

# $\log(M_{\text{Pl}}/m_{3/2})$

Oscar Loaiza-Brito\*, Johannes Martin\*, Hans Peter Nilles\* and Michael Ratz\*

\*Physikalisches Institut, Universität Bonn  
Nussallee 12, D-53115 Bonn, Germany

**Abstract.** Flux compactifications of string theory seem to require the presence of a fine-tuned constant in the superpotential. We discuss a scheme where this constant is replaced by a dynamical quantity which we argue to be a ‘continuous Chern–Simons term’. In such a scheme, the gaugino condensate generates the hierarchically small scale of supersymmetry breakdown rather than adjusting its size to a constant. A crucial ingredient is the appearance of the hierarchically small quantity  $\exp(-\langle X \rangle)$  which corresponds to the scale of gaugino condensation. Under rather general circumstances, this leads to a scenario of moduli stabilization, which is endowed with a hierarchy between the mass of the lightest modulus, the gravitino mass and the scale of the soft terms,  $m_{\text{modulus}} \sim \langle X \rangle m_{3/2} \sim \langle X \rangle^2 m_{\text{soft}}$ . The ‘little hierarchy’  $\langle X \rangle$  is given by the logarithm of the ratio of the Planck scale and the gravitino mass,  $\langle X \rangle \sim \log(M_{\text{Pl}}/m_{3/2}) \sim 4\pi^2$ . This exhibits a new mediation scheme of supersymmetry breakdown, called mirage mediation. We highlight the special properties of the scheme, and their consequences for phenomenology and cosmology.

## INTRODUCTION

Superstring theories are the most attractive candidates for a unified description of all observed phenomena. They provide all structures necessary to accommodate the matter content of the standard model as well as all known interactions. However, a commonly accepted stringy extension of the standard model has not yet emerged. Apart from the obvious problem to obtain the correct spectrum there are further, more fundamental questions, which have to be answered if one wants to relate superstring theory to observation. These questions include:

- (i) Why is the scale of weak interactions so much lower than the scale of gravity?
- (ii) Why do we observe four space-time dimensions?
- (iii) Why do we live in de Sitter (or Minkowski) space?

The first question concerns the appearance of the weak scale  $m_{\text{weak}}$  while string and Planck scale,  $M_{\text{St}}$  and  $M_{\text{Pl}}$ , are of similar size, and  $m_{\text{weak}} \ll M_{\text{Pl}}$ . To address the second question, one usually confines oneself to the problem of finding a self-consistent compactification from ten to four dimensions. This includes, in particular, the stabilization of the moduli, which parametrize the size and shape of the internal space. The last question is highly non-trivial since string compactifications admit anti-de Sitter (adS) minima, i.e. vacua with negative vacuum energy. It is challenging to understand in such a framework why the vacuum chosen by nature has positive (or zero) energy.

These questions are not unrelated. In four dimensions, hierarchically small scales can be obtained by dimensional transmutation. The conventional approach to address the hierarchy problem consists in generating a hierarchically small scale of supersymmetry (SUSY) breakdown by a non-perturbative effect, such as a gaugino condensate [1]. This leads to the appearance of the scale  $M_{\text{SUSY}} \sim M_{\text{St}} \exp(-X)$  with  $X$  being a moderately large field-dependent quantity. Once SUSY is broken, the moduli get a non-trivial potential, which might result in their stabilization. However, it is rather difficult to obtain consistent scenarios where a stabilization of all moduli occurs at realistic values. Furthermore, in this picture, one would encounter a situation where the mass of (most of) the moduli is of the order of the weak scale. It is, however, known that such moduli masses lead to severe problems for cosmology.

More recently, the picture has changed due to significant progress in understanding the role of fluxes for moduli stabilization [2]. The main new feature is that some of the moduli can be fixed at realistic values while attaining masses of the order  $M_{\text{St}}$ .

Some important aspects of the ‘flux compactification’ scheme are nicely illustrated by the toy example of KKLT [3] in type IIB string theory. Here, in a first step the complex structure moduli ( $Z^a$ ) and the dilaton ( $S$ ) get stabilized by fluxes. This results in the appearance of a (fine-tuned) constant in the superpotential which, at this stage, breaks SUSY with the scale of SUSY breaking being set by the size of the constant. In the second step, a gaugino condensate is included, which adjusts its

size to this constant, thereby fixing the Kähler modulus ( $T$ ) and restoring SUSY. At this stage, the vacuum energy is negative. This has to be rectified in the third step where an (ad-hoc) ‘uplifting’ is introduced, which renders the vacuum energy positive, and then (again) breaks SUSY. Clearly, in the context of flux compactifications one usually loses the explanation of the hierarchy  $m_{\text{weak}} \sim m_{\text{soft}} \ll M_{\text{Pl}}$ . This hierarchy is now related to the appearance of a small (quantized) constant in the superpotential, which requires a severe fine-tuning.

In Sec. 2, we propose a modification of the ‘flux compactification’ scheme where the positive features of the latter are retained while (some of) the problematic aspects are avoided. The main novelty is that the constant in the superpotential gets replaced by a dynamical quantity. This means that after the first step of moduli stabilization SUSY is unbroken and the superpotential vanishes, leading to zero vacuum energy at this stage. As a consequence, the non-perturbative effect (gaugino condensate) sets the scale of SUSY breaking rather than adjusting to a constant. We argue that the above-mentioned dynamical quantity should be given by the continuous, i.e. non-quantized, part of the Chern–Simons term appearing in dimensional reduction from ten to four dimensions.

In Sec. 3, we discuss the phenomenological consequences of the appearance of the small quantity  $\exp(-X)$ , which leads under rather general circumstances to the ‘little mass hierarchy’

$$m_{\text{moduli}} \sim \langle X \rangle m_{3/2} \sim \langle X \rangle^2 m_{\text{soft}}, \quad (1)$$

where  $\langle X \rangle \sim \log(M_{\text{Pl}}/m_{3/2}) \sim 4\pi^2$ . As we shall see, this results in a scheme with distinct properties. These properties solve (or, at least, help to solve) several problems of supersymmetric extensions of the standard model.

## AVOIDING QUANTIZED CONSTANTS

The scenario of hidden sector gaugino condensation yields a very plausible explanation of the hierarchy  $m_{\text{weak}} \ll M_{\text{Pl}}$ . Here, strong dynamics leads to the non-trivial expectation value of the gaugino bilinear [1],  $\langle \lambda \lambda \rangle = \Lambda^3$ , where  $\Lambda$  is of the order of the renormalization group (RG) invariant scale,

$$\Lambda \sim \mu \exp\{-1/[b_0 g^2(\mu)]\} \ll M_{\text{Pl}}. \quad (2)$$

$b_0$  is the coefficient of the  $\beta$ -function. This strong dynamics triggers a breakdown of SUSY that is parametrized by the gravitino mass

$$m_{3/2} \sim \frac{\Lambda^3}{M_{\text{Pl}}^2} \quad \text{and} \quad M_{\text{SUSY}} \simeq \sqrt{m_{3/2} M_{\text{Pl}}}. \quad (3)$$

Notice that SUSY breakdown requires non-trivial gauge-kinetic function [4],

$$F_i = \exp(-K) D_i W + f_i \langle \lambda \lambda \rangle + \dots \quad (4)$$

In other words, the gauge coupling has to be field-dependent,  $g^{-2} = f$ . It is further possible to include the gaugino condensate in the superpotential [5],

$$W = W_{\text{perturbative}} + C \exp(-af). \quad (5)$$

Early attempts to incorporate gaugino condensation in (heterotic) string theory [6, 5] revealed the importance of a background flux of the field strength  $H$  of the 2-index antisymmetric tensor field  $B$  to avoid a run-away behavior of the dilaton field. This is obvious from the ‘perfect square’ structure of supergravity [7],

$$S_{\text{SUGRA}} \supset (H - \alpha' \langle \lambda \lambda \rangle)^2 \quad (6)$$

Here,  $H$  is the 3-form field strength of the two-index antisymmetric tensor field  $B$  that appears in the 10d supergravity multiplet. It is important to note that the naive field strength  $H_{\text{naive}} = dB$  has to be amended by Chern–Simons terms [7, 8],

$$H = dB - \frac{1}{\sqrt{2}} (\omega^{(\text{YM})} - \omega^{(\text{L})}), \quad (7)$$

where the Yang–Mills Chern–Simons term is given by

$$\omega_{MNP}^{(\text{YM})} = \text{Tr} \left( A_{[M} F_{NP]} - \frac{2}{3} A_{[M} A_N A_P] \right), \quad (8)$$

and an analogous expression exists for the Lorentz Chern–Simons term.

The ‘perfect square structure’ (6) leads to the possibility that the flux  $H$  stabilizes the gaugino condensate (or vice versa) [6, 5]. However, generically not all moduli are stabilized. Moreover, the gaugino condensate can no longer account for the hierarchy, since it balances the value of the quantized  $H$  [9]. If one were to set  $H = 0$ , the latter problem would not arise. However, in such a scenario a non-trivial value of the gaugino condensate would now lead to a vacuum energy of order  $\Lambda^6/M_{\text{Pl}}^2$ , which is inconsistent with observation. One would therefore need a small, non-quantized piece of  $H$  to conspire with the gaugino condensate such that the vacuum energy (almost) vanishes.

Our statements rely on the non-quantized nature of the Chern–Simons terms. The quantization of  $H$  was shown in [9] for the case  $H = dB$  and did not take into account the appearance of the Chern–Simons terms.<sup>1</sup> If  $dB = 0$  to

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<sup>1</sup> When compactifying on a compact space  $K$  with  $\pi_1(K) = \mathbb{Z}_n$  there are fractional contributions to  $H$ ,  $\delta H = 1/n$  [9]. It has been argued that this might be used to stabilize the dilaton [10]. However, we find it difficult to imagine that  $1/n$  could explain the hierarchy between string and weak scale.

leading order, it was argued in [11] that the cancellation should take place between the gaugino condensate and the Chern–Simons terms, which avoid the quantization constraint.

Let us spell out these arguments in more detail. We are interested in the  $A_m A_n A_p$  part of the Chern–Simons term (8) where  $m, n, p$  are indices w.r.t. the internal dimensions. As is well known, those internal components of the gauge fields can come in two different types:

- (i) On the one hand, the ‘discrete Wilson lines’ [12] correspond to quantized background values of  $A_n^a t_a$  ( $t_a$  denotes the generator) with support on non-contractible loops in the internal space. They take values in the adjoint representation so that switching them on does not reduce the rank of the gauge group. To understand the quantization of the ‘discrete Wilson lines’  $A_n^a t_a$ , observe that an adjoint expectation value does not break the  $U(1)$  generated by  $t_a$ . Consequently, the expectation values of ‘discrete Wilson lines’ are quantized to ensure that the zero-modes living on the above-mentioned non-contractible loops are single-valued.
- (ii) ‘Continuous Wilson lines’ [13], on the other hand, transform in the coset of the gauge group (which is present before they are switched on). A generic expectation value of a ‘continuous Wilson line’ does reduce the rank. Since the  $U(1)$  generated by  $t_a$  is (generically) broken, there is no quantization constraint for ‘continuous Wilson lines’. Let us finally mention that in orbifold compactifications [14, 15] ‘continuous Wilson lines’ emerge from the untwisted sector, and can be interpreted as ‘matter fields’ in the massless spectrum [16].

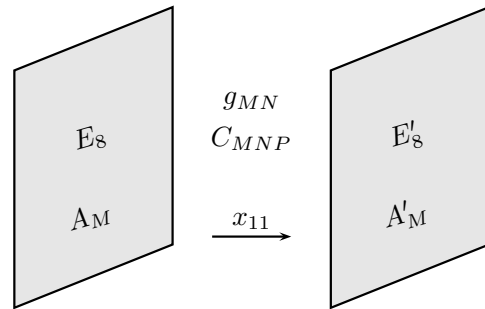
It is now clear that the trilinear term of three continuous Wilson lines can attain arbitrary values, and does in particular not suffer from quantization. It is precisely this term, which can adjust to a gaugino condensate, thus cancelling the corresponding potential energy. In the following, we will refer to such terms as ‘continuous Chern–Simons terms’. We observe that  $\omega^{(YM)}$  and  $\langle \lambda \lambda \rangle$  are both ‘ $\alpha'$  corrections’, thus suggesting their alignment without involving the quantized flux.

Further support for the cancellation between the gaugino condensate and the continuous Chern–Simons term is provided within the framework of heterotic M-theory of Hořava and Witten [17, 18]. In M-theory, gravity lives in the 11d bulk whereas the gauge fields reside on the two 10d boundaries. In the 11d bulk supergravity multiplet we find a 3-index tensor field  $C_{MNP}$  with the four index field strength  $G = dC + \text{Chern–Simons terms}$ . Dimensionally reducing to 10 dimensions one finds that  $B_{MN}$  descends from  $C_{MN,11}$  with the corresponding relation between  $H$  and  $G$ . It is now clear that the Chern–

Simons terms are located on the boundaries,

$$G = dC + \alpha' \sum_i \delta(x_{11} - x_{11}^i) \left( \omega_i^{(YM)} - \frac{1}{2} \omega_i^{(L)} \right) \quad (9)$$

with  $dG = \text{Tr} F_1^2 + \text{Tr} F_2^2 - \text{Tr} R^2$  where  $F_1$  and  $F_2$  represent the field strengths of the two  $E_8$  factors. It has further been shown in Ref. [19] that the perfect square structure between the flux  $G$  and the gaugino condensate generalizes to this case. Since the gauginos are also fields confined to the boundaries, we consider this as a further argument for a cancellation between the gaugino condensate and the Chern–Simons terms [20, 21], while the quantized bulk contribution  $dC$  should not contribute to this cancellation, thus avoiding any known quantization constraint. We sketch this local cancellation in Fig. 1.



**FIGURE 1.** M-theory set-up.  $C$  lives in the bulk whereas both the gaugino condensate and the (continuous) Chern–Simons terms live on the branes.

Let us now summarize the outcome of the discussion so far. We have argued that the appearance of quantized constants in the superpotential should be avoided in order to explain the hierarchy between the Planck and the weak scale by a natural mechanism such as gaugino condensation. We have also discussed that the field strength  $H$  and an expectation value of the gaugino bilinear combine into a perfect square (6), and that  $H$  contains both quantized and continuous parts. This leads to the possibility that the continuous Chern–Simons term cancels the contribution of the gaugino condensate to the vacuum energy.

We have now provided all the ingredients of a modification of the ‘flux compactification’ scheme, which can be outlined as follows: In a first step, before the inclusion of gaugino condensate and continuous Chern–Simons terms, the moduli are fixed by quantized 3-form fluxes without breaking SUSY—i.e.,  $D_i W = 0$  for all fields and  $W = 0$  in the vacuum. The inclusion of the gaugino condensate, in the second step, leads to SUSY breakdown. As a consequence,  $M_{\text{SUSY}}$  is now explained by dimensional transmutation (cf. Eq. (3)). At this stage, the vacuum energy is of order  $M_{\text{SUSY}}^2$ , i.e. unacceptably large. This vacuum energy can be compensated for by the continuous Chern–Simons term. A more detailed discussion and examples will be presented elsewhere [22].

## MIRAGE MEDIATION

Let us now investigate the phenomenological properties of the scheme where all but one modulus are fixed by fluxes and the last one gets stabilized through non-perturbative effects such as a gaugino condensate.

### General structure of the scheme

The stabilization of the last modulus is described by the following model-independent structure of the (effective) superpotential:

$$W = A + C \exp(-X). \quad (10)$$

Here,  $A$  and  $X$  represent vacuum expectation values of field-dependent quantities, and  $C \sim M_{\text{Pl}}^3$ . The value of  $A$  has to be small compared to the string/Planck scale, which can be achieved either through a natural mechanism, or through an explicit fine-tuning. The gravitino mass  $m_{3/2}$  will appear as

$$W \simeq \frac{m_{3/2}}{M_{\text{Pl}}} \times M_{\text{Pl}}^3, \quad (11)$$

where the Planck scale is assumed to be of order of the string scale. The Kähler potential is (up to a constant)

$$K = -n \log(X + \bar{X}) + \dots \quad (12)$$

where the omission denotes the Kähler potential for matter fields, and  $n$  is an order one constant. The scalar potential is given by

$$V = e^K \left[ K^{\alpha\bar{\beta}} (D_\alpha W)(D_{\bar{\beta}} \bar{W}) - 3|W|^2 \right], \quad (13)$$

with the Kähler derivative  $D_\alpha W = \partial_\alpha W + K_\alpha W$ , and where we set  $M_{\text{Pl}} = 1$ . The minimum occurs for  $D_X W \sim 0$ , i.e. by Eq. (11) for

$$X \sim \log(M_{\text{Pl}}/m_{3/2}) \sim 4\pi^2, \quad (14)$$

where we indicate the approximate numerical value of the logarithm of the hierarchy between  $m_{3/2}$  and  $M_{\text{Pl}}$ .

To arrive at zero vacuum energy, we have to arrange a cancellation between the terms in the brackets of (13) which are, when multiplied by  $e^K$ , both of the order of the square of the gravitino mass. We therefore have

$$V''|_{X=\langle X \rangle} \sim m_{3/2}^2, \quad (15)$$

where the prime indicates the derivative w.r.t.  $\text{Re}X$ , and  $\langle X \rangle$  denotes the position of the minimum. To evaluate the physical mass, one has to make sure that the kinetic term

of the fluctuations  $\delta X$  around the minimum is canonical,

$$K = -n \log(\langle X \rangle + \langle \bar{X} \rangle) - \frac{n}{\langle X \rangle + \langle \bar{X} \rangle} (\delta X + \delta \bar{X}) + \frac{n}{(\langle X \rangle + \langle \bar{X} \rangle)^2} (\delta X \delta \bar{X}) + \dots \quad (16)$$

This amounts to a rescaling  $\delta X \rightarrow \delta X_{\text{can}} = \delta X \times (\sqrt{n}/\text{Re}\langle X \rangle)$ . In particular, one finds for the physical mass of the modulus (taking  $\langle X \rangle$  to be real and positive)

$$m_X \sim m_{3/2} \times \langle X \rangle. \quad (17)$$

This enhancement of moduli masses is known to be a rather generic feature of the non-perturbative moduli stabilization mechanisms [23, 24]. We have sharpened the statement, and in particular shown that this enhancement occurs when (i) the Kähler potential for  $X$  is logarithmic, and (ii) the dependence of the superpotential contains the exponential term such that  $\exp(-\langle X \rangle) \sim \Lambda^3$  (with  $\Lambda$  as in Eqs. (2) and (3)). Using (14) we can recast (17) as

$$m_X \sim m_{3/2} \times \log(M_{\text{Pl}}/m_{3/2}) \sim m_{3/2} \times 4\pi^2. \quad (18)$$

### An example

Let us now discuss specifically the outcome in the simple model of KKL<sub>T</sub> [3] with matter fields on D7-branes as analyzed in [25, 26]. We concentrate on the case with the dilaton  $S$ , a Kähler modulus  $T$  and complex structure moduli  $Z_\alpha$ . Matter superfields are denoted by  $Q_7$ . We assume to be in a region of large  $S$  and  $T$ . Let us start with the D7-system [27, 28, 29, 30, 31, 32, 33, 34, 35]. The Kähler potential is assumed to be

$$K = -\log(S + \bar{S} - |Q_7|^2) - 3 \log(T + \bar{T}) + \tilde{K}(Z_\alpha, \bar{Z}_\alpha), \quad (19)$$

where  $Q_7$  denote matter multiplets on the D7 branes. The gauge kinetic function is

$$f_7 = T \quad (20)$$

for gauge bosons on the D7 branes. The inclusion of fluxes leads to a superpotential for the moduli  $S$  and  $Z_\alpha$  [3]. As a consequence, one can eliminate ('integrate out') these fields [25, 36, 37]. This leads to an effective superpotential which is given by

$$W = W(S, Z_\alpha) + C \exp(-aT) + W(Q_7), \quad (21)$$

where  $C \sim M_{\text{Pl}}^3$  and  $a$  are constants. The term  $C \exp(-aT)$  represents gaugino condensation on the D7-branes. When analyzing the potential, we look for minima where the  $Q_7$  scalars (and therefore  $W(Q_7)$  as

well) do not receive non-trivial vacuum expectation values. Extremizing the scalar potential w.r.t.  $T$  leads to an anti-de Sitter vacuum with energy  $\sim |W(S, Z_\alpha)/M_{\text{Pl}}|^2$ . To render this vacuum realistic, one introduces an (ad hoc) uplifting, which may be parametrized as

$$V_{\text{lift}} = \frac{D}{(T + \bar{T})^{n_T}}. \quad (22)$$

By tuning  $D$  it is possible to obtain local de Sitter vacua with energy consistent with observation. The relevant scales appearing in such a vacuum have been calculated in [26], and they are given by:

$$\begin{aligned} M_{\text{St}} &\sim 5 \times 10^{17} \text{ GeV}, \\ 1/R &\sim 10^{17} \text{ GeV}, \\ m_{Z,S} &\sim \frac{1}{M_{\text{St}}^2 R^3} \sim 10^{16} \text{ GeV}, \\ \Lambda_{GC} &= M_{\text{St}} e^{-\langle aT \rangle/3} \sim 10^{13} \text{ GeV}, \\ M_{\overline{\text{D3}}} &\sim e^{A_{\text{min}}} M_{\text{St}} \sim 10^{11} \text{ GeV}, \\ m_T &\sim \langle aT \rangle m_{3/2} \sim 10^6 \text{ GeV}, \\ m_{3/2} &\sim \frac{1}{M_{\text{St}}^2 R^3} \left( \frac{G_{(0,3)}}{G_{(2,1)}} \right) \sim 10^4 \text{ GeV}, \\ m_{\text{soft}} &\sim m_{\text{weak}} \sim \frac{m_{3/2}}{\langle aT \rangle} \sim 10^2 \text{ GeV}, \end{aligned} \quad (23)$$

where  $\Lambda_{GC}$  is the dynamical scale of D7 gaugino condensation,  $M_{\overline{\text{D3}}}$  is the red-shifted cutoff scale on  $\overline{\text{D3}}$ ,  $e^{-\langle aT \rangle} \sim m_{3/2}/M_{\text{St}}$  and  $e^{A_{\text{min}}} \sim \sqrt{m_{3/2}/M_{\text{St}}}$ .  $G_{(0,3)}$  and  $G_{(2,1)}$  denote the (0,3) and (2,1) components of the flux  $G$ . As is obvious from the above expressions, the SUSY breaking component  $G_{(0,3)}$  is substantially suppressed against  $G_{(2,1)}$ , which preserves SUSY. In this case we have  $X = aT$  and it will have a vacuum expectation value of order  $\langle X \rangle \sim \log(M_{\text{Pl}}/m_{3/2}) \sim 4\pi^2$  as we have discussed earlier.

## SUSY mediation

SUSY is broken by the uplifting (cf. Eq. (22)). To describe the SUSY breakdown in the usual language, one attributes the associated  $F$ -term expectation value to the so-called chiral compensator field  $\tilde{C}$  [26]. To see what this means, recall the usual supergravity relation (in the absence of  $D$ -term expectation values)

$$m_{3/2}^2 \sim \sum_i \frac{F_i^2}{M_{\text{P}}^2}, \quad (24)$$

where the sum extends over the  $F$ -term expectation values of all chiral fields. Here, the dominant  $F$ -term is the

one of the chiral compensator which is adjusted such that (24) holds. On the other hand, the  $F$ -term of the  $T$ -modulus is suppressed (cf. Eq. (23)).

Let us now explain how the suppressed  $F_T$  term emerges. Before uplifting,  $F_T$  vanishes, and  $T$  is stabilized with a mass  $\sim \langle aT \rangle m_{3/2}$  where  $m_{3/2} = e^{K/2} |W|$  is the (adS) gravitino mass. Uplifting does (practically) not change  $m_{3/2}$  but depends on  $T$  (cf. Eq. (22)). As a consequence  $T$  is slightly moved against its original minimum after uplifting. The shift in  $T$  is easily calculated in terms of the canonically normalized fluctuations around the minimum (cf. Eq. (16))

$$\frac{d}{d\delta\bar{X}_{\text{can}}} m_T^2 |\delta X_{\text{can}}|^2 \stackrel{!}{=} - \frac{d}{d\delta\bar{X}_{\text{can}}} V_{\text{lift}} \sim m_{3/2}^2 M_{\text{Pl}}^2,$$

where we used in the last relation that  $V_{\text{lift}}$  is tuned such as to cancel the negative energy of the adS minimum,  $V_{\text{adS}} = -3|W|^2 e^K = -3m_{3/2}^2 M_{\text{Pl}}^2$ . This leads to  $\delta X_{\text{can}} \sim M_{\text{Pl}}/\langle X \rangle^2$  so that in the shifted de Sitter minimum

$$F_T^2 \sim \frac{m_{3/2}^2 M_{\text{Pl}}^2}{\langle X \rangle^2}. \quad (25)$$

This implies that soft terms induced by  $F_T$  are suppressed against  $m_{3/2}$  by a factor  $\sim \langle X \rangle$ . In particular, it is the same factor  $\langle X \rangle$ , which both enhances the modulus mass and suppresses the modulus  $F$ -term.

Hence, the ‘gravity mediated’ (or ‘modulus mediated’) soft terms, being controlled by  $F_T/T$ , are suppressed against the gravitino mass, with the suppression factor  $(F_T/T)/m_{3/2} \sim F_T/F_C \sim 1/\langle X \rangle$ . This suppression is comparable to a loop-factor, and therefore anomaly mediation [38, 39] becomes competitive. As a consequence, the soft mass terms receive comparable contributions both from the  $F$ -term of the  $T$ -modulus (‘gravity mediation’) and from the super-conformal anomaly (‘anomaly mediation’). In general, one might hence expect that such a mix is a generic property of ‘sequestered’ models where the communication of SUSY breakdown can be more suppressed than by the Planck scale. We will call this scheme ‘**mirage mediation**’ in the following.

Let us emphasize the two features of mirage mediation that are most important for cosmology and phenomenology:

- The mass of  $T$  is governed by SUSY breakdown. Yet this mass is enhanced with respect to the value of the gravitino mass (cf. Eq. (17)),  $m_T = \langle X \rangle m_{3/2} \sim 4\pi^2 m_{3/2}$ , and thus becomes quite heavy.
- The soft mass terms of the matter fields are suppressed with that same factor  $m_{\text{soft}} \sim m_{3/2}/\langle X \rangle \sim m_{3/2}/4\pi^2$ . If we thus assume that the soft terms are in the region of the weak scale,  $m_{3/2}$  will be in the multi TeV region and thus heavy as well.

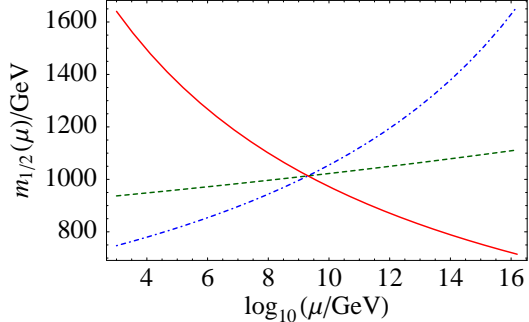
The general mass pattern of the scheme is thus determined by this little hierarchy  $\langle X \rangle = \log(M_{\text{Pl}}/m_{3/2}) \sim 4\pi^2$  with

$$m_T \sim \langle X \rangle m_{3/2} \sim \langle X \rangle^2 m_{\text{soft}}. \quad (26)$$

## Phenomenological aspects

The above-mentioned mix of gravity and anomaly mediation, i.e. the ‘mirage mediation’ scheme allows, at least in substantial regions of the parameter space, to retain the attractive features of these mediation mechanisms while discarding the problematic aspects. The most important issues are the following:

- Anomaly mediation has the notorious problem of negative mass squares for some matter fields, in particular for the sleptons. In **mirage mediation**, the ‘gravity mediated’ contribution can render the slepton mass squares positive thus leading to a consistent framework.
- We have a partial solution of the **flavour problem**. First of all, anomaly mediation is flavour-blind and thus does not cause the usual flavour problems. If, in addition, all the fields live on the D7 branes we have a common scalar mass from the modulus mediation. This additional feature is not a result of the scheme itself, but a consequence of the assumption concerning the origin of matter fields. Nevertheless, it is worthwhile to stress that in mirage mediation the flavour problem get ameliorated, and that the scheme is flexible enough to allow for the implementation of a mechanism that solves the flavour problem.
- There is also a partial solution to the **SUSY CP-problem** coming from the special property of the superpotential [40]. That is, the phases of the A-terms and gaugino masses are aligned. However, the extreme smallness of the various electric dipole moments might require further alignment of phases (see, e.g., [41]).
- The scheme leads to a distinct pattern for the spectrum of the low-energy effective theory. For example, it has been observed that the spectrum exhibits a **mirage unification** scale [42, 43]—i.e., the gaugino and scalar masses meet at an intermediate scale (an energy scale well below the GUT scale). However, this mirage unification scale does not correspond to a physical scale. It has also been argued that the partial cancellation of the RG evolution of the soft masses may ameliorate the SUSY fine-tuning problem [44, 45].
- In contrast to most of the other schemes of SUSY breakdown, in mirage mediation the lightest super-



**FIGURE 2.** Mirage unification for  $m_{3/2} = 40\text{TeV}$ ,  $M_0 := F_T/T = m_{3/2}/(4\pi)^2$  and  $\alpha := m_{3/2}/[M_0 \times \log(M_{\text{Pl}}/m_{3/2})] = 1$ . The red (solid)/green (dashed)/blue (dash-dotted) curve shows the evolution of the gluino/Wino/Bino mass.

partner (LSP) is dominated by the **Higgsino** component in large regions of the parameter space [43, 41, 44].

## Cosmological aspects

(Locally) supersymmetric theories are often in conflict with cosmology because they predict long-lived particles whose decays spoil the successful predictions from nucleosynthesis. The most prominent examples for these long-lived particles are the gravitino and the moduli. In the mirage mediation scheme, the latter are so heavy that they decay early enough not to affect nucleosynthesis. This means that the mirage mediation scheme does not suffer from the traditional gravitino and moduli problems.

Let us mention that there are further challenges for moduli cosmology, which persist even if moduli are rather heavy. These remaining problems include: moduli may not find the minimum of their effective potential at all; some of them might run to the phenomenologically unacceptable run-away minimum [46] due to a large initial velocity [47] or get destabilized by thermal effects [48, 49]. Nevertheless, there exist a few promising proposals to solve at least some of these problems (see, e.g., [50, 51, 52, 53]), but these solutions may require some further ingredients.

## SUMMARY

We presented a scheme that combines the advantages of the new ‘flux compactification’ scenarios with the traditional lore of moduli stabilization. Different from the usual models of ‘flux compactification’, a crucial feature of this scenario is that the gaugino condensate does not

adjust its size to a quantized constant. Rather it sets the scale of SUSY breakdown, and thus yields the explanation of the observed hierarchy  $m_{\text{weak}} \ll M_{\text{Pl}}$  without the need of fine-tuning. We argued that the ‘continuous Chern–Simons term’, which is comprised of ‘continuous Wilson lines’, should adjust its size such as to cancel the vacuum energy. In particular, in this scheme there is no need for the (ad hoc) uplifting procedure.

We further discussed the consequences of the scheme for phenomenology and cosmology. Most importantly, there is the little hierarchy between moduli, gravitino and soft masses,  $m_{\text{modulus}} \sim \langle X \rangle m_{3/2} \sim \langle X \rangle^2 m_{\text{soft}}$  with  $\langle X \rangle \sim \log(M_{\text{Pl}}/m_{3/2}) \sim 4\pi^2$ . The pattern of supersymmetry breakdown combines the features of gravity/moduli and anomaly mediation. As we have discussed, this leads to an attractive scenario where several problems of supersymmetric extensions of the standard model are ameliorated or even solved.

## REFERENCES

1. H. P. Nilles, *Phys. Lett.* **B115**, 193 (1982).
2. S. B. Giddings, S. Kachru, and J. Polchinski, *Phys. Rev.* **D66**, 106006 (2002), hep-th/0105097.
3. S. Kachru, R. Kallosh, A. Linde, and S. P. Trivedi, *Phys. Rev.* **D68**, 046005 (2003), hep-th/0301240.
4. S. Ferrara, L. Girardello, and H. P. Nilles, *Phys. Lett.* **B125**, 457 (1983).
5. M. Dine, R. Rohm, N. Seiberg, and E. Witten, *Phys. Lett.* **B156**, 55 (1985).
6. J. P. Derendinger, L. E. Ibáñez, and H. P. Nilles, *Phys. Lett.* **B155**, 65 (1985).
7. A. H. Chamseddine, *Nucl. Phys.* **B185**, 403 (1981).
8. M. B. Green, and J. H. Schwarz, *Phys. Lett.* **B149**, 117–122 (1984).
9. R. Rohm, and E. Witten, *Ann. Phys.* **170**, 454 (1986).
10. S. Gukov, S. Kachru, X. Liu, and L. McAllister, *Phys. Rev.* **D69**, 086008 (2004), hep-th/0310159.
11. J. P. Derendinger, L. E. Ibáñez, and H. P. Nilles, *Nucl. Phys.* **B267**, 365 (1986).
12. L. E. Ibáñez, H. P. Nilles, and F. Quevedo, *Phys. Lett.* **B187**, 25–32 (1987).
13. L. E. Ibáñez, H. P. Nilles, and F. Quevedo, *Phys. Lett.* **B192**, 332 (1987).
14. L. J. Dixon, J. A. Harvey, C. Vafa, and E. Witten, *Nucl. Phys.* **B261**, 678–686 (1985).
15. L. J. Dixon, J. A. Harvey, C. Vafa, and E. Witten, *Nucl. Phys.* **B274**, 285–314 (1986).
16. A. Font, L. E. Ibanez, H. P. Nilles, and F. Quevedo, *Nucl. Phys.* **B307**, 109 (1988).
17. P. Hořava, and E. Witten, *Nucl. Phys.* **B460**, 506–524 (1996), hep-th/9510209.
18. P. Hořava, and E. Witten, *Nucl. Phys.* **B475**, 94–114 (1996), hep-th/9603142.
19. P. Hořava, *Phys. Rev.* **D54**, 7561–7569 (1996), hep-th/9608019.
20. H. P. Nilles, M. Olechowski, and M. Yamaguchi, *Phys. Lett.* **B415**, 24–30 (1997), hep-th/9707143.
21. H. P. Nilles, M. Olechowski, and M. Yamaguchi, *Nucl. Phys.* **B530**, 43–72 (1998), hep-th/9801030.
22. O. Loaiza-Brito, J. Martin, H. P. Nilles, and M. Ratz (2005).
23. T. Banks, and M. Dine, *Phys. Rev.* **D50**, 7454–7466 (1994), hep-th/9406132.
24. T. Banks, and M. Dine, *Nucl. Phys.* **B505**, 445–460 (1997), hep-th/9608197.
25. K. Choi, A. Falkowski, H. P. Nilles, M. Olechowski, and S. Pokorski, *JHEP* **11**, 076 (2004), hep-th/0411066.
26. K. Choi, A. Falkowski, H. P. Nilles, and M. Olechowski, *Nucl. Phys.* **B718**, 113–133 (2005), hep-th/0503216.
27. P. G. Camara, L. E. Ibáñez, and A. M. Uranga, *Nucl. Phys.* **B689**, 195–242 (2004), hep-th/0311241.
28. M. Graña, T. W. Grimm, H. Jockers, and J. Louis, *Nucl. Phys.* **B690**, 21–61 (2004), hep-th/0312232.
29. T. W. Grimm, and J. Louis, *Nucl. Phys.* **B699**, 387–426 (2004), hep-th/0403067.
30. D. Lüst, P. Mayr, R. Richter, and S. Stieberger, *Nucl. Phys.* **B696**, 205–250 (2004), hep-th/0404134.
31. D. Lüst, S. Reffert, and S. Stieberger, *Nucl. Phys.* **B706**, 3–52 (2005), hep-th/0406092.
32. P. G. Camara, L. E. Ibáñez, and A. M. Uranga, *Nucl. Phys.* **B708**, 268–316 (2005), hep-th/0408036.
33. L. E. Ibáñez, *Phys. Rev.* **D71**, 055005 (2005), hep-ph/0408064.
34. H. Jockers, and J. Louis, *Nucl. Phys.* **B705**, 167–211 (2005), hep-th/0409098.
35. D. Lüst, S. Reffert, and S. Stieberger (2004), hep-th/0410074.
36. S. P. de Alwis (2005), hep-th/0506267.
37. S. P. de Alwis (2005), hep-th/0506266.
38. L. Randall, and R. Sundrum, *Nucl. Phys.* **B557**, 79–118 (1999), hep-th/9810155.
39. G. F. Giudice, M. A. Luty, H. Murayama, and R. Rattazzi, *JHEP* **12**, 027 (1998), hep-ph/9810442.
40. K. Choi, *Phys. Rev. Lett.* **72**, 1592–1595 (1994), hep-ph/9311352.
41. A. Falkowski, O. Lebedev, and Y. Mambrini (2005), hep-ph/0507110.
42. K. Choi, K. S. Jeong, and K.-i. Okumura (2005), hep-ph/0504037.
43. M. Endo, M. Yamaguchi, and K. Yoshioka, *Phys. Rev.* **D72**, 015004 (2005), hep-ph/0504036.
44. K. Choi, K. S. Jeong, T. Kobayashi, and K.-i. Okumura (2005), hep-ph/0508029.
45. R. Kitano, and Y. Nomura (2005), hep-ph/0509039.
46. M. Dine, and N. Seiberg, *Phys. Lett.* **B162**, 299 (1985).
47. R. Brustein, and P. J. Steinhardt, *Phys. Lett.* **B302**, 196–201 (1993), hep-th/9212049.
48. W. Buchmüller, K. Hamaguchi, O. Lebedev, and M. Ratz, *Nucl. Phys.* **B699**, 292–308 (2004), hep-th/0404168.
49. W. Buchmüller, K. Hamaguchi, O. Lebedev, and M. Ratz, *JCAP* **01**, 004 (2005), hep-th/0411109.
50. J. H. Horne, and G. W. Moore, *Nucl. Phys.* **B432**, 109–126 (1994), hep-th/9403058.
51. M. Dine, *Phys. Lett.* **B482**, 213–221 (2000), hep-th/0002047.
52. L. Kofman, et al., *JHEP* **05**, 030 (2004), hep-th/0403001.
53. R. Brustein, S. P. de Alwis, and P. Martens (2004), hep-th/0408160.