting M_3/M_1 significantly smaller than 6 for a positive $\alpha = \mathcal{O}(1)$.

With more data, any conclusion based on the gaugino masses can then be checked, once candidates for squark and sleptons have been identified. Mass predictions for squarks and sleptons, however, show a stronger model dependence compared to the rather robust gaugino mass patterns. Of course, even in the case of gaugino masses we have to worry about the possibility of a strong influence of physics at a high scale (e.g. string threshold effects) that might be present in some of the schemes [20, 21] and obscure the 3 simple

patterns identified above. We shall describe these uncertainties once we discuss the explicit scenarii in section III. Still we think that an analysis of the gaugino masses is a most promising first step to reveal the nature of the underlying scheme.

The organization of this paper is as follows. In the next section, we discuss the general expression of M_a/g_a^2 in the context of effective SUGRA. In section III, we consider various specific examples including the scenarii which might be realized in the context of Type IIB string theory, heterotic string/M theory, and also M theory compactified on a manifold with G_2 holonomy. Section IV contains the summary and describes future strategies to identify the underlying mechanism of supersymmetry breakdown.

2. Gaugino masses in 4D effective supergravity

To start with, let us consider 4D effective SUGRA defined at the cutoff scale Λ for the visible sector physics. This 4D SUGRA might correspond to the low energy limit of compactified string theory or brane model. The Wilsonian effective action of the model at Λ can be written as

$$\int d^4\theta \, CC^* \Big(-3e^{-K/3} \Big) + \left[\int d^2\theta \, \Big(\frac{1}{4} f_a W^{a\alpha} W_\alpha^a + C^3 W \Big) + \text{h.c} \right], \qquad (2.1)$$

where C is the chiral compensator of 4D SUGRA, K is the Kähler potential, W is the superpotential, and f_a are holomorphic gauge kinetic functions. As usual, K can be expanded as

$$K = K_0(X_I, X_I^*) + Z_i(X_I, X_I^*) Q_i^* Q_i, \qquad (2.2)$$

where Q_i are chiral matter superfields which have a mass lighter than Λ and are charged under the visible sector gauge group, and X_I are SUSY breaking (moduli or matter) fields which have nonzero F-components F^I.

The running gauge couplings and gaugino masses at a scale μ below Λ but above the next threshold scale M_{th} can be determined by the 1PI gauge coupling superfield $\mathcal{F}_a(p^2)$ $(M_{\text{th}}^2 < p^2 < \Lambda^2)$ which corresponds to the gauge kinetic coefficient in the 1PI effective action on superpspace:

$$\Gamma_{1PI} = \int d^4p \, d^4\theta \, \Big(\frac{1}{4} \mathcal{F}_a(p^2) W^a \frac{\mathcal{D}^2}{16p^2} W^a + \text{h.c} \Big) \qquad (2.3)$$

At one-loop approximation, $\mathcal{F}_a$ is given by [22–24]

$$\mathcal{F}_a(p^2) = \text{Re}(f_a^{(0)}) - \frac{1}{16\pi^2}(3C_a - \sum_i C_a^i) \ln\left(\frac{CC^*\Lambda^2}{p^2} \right)$$

$$-\frac{1}{8\pi^2} \sum_i C_a^i \ln\left(e^{-K_0/3} Z_i \right) + \frac{1}{8\pi^2} \Omega_a \qquad (2.4)$$

where $f_a^{(0)}$ are the tree-level gauge kinetic function, C_a and C_a^i are the quadratic Casimir of the gauge multiplets and the matter representation Q_i, respectively. Here Ω_a contains

the string and/or KK threshold corrections from heavy fields at scales above Λ as well as the (regularization scheme-dependent) field-theoretic one-loop part: $\frac{1}{8\pi^2} C_a \ln[\text{Re}(f_a^{(0)})]$. In the one-loop approximation, Ω_a are independent of the external momentum p^2, thus independent of C as a consequence of the super-Weyl invariance. However Ω_a generically depend on SUSY breaking fields X_I, and a full determination of their X_I-dependence requires a detailed knowledge of the UV physics above Λ.

The running gauge couplings and gaugino masses at a renormalization point μ ($M_{\text{th}} < \mu < \Lambda$) are given by

$$\frac{1}{g_a^2(\mu)} = \mathcal{F}_a|_{C=e^{K_0/6}, \, p^2=\mu^2},$$
$$M_a(\mu) = F^A \partial_A \ln(\mathcal{F}_a)|_{C=e^{K_0/6}, \, p^2=\mu^2}, \tag{2.5}$$

where $F^A = (F^C, F^I)$, $\partial_A = (\partial_C, \partial_I)$, $C = e^{K_0/6}$ corresponds to the Einstein frame condition, and

$$\frac{F^C}{C} = m_{3/2}^* + \frac{1}{3} F^I \partial_I K_0. \tag{2.6}$$

One then finds [22, 23]

$$\frac{1}{g_a^2(\mu)} = \text{Re}(f_a^{(0)}) - \frac{1}{16\pi^2} \left[(3C_a - \sum_i C_a^i) \ln\left(\frac{\Lambda^2}{\mu^2}\right) \right.$$
$$\left. + (C_a - \sum_i C_a^i) K_0 + 2 \sum_i C_a^i \ln Z_i \right] + \frac{1}{8\pi^2} \Omega_a,$$
$$\frac{M_a(\mu)}{g_a^2(\mu)} = F^I \partial_I \mathcal{F}_a + F^C \partial_C \mathcal{F}_a \tag{2.7}$$
$$= F^I \left[\frac{1}{2} \partial_I f_a^{(0)} - \frac{1}{8\pi^2} \sum_i C_a^i F^I \partial_I \ln(e^{-K_0/3} Z_i) + \frac{1}{8\pi^2} \partial_I \Omega_a \right]$$
$$- \frac{1}{16\pi^2} \left(3C_a - \sum_i C_a^i \right) \frac{F^C}{C}.$$

Note that M_a/g_a^2 do *not* run at one loop level, i.e. are independent of μ, as M_a and g_a^2 have the same running behavior in the one-loop approximation.

However, depending upon the SUSY breaking scenario, the ratios M_a/g_a^2 can receive important threshold corrections at lower intermediate threshold scales M_{th}. In fact, the expression for M_a/g_a^2 in (2.7) is valid only for the renormalization point between the high scale Λ and and the intermediate scale M_{th} where some of the particles decouple. Let us now consider how M_a/g_a^2 are modified by such threshold effects at lower scale. To see this, we assume

$$\{Q_i\} \equiv \{\Phi + \Phi^c, Q_x\}, \tag{2.8}$$

and $\Phi + \Phi^c$ get a supersymmetric mass of the order of M_{th}, while Q_x remain to be massless at M_{th}. Then $\Phi + \Phi^c$ can be integrated out to derive the low energy parameters at scales

below M_{th}. The relevant couplings of $\Phi + \Phi^c$ at M_{th} can be written as

$$\int d^4\theta C C^* e^{-K_0/3} \left(Z_\Phi \Phi^* \Phi + Z_{\Phi^c} \Phi^{c*} \Phi^c \right) + \left(\int d^2\theta\, C^3 \lambda_\Phi X_\Phi \Phi^c \Phi + \mathrm{h.c} \right), \qquad (2.9)$$

where X_Φ is assumed to have a vacuum value[1]

$$\langle X_\Phi \rangle = M_\Phi + \theta^2 F^{X_\Phi}. \qquad (2.10)$$

Then the physical mass of $\Phi + \Phi^c$ are given by

$$\mathcal{M}_\Phi = \lambda_\Phi \frac{C X_\Phi}{\sqrt{e^{-2K_0/3} Z_\Phi Z_{\Phi^c}}}, \qquad (2.11)$$

yielding a threshold correction to the gauge coupling superfield $\mathcal{F}_a$ as

$$\Delta \mathcal{F}_a(M_{\mathrm{th}}) = -\frac{1}{8\pi^2} \sum_\Phi C_a^\Phi \ln\left(\frac{\mathcal{M}_\Phi \mathcal{M}_\Phi^*}{M_{\mathrm{th}}^2} \right). \qquad (2.12)$$

For $M_{\mathrm{th}} \sim \mathcal{M}_\Phi$, this gives rise to a threshold correction of $\mathcal{O}(1/8\pi^2)$ to $1/g_a^2$. In the leading log approximation for gauge couplings, such threshold corrections can be ignored, therefore $1/g_a^2$ obeys the continuity condition at M_{th}:

$$\frac{1}{g_a^2(M_{\mathrm{th}}^+)} = \frac{1}{g_a^2(M_{\mathrm{th}}^-)}, \qquad (2.13)$$

where $M_{\mathrm{th}}^\pm$ denote the scale just above/below M_{th}. On the other hand, because F^I, F^C and F^{X_Φ} can be quite different from each other, the threshold correction to gaugino masses at M_{th} can provide an important contribution to low energy gaugino masses. For $\Delta \mathcal{F}_a$ given above, one easily finds that the threshold correction to gaugino masses at M_{th} is given by

$$M_a(M_{\mathrm{th}}^-) - M_a(M_{\mathrm{th}}^+) = g_a^2(M_{\mathrm{th}}) F^A \partial_A \Delta \mathcal{F}_a \qquad (2.14)$$

$$= -\frac{g_a^2(M_{\mathrm{th}})}{8\pi^2} \sum_\Phi C_a^\Phi \left(\frac{F^C}{C} + \frac{F^{X_\Phi}}{M_\Phi} - F^I \partial_I \ln(e^{-2K_0/3} Z_\Phi Z_{\Phi^c}) \right)$$

Adding this threshold correction to the result of (2.7), we find

$$\left(\frac{M_a}{g_a^2} \right)_{M_{\mathrm{th}}^-} = F^I \left[\frac{1}{2} \partial_I f_a^{(0)} - \frac{1}{8\pi^2} \sum_x C_a^x F^I \partial_I \left(e^{-K_0/3} Z_x \right) + \frac{1}{8\pi^2} \partial_I \Omega_a \right]$$

$$- \frac{1}{8\pi^2} \sum_\Phi C_a^\Phi \frac{F^{X_\Phi}}{M_\Phi} - \frac{1}{16\pi^2} (3 C_a - \sum_x C_a^x) \frac{F^C}{C}, \qquad (2.15)$$

where $\sum_x$ denotes the summation over $\{Q_x\}$ which remain as light matter fields at M_{th}^-.

One can repeat the above procedure, i.e. run down to the lower threshold scale, integrate out the massive fields there, and then include the threshold correction to gaugino masses until one arrives at TeV scale. Then one finally finds

$$\left(\frac{M_a}{g_a^2} \right)_{\mathrm{TeV}} = \tilde{M}_a^{(0)} + \tilde{M}_a^{(1)}|_{\mathrm{anomaly}} + \tilde{M}_a^{(1)}|_{\mathrm{gauge}} + \tilde{M}_a^{(1)}|_{\mathrm{string}} \qquad (2.16)$$

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[1]If X_Φ is not a superfield, but a parameter, then F^{X_Φ} is obviously zero.

where

$$\tilde{M}_a^{(0)} = \frac{1}{2} F^I \partial_I f_a^{(0)},$$

$$\tilde{M}_a^{(1)}|_{\text{anomaly}} = \tilde{M}_a^{(1)}|_{\text{conformal}} + \tilde{M}_a^{(1)}|_{\text{Konishi}}$$

$$= \frac{1}{16\pi^2} b_a \frac{F^C}{C} - \frac{1}{8\pi^2} \sum_m C_a^m F^I \partial_I \ln(e^{-K_0/3} Z_m),$$

$$\tilde{M}_a^{(1)}|_{\text{gauge}} = -\frac{1}{8\pi^2} \sum_\Phi C_a^\Phi \frac{F^{X_\Phi}}{M_\Phi},$$

$$\tilde{M}_a^{(1)}|_{\text{string}} = \frac{1}{8\pi^2} F^I \partial_I \Omega_a. \tag{2.17}$$

Here $\sum_m$ denotes the summation over the light matter multiplets $\{Q_m\}$ at the TeV scale, $\sum_\Phi$ denotes the summation over the gauge messenger fields $\Phi + \Phi^c$ which have a mass lighter than Λ but heavier than TeV, and

$$b_a = -3C_a + \sum_m C_a^m \tag{2.18}$$

are the one-loop beta-function coefficients at TeV. Obviously, $\tilde{M}_a^{(0)}$ is the tree level value of M_a/g_a^2, $\tilde{M}_a^{(1)}|_{\text{conformal}}$ is the SUGRA compensator-mediated one-loop contribution determined by the conformal anomaly of the effective theory at TeV scale [9], $\tilde{M}_a^{(1)}|_{\text{Konishi}}$ is a piece determined by the Konishi anomaly [25], $\tilde{M}_a^{(1)}|_{\text{gauge}}$ are field theoretic gauge thresholds due to massive particles between Λ and the TeV scale, and finally $\tilde{M}_a^{(1)}|_{\text{string}}$ includes the (UV-sensitive) string and/or KK thresholds at scales above Λ as well as the (scheme-dependent) field theoretic one-loop piece $\frac{1}{8\pi^2} C_a F^I \partial_I \ln[\text{Re}(f_a^{(0)})]$.

Depending upon the SUSY breaking scenario, M_a/g_a^2 are dominated by some of these five contributions. The stringy and KK thresholds encoded in $\frac{1}{8\pi^2} \Omega_a$ are most difficult to compute and highly model-dependent. In fact, this represents a potentially uncontrollable contribution from high energy modes. If this part gives an important contribution to M_a/g_a^2, no model independent statements about the gaugino masses can be made. On the other hand, the other parts can be reliably computed within the framework of 4D effective theory under a reasonable assumption in many SUSY breaking scenarios. Note that the anomaly-related contribution $\tilde{M}_a^{(1)}|_{\text{anomaly}}$ is determined by the matter contents at the TeV scale, while the other pieces require a knowledge of physics at scales higher than TeV.

We stress that (2.16) is valid independently of the matter content at scales above TeV, and thus is valid irrespective of whether gauge couplings are unified or not at the initial cutoff scale Λ. Also Λ does not have to be close to $M_{\text{GUT}} \simeq 2 \times 10^{16}$ GeV. The result (2.16) can be applied for the models with Λ hierarchically lower than M_{GUT}.

Formulae (2.16) and (2.17) give the most general description of gaugino masses and its origin from the underlying schemes. This is our basic tool to analyse potential candiates from future collider experiments. The SM gauge coupling constants at TeV have been measured with the (approximate) result:

$$g_1^2 : g_2^2 : g_3^2 \simeq 1 : 2 : 6. \tag{2.19}$$

Once the gaugino mass ratios at TeV are measured, the ratios of M_a/g_a^2 at TeV can be experimentally determined, which will allow us to rule out many SUSY breaking scenarios yielding M_a/g_a^2 different from the experimental values. In the next section, we list the result of low energy gaugino mass ratios for a variety of specific SUSY breaking scenarii.

3. Specific examples

3.1 mSUGRA pattern

3.1.1 Gravity mediation

The scheme of gravity mediation [2] with a universal gaugino mass at M_{GUT} is one of the most popular scenarii whose phenomenological consequences have been studied extensively under the name of mSUGRA scenario. In this scheme, M_a/g_a^2 are assumed to be universal at M_{GUT}, leading to the mSUGRA pattern of gaugino masses at the TeV scale:

$$M_1 : M_2 : M_3 \simeq 1 : 2 : 6. \tag{3.1}$$

In the language of 4D effective SUGRA discussed in the previous section, this amounts to assuming that the cutoff scale Λ of 4D effective SUGRA is close to M_{Pl} or M_{GUT}, and M_a/g_a^2 of eq. (2.16) are dominated by the contribution determined by the tree-level gauge kinetic function: $\tilde{M}_a^{(0)} = \frac{1}{2}F^I\partial_I f_a^{(0)}$, which is assumed (or predicted) to be universal. Some interesting examples of such scenario include dilaton/moduli mediation in heterotic string/M-theory [4], flux-induced SUSY breakdown in Type IIB string theory [5], and gaugino mediation realized in brane models [6]. In the following, we provide a brief sketch of dilaton/moduli mediation in heterotic string/M-theory and flux-induced SUSY breakdown in Type IIB string theory.

Dilaton/moduli-mediated SUSY breakdown in heterotic string/M-theory: the underlying UV theory of this scheme is the 11D Horava-Witten theory [26]. At tree-level of Horava-Witten theory compactified on $CY \times S^1/Z_2$, the gauge kinetic functions of the visible gauge fields take a universal form:

$$f_a^{(0)} = S + \sum_i \beta_i T_i, \tag{3.2}$$

where $\mathrm{Re}(S)$ and $\mathrm{Re}(T_i)/[\mathrm{Re}(S)]^{1/3}$ $(i = 1, \ldots, h_{1,1})$ are proportional to the CY volume and the length of the 11-th interval, respectively, measured in 11D SUGRA unit, and β_i are topological numbers of order unity given by [27, 28]

$$\beta_i = \frac{1}{8\pi^2} \int \omega_i \wedge (F \wedge F - \frac{1}{2}R \wedge R), \tag{3.3}$$

where ω_i denote the basis of harmonic two-forms on CY. In the region of moduli space in which

$$\mathrm{Re}(S) = \mathcal{O}(1), \quad \mathrm{Re}(T_i) = \mathcal{O}(1) \tag{3.4}$$

for the normalization of S and T determined by the above form of $f_a^{(0)}$, the 11D SUGRA description provides a reliable approximation for the UV theory and the corresponding compactification scale is close to $M_{\rm GUT} \sim 2 \times 10^{16}$ GeV [28]. Under the assumption of dilaton/moduli domination, $F^S/(S+S^*)$ and/or $F^i/(T_i+T_i^*)$ have a vacuum value of $\mathcal{O}(m_{3/2})$. Then M_a/g_a^2 in eq. (2.16) are dominated by the *universal* tree-level contribution [4]:

$$\tilde{M}_a^{(0)} = \frac{1}{2} F^I \partial_I f_a^{(0)} = \frac{1}{2} \left(F^S + \sum_i \beta_i F^i \right) \tag{3.5}$$

since the other parts give a subleading contribution of $\mathcal{O}(m_{3/2}/8\pi^2)$. Obviously, then the resulting low energy gaugino masses take the mSUGRA pattern. Note that in this scenario the gauge coupling unification at $M_{\rm GUT}$ is predicted by the universal form of $f_a^{(0)}$, and the universal gaugino masses at $M_{\rm GUT}$ is also an automatic consequence of the dilaton/moduli-dominated mediation scheme.

We note that as long as $F^S/(S+S^*) = \mathcal{O}(m_{3/2})$, the same conclusion applies also for the perturbative heterotic string limit in which

$$\mathrm{Re}(S) \simeq 2, \quad \mathrm{Re}(T_i) = \mathcal{O}\left(\frac{1}{4\pi}\right). \tag{3.6}$$

In this region of moduli space, the underlying UV theory corresponds to the weakly coupled 10D heterotic string theory for which the tree level gauge kinetic functions are given by $f_a^{(0)} = S$. For M_a/g_a^2 in eq. (2.16), the contribution $\frac{1}{2}\sum_i \beta_i F^i$ which was identified as a part of the tree-level contribution $\tilde{M}_a^{(0)} = \frac{1}{2} F^I \partial_I f_a^{(0)}$ in the heterotic M-theory limit should be considered as a part of the string loop contribution $\tilde{M}_a^{(1)}|_{\rm string} = \frac{1}{8\pi^2} F^I \partial_I \Omega_a$ in the perturbative heterotic string limit. For the case that $F^S/(S+S^*) = \mathcal{O}(m_{3/2})$, M_a/g_a^2 in the perturbative heterotic string limit are dominated again by the *universal* tree-level contribution $\tilde{M}_a^{(0)} = \frac{1}{2} F^S$ since

$$\sum_i \beta_i F^i \lesssim \frac{\mathrm{Re}(T_i)}{\mathrm{Re}(S)} F^S \sim \frac{F^S}{8\pi}, \tag{3.7}$$

and as a result the scheme leads to the mSUGRA pattern of low energy gaugino masses [3]. On the other hand, in case that $F^S/(S+S^*) = 0$ while $F^i/(T_i+T_i^*) = \mathcal{O}(m_{3/2})$, i.e. moduli domination scenario, M_a/g_a^2 are determined mainly by $\tilde{M}_a^{(1)}|_{\rm anomaly}$ and $\tilde{M}_a^{(1)}|_{\rm string}$ which are generically non-universal, thus should be discussed separately.

Flux-induced SUSY breakdown in Type IIB string theory: in Type IIB string theory, the 3-form flux takes an imaginary self-dual (ISD) value as a consequence of the equations of motion. Such ISD flux contributes to the F-components of Kähler moduli T_i, while giving a vanishing F-component for the dilaton and complex structure moduli [5]. More explicitly, one finds

$$\begin{aligned}
\frac{F^i}{T_i + T_i^*} &= -\frac{e^{K/2}}{T_i + T_i^*} \sum_j K^{i\bar{j}} (D_j W_{\rm flux})^* \\
&= -e^{K/2} W_{\rm flux}^* \frac{\sum_j K^{i\bar{j}} \partial_{\bar{j}} K}{T_i + T_i^*}
\end{aligned} \tag{3.8}$$

for the flux-induced superpotential

$$W_{\text{flux}} = \int (F_3 - iSH_3) \wedge \Omega, \tag{3.9}$$

where F_3 and H_3 are the RR and NS-NS fluxes, respectively, and Ω is the holomorphic $(3,0)$-form of CY. At leading order in the α'-expansion, the Kähler potential of Kähler moduli obeys the no-scale relation

$$\sum_i (T_i + T_i^*)\partial_i K = -3, \quad \sum_j (T_j + T_j^*)K_{i\bar{j}} = -\partial_i K \tag{3.10}$$

leading to *universal* Kähler moduli F-components together with a vanishing F-component of the SUGRA compensator:

$$\frac{F^i}{T_i + T_i^*} = m_{3/2}^*, \quad \frac{F^C}{C} = m_{3/2}^* + \frac{1}{3}F^i\partial_i K = 0. \tag{3.11}$$

In fact, the above results of flux-induced SUSY breakdown have been obtained in the limit in which T_i are not stabilized yet. It has been suggested that the overall volume modulus T can be stabilized by a competition between two small perturbative corrections to K, while keeping (approximately) the no-scale structure [29]. Including the relevant higher order corrections, the Kähler potential of T takes the form:

$$K = -3\ln(T + T^*) + \frac{\xi_1}{(T + T^*)^{3/2}} - \frac{\xi_2}{(T + T^*)^2}, \tag{3.12}$$

where ξ_1 is the coefficient of higher order α' correction and ξ_2 is the coefficient of string loop corrections.[2] If $\xi_1 > 0$ and $\xi_2 > 0$, one finds that T is stabilized with $F^T/(T + T^*) \simeq m_{3/2}^*$ and $|F^C/C| \ll |m_{3/2}|$ [29, 30]. However $\xi_1 > 0$ requires the Euler number $\chi = 2(h_{1,1} - h_{2,1})$ of the underlying CY orientifold to be positive as well. On the other hand, most of interesting CY compactifications have nonzero $h_{2,1}$. In particular, if one wishes to have a landscape of flux vacua which might contain a state with nearly vanishing cosmological constant, one typically needs $h_{2,1} = \mathcal{O}(100)$ to accommodate a sufficient number of independent 3-form fluxes. At the moment, it is unclear how the above perturbative scheme of volume modulus stabilization can be extended to the case with $h_{1,1} > 1$, while keeping the flux-induced pattern of SUSY breakdown maintained. It remains to be seen whether such a scheme can be realized.

If such scheme of Kähler moduli stabilization exists, an interesting feature of this flux-induced SUSY breakdown is that the resulting $F^i/(T_i + T_i^*)$ have *universal* vacuum values at leading order in the α'-expansion. To proceed, let us suppose that the visible sector gauge fields live on $D7$ branes. Then at leading order in the α'-expansion, the visible sector gauge kinetic functions are generically given by

$$f_a^{(0)} = \sum_i k_{ai}T_i, \tag{3.13}$$

---

[2]In fact, string loop corrections to K includes also a term of the form $X(S + S^*, \mathcal{Z}, \mathcal{Z}^*)/(T + T^*)$, where X is a function of the string dilaton S and complex structure moduli $\mathcal{Z}$. After S and $\mathcal{Z}$ are fixed by flux, X can be treated as a constant in the effective theory of T. Then this correction of $\mathcal{O}(1/(T + T^*))$ can be absorbed into a field redefinition $T + T^* \to T + T^* - \frac{1}{3}X$, after which K takes the form of (3.12).

where k_{ai} are discrete numbers of order unity. Applying the universality of $F^i/(T_i + T_i^*)$ and the vanishing F^C to M_a/g_a^2 in eq. (2.16) for this form of $f_a^{(0)}$, one easily finds

$$\left(\frac{M_a}{g_a^2}\right)_{\text{TeV}} \simeq \frac{m_{3/2}}{g_a^2(\Lambda)}, \tag{3.14}$$

where $1/g_a^2(\Lambda) = \sum_i k_{ai}\text{Re}(T_i)$ for the cutoff scale Λ of 4D effective theory. Note that the gaugino masses at Λ are universal irrespective of the values of k_{ai}, i.e. irrespective of whether or not the gauge couplings are unified at Λ. For $\text{Re}(T_i) = \mathcal{O}(1)$, Λ is close to $M_{\text{GUT}} \sim 2 \times 10^{16}\,\text{GeV}$, and then it is reasonable to assume that g_a^2 are unified at Λ: $g_a^2(\Lambda) \simeq g_{\text{GUT}}^2$. Under this assumption of gauge coupling unification, M_a/g_a^2 at the TeV scale are (approximately) universal, and thereby the low energy gaugino masses take the mSUGRA pattern.

3.1.2 Gauge mediation

In the gauge mediation scenario [7], SUSY breakdown is assumed to be mediated dominantly by the loops of gauge-charged messenger fields $\Phi + \Phi^c$ at a threshold scale M_Φ well above TeV, so M_a/g_a^2 in eq. (2.16) are dominated by the gauge-threshold contribution:

$$\left(\frac{M_a}{g_a^2}\right)_{\text{TeV}} \simeq \tilde{M}_a^{(1)}|_{\text{gauge}} = -\frac{1}{8\pi^2} \sum_\Phi C_a^\Phi \frac{F^{X_\Phi}}{M_\Phi}. \tag{3.15}$$

This gauge threshold correction to M_a^2/g_a^2 accompanies additional running of gauge coupling constants over the scales between M_{GUT} and M_Φ:

$$\Delta\left(\frac{1}{g_a^2}\right) = -\frac{1}{8\pi^2} \sum_\Phi C_a^\Phi \ln\left(\frac{M_{\text{GUT}}^2}{M_\Phi^2}\right). \tag{3.16}$$

In order to maintain the successful gauge coupling unification of the MSSM, one usually assumes that the messenger fields $\Phi + \Phi^c$ form a full SU(5) multiplet, for which $\tilde{M}_a^{(1)}|_{\text{gauge}}$ are universal. The resulting low energy gaugino masses then take the mSUGRA pattern, as a result of the assumption of gauge coupling unification at M_{GUT}.

3.1.3 Large volume compactification of Type IIB string theory

In models with an exponentially large compactification volume, the messenger scale of SUSY breakdown is around the string scale M_{st} which is hierarchically lower than the 4D Planck scale. At this moment, the only known example of moduli stabilization giving a large compactification volume is the model of [8] based on the following form of moduli Kähler potential and superpotential:

$$K = -2\ln\left((T_b + T_b^*)^{3/2} - (T_s + T_s^*)^{3/2} + \xi\right),$$
$$W = w_0 + Ae^{-aT_s}, \tag{3.17}$$

where T_b is the large 4-cycle Kähler modulus for which the bulk CY volume is given by $V_{\text{CY}} \sim (T_b + T_b^*)^{3/2}$ in the string unit with $M_{\text{st}} = 1$, T_s is the Kähler modulus of small

4-cycle wrapped by D7 branes on which the visible fields are assumed to live, and w_0 is assumed to be of order unity. In the limit $\text{Re}(T_b) \gg \text{Re}(T_s)$, one then finds

$$e^{aT_s} \sim (T_b + T_b)^{3/2}, \tag{3.18}$$

thus an exponentially large volume when $aT_s \gg 1$, and also the following pattern of F-components [31]:

$$
\begin{aligned}
\frac{F^{T_b}}{T_b + T_b} &= m_{3/2}\left[1 + \mathcal{O}\left(\frac{1}{(T_b + T_b^*)^{3/2}}\right)\right], \\
\frac{F^{T_s}}{T_s + T_s^*} &\simeq \frac{m_{3/2}}{\ln(M_{\text{Pl}}/m_{3/2})}, \\
\frac{F^C}{C} &= m_{3/2} + \frac{1}{3}F^I \partial_I K_0 = \mathcal{O}\left(\frac{m_{3/2}}{(T_b + T_b^*)^{3/2}}\right), \\
F^S &= F^U = 0,
\end{aligned}
\tag{3.19}
$$

where S and U denote the string dilaton and complex structure moduli, respectively. The string scale, 4D Planck scale and $m_{3/2}$ are related as [8]

$$\frac{M_{\text{st}}^2}{M_{\text{Pl}}^2} \sim \frac{m_{3/2}}{M_{\text{Pl}}} \sim \frac{1}{(T_b + T_b^*)^{3/2}}, \tag{3.20}$$

and $m_{3/2} \sim 10\,\text{TeV}$ (which would give the sparticle mass $m_{\text{soft}} \sim m_{3/2}/\ln(M_{\text{Pl}}/m_{3/2}) \sim 1\,\text{TeV}$) can be obtained for $(T_b + T_b^*)^{3/2} \sim 10^{14}$ giving $M_{\text{st}} \sim 10^{11}\,\text{GeV}$ [8].

In the above F^{T_b} originates mostly from the 3-form flux generating w_0 in the superpotential. On the other hand, F^{T_s} receives contributions from both the flux and the non-perturbative dynamics generating Ae^{-aT_s}, which dynamically cancel each other, making F^s suppressed by $1/\ln(M_{\text{Pl}}/m_{3/2})$. As the Kähler potential of the large volume modulus T_b takes the no-scale form:

$$K_0 = -3\ln(T_b + T_b^*) + \mathcal{O}(1/(T_b + T_b)^{3/2}), \tag{3.21}$$

$m_{3/2}$ in F^C/C is cancelled by $\frac{1}{3}F^{T_b}\partial_{T_b}K_0 = -m_{3/2}$, making F^C/C negligibly small.

Obviously, the visible sector gauge fields can not originate from $D7$ branes wrapping the large 4-cycle since the corresponding gauge couplings are too weak. For the gauge fields on $D7$ branes wrapping the small 4-cycle, the tree-level gauge kinetic functions[3] are given by [31]

$$f_a^{(0)} = T_s + h_a S. \tag{3.22}$$

The corresponding gauge coupling superfields $\mathcal{F}_a$ (see eq. (2.4)) should be independent of the large 4-cycle volume $\text{Re}(T_b) \sim 10^9$ [31], i.e.

$$\partial_{T_b}(e^{K_0/3}Z_i) = \partial_{T_b}\Omega_a = 0, \tag{3.23}$$

---

[3]Such a scheme with $M_{\text{st}} \sim 10^{11}$ GeV cannot accommodate (and is thus not constrained by) conventional gauge coupling unification. The normalization of $\text{U}(1)_Y$ charge is model-dependent and the $\text{U}(1)_Y$ gauge kinetic function can be generalized to $f_Y = k_Y T_s + h_Y S$, where k_Y is a generic rational number. In such a case, the predicted Bino mass M_1 should be multiplied by k_Y.

so F^{T_b} does not participate in the generation of the visible sector gaugino masses. Then, for the SUSY breaking pattern of large volume compactification summarized above, M_a/g_a^2 are dominated by the contribution from the *universal* tree-level gauge kinetic function:

$$\left(\frac{M_a}{g_a^2}\right)_{\text{TeV}} = \frac{1}{2}F^I \partial_I f_a^{(0)} = \frac{1}{2}F^{T_s}, \qquad (3.24)$$

leading to the mSUGRA pattern of low energy gaugino masses:

$$M_1 : M_2 : M_3 \simeq 1 : 2 : 6. \qquad (3.25)$$

We stress that this result is obtained independently of the value of the 4D cutoff scale $\Lambda \sim M_{\text{st}}$ at which $\text{Re}(f_a^{(0)}) \simeq 1/g_a^2(\Lambda)$, and also of whether g_a^2 are unified or not at M_{st}. For instance, for the case with $M_{\text{st}} \sim 10^{11}$ GeV, $g_a^2(M_{\text{st}})$ might be unified or not, depending upon whether or not there exist additional matter fields other than those in the MSSM at scales between TeV and M_{st}. The large volume compactification predicts that M_a/g_a^2 are (approximately) universal regardless of the gauge coupling unification near M_{st} and also of the matter contents above TeV. We also have

$$g_1^2 : g_2^2 : g_3^2 \simeq 1 : 2 : 6, \qquad (3.26)$$

and thus the mSUGRA pattern of gaugino masses at TeV, regardless of whether there exist extra matter states and also of whether g_a^2 are unified at M_{st}.

3.2 Mirage pattern

In many compactifications of Type IIB theory with an explicit scheme of Kähler moduli stabilization, SUSY is broken *not* dominantly by flux, but by other effects such as the uplifting mechanism. For instance, in KKLT stabilization of Kähler moduli T_i [13], the flux-induced SUSY breaking $\propto D_i W_{\text{flux}} = (\partial_i K)W_{\text{flux}}$ is *dynamically cancelled* by the non-perturbative SUSY breaking $\propto D_i W_{\text{np}}$ stabilizing T_i, thereby yielding a SUSY-preserving solution $D_i(W_{\text{flux}} + W_{\text{np}}) = 0$ before the uplifting potential is taken into account. Such a dynamical cancellation of flux-induced SUSY breaking by moduli-stabilizing dynamics appears to be a somewhat generic feature of moduli stabilization which admits a supersymmetric configuration satisfying $D_i W = 0$ since the admitted supersymmetric configuration is always a stationary solution of the SUGRA scalar potential.[4]

Here we present three examples of SUSY breaking by uplifting dynamics, the KKLT stabilization [13], a partial-KKLT stabilization [15] and an uplift by F-terms of matter superpotentials [14], all of which give rise to the mirage pattern of low energy gaugino masses [10, 16] (but not necessarily a mirage pattern for sfermion masses) if the visible sector is assumed to live on $D7$ branes. We also present another example of SUSY-breaking scenario, the deflected anomaly mediation [17], which gives the same gaugino mass pattern although it is quite different from the above three schemes in other aspects.

---

[4] Of course, the perturbative stabilization of T by the Kähler potential (3.12) is an exception since it does not admit supersymmetric configuration in a region where the Kähler potential (3.12) is reliable [29]. Note however that such perturbative stabilization can be applied only for a quite limited situation as we have noticed in the previous subsection.

3.2.1 KKLT with visible sector on D7 branes

KKLT stabilization of Kähler moduli T_i in Type IIB string theory is based on the following features of fluxed compactification: i) bulk geometry contains a warped throat generated by flux, ii) SUSY is broken by a brane-localized source stabilized at the IR end of warped throat, iii) there exist non-perturbative dynamics, e.g. hidden gaugino condensation or D-brane instantons, generating a non-perturbative superpotential of the form $W_{\mathrm{np}} = \sum_i A_i e^{-a_i T_i}$. In this set-up, SUSY-breaking at the IR end of the throat is sequestered from the Kähler moduli and visible sector fields living in the bulk CY located at the UV end of the throat, realizing the gravity dual of the conformal sequestering scenario [11, 12]. After integrating out the heavy dilaton and complex structure moduli as well as the massive degrees of freedom on SUSY-breaking brane, the effective action of Kähler moduli is given by

$$\mathcal{L}_{\mathrm{eff}} = \int d^4\theta \left[-3CC^* \exp\left(-\frac{1}{3}K(T_i + T_i^*)\right) - \theta^2\bar{\theta}^2 C^2 C^{*2}\mathcal{P}_0 \right] + \left(\int d^2\theta C^3 W + \mathrm{h.c.} \right), \tag{3.27}$$

where C is the SUGRA compensator, the moduli Kähler potential K obeys the no-scale condition (3.10) at leading order in the α'-expansion, and

$$W = w_0 + \sum_i A_i e^{-a_i T_i},$$
$$\mathcal{P}_0 = \mathrm{constant}, \tag{3.28}$$

where w_0 is a constant including the flux-induced contribution. The uplifting operator $\theta^2\bar{\theta}^2\mathcal{P}_0$ represents the low energy consequence of the *sequestered* SUSY-breaking brane, and thereby $\mathcal{P}_0$ is independent of T_i and visible matter fields.

In the Einstein frame with $C = e^{K/6}$, the above effective lagrangian gives the Kähler moduli potential:

$$V_{\mathrm{TOT}} = V_F + V_{\mathrm{lift}} = e^K \left(K^{i\bar{j}} D_i W (D_j W)^* - 3|W|^2 \right) + e^{2K/3}\mathcal{P}_0, \tag{3.29}$$

where V_F is the conventional $N = 1$ SUGRA potential and $V_{\mathrm{lift}} = e^{2K/3}\mathcal{P}_0$ is the uplifting potential induced by SUSY-breaking brane. The Kähler moduli vacuum values and their F-components can be computed by minimizing V_{TOT} under the condition $\langle V_{\mathrm{TOT}} \rangle = 0$. This can be done in two steps: first minimize V_F yielding a supersymmetric AdS solution of $D_i W = 0$, and next compute the small vacuum shifts δT_i caused by V_{lift} in a perturbative expansion in powers of $1/\ln(M_{\mathrm{Pl}}/m_{3/2})$ [10, 16, 32]. The first step yields the following moduli vacuum values:

$$a_i T_i = \ln(M_{\mathrm{Pl}}/m_{3/2}) \left[1 + \mathcal{O}\left(\frac{1}{\ln(M_{\mathrm{Pl}}/m_{3/2})} \right) \right], \tag{3.30}$$

where the A_i of (3.28) have been assumed to be of order unity. As the Kähler moduli masses at this supersymmetric solution are significantly heavier than $m_{3/2}$,

$$m_{T_i} = \mathcal{O}\left(m_{3/2} \ln(M_{\mathrm{Pl}}/m_{3/2}) \right), \tag{3.31}$$

while V_{lift} has a vacuum value of $\mathcal{O}(m_{3/2}^2)$ (in the unit with $M_{\text{Pl}} = 1$), the vacuum shift induced by V_{lift} is as small as

$$\delta T_i \equiv \delta\phi_i + i\delta a_i = \mathcal{O}\left(\frac{1}{[\ln(M_{\text{Pl}}/m_{3/2})]^2}\right). \tag{3.32}$$

This makes the computation of the moduli F-components in KKLT stabilization rather straightforward.

To proceed, let us choose the field basis for which w_0 and A_i take a real value. Since V_{lift} is a real function of real variables $T_i + T_i^*$, it is obvious that V_{lift} does not induce a tadpole of the axion component: $\delta a_i = 0$. To compute $\delta\phi_i$, let us expand V_{TOT} around the supersymmetric solution $\vec{T}_0$ of $D_i W = 0$:

$$V_{\text{TOT}} = -3m_{3/2}^2(\vec{T}_0) + V_{\text{lift}}(\vec{T}_0) + 2\partial_i V_{\text{lift}}(\vec{T}_0)\delta\phi_i + \left(m_\phi^2\right)_{ij}\delta\phi_i\delta\phi_j + \ldots, \tag{3.33}$$

where $\partial_i V_{\text{lift}} = \partial V_{\text{lift}}/\partial T_i$, and

$$\left(m_\phi^2\right)_{ij} = m_{3/2}^2 a_i a_j K^{i\bar{j}}\partial_i K\partial_{\bar{j}} K\left[1 + \mathcal{O}\left(\frac{1}{\ln(M_{\text{Pl}}/m_{3/2})}\right)\right]. \tag{3.34}$$

Upon ignoring the small corrections further suppressed by $1/\ln(M_{\text{Pl}}/m_{3/2})$, minimizing V_{TOT} under the condition $\langle V_{\text{TOT}}\rangle = 0$ leads to

$$\sum_j \left(m_\phi^2\right)_{ij}\delta\phi_j = -\partial_i V_{\text{lift}} = -\frac{2}{3}V_{\text{lift}}\partial_i K = -2m_{3/2}^2\partial_i K, \tag{3.35}$$

and thus

$$\sum_j a_i a_j K^{i\bar{j}}\partial_{\bar{j}} K\delta\phi_j = -2, \tag{3.36}$$

where we have used that $V_{\text{lift}} = e^{2K/3}\mathcal{P}_0$ for a constant $\mathcal{P}_0$, and also $\langle V_{\text{lift}}\rangle \simeq 3m_{3/2}^2$ as required for $\langle V_{\text{TOT}}\rangle \simeq 0$. Applying the above result and (3.30) to

$$\begin{aligned}F^i &= -e^{K/2}K^{i\bar{j}}(D_j W)^* \\ &= -m_{3/2}\sum_j a_j K^{i\bar{j}}\partial_{\bar{j}} K\delta\phi_j\left[1 + \mathcal{O}\left(\frac{1}{\ln(M_{\text{Pl}}/m_{3/2})}\right)\right],\end{aligned} \tag{3.37}$$

one finds that $F^i/(T_i + T_i^*)$ are *universal* [32] at leading order in the expansion in $1/\ln(M_{\text{Pl}}/m_{3/2})$:

$$\frac{F^i}{T_i + T_i^*} = \frac{m_{3/2}}{\ln(M_{\text{Pl}}/m_{3/2})} \equiv M_0. \tag{3.38}$$

Note that the universality of $F^i/(T_i + T_i^*)$ is obtained independently of the detailed form of the moduli Kähler potential. The F-component of the SUGRA compensator C in KKLT stabilization is given by

$$\frac{F^C}{C} = m_{3/2} + \frac{1}{3}\sum_i F^i\partial_i K \simeq m_{3/2}. \tag{3.39}$$

Applying the above pattern of F-components to M_a/g_a^2 in eq. (2.16) for the case that visible gauge fields live on $D7$ branes, i.e. for the case with

$$f_a^{(0)} = \sum_i k_{ai} T_i, \tag{3.40}$$

we find that M_a/g_a^2 are dominated by the tree-level contribution $\tilde{M}_a^{(0)} = \frac{1}{2} F^i \partial_i f_a^{(0)}$ and the one-loop conformal anomaly contribution $\tilde{M}_a^{(1)}|_{\text{conformal}}$ which are comparable to each other:

$$
\begin{aligned}
\left(\frac{M_a}{g_a^2}\right)_{\text{TeV}} &= \frac{1}{2} \sum_i k_{ai} F^i + \frac{b_a}{16\pi^2} m_{3/2} \\
&= \left(1 + \frac{\ln(M_{\text{Pl}}/m_{3/2})}{16\pi^2} g_{\text{GUT}}^2 b_a\right) \frac{M_0}{g_{\text{GUT}}^2},
\end{aligned}
\tag{3.41}
$$

where we have assumed that g_a^2 are unified at M_{GUT}: $\sum_i k_{ai} \text{Re}(T_i) = 1/g_a^2(M_{\text{GUT}}) = 1/g_{\text{GUT}}^2$, and again the subleading parts suppressed by $1/\ln(M_{\text{Pl}}/m_{3/2})$ are ignored. This corresponds to the mirage pattern (1.6) with $\alpha = 1$, for which the low energy gaugino mass ratios are given by

$$M_1 : M_2 : M_3 \simeq 1 : 1.3 : 2.5. \tag{3.42}$$

We note that a more general form of mirage pattern can be obtained for instance when the gauge kinetic functions are generalized to include higher order correction in the α'-expansion [33], e.g.

$$f_a^{(0)} = \sum_i k_i T_i + hS, \tag{3.43}$$

where we have assumed a *universal* form of f_a for the gauge coupling unification at M_{GUT}. The higher order term hS in $f_a^{(0)}$ can be induced by a magnetic flux on $D7$ branes. Even when the compactification radius R is significantly bigger than the string length scale $l_s = \sqrt{\alpha'}$, so $\text{Re}(S)/\text{Re}(T) \sim \alpha'^2/R^4 \ll 1$, the magnetic flux $h = \frac{1}{8\pi^2} \int F \wedge F$ might have a large value, which would make hS in $f_a^{(0)}$ non-negligible. The inclusion of hS in $f_a^{(0)}$ suggests that the non-perturbative superpotential should be generalized also as

$$W = w_0 + \sum_i A_i e^{-(a_i T_i + b_i S)}, \tag{3.44}$$

where $A_i = \mathcal{O}(1)$, and $b_i S/a_i T_i$ might have a sizable value. For S fixed at $\langle S \rangle = S_0$ by flux and a sequestered uplifting operator with $\mathcal{P}_0 = \text{constant}$, it is straightforward to repeat the previous analysis to compute the vacuum values of F^i and T_i determined by the above generalized form of moduli superpotential. One then finds [33]

$$
\begin{aligned}
a_i T_i + b_i S_0 &= \ln(M_{\text{Pl}}/m_{3/2}) \left[1 + \mathcal{O}\left(\frac{1}{\ln(M_{\text{Pl}}/m_{3/2})}\right)\right], \\
a_i F^i &= 2m_{3/2} \left[1 + \mathcal{O}\left(\frac{1}{\ln(M_{\text{Pl}}/m_{3/2})}\right)\right],
\end{aligned}
\tag{3.45}
$$

which yields

$$\left(\frac{M_a}{g_a^2}\right)_{\text{TeV}} = \frac{1}{2}\sum_i k_i F^i + \frac{1}{16\pi^2} b_a m_{3/2}$$

$$= \left(1 + \frac{\ln(M_{\text{Pl}}/m_{3/2})}{16\pi^2} g_{\text{GUT}}^2 b_a \alpha\right)\frac{M_0}{g_{\text{GUT}}^2}, \tag{3.46}$$

where

$$\frac{1}{g_{\text{GUT}}^2} = \sum_i k_i \text{Re}(T_i) + h\text{Re}(S_0),$$

$$M_0 = F^i \partial_i \ln(\text{Re}(f_a^{(0)})) = \frac{1}{2} g_{\text{GUT}}^2 \sum_i k_i F^i,$$

$$\alpha = \frac{m_{3/2}}{M_0 \ln(M_{\text{Pl}}/m_{3/2})} = \frac{h\text{Re}(S_0) + \sum_i k_i \text{Re}(T_i)}{\sum_i [\frac{k_i b_i}{a_i}\text{Re}(S_0) + k_i \text{Re}(T_i)]}. \tag{3.47}$$

Note that $\alpha \to 1$ in the limit $\text{Re}(S_0)/\text{Re}(T_i) \to 0$, i.e. in the limit when the magnetic flux-induced $\mathcal{O}(\alpha'^2)$ correction in $f_a^{(0)}$ is negligible. At any rate, the above result gives rise to

$$M_1 : M_2 : M_3 \simeq (1 + 0.66\alpha) : (2 + 0.2\alpha) : (6 - 1.8\alpha), \tag{3.48}$$

where now α can take a generic value of order unity.

3.2.2 Partial KKLT stabilization [15]

In the partial-KKLT stabilization scenario, some Kähler moduli (T_i) are stabilized by a non-perturbative superpotential, while the remaining Kähler moduli (X_p) are stabilized by the sequestered uplifting potential. This scenario is an interesting generalization of KKLT stabilization in which one combination of $\text{Im}(X_p)$ can be identified as the QCD axion solving the strong CP problem [15].

Like the case of KKLT stabilization, the SUSY-breaking brane is assumed to be stabilized at the IR end of a warped throat, and thus the resulting uplifting operator $\theta^2\bar{\theta}^2\mathcal{P}_0$ is independent of the entire set of Kähler moduli $T_I = (T_i, X_p)$:

$$K = K_0(T_i + T_i^*, X_p + X_p^*),$$
$$W = w_0 + \sum_i A_i e^{-a_i T_i},$$
$$\mathcal{P}_0 = \text{constant}. \tag{3.49}$$

It is also assumed that the model allows a supersymmetric configuration satisfying

$$D_i W = 0, \quad D_p W = W\partial_p K = 0. \tag{3.50}$$

One simple such example would be

$$K = -2\ln[(\Phi_1 + \Phi_1^*)^{3/2} - (\Phi_2 + \Phi_2^*)^{3/2} - (\Phi_3 + \Phi_3^*)^{3/2}],$$
$$W = w_0 + A_1 e^{-a_1\Phi_1} + A_2 e^{-a_2(\Phi_2+\Phi_3)}, \tag{3.51}$$

for which one can rewrite the effective SUGRA in terms of $T_1 = \Phi_1$, $T_2 = \Phi_2 + \Phi_3$, and $X_1 = \Phi_2 - \Phi_3$. Although X_p are stabilized by the uplifting potential $V_{\text{lift}} = e^{2K/3}\mathcal{P}_0$, while T_i are stabilized by non-perturbative superpotential, it turns out that the $F^I/(T_I + T_I^*)$ are again *universal* for the entire Kähler moduli $T_I = (T_i, X_p)$ as long as the moduli Kähler potential obeys the no-scale relation [15]:

$$\sum_J K^{I\bar{J}}\partial_{\bar{J}}K = -(T_I + T_I^*), \tag{3.52}$$

which is indeed satisfied at leading order in the α'-expansion. As a result, the partial KKLT stabilization also gives rise to the mirage pattern of gaugino masses as the KKLT stabilization, although the details of moduli stabilization are quite different.

Similarly to the KKLT case, one can compute F^I in two steps: first start with the supersymmetric solution and then compute the vacuum shift:

$$\delta T_I = (\delta T_i, \delta X_p) = \delta\phi_I + i\delta a_I \tag{3.53}$$

caused by V_{lift}. Again, for $A_i = \mathcal{O}(1)$, $D_iW = 0$ in the first step determines T_i as

$$a_iT_i = \ln(M_{\text{Pl}}/m_{3/2})\left[1 + \mathcal{O}\left(\frac{1}{\ln(M_{\text{Pl}}/m_{3/2})}\right)\right], \tag{3.54}$$

regardless of the detailed form of K. It is also obvious that V_{lift} does not induce any vacuum shift in the axion direction: $\delta a_I = 0$. As they are stabilized by a non-perturbative super-potential, T_i get a mass of $\mathcal{O}(m_{3/2}\ln(M_{\text{Pl}}/m_{3/2}))$, and thus $\delta\phi_i = \mathcal{O}(1/[\ln(M_{\text{Pl}}/m_{3/2})]^2)$. On the other hand, since W is independent of X_p, $\text{Re}(X_p)$ has a mass of $\mathcal{O}(m_{3/2})$. One might then expect that V_{lift} causes a large vacuum shift in the direction of $\text{Re}(X_p)$, which would result in a breakdown of our perturbative expansion. However, this is not the case since the supersymmetric configuration satisfying $D_pW = W\partial_pK = 0$ is a stationary point of $V_{\text{lift}} = e^{2K/3}\mathcal{P}_0$ for a constant $\mathcal{P}_0$. In fact, $\delta\phi_p$ is induced by a Kähler mixing between $\text{Re}(X_p)$ and $\text{Re}(T_i)$, and as a result $\delta\phi_p$ has the same order of magnitude as $\delta\phi_i = \mathcal{O}(1/[\ln(M_{\text{Pl}}/m_{3/2})]^2)$.

Expanding V_{TOT} around the supersymmetric configuration $\vec{T}_0$, we find

$$V_{\text{TOT}} = -3m_{3/2}^2(\vec{T}_0) + V_{\text{lift}}(\vec{T}_0) + 2\partial_I V_{\text{lift}}(\vec{T}_0)\delta\phi_I + \left(m_\phi^2\right)_{IJ}\delta\phi_I\delta\phi_J + \ldots, \tag{3.55}$$

where the moduli mass matrix is given by

$$\left(m_\phi^2\right)_{ij} = a_ia_jK^{i\bar{j}}\partial_iK\partial_{\bar{j}}Km_{3/2}^2\left[1 + \mathcal{O}\left(\frac{1}{\ln(M_{\text{Pl}}/m_{3/2})}\right)\right],$$

$$\left(m_\phi^2\right)_{pi} = 2m_{3/2}^2K_{p\bar{i}}, \quad \left(m_\phi^2\right)_{pq} = 2m_{3/2}^2K_{p\bar{q}}, \tag{3.56}$$

where we have used $V_{\text{lift}}(\vec{T}_0) \simeq 3m_{3/2}^2$. Minimizing V_{TOT} determines $\delta\phi_I = (\delta\phi_i, \delta\phi_p)$ as

$$\sum_j a_ia_jK^{i\bar{j}}\partial_{\bar{j}}K\delta\phi_j = -2,$$

$$\sum_q K_{p\bar{q}}\delta\phi_q = -\sum_j K_{p\bar{j}}\delta\phi_j, \tag{3.57}$$

– 19 –

which show that $\delta\phi_i$ and $\delta\phi_p$ are all of the order of $1/[\ln(M_{\rm Pl}/m_{3/2})]^2$. For such small vacuum shifts, the moduli F-components are obtained to be

$$
\begin{aligned}
F^i &= -e^{K/2}\sum_J K^{i\bar{J}}(D_J W)^* \\
&= -m_{3/2}\sum_j a_j K^{i\bar{j}}\partial_{\bar{j}}K\delta\phi_j\left[1+\mathcal{O}\left(\frac{1}{\ln(M_{\rm Pl}/m_{3/2})}\right)\right], \\
F^p &= -e^{K/2}\sum_J K^{p\bar{J}}(D_J W)^* \\
&= -m_{3/2}\sum_j a_j K^{p\bar{j}}\partial_{\bar{j}}K\delta\phi_j\left[1+\mathcal{O}\left(\frac{1}{\ln(M_{\rm Pl}/m_{3/2})}\right)\right].
\end{aligned}
\tag{3.58}
$$

Then, using (3.54) and (3.57), we find

$$
\begin{aligned}
\frac{F^i}{T_i+T_i^*} &= \frac{m_{3/2}}{\ln(M_{\rm Pl}/m_{3/2})} \equiv M_0, \\
\sum_q K_{p\bar{q}}F^{q*} &= -\sum_j K_{p\bar{j}}F^{j*},
\end{aligned}
\tag{3.59}
$$

So far, we have *not* used any property of the moduli Kähler potential K. If K obeys the no-scale condition (3.52), the result on F^p can be further simplified. Combining the no-scale condition (3.10) with $\partial_p K(\vec{T}_0)=0$, one easily finds

$$
\sum_j (T_j+T_j^*)K_{p\bar{j}} = -\sum_q (X_q+X_q^*)K_{p\bar{q}}.
\tag{3.60}
$$

Combining this with (3.59), one finds also

$$
\sum_q K_{p\bar{q}}\Big[F^q - (X_q+X_q^*)M_0\Big] = 0,
\tag{3.61}
$$

which finally leads to

$$
\begin{aligned}
\frac{F^p}{X_p+X_p^*} = \frac{F^i}{T_i+T_i^*} &= \frac{m_{3/2}}{\ln(M_{\rm Pl}/m_{3/2})}, \\
\frac{F^C}{C} = m_{3/2}+\frac{1}{3}F^I\partial_I K &\simeq m_{3/2}.
\end{aligned}
\tag{3.62}
$$

Thus, although the details of moduli stabilization in the partial-KKLT scenario are quite different from the KKLT scenario, the resulting pattern of SUSY breakdown is the same, and gives the mirage pattern of low energy gaugino masses for the MSSM gauge fields living on $D7$ branes.

3.2.3 Uplifting via matter superpotentials [14]

The KKLT and partial KKLT scheme assume that the uplifting of the AdS vacuum is achieved by a brane-localized source of SUSY breaking which is sequestered from the volume modulus and visible matter fields. In Type IIB string theory, such sequestered

SUSY breakdown can be naturally realized if the SUSY-breaking brane is stabilized at the IR end of throat, while the visible sector and Kähler moduli live on the bulk CY which corresponds to the UV end of throat.

On the other hand, it has been pointed out that a successful uplifting can be achieved also by a hidden sector realizing the conventional spontaneous breakdown of $N = 1$ SUGRA [14, 34]. In such a scheme the hidden sector for uplifting is not necessarily sequestered from the volume modulus and visible matter fields.

More specifically, the model proposed in [14] is given by

$$
\begin{aligned}
K &= -3\ln(T + T^*) + ZZ^*, \\
W &= \phi(Z) + A(Z)e^{-aT}, \\
f_a &= T,
\end{aligned}
\tag{3.63}
$$

where T is the volume modulus and Z is a hidden matter field with nonzero F^Z induced by a proper matter superpotential $\phi(Z)$. Since the superpotential of T has a non-perturbative origin, there is still the natural possibility that

$$
|\partial_T^2 W| \gg |W|,
\tag{3.64}
$$

realizing

$$
m_T \gg m_{3/2}, \quad \frac{F^T}{T + T^*} \ll m_{3/2}.
\tag{3.65}
$$

On the other hand, the condition of nearly vanishing cosmological constant requires [34]

$$
F^Z = \mathcal{O}(m_{3/2}), \quad \frac{F^C}{C} = \mathcal{O}(m_{3/2}).
\tag{3.66}
$$

In such a case, the conformal anomaly mediation from $F^C/C = \mathcal{O}(m_{3/2})$ and also the one-loop contributions from $F^Z = \mathcal{O}(m_{3/2})$ can become equally important to the tree level contribution to M_a/g_a^2 from $|F^T| \ll m_{3/2}$. In particular, string thresholds might depend on Z, and thus give a UV sensitive contribution to M_a/g_a^2 through the term $\frac{1}{8\pi^2}F^Z\partial_Z\Omega_a$ in (2.16), which could spoil the predictability of the scheme.

However, this potential difficulty can be avoided by making a simple assumption on the property of the hidden matter Z, for instance by assuming a discrete symmetry $Z \to -Z$ which is broken only by non-perturbative dynamics that is responsible for SUSY breakdown. This (approximate) symmetry ensures that the vacuum values of $\partial_Z \ln(e^{-K_0/3}Z_m)$ and $\partial_Z\Omega_a$ are negligibly small, and therefore the one-loop contribution to M_a/g_a^2 is dominated by the conformal anomaly mediation. Note that the above model does not necessarily take a sequestered form: $e^{-K/3} = \Omega_1(T, T^*) + \Omega_2(Z, Z^*)$ and $W = W_1(T) + W_2(Z)$. This has several interesting consequences. First of all, it allows more freedom for the relative importance of anomaly mediation, i.e. a much larger range of values of α is allowed. Secondly, while a mirage pattern is expected for the gaugino masses, this is not necessarily true for the squark and slepton masses as well [14, 34].

To see the pattern of low energy gaugino masses obtainiable by the matter uplifting scenario, let us consider an example with a rather large value of α. It is the example of ref. [14] with a superpotential W expanded around the vacuum configuration $\langle T \rangle = 2$ and $\langle Z \rangle = 0$:

$$W = \epsilon \left[0.577 + Z + 0.441(T-2) + 0.592Z^2 + 9.595(T-2)^2 \right.$$
$$\left. + 0.114Z^3 + 0.220Z^2(T-2) + 46.451(T-2)^3 + \text{higher order terms} \right], \quad (3.67)$$

where ϵ is a small parameter of $\mathcal{O}(m_{3/2})$ which might be generated by nonperturbative dynamics, like e.g. gaugino condensation. This particular example of W stabilizes T and Z at a nearly Minkowski vacuum with

$$F^Z = \frac{\epsilon}{8}, \quad F^T = \frac{\epsilon}{200}, \quad m_{3/2} = \frac{\epsilon}{14}. \quad (3.68)$$

The resulting M_a/g_a^2 take a mirage pattern:

$$\left(\frac{M_a}{g_a^2} \right)_{\text{TeV}} = \frac{1}{2}F^T + \frac{b_a}{16\pi^2}m_{3/2}$$
$$= \left(1 + \frac{\ln(M_{\text{Pl}}/m_{3/2})}{16\pi^2} g_{\text{GUT}}^2 b_a \alpha \right) \frac{M_0}{g_{\text{GUT}}^2}, \quad (3.69)$$

with

$$\alpha \simeq 1.7, \quad (3.70)$$

and thus

$$M_1 : M_2 : M_3 \simeq 1 : 1.1 : 1.4. \quad (3.71)$$

Although the above specific example of W leads to the mirage parameter $\alpha \simeq 1.7$, one can construct different superpotentials giving a different value of α within the range of order unity and smaller. This is in contrast to the KKLT and partial KKLT schemes with minimal gauge kinetic function $f_a = T$, where we obtain $\alpha \simeq 1$. This difference is due to the fact that the SUSY breaking sector is sequestered from T in KKLT and partial KKLT, while it is not in the matter uplifting scenario. As we have noticed, $\alpha \neq 1$ can be obtained also in KKLT and partial KKLT while keeping the SUSY breaking sector sequestered from T if the magnetic flux-induced corrections to f_a are sizable. Note also that a non-sequestered coupling between Z and the visible matter fields can give a contribution of $\mathcal{O}(m_{3/2}^2)$ to the sfermion mass-squared in the matter uplifting scenario [14].

3.2.4 Deflected anomaly mediation [17]

Deflected anomaly mediation [17] has been proposed as a solution to the tachyonic slepton problem of the original anomaly mediation scenario [9]. The scheme assumes a gauge-charged messenger sector which experiences a non-decoupled SUSY breakdown triggered by the F-component of the SUGRA compensator. Such a gauge messenger sector then provides a gauge-mediated contribution to soft masses comparable to anomaly mediation

at the threshold scale of gauge-charged messenger fields. This deflects the soft masses from the anomaly mediation trajectory at scales below the gauge threshold scale, thereby avoiding tachyonic sleptons.

A simple example of such model is provided by the following messenger sector superpotential containing a scale *non-invariant* term:

$$W_{\text{mess}} = \lambda_\Phi X_\Phi \Phi^c \Phi + M_*^{3-n} X_\Phi^n, \tag{3.72}$$

where $n \neq 3$, M_* is a model-dependent mass parameter, X_Φ is a singlet chiral superfield, and finally the gauge-charged $\Phi + \Phi^c$ are assumed to form a GUT representation to maintain the gauge coupling unification at $M_{\text{GUT}} \sim 2 \times 10^{16}$ GeV. The chiral compensator C couples to X_Φ at tree level through the scale non-invariant term $M_*^{3-n} X_\Phi^n$ ($n \neq 3$), thereby fixing the vacuum value of X_Φ. One then finds that X_Φ is stabilized at $\langle X_\Phi \rangle \gg m_{3/2}$ for $M_* \gg m_{3/2}$ if $n < 0$ or $n > 3$ [17], explicitly

$$\langle X_\Phi \rangle = M_\Phi + \theta^2 F^{X_\Phi}, \tag{3.73}$$

where

$$M_\Phi \sim m_{3/2}^{1/(n-2)} M_*^{(n-3)/(n-2)},$$
$$\frac{F^{X_\Phi}}{M_\phi} = -\frac{2}{n-1} \frac{F^C}{C}. \tag{3.74}$$

In fact, even in the absence of the scale non-invariant term $M_*^{n-3} X_\Phi^n$ ($n > 3$ or $n < 0$) in the superpotential, X_Φ still can be stabilized by radiative corrections to its Kähler potential, yielding [17]

$$\frac{F^{X_\Phi}}{M_\phi} \simeq -\frac{F^C}{C}. \tag{3.75}$$

As $F^{X_\Phi}/M_\Phi = \mathcal{O}(F^C/C)$ in the above case, the resulting gauge-mediated contribution to soft masses is comparable to the anomaly-mediated ones. Applying the result of eq. (2.16) to this case, one easily finds

$$\begin{aligned}
\left(\frac{M_a}{g_a^2} \right)_{\text{TeV}} &= \frac{b_a}{16\pi^2} \frac{F^C}{C} - \frac{1}{8\pi^2} \sum_\Phi C_a^\Phi \frac{F^{X_\Phi}}{M_\Phi} \\
&= \frac{1}{16\pi^2} \frac{F^C}{C} \left(b_a + \frac{2}{n-1} N_\Phi \right),
\end{aligned} \tag{3.76}$$

where $n \geq 3$ or $n < 0$, and $N_\Phi = \sum_\Phi (C_a^\Phi + C_a^{\Phi^c})$ is the number of messenger pairs $\Phi + \Phi^c$. This is the mirage pattern of gaugino masses with

$$\alpha = \frac{16\pi^2}{g_{\text{GUT}}^2 \ln(M_{\text{Pl}}/m_{3/2})} \frac{n-1}{2N_\Phi}. \tag{3.77}$$

We note that, for a given value of α, the sfermion spectrum in deflected anomaly mediation [17] takes a different pattern than that of mirage mediation [16], although the gaugino masses share the same mirage pattern.

3.3 Anomaly pattern

In anomaly mediation [9], SUSY breaking fields X^I are assumed to be (effectively) sequestered from the visible sector fields, which means that $F^I \partial_I f_a^{(0)}$, $\frac{1}{8\pi^2} F^I \partial_I (e^{-K_0/3} Z_i)$ and $\frac{1}{8\pi^2} F^I \partial_I \Omega_a$ in eq. (2.16) are all subdominant compared to the SUGRA compensator mediated contribution [11, 12]. Then M_a/g_a^2 are determined as

$$\left(\frac{M_a}{g_a^2} \right)_{\text{TeV}} \simeq \tilde{M}_a^{(1)}|_{\text{conformal}} = \frac{b_a}{16\pi^2} \frac{F^C}{C}. \tag{3.78}$$

If the effective theory around TeV is given by the MSSM, $b_a = (\frac{33}{5}, 1, -3)$ and the low energy gaugino masses take the anomaly pattern:

$$M_1 : M_2 : M_3 \simeq 3.3 : 1 : 9. \tag{3.79}$$

One example of string-based scenario which can give the anomaly pattern of gaugino masses is the fluxed compactification of Type IIB string theory with visible sector living on $D3$ branes [10].

3.3.1 Visible sector on $D3$ branes in Type IIB string theory

In fluxed Type IIB compactification, the ISD 3-form flux fixes the dilaton S and complex structure moduli U at a supersymmetric solution of $D_S W_{\text{flux}} = D_U W_{\text{flux}} = 0$. As the flux-induced masses of S and U are hierarchically heavier than $m_{3/2}$, the F-components of S and U remain to be negligible even after the subsequent SUSY-breaking effect is taken into account:[5]

$$|F^{S,U}| \sim \frac{m_{3/2}^2}{m_{S,U}} \ll \frac{m_{3/2}}{8\pi^2}, \tag{3.80}$$

while the F-components of Kähler moduli T_i can be $\mathcal{O}(m_{3/2}/8\pi^2)$ or bigger, depending upon how T_i are stabilized.

To be specific, let us first consider the KKLT-type stabilization of T_i, yielding the following pattern of F-terms [10]:

$$\frac{F^i}{T_i + T_i^*} = \mathcal{O} \left(\frac{m_{3/2}}{8\pi^2} \right), \quad \frac{F^C}{C} \simeq m_{3/2}. \tag{3.81}$$

For a gauge field on $D3$, the corresponding gauge kinetic function is given by

$$f_{3a} = S \tag{3.82}$$

at leading order in the α' and string loop expansions. The only possible correction to f_{3a} allowed by the axionic shift symmetries takes the form: $\Delta f_{3a} = \epsilon_a T$, where ϵ_a are real

---

[5]In fact, this is not a precise statement for the KKLT scenario. In the KKLT scenario, the complex structure modulus describing the collapsing 3-cycle of the warped throat can develop a large F-component through its direct coupling to the SUSY breaking brane at the IR end of throat. However, such particular complex structure modulus is sequestered from the visible sector $D3$ branes which are assumed to be stabilized somewhere in the bulk CY, and thus does not affect the soft terms of visible fields.

(discrete) constants. However any nonzero value of $\epsilon_a = 0$ can not give a sensible behavior of $D3$ brane gauge coupling in the large volume limit $\mathrm{Re}(T)/\mathrm{Re}(S) \gg 1$.[6] Thus the above form of $D3$ gauge kinetic function is exact up to non-perturbative effects of $\mathcal{O}(e^{-8\pi^2 S})$ or $\mathcal{O}(e^{-8\pi^2 T})$. Applying the F-components in KKLT scenario to M_a/g_a^2 in (2.16) with $f_a = S$, one easily finds that M_a/g_a^2 are dominated by the conformal anomaly contribution $b_a F^C/16\pi^2 C = \mathcal{O}(m_{3/2}/8\pi^2)$, and thereby the resulting low energy gaugino masses take the anomaly pattern (3.79). Note that the contributions from $\frac{1}{8\pi^2} F^i \partial_i \ln(e^{-K_0/3} Z_m)$ and $\frac{1}{8\pi^2} F^i \partial_i \Omega_a$ are of the order of $m_{3/2}/(8\pi^2)^2$, *independently of* how the matter Kähler metric Z_m and the string threshold correction Ω_a depend on T_i.

In KKLT scenario, $F^T/(T + T^*) = \mathcal{O}(m_{3/2}/8\pi^2)$ and $F^C/C \simeq m_{3/2}$ guarantees that M_a/g_a^2 are dominated by the conformal anomaly mediation irrespective of the T-dependence of $e^{-K_0/3} Z_m$ and Ω_a. However, in flux-induced SUSY-breaking scenario with T stabilized by the perturbative Kähler corrections in (3.12), we have an opposite hierarchy of F-terms [30]:

$$\frac{F^T}{T + T^*} \simeq m_{3/2}, \quad \frac{F^C}{C} \simeq \frac{0.5\xi_1}{(T + T^*)^{3/2}} m_{3/2} \ll m_{3/2}. \tag{3.83}$$

As a result, one needs a higher degree of sequestering, i.e. a strong suppression of $\partial_T \ln(e^{-K_0/3} Z_m)$ and $\partial_T \Omega_a$, in order to achieve a conformal anomaly domination in flux-induced SUSY breaking scenario. In fact, we have $\partial_T \ln(e^{-K_0/3} Z_m) = 0$ for the matter Kähler metric Z_m on $D3$ at leading order in the α' and string loop expansions. It is also expected that the tree-level masses of open string modes on $D3$ are determined by the local physics on the $D3$ world volume, thus are independent of the volume modulus T. This would result in $\partial_T \Omega_a = 0$ for one-loop string threshold corrections. These indicate that the anomaly pattern of $D3$ gaugino masses might be possible in flux-induced SUSY breaking scenario also, although it requires that the sequestering persists for higher order matter Kähler metric and string loop threshold corrections.

As is well known, pure anomaly mediation for sfermion masses is problematic, as it leads to tachyonic sleptons at the TeV scale. However, it might be possible to avoid tachyonic sleptons while keeping the anomaly pattern of gaugino masses maintained [35], for instance by introducing a D-term contribution to sfermion masses.

3.4 Schemes with strong model dependence

In SUSY-breaking scenarii leading to the above three patterns of gaugino masses, the UV sensitive string and/or KK threshold contributions $\tilde{M}_a^{(1)}|_{\mathrm{string}}$ to the ratios M_a/g_a^2 are always subleading effects compared to the UV insensitive contributions calculable within the 4D effective theory. This is the reason that we can make reliable predictions for the low energy gaugino mass ratios. However, in certain cases, $\tilde{M}_a^{(1)}|_{\mathrm{string}}$ becomes one of the dominant contributions, and then we loose the predictive power as the low energy gaugino masses become sensitive to the UV physics above the 4D cutoff scale Λ. Here

---

[6]Note that $\mathrm{Re}(S) \propto 1/g_{\mathrm{st}}$ and $\mathrm{Re}(T) \propto R^4/g_{\mathrm{st}}$, where g_{st} and R denote the string coupling and the compactification radius, respectively.

we present two such examples; one is the volume moduli-dominated SUSY breakdown in perturbative heterotic string theory [3, 20] and the other is the M theory compactification on G_2 manifolds with moduli stabilization by non-perturbative dynamics [21, 37].

3.4.1 Moduli dominated scenario in heterotic orbifold compactification

As an example of the scheme in which $\tilde{M}_a^{(1)}|_{\text{string}}$ is one of the dominant contributions to M_a/g_a^2, let us consider the heterotic string compactification on orbifolds for which the tree-level Kähler potential and gauge kinetic function are given by

$$K = -\ln(S + S^*) - \sum_i \ln(T_i + T_i^*) + \prod_i (T_i + T_i^*)^{n_i^m} Q_m^* Q_m,$$

$$f_a^{(0)} = S, \tag{3.84}$$

where T_i $(i = 1, 2, 3)$ are the volume moduli of the underlying toroidal compactification. Using the constraint from the $SL(2, Z)$ modular invariance for each T_i, the one-loop gaugino masses including string threshold correction have been obtained in [20]. Applying the result of [20] for the following pattern of volume moduli-dominated SUSY breakdown:

$$\left| \frac{F^S}{S + S^*} \right| \lesssim \mathcal{O}\left(\frac{m_{3/2}}{8\pi^2}\right), \quad \frac{F^i}{T_i + T_i^*} = \mathcal{O}(m_{3/2}), \tag{3.85}$$

and ignoring the subdominant pieces, we find

$$\left(\frac{M_a}{g_a^2}\right)_{\text{TeV}} = \tilde{M}_a^{(0)} + \tilde{M}_a^{(1)}|_{\text{conformal}} + \tilde{M}_a^{(1)}|_{\text{Konishi}} + \tilde{M}_a^{(1)}|_{\text{string}} \tag{3.86}$$

where

$$\tilde{M}_a^{(0)} = \frac{1}{2} F^I \partial_I f_a^{(0)} = \frac{1}{2} F^S,$$

$$\tilde{M}_a^{(1)}|_{\text{conformal}} = \frac{1}{16\pi^2} b_a \frac{F^C}{C}$$

$$= -\frac{1}{16\pi^2}(3C_a - \sum_m C_a^m)\left(m_{3/2} - \frac{1}{3}\sum_i \frac{F^i}{T_i + T_i^*}\right)$$

$$\tilde{M}_a^{(1)}|_{\text{Konishi}} = -\frac{1}{8\pi^2} \sum_m C_a^m F^I \partial_I \ln(e^{-K_0/3} Z_m)$$

$$= -\frac{1}{8\pi^2} \sum_{i,m} C_a^m \left(\frac{1}{3} + n_i^m\right) \frac{F^i}{T_i + T_i^*},$$

$$\tilde{M}_a^{(1)}|_{\text{string}} = \frac{1}{16\pi^2} \sum_i F^i \left[\frac{1}{3} \delta_{\text{GS}} \partial_i \ln\left((T_i + T_i^*)\eta^2(T_i)\right) \tag{3.87} \right.$$

$$\left. + \left(C_a - \sum_m C_a^m(1 + 2n_i^m)\right)\partial_i \ln \eta^2(T_i) \right],$$

where $\eta(T) = e^{-\pi T/12} \prod_{n=1}^{\infty}(1 - e^{-2\pi nT})$ is the Dedekind function, and δ_{GS} is the coefficient of the Green-Schwarz counterterm. Rearranging the above result, we finally find

$$\left(\frac{M_a}{g_a^2}\right)_{\text{TeV}} = \frac{1}{2} F^S + \frac{b_a}{16\pi^2} m_{3/2} \tag{3.88}$$

$$+\frac{1}{16\pi^2}\sum_i\left(\frac{F^i}{T_i+T_i^*}\right)\kappa(T_i)\left(\frac{\delta_{\mathrm{GS}}}{3}+C_a-\sum_m(1+2n_i^m)C_a^m\right),$$

where

$$\kappa(T_i)=1+(T_i+T_i^*)\partial_i\ln\eta^2(T_i). \tag{3.89}$$

The above result shows that the low energy gaugino masses in the volume moduli-dominated SUSY-breaking scenario in heterotic string theory are indeed quite sensitive to the string threshold $\tilde{M}_a^{(1)}|_{\mathrm{string}}$. Although, in this case, one could determine $\tilde{M}_a^{(1)}|_{\mathrm{string}}$ using the $SL(2,Z)$ modular invariance [20], usually this is not possible in more generic compactifications, and then we loose the predictive power.

The volume moduli domination in heterotic string might be achieved in the racetrack scenario of dilaton stabilization [36]. However, the racetrack stabilization typically leads to an AdS vacuum and still lacks a mechanism lifting this AdS vacuum to a dS vacuum while keeping the volume-moduli domination.

We note that the gaugino mass pattern (3.86), although generically not predictive since it depends on many model parameters, corresponds to the mirage pattern in a special circumstance in which both T_i and F^i have *universal* vacuum values,

$$T_1=T_2=T_3=T,$$
$$F^{T_1}=F^{T_2}=F^{T_3}=F^T, \tag{3.90}$$

and also all gauge charged matter fields originate from the untwisted sector, so that

$$\sum_i n_i^m=-1. \tag{3.91}$$

In such case, (3.86) gives

$$\left(\frac{M_a}{g_a^2}\right)_{\mathrm{TeV}}=\frac{1}{2}F^S+\frac{1}{16\pi^2}\kappa(T)\delta_{\mathrm{GS}}\frac{F^T}{T+T^*}+\frac{b_a}{16\pi^2}\left(m_{3/2}-\kappa(T)\frac{F^T}{T+T^*}\right), \tag{3.92}$$

leading to

$$M_1:M_2:M_3\simeq(1+0.66\alpha):(2+0.2\alpha):(6-1.8\alpha) \tag{3.93}$$

with

$$\alpha=\left(\frac{16\pi^2}{g_{\mathrm{GUT}}^2\ln(M_{\mathrm{Pl}}/m_{3/2})}\right)\left(\frac{(T+T^*)m_{3/2}-\kappa(T)F^T}{8\pi^2(T+T^*)F^S+\kappa(T)\delta_{\mathrm{GS}}F^T}\right). \tag{3.94}$$

3.4.2 M theory on G_2 manifolds

Another scheme in which the string thresholds can significantly affect the gaugino mass ratios is the recently studied M theory compactification on G_2 manifolds in which the

moduli are stabilized by non-perturbative dynamics [21, 37]. The Kähler potential and gauge kinetic functions of 4D effective theory are given by

$$K = -\sum_i n_i \ln(T_i + T_i^*) + Z_\phi(T_i + T_i^*)\phi\phi^*$$

$$+ Z_m(T_i + T_i^*, \phi, \phi^*)Q_m^* Q_m + \dots$$

$$f_a = \sum_i k_i T_i, \tag{3.95}$$

where n_i are positive rational numbers satisfying $\sum_i n_i = 7/3$, k_i are integers, and the ellipsis stands for the terms higher order in ϕ and Q_m. Here $\text{Re}(T_i)$ correspond to the 3-cycle volume moduli, ϕ is a composite hidden matter whose F-component is crucial for the model to have a phenomenologically viable dS or Minkowski vacuum, and Q_m are the visible matter superfields. Under the assumption of non-perturbative dynamics generating for instance a superpotential of the (racetrack) form:

$$W = A_1 \phi^k e^{-a_1 f_h} + A_2 e^{-a_2 f_h} \tag{3.96}$$

for a hidden sector gauge kinetic function $f_h = \sum_i \tilde{k}_i T_i$, it has been noticed that ϕ and T_i can be stabilized (except for some axion components which are harmless) at a SUSY-breaking dS vacuum with the following features [37]:

$$\phi = \mathcal{O}(1),$$

$$\frac{F^\phi}{\phi} = \mathcal{O}(m_{3/2}),$$

$$\frac{F^C}{C} = \mathcal{O}(m_{3/2}),$$

$$\frac{F^i}{T_i + T_i^*} = \mathcal{O}\left(\frac{m_{3/2}}{\ln(M_{\text{Pl}}/m_{3/2})}\right). \tag{3.97}$$

Note that $F^\phi = \mathcal{O}(m_{3/2})$ is crucial in order for the vacuum energy density $\langle V \rangle = K_{I\bar{J}}F^I F^{J*} - 3|m_{3/2}|^2$ to be nearly vanishing.

With the above pattern of SUSY-breaking F-components, the one-loop contributions to M_a/g_a^2 are generically comparable to the tree-level contribution, thus should be carefully taken into account. Explicitly, we have

$$\left(\frac{M_a}{g_a^2}\right)_{\text{TeV}} = \tilde{M}_a^{(0)} + \tilde{M}_a^{(1)}|_{\text{conformal}} + \tilde{M}_a^{(1)}|_{\text{Konishi}} + \tilde{M}_a^{(1)}|_{\text{string}}, \tag{3.98}$$

where

$$\tilde{M}_a^{(0)} = \frac{1}{2}\sum_i k_i F^i = \mathcal{O}\left(\frac{m_{3/2}}{\ln(M_{\text{Pl}}/m_{3/2})}\right),$$

$$\tilde{M}_a^{(1)}|_{\text{conformal}} = \frac{1}{16\pi^2}b_a\left(m_{3/2} + \frac{1}{3}F^\phi\partial_\phi K\right),$$

$$\tilde{M}_a^{(1)}|_{\text{Konishi}} = -\frac{1}{8\pi^2}\sum_m C_a^m F^\phi \partial_\phi \ln(e^{-K_0/3}Z_m),$$

$$\tilde{M}_a^{(1)}|_{\text{string}} = \frac{1}{8\pi^2}F^\phi \partial_\phi \Omega_a, \tag{3.99}$$

where K_0 is the Kähler potential of T_i and ϕ, $\frac{1}{8\pi^2}\Omega_a$ are the M-theory thresholds for the visible gauge coupling superfields, and we have ignored the subleading part of $\mathcal{O}(F^i/8\pi^2)$. As for the relative importance of each of the one-loop contributions, one can consider two distinct possibilities. The first possibility is that the hidden matter ϕ is sequestered from the visible sector, which would mean

$$\Gamma_\phi \equiv \partial_\phi \ln(e^{-K_0/3} Z_m) \simeq 0, \qquad \Omega_{a\phi} \equiv \partial_\phi \Omega_a \simeq 0. \tag{3.100}$$

This is in principle an open possibility as ϕ lives at a point-like conical singularity spatially separated from the 3 cycle (and the conical singularities in it) of the visible gauge fields [37]. In such sequestered case, M_a/g_a^2 are determined by the universal $\tilde{M}_a^{(0)}$ and the comparable conformal anomaly contribution $\tilde{M}_a^{(1)}|_{\text{conformal}}$, and thus take the mirage pattern.

However, sequestering is not a generic consequence of geometric separation, but arises only in special circumstances such as the case of 5D bulk geometry or a geometric separation by warped throat [38]. In case that ϕ is not sequestered from the visible sector, which is actually the case assumed in [21, 37], both Γ_ϕ and $\Omega_{a\phi}$ are expected to be of order unity. If one assumes (as is done explicitly in [37]) that the hidden matter Kähler metric takes the minimal form and the visible matter Kähler metrics are independent of ϕ, i.e.

$$K = -\sum_i n_i \ln(T_i + T_i^*) + \phi\phi^* + Z_m(T_i + T_i^*)Q_m Q_m^*, \tag{3.101}$$

the resulting value of Γ_ϕ is indeed of order unity:

$$\Gamma_\phi = -\frac{1}{3}\phi^* = \mathcal{O}(1). \tag{3.102}$$

Although the order of magnitude is unchanged, including the terms higher order in ϕ can significantly change the size of Γ_ϕ, and thus the size of the Konishi anomaly contribution $\tilde{M}_a^{(1)}|_{\text{Konishi}}$. Furthermore, the masses of superheavy gauge-charged M-theory matter fields $Q_H + Q_H^c$ [39] living at the same conical singularity as Q_m can have a unsuppressed ϕ-dependence through for instance $e^{-K_0/3} Z_H$ where Z_H denotes the Kähler metric of Q_H or through the higher-dimensional superpotential coupling between $Q_H + Q_H^c$ and the hidden matter ϕ such as $\phi^l Q_H Q_H^c$. Since ϕ has a vacuum value of order unity in the unit with $M_{\text{Pl}} = 1$, this eventually yields a sizable M-theory threshold to gaugino masses:

$$\Omega_{a\phi} = \mathcal{O}(1). \tag{3.103}$$

In this case, the gaugino mass ratios can be determined only when one can reliably compute the values of the highly UV sensitive Γ_ϕ and $\Omega_{a\phi}$, which is not available with our present understanding of M theory compactification.

4. Summary and search strategy

In the search for supersymmetry at the LHC, the identification of gauginos will play a crucial role. Predictions for the masses of gauginos are rather robust and seem to favor few distinctive patterns. Of those, the mSUGRA pattern

$$\text{mSUGRA pattern:} \quad M_1 : M_2 : M_3 \simeq 1 : 2 : 6 \tag{4.1}$$

is shared by many schemes, such as

- gravity mediation [2],

- various schemes of dilaton/moduli mediation in string and M-theory including the flux-induced SUSY breakdown [3 – 5],

- gaugino mediation [6],

- gauge mediation [7],

- large volume compactification in Type IIB string theory [8].

The mSUGRA pattern arises if the ratios M_a/g_a^2 in (2.16) are dominated by universal tree level contribution $\tilde{M}_a^{(0)}$ or by universal gauge threshold contribution $\tilde{M}_a^{(1)}|_{\text{gauge}}$. This pattern is closely related to the gauge coupling constants in the TeV range and rather independent of the ultraviolet properties of the underlying scheme. It can appear independently of gauge coupling unification at a large scale, although in some cases (like gauge mediation) such an assumption seems to be required.

In the anomaly pattern

$$\text{Anomaly pattern:} \quad M_1 : M_2 : M_3 \simeq 3.3 : 1 : 9, \qquad (4.2)$$

M_a/g_a^2 are dominated by the conformal anomaly contribution $\tilde{M}_a^{(1)}|_{\text{conformal}}$ related to the $SU(3) \times SU(2) \times U(1)$ β-functions [9]. Schemes that realize this pattern need a very strict separation of the hidden SUSY-breakdown sector from the visible sector of the supersymmetric standard model (MSSM), i.e. a strong sequestering [11, 12]. Because of this restriction its appearance is rather rare and delicate. One possibility can be found in Type IIB string theory with visible sector on D3 branes. In its pure form, anomaly mediation is problematic, as it predicts tachyonic sleptons. This problem has to be removed without disturbing the gaugino mass pattern.

The mirage pattern for gaugino masses

$$\text{Mirage pattern:} \quad M_1 : M_2 : M_3 \simeq (1 + 0.66\alpha) : (2 + 0.2\alpha) : (6 - 1.8\alpha) \qquad (4.3)$$

with $\alpha = \mathcal{O}(1)$ arises if M_a/g_a^2 are dominated by $\tilde{M}_a^{(0)}$ and $\tilde{M}_a^{(1)}|_{\text{conformal}}$ which are comparable to each other. Schemes yielding the mirage pattern have recently been identified in various versions of string and M theory [10, 16]:

- KKLT moduli stabilization [13] in Type IIB string theory with visible sector on D7 branes,

- partial KKLT moduli stabilization [15],

- uplifting via matter superpotentials [14],

- deflected anomaly mediation [17].

In these schemes, the leading contribution of moduli mediation is suppressed by a factor $\log(M_{\mathrm{Planck}}/m_{3/2})$ such that the contribution of the conformal anomaly mediation (suppressed by a loop factor) becomes competitive, while the other (UV sensitive) one-loop contributions, i.e. the Konishi anomaly contribution $\tilde{M}_a^{(1)}|_{\mathrm{Konishi}}$ and the string threshold correction $\tilde{M}_a^{(1)}|_{\mathrm{string}}$ in the formulae (2.16) and (2.17), remain to be subleading. In such a scheme of mixed modulus-anomaly mediation the predictions for gaugino masses are again pretty robust and reliable, while the patterns of squark and slepton masses show stronger model dependence.

Besides the schemes leading to the above three patterns of gaugino masses, one can imagine other scenario in which $\tilde{M}_a^{(1)}|_{\mathrm{Konishi}}$ and/or $\tilde{M}_a^{(1)}|_{\mathrm{string}}$ become important, thus give a different gaugino mass pattern. If $\tilde{M}_a^{(1)}|_{\mathrm{Konishi}}$ from light matter fields is important, $\tilde{M}_a^{(1)}|_{\mathrm{string}}$ from heavy string or M theory modes is expected to become important also. We then loose the predictive power as the low energy gaugino masses become sensitive to the UV physics above the 4D cutoff scale Λ. Examples of such scheme include the volume moduli-dominated SUSY breakdown in perturbative heterotic string theory [20] and the M theory compactification on G_2 manifolds with moduli stabilization by non-perturbative dynamics [21, 37].

Of course, the mass patterns identified so far correspond to the parameters M_1, M_2 and M_3 of the MSSM and not yet to the mass eigenstates. The challenge for phenomenological analyses will be the connection of the M_a ($a = 1, 2, 3$) to the physical masses. Sample spectra for the mirage pattern have been worked out in [40]. The gluino mass is directly related to M_3, whereas the LSP-neutralino can be a mixture of bino, wino as well as Higgsino. The mSUGRA pattern would favor a bino-LSP with a ratio 1/6 compared to the gluino while in the anomaly mediation scheme we might have a wino-LSP with ratio 1/9 to the gluino mass. Of course, we expect mixed states and have to take into account a possibly sizable Higgsino component. If the gluino to LSP ratio turns out to be anomalously large, this could then be a signal for a Higgsino-LSP. On the other hand, a small ratio of the gluino to the LSP-mass (less than 6) might be a hint towards the mirage pattern. Thus even with the knowledge of two of the M_a parameters we might already distinguish between the various different patterns.

If we have identified the LSP, most probably we shall also have some information on the heavier neutralinos as the LSP might be the end-product of a cascade decay. This would allow us to formulate sum rules including all three M_a to further check the patterns. Note that the combination

$$r = \frac{1}{M_3}\left(2(M_1 + M_2) - M_3\right) \qquad (4.4)$$

will approximately vanish both in the mSUGRA and anomaly pattern. A nonvanishing r would thus be a sign of the mirage scheme and allow a determination of the mirage mediation parameter α. For the benchmark scenario with $\alpha = 1$ (intermediate mirage messenger scale) discussed earlier, one would obtain $r \approx 0.84$, while for the case with $\alpha = 2$ (TeV mirage scale) we get $r \approx 2.93$.

Having determined the pattern of gaugino masses we would then include information on squark and slepton masses to break possible degeneracies. Unfortunately, sfermion

masses show a stronger model dependence. Still, as we have seen in our discussion in section 3, useful information could be extracted once the pattern has been identified via the gaugino masses. Such a discussion is beyond the scope of this paper and will be the subject of future investigations.[7] Still we think that the values of the gaugino masses will give us the first hint to unravel the underlying structure of supersymmetry breakdown.

Acknowledgments

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