Mirage Pattern from the Heterotic String

Valéri Löwen* and Hans Peter Nilles†

Physikalisches Institut der Universität Bonn
Nussallee 12, 53115 Bonn, Germany

Abstract

We provide a simple example of dilaton stabilization in the framework of heterotic string theory. It requires a gaugino condensate and an up-lifting sector similar to the one postulated in type IIB string theory. Its signature is a hybrid mediation of supersymmetry breakdown with a variant of a mirage pattern for the soft breaking terms. The set-up is suited for the discussion of heterotic MSSM candidates.

*E-mail: loewen@th.physik.uni-bonn.de
†E-mail: nilles@th.physik.uni-bonn.de
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1 Introduction

With the recent success of model building in the framework of the heterotic $E_8 \times E_8$ string theory \cite{1–3} it becomes important to reconsider the questions of moduli stabilization and supersymmetry breakdown. Early attempts considered fluxes and gaugino condensates \cite{4–6}, race track superpotentials \cite{7, 8} and/or Kähler stabilization \cite{9–11}.

More recently, these questions have been studied in the framework of type IIB string theory. While explicit model building towards the MSSM (the minimal supersymmetric extension of the standard model) is more difficult in this framework, it is usually argued that the stabilization of moduli can be achieved quite easily. The reason is the appearance of two types of fluxes [Neveu-Schwarz–Neveu-Schwarz (NS-NS) and Ramond–Ramond (R-R)] \cite{12, 13} while only one of them is present in the heterotic theory\footnote{More fluxes can appear if we go beyond Calabi-Yau compactification \cite{14–16}.}. With the inclusion of nonperturbative effects like gaugino condensation one might then stabilize all moduli \cite{17}. Still, there is a problem with the adjustment of the vacuum energy $E_{\text{VAC}}$ to a small positive value. One has to postulate a so-called “up-lifting” sector that adjusts $E_{\text{VAC}}$ to the desired value \cite{17}. It turns out that this up-lifting sector has important consequences for the explicit pattern of supersymmetry breakdown \cite{18}. Instead of modulus mediation one is led to a hybrid mediation scheme and a so-called “mirage pattern” of the soft breaking terms emerges \cite{19–23}.

Within the framework of the heterotic string, the importance of this up-lifting sector has not been fully appreciated so far. The present paper is an attempt to fill this gap. To illustrate the importance of this sector, we consider a simple example: a gaugino condensate in the absence of a flux background. This is known to lead to a so-called run-away potential with a supersymmetric vacuum at $S \to +\infty$ (where $S$ denotes the dilaton field). Attempts to stabilize the dilaton at a finite value include the consideration of nonperturbative correction to the Kähler potential \cite{9–11}. The resulting scheme is believed to provide a mediation of supersymmetry breakdown that is dominated by the dilaton. Still the question of the fine-tuning of the vacuum energy has to be addressed. In the existing examples it seems again that an additional sector is needed to adjust the vacuum energy to the desired value \cite{9}, just like the up-lifting sector in the type IIB case. To be explicit, we consider the model of Barreiro, Carlos and Copeland (BCC) ref. \cite{11} that appears to be particularly suited for a class of realistic heterotic models \cite{24}. A stabilization of the dilaton can be achieved, but the actual value of $E_{\text{VAC}}$ turns out to be large and positive so that the desired “up-lifting” mechanism should provide a “down-lift” of $E_{\text{VAC}}$. We here consider the mechanism of Lebedev, Nilles and Ratz (LNR) \cite{25} to adjust $E_{\text{VAC}}$. In the framework of type IIB theory this is known as $F$-term uplifting \cite{26–31} and could be originated in dynamical schemes as considered in \cite{32} or schemes that might require the existence of additional branes and/or anti-branes \cite{17}. An explicit discussion of supersymmetry breakdown leads to similar conclusions as in the type IIB case, the uplifting sector is important for the mechanism of supersymmetry (SUSY) breakdown and its mediation such that a variant of a mirage pattern emerges \cite{25}. There are, however, some quantitative differences between the heterotic and the type IIB case that will be discussed in detail.

So far it seems that the stabilization of the dilaton requires sizeable nonperturbative corrections to the Kähler potential while the uplifting sector adjusts $E_{\text{VAC}}$. A closer inspection, however, reveals the surprising fact that one can switch off the corrections to the Kähler potential and still remain with a stabilized dilaton! The uplifting sector alone is responsible for both, modulus stabilization and adjustment of the vacuum energy. Stabilization of the dilaton neither requires nonvanishing flux nor corrections to the Kähler potential. The appearance of a gaugino...
condensate, however, remains crucial.

This remarkable fact shows that modulus stabilization in heterotic string theory is not as difficult as usually assumed. One just needs a suitable uplifting sector very similar to the one postulated in the framework of type IIB string theory. It also shows the importance of this uplifting sector for moduli stabilization and supersymmetry breakdown. Moreover it leads to a hybrid mediation scheme and its signature is a mirage pattern of the soft supersymmetry breaking terms.

The paper is organized as follows. In section 2 we shall discuss the model of ref [11] with a gaugino condensate and nonperturbative corrections to the Kähler potential. We present the heterotic analogy of uplifting in the spirit of LNR [25], the adjustment of the vacuum energy of the metastable vacuum and the discussion of supersymmetry breakdown. This will be followed by a discussion of the soft supersymmetry breaking parameters, the appearance of a mirage pattern and the phenomenological properties of the set-up. Section 4 will contain some concluding remarks.

2 Dilaton stabilization in heterotic string theory

Dilaton stabilization in the context of heterotic string theory can occur via a racetrack mechanism [7, 8] or by the means of nonperturbative corrections to the Kähler potential [9–11]. The former option, consisting of at least two gaugino condensates, leads to a scenario where both the auxiliary field of the dilaton and the $T$ modulus are nonzero with typically $F_T > F_S$. The latter option might require just one gaugino condensate, but sizeable nonperturbative corrections to the tree level Kähler potential. Under certain circumstances this leads to a scenario with $F_T = 0$ which is also known as the dilaton domination scenario.

Casas [9] has investigated the impact of an arbitrary Kähler potential on the low energy theory and also introduced possible string motivated nonperturbative corrections to the Kähler potential. The distinct feature of this ansatz was the fact that it became possible to stabilize the dilaton at phenomenologically acceptable values $\Re S \simeq 2$, although the minimum of the scalar potential turned out to be de Sitter.

In this work we shall investigate the impact of the presence of an up-lifting sector (consisting of hidden sector matter fields [25]) on a set-up with just one gaugino condensate and the tree-level Kähler potential. But let us first repeat the arguments of [9] and present the explicit model of BCC [11].

2.1 Problems with a single gaugino condensate

Consider [9, 11]

$$W(S, T) = -A \frac{1}{\eta^6(iT)} e^{-aS},$$

where $\eta(iT)$ is the Dedekind eta function (convention as in [33]), insuring the correct transformation under the $SL(2, \mathbb{Z})$ target space modular invariance and $a = \frac{8\pi^2}{N}$ if the condensing gauge group is $SU(N)$. The Kähler potential at tree-level is given by

$$K = -3 \log (T + \overline{T}) - \log (S + \overline{S}).$$
In the supergravity language the scalar potential and the $F$ terms are expressed in terms of the function

$$G = K + \log W\overline{W},$$

with $K$ and $W$ being the Kähler potential and the superpotential, respectively. The scalar potential is given by

$$V = e^G \left( G^{-\frac{1}{2}} G_{\alpha} G_{\overline{\beta}} - 3 \right),$$

where $\alpha, \beta$ denote differentiation with respect to the fields and $G^{-\frac{1}{2}}_{\alpha\beta} = K^{-\frac{1}{2}}_{\alpha\beta}$ is the inverse Kähler metric. The $F$ terms are found to be

$$F^\alpha = e^{G/2} G^{-\frac{1}{2}} G_{\alpha \overline{\beta}} G_{\overline{\beta}}.$$  

Using eqs. (1, 2) we obtain

$$F^S = e^{G/2} \frac{1}{SS} \left( K - a \right),$$

$$F^T = -e^{G/2} \left( T + \overline{T} \right)^2 \mathcal{E}(T, \overline{T})$$

with $\mathcal{E}(T, \overline{T}) = (T + \overline{T})^{-1} + 2\eta^{-1}(iT)\partial\eta(iT)/\partial T$. In this work we will assume that the Kähler modulus is stabilized at one of the self-dual points of the $\mathcal{E}$ function, leading to $F^T = 0$. From now on we will drop the $T$ dependence in Kähler potential, gauge kinetic function and superpotential and rescale $A^{-6}(iT_0) \rightarrow A$.

In order to provide the formation of a minimum and so to stabilize the dilaton we look at the stationary point condition

$$V_S = G_S V + e^G \frac{\partial}{\partial S} \left( K^{-\frac{1}{2}}_{SS} G_S G_{\overline{S}} \right) \frac{1}{2} = 0,$$

which relates the derivatives of $G$ as

$$K^{-\frac{1}{2}}_{SS} G_S |G_S|^2 - K_S \overline{S} K^{-\frac{1}{2}}_{SS} |G_S|^2 - G_S = 0.$$  

If we only use the tree-level Kähler potential $\log (S + \overline{S})$, eq. (9) becomes

$$\left[ (2\text{Re}S)^2 \left| \frac{1}{2\text{Re}S} + a \right|^3 - 4\text{Re}S \left| \frac{1}{2\text{Re}S} + a \right|^2 - \left( \frac{1}{2\text{Re}S} + a \right) \right]_{\text{Re}S \approx 2} = 0,$$

which can not be solved for any reasonable value of $a$ or $N$, respectively. Therefore one obtains a runaway potential for the dilaton (fig.1). Adding nonperturbative corrections to standard tree-level Kähler potential as motivated in [9]

$$K_{\text{tree+np}} = \log \left[ \frac{1}{S + \overline{S}} + d \left( \frac{S + \overline{S}}{2} \right)^{p/2} e^{-b\sqrt{\frac{S + \overline{S}}{2}}} \right]$$

provides an extra contribution to eq. (10) which must be sizeable in order to allow for a solution. The explicit model of [11] typically obtains a potential as displayed in fig.2.
Stabilization of the dilaton depends crucially on the presence of nonperturbative corrections to the Kähler potential. Otherwise we would have a run-away behaviour of the dilaton and no stabilization. Typically the vacuum energy turns out to be too large and a “down-lift” to a smaller value is necessary. An additional contribution to the scalar potential is therefore needed to give a solution to eq. (9). The role of such an additional contribution should allow the adjustment of the vacuum energy to a zero (or small) value. We shall see that such an additional sector will stabilize the dilaton even in the absence of the nonperturbative corrections to the Kähler potential.

2.2 Down-lifting the vacuum energy

We would like to investigate the interaction of the set-up described in the previous section with a new hidden sector designed to adjust the vacuum energy. This is the heterotic analogy of [25]. The corresponding Kähler potential reads

\[ K = -\log(S + S) + C\overline{C}, \]

with \( C \) being a hidden sector matter field, assumed to be a singlet under unbroken gauge symmetries. From eq. (11) we obtain

\[ V = e^G \left( K^{-1}_{SS} G_S G_{\overline{S}} + G_C G_{\overline{C}} - 3 \right). \]

As obvious from eq. (13) there are now new contributions to the scalar potential which, by properly adjusting the parameters of the new sector, may modify the shape of the scalar potential.

Let us for concreteness and simplicity consider the Polonyi superpotential [34] and focus on the possibility of having a Minkowski vacuum. The superpotential for this choice is given by

\[ W = \omega + \mu^2 C - Ae^{-a S}, \]

where \( \omega \) and \( \mu^2 \) are constants. The scalar potential eq. (13) becomes

\[ V = e^K \left( |WS(S + S) - W|^2 + |W\overline{C} + WC|^2 - 3|W|^2 \right), \]

\[ \equiv e^K \left( |\gamma_S|^2 + |\gamma_C|^2 - 3|W|^2 \right), \]
where we have introduced the quantities
\[ \gamma^S = W_S (S + \overline{S}) - W, \]
\[ \gamma^C = W_C + W_C, \]
which denote the contribution of the dilaton and hidden sector matter field to SUSY breaking, respectively. The stationary point conditions now become
\[ \frac{\partial V}{\partial S} = -\frac{V}{S + \overline{S}} + e^K \left[ W_S \gamma^S + W_S (S + \overline{S}) + W_S C \gamma^C - 3W_S \overline{W} \right] \overset{!}{=} 0, \]
\[ \frac{\partial V}{\partial C} = V \overline{C} + e^K \left[ -W_C \gamma^S + W_C \gamma^C + W_C \overline{C} \gamma^C - 3W_C \overline{W} \right] \overset{!}{=} 0. \]
In the following we show that eqs. (19, 20) are satisfied at
\[ S = S_0 = O(1), \quad C = C_0 = O(1). \]
To prove whether the stationary point corresponds to a minimum we have to calculate the Hessian matrix for the scalar potential and subsequently evaluate the eigenvalues. In this analysis we can neglect the vacuum energy \( V_0 = V(S_0, C_0) \ll 1. \) Up to the factor \( \exp \left[ K(S_0, C_0) \right] \sim O(1) \) the second derivatives of \( V \) are given by
\[ V_{SS} \simeq W_{SS} (S_0 + \overline{S}_0) \gamma^S + 2W_{SS} (S_0 + \overline{S}_0) W_S + W_{SS} (\gamma^S + \overline{C}_0 \gamma^C - 3\overline{W}_0), \]
\[ V_{SS} \simeq \left| W_{SS} (S_0 + \overline{S}_0) \right|^2 + W_{SS} \gamma^S + \overline{W}_S \gamma^S + \left| W_S \right|^2 \left( -2 + \left| C_0 \right|^2 \right), \]
\[ V_{CC} \simeq 2\overline{W}_0 C_0 W_C, \]
\[ V_{CC} \simeq \left| W_0 + W_C C_0 \right|^2, \]
\[ V_{SC} \simeq -\overline{W}_S W_C + \overline{W}_0 W_S C_0, \]
\[ V_{SC} \simeq -\overline{W}_C W_{SS} (S_0 + \overline{S}_0) + W_S \left( \overline{W}_C \left| C_0 \right|^2 + \gamma^C - 3\overline{W}_C \right). \]
One can now consider two cases. For \( \gamma^C \ll \gamma^S \) the SUSY breaking is dominated by the dilaton. However this does not correspond to a satisfactory choice as the contribution from the Polonyi sector would then not be sufficient to solve eqs. (19, 20). Therefore it is more interesting to look at the so-called *matter dominated* SUSY breaking where \( \gamma^S \ll \gamma^C \). The leading order term in eqs. (22, 27) is \( |W_{SS}(S_0 + \overline{S}_0)|^2 \). Next to leading order terms are \( |W_{SS}(S_0 + \overline{S}_0)\overline{W}_C| \) and \( |W_0|^2 \). All other contributions are sub-leading and can be neglected in this analysis.\(^2\)
Thus we arrive at
\[ \frac{\partial^2 V}{\partial x_i \partial x_j} \sim \left( \begin{array}{ccc} \left| \Gamma \right|^2 & 0 & \Gamma \overline{\theta} \\ 0 & \left| \Gamma \right|^2 & 0 \\ \Gamma \theta & 0 & \Delta \end{array} \right), \]
with
\[ \Gamma = W_{SS} (S_0 + \overline{S}_0), \]
\[ \theta = -W_C, \]
\[ \Delta = \left| W_0 \right|^2. \]
\(^2\)The term \( W_{SS}(S_0 + \overline{S}_0)\overline{W}_C \) in the matrix element \( V_{SS} \) contributes at the next to leading order. We however neglect it here to keep the Hessian as simple as possible in order to obtain transparent equations for the eigenvalues.
where $|\Gamma| \gg |\theta|$, $\Delta$. The indices of the Hessian are defined as $(x_1, x_2, x_3, x_4) = (S, \overline{S}, C, \overline{C})$. The eigenvalues are given by
\begin{equation}
\frac{1}{2} \left( \Delta + |\Gamma|^2 - \sqrt{\Delta^2 - 2|\Delta||\Gamma|^2 + |\Gamma|^4 + 4|\Gamma|^2|\theta|^2} \right) \simeq \frac{\Delta}{2},
\end{equation}
\begin{equation}
\frac{1}{2} \left( \Delta + |\Gamma|^2 + \sqrt{\Delta^2 - 2|\Delta||\Gamma|^2 + |\Gamma|^4 + 4|\Gamma|^2|\theta|^2} \right) \simeq \frac{\Delta}{2} + |\Gamma|^2
\end{equation}
and are all positive. This proves that the stationary point is a local minimum. Moreover the spectrum consists of two heavy states with masses of order $|\Gamma| \sim \left| W_{SS} \right|$ and two light states with masses of order $\sqrt{\Delta} = |W_0| \sim \mu^2$.

2.3 Adjusting the cosmological constant

In consideration of a vanishing/small cosmological constant eq. (13) can be parameterized as
\begin{equation}
G_C G_{\overline{C}} + \lambda = 3 + \epsilon,
\end{equation}
where $\epsilon$ is a fine-tuning parameter and $\lambda = K_{SS}^{-1} G_S G_{\overline{S}} \ll 1$. Then the vacuum energy is
\begin{equation}
V(S_0) = V_0 \sim \epsilon \mu^4.
\end{equation}
The solution to first order in $\epsilon$ reads
\begin{equation}
\omega \simeq \left( 2 - \sqrt{3 - \lambda} - \frac{\epsilon}{\sqrt{3 - \lambda}} \right) \mu^2,
\end{equation}
\begin{equation}
C_0 \simeq -1 + \sqrt{3 - \lambda} + \frac{\epsilon}{2\sqrt{3 - \lambda}}.
\end{equation}

By choosing $\mu^2$, which sets the scale of the Polonyi field, one obtains $V_0/\mu^4 \sim \epsilon$, implying that the system under consideration can be used to construct a Minkowski minimum (or adjusting the vacuum energy to a small value). An example is presented in figs. 3 and 4.

2.4 Supersymmetry breaking parameters

Let us now have a look at the gravitino mass. It is given by
\begin{equation}
m_{3/2} = e^{G/2} = e^{K/2} |W(S_0, C_0)| \simeq \mu^2,
\end{equation}
and the Polonyi part in the superpotential dominates. So we see that $\mu^2$ controls not only the mass of the Polonyi field but also the gravitino mass. Furthermore, for the dilaton we obtain
\begin{equation}
F^S = e^{G/2} K_{SS}^{-1} G_{\overline{S}}.
\end{equation}
The quantity $G_{\overline{S}}$ can be estimated as follows. The stationary point conditions for the scalar potential eq. (13) read
\begin{equation}
V_S = G_S V + e^{G} \frac{\partial}{\partial S} \left( K_{SS}^{-1} G_S G_{\overline{S}} \right) + e^{G} \frac{\partial}{\partial S} \left( G_C G_{\overline{C}} \right) \overset{!}{=} 0,
\end{equation}
\begin{equation}
V_C = G_C V + e^{G} K_{SS}^{-1} \frac{\partial}{\partial C} \left( G_S G_{\overline{S}} \right) + e^{G} \frac{\partial}{\partial C} \left( G_C G_{\overline{C}} \right) \overset{!}{=} 0.
\end{equation}
The first term in each equation vanishes at the minimum and the remaining terms give relations among the derivatives of $G$. The solution is given by

$$G \sim \frac{K_{SS} G_C}{a},$$

(42)

with $G_C \sim \mathcal{O}(1)$. We then arrive at

$$F^S \sim \frac{m_{3/2}}{a}.$$

(43)

The choice of a Minkowski vacuum (or small vacuum energy) requires $G_S$ to have a value that cancels the contribution of $K_{SS}^{-1}$ in eq. (39) and furthermore acts to suppress $F^S$ by $a$, where

$$a \Re S \sim \log \left( \frac{A}{\mu^2} \right) \sim \log \left( \frac{M_{\text{Planck}}}{m_{3/2}} \right) \sim \mathcal{O}(4\pi^2).$$

(44)

For the Polonyi field we have

$$F^C = e^{G/2} G_C \sim m_{3/2},$$

(45)

since $G_C \sim \mathcal{O}(1)$. Thus $F^C$ turns out to be the dominant part in SUSY breakdown whereas $F^S$ appears as being suppressed by a factor as given in eq. (44). The mass of the dilaton field is expressed as

$$m_S^2 = \frac{V_{\overline{S}}}{K_{SS}},$$

(46)

where the second derivative of the scalar potential has the behavior $V_{SS} \sim a^2 m_{3/2}^2$ and thus

$$m_S^2 \sim a^2 m_{3/2}^2$$

(47)

and the mass of the dilaton is enhanced compared to $m_{3/2}$. The mass of the Polonyi field is

$$m_C^2 = \frac{V_{\overline{C}}}{K_{CC}} = V_{C\overline{C}} \sim m_{3/2}^2,$$

(48)

which means that it is comparable to the gravitino mass. This is in accordance with the previous discussion, namely $\mu^2$ sets the scale of the Polonyi field as well as that of the gravitino mass.

### 2.5 The little hierarchy: $\log \left( \frac{M_{\text{Planck}}}{m_{3/2}} \right)$

If we compare this to the type IIB case we, in fact, end up with very similar results (although the starting point was quite different). In type IIB one started with a supersymmetric theory in an anti-de Sitter vacuum which then gets up-lifted to a Minkowski vacuum with SUSY breakdown induced by the up-lifting sector. In the heterotic case we started with an unstabilized dilaton. The superpotential interaction involving matter fields provides the stabilization of the dilaton and also induces the breakdown of supersymmetry. Finally a Minkowski vacuum can be realized by properly adjusting the parameters of the Polonyi sector. The hierarchical structure among $m_S, m_{3/2}$ and gaugino masses $m_{1/2} \sim F^S/s+\overline{s}$ is

$$m_S \sim a \cdot m_{3/2} \sim a^2 \cdot m_{1/2}.$$  

(49)

This is similar as the result obtained in the type IIB case [18, 19]. In both cases, this little hierarchy has its origin in the appearance of the factor $\log \left( \frac{M_{\text{Planck}}}{m_{3/2}} \right)$. It suppresses the modulus contribution to the soft mass terms such that loop induced effects become competitive to the tree level ones. The result is the so-called mirage pattern of soft mass terms.
3 Phenomenological properties of the set-up

Having presented the set-up we would like to analyze in detail the pattern of the emerging soft terms in the effective low energy theory. First we derive the soft breaking terms as boundary conditions valid at the grand unified theory (GUT) scale with a suitable parameterization. We then evolve these soft terms down to the electroweak (EW) scale and impose several phenomenological constraints of theoretical and experimental nature.

3.1 Soft supersymmetry breaking terms

Our analysis in section 2.2 was done in the framework of the Polonyi model. Following the lines of [25] one can also construct generalized superpotentials for which the soft terms take a simpler form. In particular we are interested in a simple pattern for the gaugino masses. For a more general discussion see [35]. In order to simplify the discussion we thus assume that the minimum of the scalar potential emerges at $S_0 = 2$ and $C_0 \ll 1$. The full Kähler potential including interaction with observable matter fields is given by

$$K = -\log (S + \overline{S}) + C\overline{C} + Q_i \overline{Q}_i \left[ 1 + \xi_i C\overline{C} \right], \quad (50)$$

where $Q_i$ are the visible sector matter fields and $\xi_i$ describe the coupling between observable and hidden matter. Furthermore we assume that the string threshold corrections to the gauge kinetic function involving the $C$ field are negligible at $C \ll 1$. The moduli/dilaton mediated contribution to the soft terms is then given by

$$M_a = \frac{1}{2 \Re f_a} F^\alpha \partial_\alpha f_a, \quad (51)$$

$$A_{ijk} = F^\alpha \left[ K_\alpha + \partial_\alpha \log Y_{ijk} - \partial_\alpha \log (K_i K_j K_k) \right], \quad (52)$$

$$m_i^2 = m^2_{3/2} - \overline{F}_i \overline{F}^j \partial_\alpha \partial_\beta \log K_i, \quad (53)$$
where $\alpha$ and $\beta$ run over the SUSY breaking fields, $f_a = S$ are the gauge kinetic functions, $K_\alpha = \partial_\alpha K$ and $K_i$ is the K"ahler metric for the visible fields
\[
K_i = \frac{\partial^2 K}{\partial Q_i \partial \bar{Q}_i}.
\]
Assuming that $Y_{ijk}$ are independent of $S$ and $C$ we obtain
\[
M_a = \frac{F^S}{S_0 + S_0} S_0 + S_0 + b_a g_a^2 \frac{F^\phi}{16\pi^2},
\]
\[
A_{ijk} = -\frac{F^S}{S_0 + S_0} S_0 + S_0 + (\gamma_i + \gamma_j + \gamma_k) \frac{F^\phi}{16\pi^2},
\]
\[
m_i^2 = m_{3/2}^2 - \xi_i |F^C|^2,
\]
with $K_S$ being the derivative of $K$ with respect to $S$. The condition for having a Minkowski vacuum gives a relation among $F^S$ and $F^C$, namely,
\[
\frac{|F^S|^2}{(S_0 + S_0)^2} + |F^C|^2 = 3m_{3/2}^2.
\]
Under these assumptions (minimum at $C_0 \ll 1$, $f_a$ independent of $C$), the tree level soft terms for the gauginos and the $A$ parameters are independent of $F^C$. As discussed in section 2.4, $F^S$ is suppressed. Thus, we have to worry about loop suppressed contributions to the soft terms coming e.g. from the superconformal anomaly [36]. The masses squared of the scalars, however, could behave differently. They do contain contributions from $F^C$ which dominate over $F^S$. Nevertheless the $F^S$ contribution as well as the anomaly part may be of interest if we consider $\xi_i \sim O(1/3)$. Including anomaly mediated contributions into eqs. (51-53) the GUT scale boundary values are given by
\[
M_a = \frac{F^S}{S_0 + S_0} S_0 + S_0 + b_a g_a^2 \frac{F^\phi}{16\pi^2},
\]
\[
A_{ijk} = -\frac{F^S}{S_0 + S_0} S_0 + S_0 + (\gamma_i + \gamma_j + \gamma_k) \frac{F^\phi}{16\pi^2},
\]
\[
m_i^2 = \xi_i \frac{|F^S|^2}{(S_0 + S_0)^2} - \frac{|F^\phi|^2}{(16\pi^2)^2} + \frac{2F^\phi F^S}{16\pi^2} \partial S \gamma_i + (1 - 3\xi_i) m_{3/2}^2,
\]
with $F^\phi$ being the auxiliary field of the conformal compensator, $b_a$ are the beta function coefficients, $\gamma_i$ gives the anomalous dimension and $\tilde{\gamma}_i = 16\pi^2 (\partial \gamma_i/\partial \log Q)$ with $Q$ being the renormalization scale. More details can be found in [19].

For the gauginos eq. (60) we obtain a similar result as in the type IIB picture. They are split at the GUT scale according to their beta function coefficients. Since the evolution of the gauginos from the GUT to the EW scale is governed by the same beta function coefficients the splitting disappears at an intermediate scale, leading to the mirage unification of the gaugino masses.

The form of the $A$ terms and the scalar masses is very similar to the type IIB case. For $\xi_i \sim 0$, $F^C$ dominates and we obtain matter dominated SUSY breaking with $m_i^2 \sim m_{3/2}^2$. The choice $\xi_i \sim O(1/3)$ would suppress the $F^C$ contribution and make it comparable to the dilaton
and anomaly mediated contributions. The problematic feature of anomaly mediation is the potential presence of tachyonic sleptons. Due to the mixing between the dilaton and anomaly mediated contributions also the squarks might become tachyonic here. The presence of tachyons can be avoided by appropriately choosing $\xi_i$.

As already studied in the type IIB case, the phenomenology of such a mixed mediation scheme depends crucially on the ratio between modulus/dilaton and anomaly contributions. We introduce the parameterization

$$\vartheta \equiv \frac{1}{M} \frac{F_S}{S_0 + \bar{S}_0},$$

$$M \equiv \frac{m_{3/2}}{16\pi^2},$$

where $\vartheta$ measures the relative importance of dilaton and anomaly mediation and $M$ sets the scale of the soft mass terms. Note that the limit $\vartheta = 0$ corresponds to pure anomaly mediation whereas $\vartheta \gg 1$ is pure dilaton domination. The last term in eq. (61) is enhanced by $(16\pi^2)^2$ with respect to the other terms. In order to compare its contribution with the remaining terms we use

$$\eta_i^2 \equiv (1 - 3\xi_i)(16\pi^2)^2,$$

with $\eta_i = 0$ corresponding to $\xi_i = 1/3$. Then, the soft terms eqs. (69)-(61) take the form

$$M_a = M \left[ \vartheta + b_a g_a^2 \right],$$

$$A_{ijk} = M \left[ - \vartheta + (\gamma_i + \gamma_j + \gamma_k) \right],$$

$$m_i^2 = M^2 \left[ \xi_i \vartheta^2 - \gamma_i + 2\vartheta (S_0 + \bar{S}_0) \partial S \gamma_i + \eta_i^2 \right].$$

Let us close this section by pointing out that in the heterotic case the anomaly mediated contributions to the $A$ parameters and to the scalar masses squared (for $\eta_i \neq 0$) are enhanced compared to the type IIB situation, where the modulus mediated contribution contained a factor of 3 originating from $3 \log(T + \bar{T})$ as compared to $\log(S + \bar{S})$.

### 3.2 Analysis of the parameters

The soft terms in our set-up are described by two continuous parameters

$$\{ \vartheta, m_{3/2} \},$$

the three quantities

$$\{ \tan \beta, \text{sign } \mu, m_t \},$$

and the $\eta_i$ parameters from the matter sector. In the following we will assume non-universal masses for sfermions and Higgses [37,38] and denote

$$\eta_i^{(\text{sfermions})} \equiv \eta_i,$$

$$\eta_i^{(\text{Higgses})} \equiv \eta_i'.$$

We use $\tan \beta$ to fix the $B\mu$ term. The requirement of correct electroweak symmetry breaking fixes the size of $\mu^2$ so its sign remains a free parameter. We will use $m_t = 175$ GeV for the top quark mass and sign $\mu = +1$ throughout our low energy analysis. For the calculation of the low energy data we use the public codes SOFTSUSY [39] and micrOMEGAs [40].
3.2.1 Gauginos

The gaugino soft terms eq. (65) approximately read

\[ M_1 \simeq (3.3 + \varrho)M, \quad M_2 \simeq (0.5 + \varrho)M, \quad M_3 \simeq (-1.5 + \varrho)M. \]  

(72)

The non-universality of the gaugino masses arises from the anomaly mediated contributions which are proportional to the beta function coefficients \( b_a = (33/5, 1, -3) \). At the GUT scale the gauginos show a pattern \( M_1 > M_2 > M_3 \) because \( M_3 \) is suppressed by the negative contribution from anomaly mediation. Depending on the value of \( \varrho \) this negative contribution to \( M_3 \) might become more or less important. At \( \varrho \sim 1.5 \) it might even lead to vanishing gluino mass.

3.2.2 Scalars

As already analyzed in the type IIB case the scalar masses squared could become tachyonic due to the contributions from anomaly mediation. In pure anomaly mediation sleptons are tachyonic. Due to the mixing between dilaton and anomaly mediation also squarks might become tachyonic here (fig. 5). For small \( \varrho \) the dilaton mediated contribution is too weak to cancel the negative anomaly contribution. Nevertheless we can use the \( \eta \) and \( \eta' \) parameters to avoid this problem. In
oder to study a tachyon-free set-up for all values of \( \varrho \) we scan over \( \{ \eta, \eta' \} \) and exclude tachyonic regions. This implies lower bounds on \( \eta \geq 3.5 \) and \( \eta' \geq 1.7 \).

In fact, we might choose the values of \( \eta \) and \( \eta' \) in such a way that the so-called “MSSM hierarchy problem” can be avoided. Correct electroweak symmetry breakdown (EWSB) requires

\[
\frac{m_Z^2}{2} = -\mu^2 + \frac{m_{H_u}^2 - m_{H_d}^2 \tan^2 \beta}{\tan^2 \beta - 1}
\]  

(73)

at the EW scale. Using the input soft terms eqs. (65) - (67) this condition can be rewritten at one-loop level as [41]

\[
m_Z^2 \simeq -1.8\mu^2 + 5.9M_3^2 - 0.4M_Z^2 - 1.2m_{H_u}^2 + 0.9m_{Q_3}^2 + 0.7m_{U_3}^2 - 0.6A_tM_3 + 0.4M_2M_3
\]

\[
:= -1.8\mu^2 + \tilde{m}_Z^2,
\]  

(74)

where we have considered \( \tan \beta = 5 \) and neglected terms with smaller numerical coefficients. If all the parameters on the right hand side of eq. (74) are of order of magnitude of 100 GeV, no significant fine-tuning is needed. However, the soft masses are typically in the TeV region. Then, in order to obtain the correct value of \( m_Z \), cancellations in eq. (74) are required. One could, of course, adjust \( \mu \) such that \( m_Z \) has the correct value but then \( \mu \) might have to be very large. But one might also be interested in a situation where \( \mu \) has a value of order of the weak scale. This would then require cancellations inside \( \tilde{m}_Z^2 \).

The largest contribution to \( \tilde{m}_Z^2 \) comes from the gluino. In order to keep \( \tilde{m}_Z^2 \) small one would have to keep \( M_3 \) under control [42, 43]. As we have seen in section 3.2.1 the gluino mass might become small for small \( \varrho \) (at \( \varrho \sim 1.5 \) it might even vanish). This is ruled out, as the gluino would be the lightest supersymmetric particle (LSP). In addition if the gluino is light it cannot provide the necessary renormalization group (RG) contribution to \( m_{H_u}^2 \) such that eq. (73) will no longer be satisfied and consequently EWSB will not be realized around \( \varrho \sim 1.5 \). Thus larger values of \( \varrho \) are required. In order to achieve a cancellation within \( \tilde{m}_Z^2 \) for moderate values of \( \varrho \) one has to adjust the masses of the scalars (sfermions and Higgses). Here the freedom of choosing \( \eta \) and \( \eta' \) enters the game.

As is evident from eq. (74) the contribution from \( m_{H_u}^2 \) is negative and thus by increasing \( m_{H_u}^2 \) one obtains a sizeable term that could cancel the contribution of the gluino \( M_3 \). The contribution from squarks is positive and one has to keep their masses low. However, we cannot choose \( \eta \) too small, otherwise the squarks might become tachyonic at the GUT scale. The essential lesson we learn from these considerations is to keep \( \eta \) as low as possible and then adjust \( \eta' \). If we want to keep \( \tilde{m}_Z^2 \equiv (100 \text{GeV})^2 \) a relation between the values of \( \varrho, m_{3/2}, \eta \) and \( \eta' \) has to be fulfilled. In that sense the “MSSM hierarchy problem” can be avoided at the expense of a fine tuning of \( \eta' \).

### 3.3 Constraints

After having excluded tachyons, we shall see that the parameter space of our set-up is further restricted by phenomenological constraints. These include the quest for correct EWSB, mass bounds from LEP and the cosmological relic abundance of neutralino dark matter.

- **Correct EWSB**

  The minimization of the MSSM Higgs scalar potential leads to eq. (73). Here \( \mu^2 \) should be positive and \( m_{H_u}^2 \) should be negative at the EW scale. Given a positive \( m_{H_u}^2 \) at the GUT
scale, it will be driven to negative values at the EW scale by the renormalization group evolution according to

$$\frac{dm_{\tilde{H}_u}^2}{d\log Q} \simeq |y_t|^2 \left( m_{\tilde{H}_u^+}^2 + m_{Q_3}^2 + m_{U_3}^2 \right) + |A_t|^2.$$  \hspace{1cm} (75)

The RG evolution is most sensitive to the gluino mass which induces an increase of \(m_{Q_3}\) and \(m_{U_3}\). In a mirage mediation scheme as the one considered here a cancellation between dilaton and anomaly mediated contributions for the gluino mass occurs for small values of \(\varrho\) leading to an ultra-light gluino around \(\varrho \sim 1.5\). There the RG contribution from the gluino to eq. (75) is too small and a satisfactory value of \(m_{\tilde{H}_u}^2\) can not be obtained. The requirement of correct EWSB sets a lower bound on the \(\varrho\) parameter.

- **LEP mass bounds**
  Direct collider searches set lower bounds on the sparticle spectrum and Higgs masses. Most important and restrictive bounds are due to the lightest Higgs boson mass \(m_h > 114\mbox{ GeV}\), the lightest chargino \(m_{\tilde{\chi}^+} > 103.5\mbox{ GeV}\) and the lightest stop quark \(m_{\tilde{t}_1} > 95.7\mbox{ GeV}\) [44]. Regions of parameter space violating one of these bounds are called below LEP. These constraints set a lower bound on \(m_{3/2}\).

- **Neutralino Dark Matter**
  In SUSY models the weakly interacting neutralinos tend to be the LSPs and they are perfect Dark Matter candidates (under the assumption of R-parity conservation). In our model this is true throughout most of the parameter space. The four neutralinos of the MSSM \(\tilde{\chi}_1, \tilde{\chi}_2, \tilde{\chi}_3, \tilde{\chi}_4\) are superpositions of the neutral Higgs fermions \(\tilde{H}_u^0, \tilde{H}_d^0\) and the fermionic partners of the EW gauge bosons \(\tilde{B}_0, \tilde{W}_3^0\). The neutralino mass matrix can be diagonalized by an orthogonal matrix \(Z\), such that the lightest neutralino is given by

$$\tilde{\chi}_1^0 = Z_{11} \tilde{B}_0^0 + Z_{12} \tilde{W}_3^0 + Z_{13} \tilde{H}_d^0 + Z_{14} \tilde{H}_u^0.$$  \hspace{1cm} (76)

Using this decomposition one defines

$$\tilde{\chi}_1^0 = \begin{cases} \text{bino-like} & |Z_{11}|^2 + |Z_{12}|^2 > 0.9 \wedge |Z_{11}| > |Z_{12}|, \\ \text{wino-like} & |Z_{11}|^2 + |Z_{12}|^2 > 0.9 \wedge |Z_{11}| < |Z_{12}|, \\ \text{higgsino-like} & |Z_{11}|^2 + |Z_{12}|^2 < 0.1, \\ \text{mixed} & \text{otherwise}. \end{cases}$$  \hspace{1cm} (77)

We use the 3\(\sigma\) limit from the WMAP collaboration on the neutralino Dark Matter relic abundance [45]:

$$0.087 \leq \Omega_{\tilde{\chi}_1}h^2 \leq 0.138.$$  \hspace{1cm} (78)

We will require that the neutralinos annihilate efficiently enough to satisfy the bound eq. (78) and assume that the LSP abundance is thermal. In the remainder of this paper the regions of the parameter space that violate the upper WMAP bound are called forbidden, those within the bounds are called favored and those below the lower bound are denoted as allowed. In the latter case the correct cosmological abundance of Dark Matter could be achieved with additional dark matter particles and/or a nonthermal origin.
3.4 Low energy aspects of the spectrum

In models where the contributions from modulus and anomaly mediation are comparable we experience the phenomenon of mirage unification at an intermediate scale. The same happens, of course, in the present case of mixed dilaton-anomaly mediation.

- **Gauginos**
  Below the GUT scale the nonuniversality of the gaugino masses is given by the respective beta function coefficients. The renormalization group running of the gaugino masses is governed by the same coefficients, thus at an intermediate scale, the splitting disappears and the gauginos unify. Since there is no physical threshold associated to this scale it is called \textit{mirage scale}. Using the parameterization eq. (62) the mirage scale is given by

\[ M_{\text{MIR}} = M_{\text{GUT}} \frac{e^{-8\pi^2/\rho}}{\rho}. \] (79)

For \( \rho = 5 \) the mirage scale is intermediate while for \( \rho \approx 2 \) mirage unification occurs at the TeV scale. The pattern of the gaugino masses at the GUT scale is \( M_1 > M_2 > M_3 \). At the EW scale this pattern becomes inverted \( M_3 > M_2 > M_1 \) and a compressed spectrum is obtained (fig.6).

- **First and second generation scalars**
  These sparticles behave in a similar way as the gauginos. The reason for this is that they
are unaffected by the large Yukawa coupling $y_t$ and string threshold corrections as well as Kähler anomalies are negligible. The renormalization group running of the first and second generation scalars is given by

$$\frac{dm_i^2}{d\log Q} \sim \sum_a g_a^2 M_a^2 C_i.$$  

(80)

Under these circumstances the first and second generation scalars unify at the same $M_{\text{MR}}$ eq. (79) as the gauginos (fig. 6).

- **Third generation scalars**
  Here we have to distinguish between those whose RG running depends on the Yukawa coupling $y_t$ and those that are unaffected by $y_t$. The only third generation scalars that feel the effect of $y_t$ are $m_{Q_3}^2$, $m_{U_3}^2$ and $m_{H_u}^2$. Consequently they do not unify at the mirage scale eq. (79). All other third generation scalars and $m_{H_d}^2$ are only affected by the smaller

\[^3\text{This is strictly true only for small values of } \tan \beta. \text{ For large values of } \tan \beta \text{ we also have to take into account the bottom quark Yukawa coupling } y_b.\]
Yukawa couplings $y_b, y_\tau \ll y_t$ and the structure of their RG running (fig. 7) is very similar to eq. (80), thus they (partially) share the mirage unification feature.

3.4.1 Small $\tan \beta$ regime

As an illustrative example let us consider $\eta = 4$ and $\eta' = 6$. The corresponding parameter space is shown in fig. 8. In this scenario there are no tachyons. However, correct EWSB and current LEP bounds for $m_h$, $m_{\tilde{\chi}^\pm}$ and $m_{\tilde{\tau}}$ put severe constraints on $\varrho$ and $m_{3/2}$. Particularly we find that for $\tan \beta = 5$ only $\varrho \geq 5$ and $m_{3/2} \geq 8$ TeV are allowed.

The presence of the no-EWSB region appears because at $\varrho \sim 1.5$ the gluino contribution to the RG is too small to make $m_{H_u}^2$ negative at the EW scale. In contrast to type IIB, in the heterotic case we find a region in the parameter space where a chargino is the LSP. This happens because the $A$ terms in the heterotic case are smaller than in type IIB\textsuperscript{[4]} Additionally, these reduced $A$ terms lead to a smaller intra-generational mixing and increase the masses of the stop quark and the stau lepton. Therefore, stop or stau (N)LSP are not realized in this scenario.

\textsuperscript{[4]}More precisely, the contribution from modulus mediation in the $A$ terms is reduced by a factor of 3.
Fig. 9: Same as fig. 8 but with $\tan \beta = 30$. The LEP constraints are less restrictive whereas the chargino-LSP and the No-EWSB regions increase with larger $\tan \beta$. Compared to fig. 8 larger portions of the parameter space are attractive for the discussion of the relic abundance.

For $\varrho$ values close to the no-EWSB region we have $|\mu| < M_1$ and the neutralino LSP is higgsino-like. Going to larger $\varrho$ values the LSP becomes a mixed higgsino-bino state. From $\varrho \sim 7$ we have a mostly bino-like LSP and also $|\mu| \gg M_1$. Thus, in most of the parameter space the LSP is bino-like.

The shaded/strip in fig. 8 shows the region of the parameter space which is favoured by the WMAP results eq. (78). The region below the shaded/strip is allowed (lower abundance) and that above (and to the right) of the shaded/strip is forbidden (too large relic abundance). For $m_{3/2} = 40$ TeV and $\varrho \sim 5$ we are close to the $\tilde{\chi}^+$ LSP region and therefore we have $\tilde{\chi}^0\tilde{\chi}^+$ coannihilation, which enhances the annihilation cross section and lowers the relic abundance. The neutralino in this region is higgsino-like. As $\varrho$ increases, the neutralino becomes mixed higgsino-bino and the $\mu$ term increases. The coannihilation with the chargino gets reduced and the annihilation cross section decreases leading to a higher relic abundance, so that there ($\varrho \sim 6$) we reach the shaded/strip. When we proceed to increase $\varrho$, the $\tilde{\chi}^0$ becomes bino-like. However, we then reach $m_A/2 \sim m_{\tilde{\chi}^0}$ and there the annihilation proceeds efficiently through the pseudo-scalar Higgs exchange $\tilde{\chi}^0\tilde{\chi}^0 \rightarrow A \rightarrow f\bar{f}$. This enhances the annihilation cross section and reduces the relic abundance. For $\varrho > 7$ the mass gap $|m_A/2 - m_{\tilde{\chi}^0}|$ grows and the efficiency of the $A$-channel reduces. As there are no other coannihilation channels available the cross section decreases and the relic abundance becomes too large.
3.4.2 Large tan β regime

For large values of tan β, the LEP mass constraint becomes less restrictive (fig. 9). However, the no-EWSB region gets slightly bigger and the \( \tilde{\chi}^+ \) LSP region covers a larger part of the parameter space (compared to the case of small tan β). The composition of the neutralino LSP is similar to the tan β = 5 situation. For low \( \rho \) values (close to the no-EWSB region) the neutralino is higgsino-like. Then, for larger \( \rho \) it becomes more and more bino-like.

The shaded/strip, satisfying the WMAP limits, differs significantly from the one discussed above. Now, a larger part of the parameter space is consistent with the correct amount of Dark Matter. This is because for large tan β the pseudo-scalar Higgs exchange provides a sizable contribution to the annihilation cross section. For \( m_3/2 = 30 \) TeV and \( \rho \sim 5.5 \) we have \( m_{\tilde{\chi}^0} \sim m_{\tilde{\chi}^0} \sim \mu \) and chargino coannihilation enhances the annihilation cross section and lowers the relic abundance. When \( \rho \) increases, \( \mu \) gets larger (the whole sparticle spectrum becomes heavier) and the \( \tilde{\chi}^0 \) becomes bino-like. For \( \rho > 7 \) the mass gap \( |m_{\tilde{\chi}^+} - m_{\tilde{\chi}^0}| \) grows, thus the \( \tilde{\chi}^+ \tilde{\chi}^0 \) coannihilation channel no longer provides a sizable effect. As a result, the annihilation cross section decreases and the relic abundance grows above the upper WMAP bound. At the same time the mass of the \( \tilde{\chi}^0 \) approaches the value \( m_A/2 \) and thus the pseudo-scalar Higgs exchange begins to contribute. The cross section \( \sigma(\tilde{\chi}^0\tilde{\chi}^0 \rightarrow A \rightarrow b\bar{b}) \) grows with tan² β and so this channel overcomes the decrease of the annihilation cross section caused by the bino component of the neutralino. For \( \rho > 9.5 \) the production of the relic abundance is in the allowed range. For still larger \( \rho \) finally this effect dies out and the relic abundance becomes too large.

3.4.3 Numerical results

Some points are selected from the the allowed parameter space in figs. 8 and 9 and the spectrum is analysed in detail. Examples of the spectra are displayed in table 1.

4 Conclusions

As we have seen, dilaton stabilization in the framework of the heterotic string can be achieved quite easily, if we accept the existence of an up-lifting sector (as postulated previously in type IIB theory) which in any case is needed to adjust the vacuum energy to an acceptable value. One just needs a gaugino condensate, while nontrivial background flux and/or nonperturbative corrections to the Kähler potential are not necessarily required. In that sense the heterotic mechanism of dilaton stabilization is somewhat similar to the one conjectured in the framework of M theory on \( G_2 \)-manifolds (which requires a racetrack scenario and an up-lifting sector [47,48]). A comparison of these two scenarii will be the subject of future work, where we shall also investigate heterotic M theory with a gaugino condensate [49,50].

Such a (partially sequestered) up-lifting sector is thus common to many of the string schemes considered so far. While its exact origin has to be clarified (branes, antibranes etc.), its importance for the resulting phenomenology cannot be overestimated. It is the dominant source of supersymmetry breakdown, but as it is (partially) sequestered it leads to mediation schemes where tree-level moduli contributions compete with loop-effects from the up-lifting sector. The gravitino and the moduli fields become rather heavy, but the soft terms of the MSSM particles are suppressed by a factor of the order of \( \log(M_{\text{Planck}}/m_{3/2}) \) [19,51]. These soft masses often show a characteristic pattern, known as the mirage pattern. It leads to a rather compressed spectrum of masses and seems to be especially robust in the case of gauginos [35]. This specific pattern...
Tab. 1: Three sample spectra. All masses are given in TeV.

|                | A       | B       | C       |
|----------------|---------|---------|---------|
| $\tan \beta$  | 5       | 30      | 10      |
| $\varrho$     | 6       | 10      | 6.5     |
| $m_{3/2}$     | 40      | 6       | 25      |
| $\eta$        | 4       | 4       | 5       |
| $\eta'$       | 6       | 6       | 5       |
| $M_1$         | 1.040   | 0.211   | 0.676   |
| $M_2$         | 1.317   | 0.311   | 0.881   |
| $M_3$         | 2.391   | 0.743   | 1.697   |
| $m_h$         | 0.119   | 0.115   | 0.120   |
| $m_A$         | 2.118   | 0.468   | 1.408   |
| $\mu$         | 0.860   | 0.413   | 0.885   |
| $m_{\tilde{\chi}_i^0}$ | 0.850 | 0.204 | 0.665 |
| $m_{\tilde{\chi}_i^\pm}$ | 0.870 | 0.296 | 0.838 |
| $m_{\tilde{\chi}_i^0}$ | 0.855 | 0.296 | 0.836 |
| $m_{\tilde{t}_1}$ | 1.610 | 0.488 | 1.236 |
| $m_{\tilde{t}_2}$ | 2.110 | 0.694 | 1.578 |
| $m_{\tilde{\chi}_1}$ | 1.398 | 0.233 | 1.021 |
| $m_{\tilde{\chi}_2}$ | 1.522 | 0.347 | 1.112 |
| $\Omega_{\chi^0 h^2}$ | 0.088 | 0.115 | 0.092 |

has been investigated thoroughly in the framework of the type IIB string. Given the particularly successful attempts of realistic MSSM model building in the heterotic theory [1, 2, 52–55], it is reassuring to see that a similar patterns seems to emerge here as well.

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