Phenomenological Implications of Supersymmetry Breaking by the Dilaton

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Abstract

We investigate the low energy properties of string vacua with spontaneously broken $N = 1$ supersymmetry by a dilaton $F$-term. As a consequence of the universal couplings of the dilaton, the supersymmetric mass spectrum is determined in terms of only three independent parameters and more constrained than in the minimal supersymmetric Standard Model. For a $\mu$-term induced by the Kähler potential the parameter space becomes two-dimensional; in the allowed regions of this parameter space we find that most supersymmetric particles are determined solely by the gluino mass. The Higgs is rather light and the top-quark mass always lower than 180 GeV.
The leading candidate for consistently incorporating quantum gravity into the standard interactions of particle physics is a heterotic superstring theory. However, despite its numerous attractions it has not been possible to identify quantitatively the Standard Model (SM) as the low energy limit of string theory. In part this is due to our lack of conceptual understanding of the theory; so far we only enjoy control over (some of) its perturbative regime. Unfortunately, low energy string phenomenology does seem to depend crucially on non-perturbative properties of string theory. The mechanism for supersymmetry breaking, the choice of the string vacuum, or the determination of the gauge couplings are believed to be governed by (possibly ‘stringy’) non-perturbative effects; our current techniques are inappropriate to incorporate such effects into the low energy effective Lagrangian.

Ultimately, we have to come to terms with this deficiency; in the mean time various strategies have been employed in order to investigate and/or constrain the low energy limit of the string. We refrain here from systematically reviewing the subject, instead we briefly outline the method that we are going to follow in this letter. Based on work in the context of gaugino condensation [1, 2] and duality-invariant effective Lagrangians [2–4] it was recently suggested [5] to simply parametrize the unknown non-perturbative physics. All relevant low energy interactions are expressed in terms of couplings calculable in string perturbation theory and couplings encoding the non-perturbative dynamics; the latter then appear as arbitrary parameters in the low energy effective theory. Surprisingly, even in such a general framework this effective theory can display rather distinct properties. In ref. [5] the non-perturbative couplings are constrained by some assumptions about the nature of the non-perturbative dynamics and the nature of supersymmetry breaking. In particular, it was assumed that supersymmetry is spontaneously broken in the moduli/dilaton sector of string theory. Such scalar multiplets are always present in the massless string spectrum and the couplings of the low energy effective Lagrangian are determined by their vacuum expectation values (VEVs). In string perturbation theory, both the moduli and the dilaton are exact flat directions of the effective potential, leaving their VEVs undetermined. In ref. [5] this perturbative degeneracy is assumed to be completely lifted by the non-perturbative dynamics and VEVs for moduli and dilaton to be induced. In addition, supersymmetry is assumed to be spontaneously broken by the auxiliary $F$-terms of the moduli/dilaton supermultiplets. Indeed, in the context of gaugino condensation, such a scenario can occur [1, 6].
The dilaton plays a distinct role in the low energy theory; all its couplings are universal (at the tree level), that is, they are identical for all $N = 1$ heterotic string vacua or equivalently they do not depend on the details of the internal conformal field theory. As a consequence supersymmetry breaking dominated by the dilaton $F$-term leads to very specific and model-independent low energy properties. It is the purpose of this letter to study the phenomenological implications of supersymmetry breaking in the dilaton sector. Such an analysis has not been done previously since supersymmetry, in the context of gaugino condensation, usually breaks in the moduli direction. However, if one does not specify the non-perturbative physics (in the spirit of ref. [5]), supersymmetry breaking by a dilaton $F$-term is a conceivable scenario. As we will see from a phenomenological point of view it automatically leads to some desired features.

In string perturbation theory there is an enormous vacuum degeneracy and out of this plethora of possibilities one chooses (by hand) phenomenologically promising candidate vacua. For the purpose of this article we require the class of string vacua under consideration to satisfy a few standard properties. In addition to the moduli and dilaton, the string spectrum contains families of matter multiplets that are charged under the gauge group $G$. Part of this gauge group has to contain the standard $SU(3) \times SU(2) \times U(1)$, and we denote all light $N = 1$ chiral multiplets in this ‘observable sector’ by $Q^I$. For simplicity we assume that the $Q^I$ coincide with the multiplets present in the minimal supersymmetric Standard Model (MSSM), that is, all particles of the SM occur in chiral superfields with one additional Higgs doublet [8]. (This assumption is not crucial for the structure of the soft supersymmetry-breaking terms themselves; however, some of the low energy properties do depend for example on the Higgs sector.) Thus, the index $I$ labels collectively the quark multiplets $(Q_L, U_R, D_R)$, the leptons $(L_L, E_R)$ and the two Higgs doublets $(H_1, H_2)$, and we suppress their gauge quantum numbers.

The low energy interactions of the observable fields consist of supersymmetric couplings (encoded in a superpotential $W$) and a set of soft supersymmetry-breaking

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* The dilaton VEV also determines the tree-level gauge couplings.
† See however ref. [5].
* Note that $Q_L$ denotes the left-handed quark supermultiplet and should not be confused with $Q^I$, which stands for all matter multiplets.
parameters. The masses and Yukawa couplings of the chiral (matter) fermions are summarized by an effective superpotential of the form

\[ W^{(\text{eff})} = \frac{1}{2} \mu_{IJ} \, Q^I Q^J + \frac{1}{3} y_{IJK} \, Q^I Q^J Q^K , \]

where we have chosen a basis for \( Q^I \) with canonically normalized kinetic energy terms. Thus, \( \mu_{IJ} \) are the physical (supersymmetric) mass terms and \( y_{IJK} \) denote the physical Yukawa couplings.\(^\dagger\) These can be computed as functions of the moduli in string perturbation theory, and we assume here that they reproduce the known Yukawa couplings of the MSSM. On the other hand \( \mu_{IJ} \) is generated either in string perturbation theory or by non-perturbative effects \(^\ddagger\); because of gauge invariance it has only one non-vanishing entry in the direction of the two Higgs doublets \( \mu_{12} = \mu \). Thus, in the standard notation of the MSSM, eq. (1) reads

\[ W^{(\text{eff})} = \mu H_1 H_2 + \sum_{\text{generations}} (y_U Q_L U_R H_2 + y_D Q_L D_R H_1 + y_L L_L E_R H_1) . \] (2)

In addition to the supersymmetric interactions (eqs. (1), (2)) supersymmetry breaking induces soft breaking parameters in the observable sector. The general structure of these soft terms in string theory was analysed in refs. \[1\text{-}3\] and will not be repeated here. Instead, we just recall their form for the particular case of supersymmetry breaking induced by a dilaton \( F \)-term. Because of the universal couplings of the dilaton, the entire effect of the breaking can be parametrized by the gravitino mass \( m_{3/2} \). Neglecting string loop corrections one finds a universal (gauge-group-independent) gaugino mass, which is determined by \( m_{3/2} \)

\[ \tilde{m}_a = \sqrt{3} \, m_{3/2} , \quad \forall a , \] (3)

where we label the different factors in the observable gauge group \( G \) according to \( G = \prod_a G_a \). (The mass term of eq. (3) is given in a basis where the gauginos are

\(^\dagger\) Our definition of \( Q^I, \mu_{IJ} \) and \( y_{IJK} \) differs from the definition in eqs. (9) and (10) of \[4\] in that here we use canonically normalized fields throughout.

\(^\ddagger\) The gaugino mass given in ref. \[4\] (eqs. (8) and (15)) incorrectly includes a factor of \( \frac{1}{2} \). We thank L. Ibáñez for pointing this out.
canonically normalized.) The potential for the scalar fields $q^I$ in the supersymmetric multiplets takes the form

$$V^{\text{(eff)}}(q, \bar{q}) = \sum_a \frac{g_a^2}{4} (q^T a q)^2 + |\partial_I W^{\text{(eff)}}|^2$$

$$+ m_{I\bar{J}}^2 q^I \bar{q}^\bar{J} + \left( \frac{1}{3} A_{I\bar{J}L} q^I q^\bar{J} q^L + \frac{1}{2} B_{IJJ} q^I q^J \right) + \text{h.c.}$$  \hspace{1cm} (4)

The first two terms are the standard supersymmetric potential, whereas the last three are soft supersymmetry-breaking interactions. Their structure for dilaton-induced supersymmetry breaking is highly constrained and given by [5] (again neglecting string loops)

$$m_{I\bar{J}}^2 = \frac{m_{3/2}}{2} \delta_{I\bar{J}}, \quad A_{I\bar{J}L} = -\sqrt{3} \frac{m_{3/2}}{2} y_{I\bar{J}L};$$  \hspace{1cm} (5)

$B_{IJJ}$ also has only one non-vanishing entry – the coefficient of the Higgs doublets ($B_{12} \equiv B$) – but in general is not restricted further. From eq. (5) we learn that all scalar masses $m_{I\bar{J}}^2$ are flavour-independent (universal) and furthermore the $A$-terms are strictly proportional to the Yukawa couplings with a universal constant of proportionality. Both features are commonly assumed in phenomenological investigations of the MSSM, but generically do not hold in string theory. As a consequence, ‘dilaton breaking’ automatically ensures the smallness of flavour-changing neutral currents (FCNC). This is not guaranteed in other scenarios of supersymmetry breaking and in general imposes strong constraints on the perturbative couplings of the string vacuum [2]. The other distinct feature displayed by eqs. (3) and (5) is the fact that gaugino and scalar masses, as well as $A$-terms, are locked in terms of $m_{3/2}$ with no free parameter to vary. This leads to significant constrains on the low-energy mass spectrum and we find part of this spectrum directly determined by $m_{3/2}$. To summarize, the entire supersymmetric mass spectrum is expressed in terms of only three independent parameters $m_{3/2}$, $\mu$ and $B$.

This three-dimensional parameter space can be further reduced if one specifies the mechanism responsible for generating the $\mu$-term. Generically, there is a danger in string theory of inducing a large $\mu$, which prohibits a light Higgs. However, if $\mu$ arises from couplings in the Kähler potential (which do occur in string theory) its size is automatically $O(m_{3/2})$ [4]. For a $\mu$-term solely generated by this mechanism, $B$ is

\begin{itemize}
  \item Note that $B$ is not necessarily proportional to $\mu$.
  \item This corresponds to $\tilde{\mu} = 0$ in eqs. (2) and (9) of ref. [3].
\end{itemize}
no longer an independent parameter but instead obeys \( B = 2 \mu m_{3/2} \). In this case the mass spectrum is determined by two independent parameters, \( m_{3/2} \) and \( \mu \).

The mass relations of eqs. (3) and (5) should be viewed as a boundary condition at the unification scale \( M_X \) before (low-energy) renormalization effects are taken into account. The mass spectrum of the supersymmetric particles at the weak scale is determined by the evolution of the couplings according to their renormalization group (RG) equations. Here, we use the standard RG analysis where only the top-quark Yukawa coupling \( y_t \) is kept [11]. As the unification scale we choose \( M_X = 3 \times 10^{16} \) GeV in order to be consistent with the unification of the gauge couplings. String theory indeed implies a unification of gauge couplings; however, it occurs at the characteristic string scale, which is approximately \( 5 \times 10^{17} \) GeV and does not coincide with \( M_X \). There exist various suggestions of how to remedy this fact and we assume here that the string scale is effectively lowered by large threshold corrections [12, 13].

Let us turn to the Higgs sector, which is responsible for the electroweak symmetry breaking [8]. From eq. (4) we learn that the potential for the two neutral components \( h_1^0, h_2^0 \) of the Higgs doublets is given by

\[
V = \frac{1}{8}(g_1^2 + g_2^2)(|h_1^0|^2 - |h_2^0|^2)^2 + m_1^2|h_1^0|^2 + m_2^2|h_2^0|^2 - m_3^2(h_1^0h_2^0 + \text{h.c.}),
\]

with the boundary conditions at \( M_X \)

\[
m_1^2 = m_2^2 = m_{3/2}^2 + \mu^2, \quad m_3^2 = -B.
\]

These mass parameters evolve according to their RG equation, which can be solved analytically, as a function of the top Yukawa coupling \( y_t \) only. At low energies one finds (assuming (3) and (5) to hold) [11]

\[
\begin{align*}
m_1^2 &= c_1 m_{3/2}^2 + \mu_R^2, \\
m_2^2 &= c_2(y_t) m_{3/2}^2 + \mu_R^2, \\
m_3^2 &= c_4(y_t) B + c_5(y_t) \mu_R m_{3/2}, \\
\mu_R^2 &= c_3(y_t) \mu^2,
\end{align*}
\]

where \( c_{2-5} \) develop a (complicated) dependence on the unknown \( y_t \); their precise functional form can be found in ref. [11]. For the following analysis we only need

\( \nabla \) Our numerical calculation also takes into account the effects of the supersymmetry threshold.

\( \dagger \) Another possibility would be to assume extra light states in the spectrum, which decouple at some intermediate scale [14], but we do not entertain this option here.
to record that $c_1$ is independent of $y_t$ whereas $c_2$ obeys $c_2(y_t) \leq c_1$ and reaches its maximum at $y_t = 0$, i.e. $c_2(0) = c_1$. Furthermore, at low energies the renormalized $y_t$ cannot grow arbitrarily, but is instead ‘attracted’ by an infrared fixed point of its RG equation $y_t \to y_t^{\text{crit}}$. At that fixed point the coefficient $c_3$ vanishes: $c_3(y_t = y_t^{\text{crit}}) = 0$.

In order to induce electroweak symmetry breaking the renormalized masses have to satisfy

$$2m_3^2 < m_1^2 + m_2^2, \quad m_1^2 m_2^2 < m_3^4,$$  \hspace{1cm} (9)

and

$$M_Z^2 = 2 \frac{m_1^2 - m_2^2 \tan^2 \beta}{\tan^2 \beta - 1},$$ \hspace{1cm} (10)

where

$$\tan \beta = \frac{\langle h_2^0 \rangle}{\langle h_1^0 \rangle}, \quad \sin 2\beta = \frac{2m_3^2}{m_1^2 + m_2^2}, \quad \frac{\pi}{4} \leq \beta \leq \frac{\pi}{2}. \hspace{1cm} (11)$$

The relations (9) do not hold automatically but constrain the initial soft parameter space as we will see shortly. Even though $y_t$ is unknown, it cannot be viewed as an independent parameter because of the constraint equation (10). Which parameter one chooses to eliminate via (10) is a matter of convenience and taste. In our analysis, we eliminate $y_t$ and determine the top-quark mass $m_t (= y_t \langle h_2^0 \rangle)$ as a function of the soft parameters. Equivalently, one could trade $y_t$ for one of the soft parameters and use instead $m_t$ as an input parameter.

One of the distinct features of the supersymmetric Higgs potential (6) is the occurrence of a light Higgs boson. At the tree level its mass is given by

$$m_h^2 = \frac{1}{2} \left[ m_A^2 + M_Z^2 - \left( (m_A^2 + M_Z^2)^2 - 4m_A^2 M_Z^2 \cos^2 2\beta \right)^{1/2} \right],$$ \hspace{1cm} (12)

where $m_A^2 = m_1^2 + m_2^2$, and $m_1, m_2$ are as defined in eq. (8). In the limit $\tan \beta \to 1$, $m_h$ approaches zero.*

After these preliminaries we are in a position to discuss the mass spectrum of the supersymmetric particles. We start our investigation with the three-parameter case and afterwards consider the two-parameter scenario. We confine our attention

* It has recently been realized that one-loop corrections can significantly raise $m_h$, because of the heavy top quark \[15\]. In our numerical evaluation, we take this into account.
to those features of the mass spectrum that differ from the typical MSSM results. At an arbitrary RG scale \( p \), the gaugino masses obey

\[
\tilde{m}_a(p) = k_a \frac{\alpha_a(p)}{\alpha_X} \tilde{m}(M_X), \quad k_2 = k_3 = 1, \quad k_1 = 5/3.
\] (13)

Equation (3), together with \( \alpha_3(M_Z) \approx 0.118 \) and \( \alpha_X \approx 1/24 \), implies at low energies

\[
\tilde{m}_a = d_a m_{3/2}; \quad d_3 \approx 5, \quad d_2 \approx 1.5, \quad d_1 \approx 0.75.
\] (14)

(\( \tilde{m}_3 \) is the gluino and should not be confused with \( m_3 \) of eq. (8).) Again because of eq. (3) all squark and slepton masses \( m_{q_i} \) (except the stop mass) are determined by \( m_{3/2} \) (or equivalently \( \tilde{m}_3 \)). One finds \[11\]

\[
m^2_{q_i} = l_i \tilde{m}^2_3 + n_i M^2_Z, \quad \text{where} \quad l_i \geq 0.3, \quad -\frac{1}{2} < n_i < \frac{1}{2}.
\] (15)

The \( l_i \) are fixed numerical coefficients with no dependence on the soft parameters, whereas \( n_i \) depend on \( \tan \beta \). Instead of listing \( l_i \) and \( n_i \), we display the squark and lepton masses as a function of the gluino mass \( \tilde{m}_3 \) in fig. 1. Owing to the small second term in eq. (15), they lie within a tiny band, which is invisible in fig. 1. The slepton masses coincide to a very good approximation with \( 0.3 \tilde{m}_3 \), whereas the squark masses are essentially degenerate with \( \tilde{m}_3 \). Because of the (large) top Yukawa coupling, the left- and right-handed stop can have a large mixing term; as a consequence the stop mass is not accurately described by eq. (15). For large \( \mu \) it can be significantly lower than the other squark masses \[8\].

The masses of the four neutralinos \( \chi^0 \) (linear combinations of Higgsinos, photino and zino) are determined by the eigenvalues of a \( 4 \times 4 \) mass matrix with input parameters \( \tilde{m}_3, \mu_R \) and \( \tan \beta \) \[8\]. The scale of the lowest eigenvalue is set either by \( \tilde{m}_1 \) (\( \approx 0.16 \tilde{m}_3 \)) or by \( \mu_R \), whichever is lower. Since the lightest slepton mass is approximately \( 0.3 \tilde{m}_3 \) we immediately conclude that the lightest supersymmetric particle (LSP) is always a neutralino. Its mass range is similar to the mass range found in the standard MSSM analysis, which can be understood from the fact that in both cases the neutralino masses are determined by three independent parameters. The masses of the charginos \( \chi^\pm \) (linear combination of the charged Higgsino and the charged wino) are determined by the exact same three input parameters and, as a consequence, we
find no significant deviation from the MSSM in the chargino sector. A similar conclusion holds in the Higgs sector. In our numerical analysis we find no restriction of \( \tan \beta \) and consequently \( m_h \) also varies within the standard ranges over the allowed parameter space. This will change in the two-parameter case to which we now turn our attention.

The supersymmetric mass spectrum only depends on two soft parameters if \( \mu \) is generated by terms in the Kähler potential as was first suggested in ref. [9]. In this case \( B \) is no longer independent but obeys \( B = 2 \mu m_{3/2} \). With respect to the three-dimensional parameter space we just discussed, the main difference arises from the fact that \( \mu \) is now constrained to lie well above \( m_{3/2} \). Numerically we find that \( \mu > 0.4 \tilde{m}_3 \) and \( \tilde{m}_3 > 225 \text{ GeV} \) (16)

has to be satisfied in order to evade the experimental bounds on the top-quark mass \( (m_t \geq 108 \text{ GeV} \ [16]) \) and the Higgs mass \([17]\). In deriving (16) we first observe that the top mass limit alone pushes the mass of the pseudoscalar Higgs \( m_A \) well above 70 GeV (this can be seen from (8) and the second equation in (12)) and as a consequence the scalar Higgs becomes effectively the SM Higg with a lower bound of \( \approx 60 \text{ GeV} \ [17] \). The combined limits of \( m_t \geq 108 \text{ GeV} \) and \( m_h \geq 60 \text{ GeV} \) then result in the constraint (16). (For gluino masses close to their lower bound \( \mu \) has to be bigger than \( \tilde{m}_3 \).) In Fig. 2 we display the top and Higgs mass (along with the lightest neutralino and chargino) as a function of \( \mu \) for a fixed gluino mass \( \tilde{m}_3 = 400 \text{ GeV} \). We clearly see that the experimental bounds imply a large \( \mu \). Analytically, the constraint (16) can be understood from the fact that for \( \mu = m_{3/2} \) at \( M_X \) the Higgs mass matrix (which we can read off from eq. (6)) has a zero eigenvalue and (almost) causes an instability after RG effects are included. In addition, for small \( \mu \) eqs. (10) and (8) force \( y_t \) to very small values, which results in a low top-quark mass.

For large \( \mu \) we find a rather different behaviour. From eqs. (10) and (11) we immediately infer that a large \( \mu \) is only accessible if at the same time \( c_3 \rightarrow 0 \), such that \( \mu_R \) stays fixed. As we already indicated, this behaviour of \( c_3 \) precisely occurs for \( y_t \) approaching its infrared fixed point. Since the physical masses depend on \( \mu_R \),

\[ \mu > 0.4 \tilde{m}_3 \quad \text{and} \quad \tilde{m}_3 > 225 \text{ GeV} \]
they display an asymptotic behaviour as a function of $\mu$ which can be observed in Fig. 2. As a consequence, the lightest neutralino, chargino and stop as well as the pseudoscalar Higgs effectively depend only $\tilde{m}_3$ in the allowed region of large $\mu$. Fig. 3 displays their masses as a function of $\tilde{m}_3$; the stop and the chargino retain a weak dependence on $\mu$ which results in the ‘spread’ seen in the plot. The squark and slepton masses (except the stop) again obey eq. (15), where $n_i$ is now a pure number. Since this second term is small, Fig. 1 adequately summarizes the squark masses also in this case; the stop is the lightest squark with a mass well above $m_t$. Thus, for most of the supersymmetric particles the original two-dimensional parameter space effectively reduces to a one-dimensional space with the masses determined solely by $\tilde{m}_3$.

For the Higgs and the top quark the parameter space remains two-dimensional. However, from eqs. (10) and (8) one infers that for large $\mu$, $\tan\beta$ is almost independent of $\tilde{m}_3$ which results in an upper bound $\tan\beta \leq 2$. This in turn, implies an upper bound on the top mass $m_t \leq 180$ GeV and leads to a relatively light Higgs boson all over the allowed parameter space (as can be seen from eq. (12)). In Fig. 4 we show the (one-loop corrected) Higgs mass as a function of $m_t$ for different values of $\tilde{m}_3$. We see that both top and Higgs are constrained and only for a large gluino and a large top mass the Higgs can be heavier than the $Z$-boson. In addition, for fixed gluino mass there is a linear correlation between Higgs and top mass.

Let us conclude. We investigated the low-energy supersymmetric mass spectrum, which arises under the assumption that supersymmetry is spontaneously broken by the dilaton $F$-term. Owing to the universal couplings of the dilaton, the structure of the soft parameters as given in eqs. (3) and (5) holds for all $N = 1$ vacua of the heterotic string. Compared with other scenarios of supersymmetry breaking in string theory, they display some simplicity and have a phenomenological appeal. The scalar and gaugino masses as well as the $A$-terms are automatically universal (a feature that generically does not hold in string theory) and determined in terms of $m_{3/2}$. Without specifying the mechanism for generating the $\mu$-term, the supersymmetric mass spectrum is determined in terms of only three independent parameters $m_{3/2}$, $\mu$, and $B$, and as a consequence the masses are slightly more constrained than in the MSSM. For a $\mu$-term induced by the Kähler potential, $B$ is related to $\mu$ and the parameter space becomes two-dimensional. Current experimental limits of the Higgs and the top-quark mass further constrain the range of $\mu$ and lead to a reduction of the parameter space for most of the supersymmetric particles. $\tan\beta$ is always small and as a consequence the Higgs is rather light whereas the top mass is bounded by 180 GeV.
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Note added

In a previous version of this article the gaugino masses (eq. (3)) were given incorrectly. We thank L. Ibáñez for correcting this error.

After completion of this paper we received a preprint by J. López, D. Nanopoulos and A. Zichichi, (CERN-TH.6903/93), which analyses supersymmetry breaking by the dilaton in the context of a flipped $SU(5)$ model.
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Figure captions

Fig. 1 Masses of squarks (except stop) $Q_L (-)$, $U_R$, $D_R (- - -)$, and sleptons $L_L (- - -)$, $E_R (\cdots)$ as a function of the gluino mass $\tilde{m}_3$.

Fig. 2 Masses of top quark ($-$), lightest scalar Higgs boson (- - -), lightest neutralino ($\cdots$) and lightest chargino (\cdots) as a function of $\mu$ for a gluino mass $\tilde{m}_3 = 400\, GeV$.

Fig. 3 Masses of pseudoscalar Higgs (- - -), lightest stop ($-$), lightest chargino (- - -) and lightest neutralino ($\cdots$) for large $\mu$ as a function of the gluino mass $\tilde{m}_3$.

Fig. 4 Higgs mass as a function of top-quark mass for gluino masses $\tilde{m}_3 = 250, 500, 1000, 1500\, GeV$. 

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