

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Dilaton domination in the MSSM and its singlet extensions

Jan Louis^{a,b,c}, Kai Schmidt-Hoberg^d, Lucila Zarate^a^a *Fachbereich Physik der Universität Hamburg, Luruper Chaussee 149, D-22761 Hamburg, Germany*^b *Zentrum für Mathematische Physik, Universität Hamburg, Bundesstrasse 55, D-20146 Hamburg, Germany*^c *Albert Einstein Minerva Center, Weizmann Institute of Science, Rehovot 76100, Israel*^d *Theory Division, CERN, 1211 Geneva 23, Switzerland*

ARTICLE INFO

Article history:

Received 21 February 2014

Received in revised form 24 April 2014

Accepted 19 May 2014

Available online 23 May 2014

Editor: A. Ringwald

ABSTRACT

We analyse the current status of the dilaton domination scenario in the MSSM and its singlet extensions, taking into account the measured value of the Higgs mass, the relic abundance of dark matter and constraints from SUSY searches at the LHC. We find that in the case of the MSSM the requirement of a dark matter relic abundance in accord with observation severely restricts the allowed parameter space, implying an upper bound on the superpartner masses which makes it fully testable at the LHC-14. In singlet extensions with a large singlet-MSSM coupling λ as favoured by naturalness arguments the coloured sparticles should again be within the reach of the LHC-14, while for small λ it is possible to decouple the MSSM and singlet sectors, achieving the correct dark matter abundance with a singlino LSP while allowing for a heavy MSSM spectrum.

© 2014 The Authors. Published by Elsevier B.V. This is an open access article under the CC BY license (<http://creativecommons.org/licenses/by/3.0/>). Funded by SCOAP³.

1. Introduction

With the discovery of the Higgs boson and no evidence for supersymmetry there is a renewed interest in non-minimal versions of the supersymmetric Standard Model. In particular in singlet extensions such as the NMSSM an additional neutral chiral multiplet S with specific couplings to the Higgs sector is added to the field content of the minimal supersymmetric Standard Model (MSSM) (see e.g. [1,2] for reviews on the subject). Originally this model was suggested as a dynamical mechanism to generate the supersymmetric μ -term with its size being naturally of order $\mathcal{O}(M_Z)$. However, due to the mixing of the scalar components of S with the Higgs sector, a Higgs mass $m_h \simeq 125$ GeV is obtained with more ease. This is due to an additional tree-level F -term contribution to the Higgs mass, which also helps alleviating the notorious fine tuning problem of the MSSM. This has been observed in the \mathbb{Z}_3 symmetric NMSSM [3] and particularly in its generalised version (GNMSSM) [4–7], see also [8,9].

It is of obvious interest to embed the NMSSM into string backgrounds. Since scalar fields are abundant in string theory, singlet extensions appear to be a promising low energy limit. The corresponding string model building has as usual two aspects. On the one hand, there is the need to construct explicit string backgrounds with the additional chiral multiplet S and its specific

couplings.¹ On the other hand, the mechanism for supersymmetry breaking determines the soft terms and thus the specific low energy particle spectrum.

In this paper we focus on the second aspect in that we do not attempt to construct a string background but instead study the effects of string motivated soft terms. Specifically we concentrate on the situation where the supersymmetry breaking occurs in the dilaton direction of the heterotic string [12]. This leads to a rather constrained set of (universal) soft terms – termed “dilaton dominated soft terms” – which were analysed within the MSSM in Refs. [13–19].

It has been argued that dilaton domination in the MSSM suffers from the existence of charge and colour breaking vacua which are deeper than the local electroweak vacuum [18]. Subsequently, however, it has been realised that, while these deeper vacua exist, the lifetime of the electroweak vacuum is almost always considerably longer than the age of the Universe [20,21] (see e.g. Fig. 1 of [21] for typical lifetimes). This renders the existence of such deeper vacua effectively harmless and the dilaton domination scenario phenomenologically viable (see also [22]).

In this paper we reconsider the MSSM with dilaton dominated soft terms and then extend the analysis to the NMSSM. In the MSSM there are only two independent parameters after electroweak symmetry breaking has been implemented. Within

E-mail addresses: jan.louis@desy.de (J. Louis), kai.schmidt-hoberg@cern.ch (K. Schmidt-Hoberg), lucila.zarate@desy.de (L. Zarate).

¹ See, for example, Refs. [10,11].

this two-dimensional parameter space one can obtain the observed Higgs mass for large $\tan\beta$ and simultaneously satisfy all other current experimental constraints. Nevertheless, the thermal relic abundance is generically too large and the requirement of a dark matter relic abundance in accord with observation severely restricts the allowed parameter space, implying an upper bound on the superpartner masses which makes it fully testable at the LHC-14. This also implies that the leading radiative corrections to the SM-like Higgs mass cannot be large enough to push the Higgs mass to its observed central value, resulting in a Higgs mass that is acceptable given theoretical and experimental uncertainties, but slightly low.

As mentioned above, singlet extensions benefit from a sizeable additional contribution to the diagonal component of the tree-level Higgs mass matrix in the small $\tan\beta$ regime if the MSSM-singlet coupling λ is large. This will result in an enhanced lightest Higgs mass if the mixing within the Higgs sector is small. While this is not possible in the usual \mathbb{Z}_3 symmetric NMSSM, we will consider very well motivated generalised versions of the NMSSM where small mixings can be natural, resolving the tension between the Higgs mass and a viable dark matter sector. The mass scale of the MSSM superpartners is however still very constrained, as for large λ the singlet states have to be rather heavy in order to avoid large mixings, which in turn implies a bino-like LSP whose relic abundance can only be kept small enough for not too heavy MSSM states. A possibility to circumvent this conclusion is the decoupling regime of small λ , where the LSP can be singlino-like with the correct relic abundance while all MSSM states can be heavy.

This paper is organized as follows. In Section 2 we introduce the MSSM and NMSSM and recall the dilaton dominated soft terms. In Section 3.2 we compute the low energy spectrum for the MSSM while in Section 3.3 we repeat the analysis for the NMSSM.

2. Dilaton dominated supersymmetry breaking

In this section we review the high scale boundary conditions for dilaton dominated supersymmetry breaking in the MSSM and NMSSM in order to set the stage for the following sections. Let us start with the MSSM.

2.1. Dilaton domination in the MSSM

The superpotential of the MSSM is given by

$$\begin{aligned} \mathcal{W}_{\text{MSSM}} = & \mu_h H_u H_d + \sum_{\text{generations}} (y_U Q_L U_R H_u + y_D Q_L D_R H_d \\ & + y_L L_L E_R H_d), \end{aligned} \quad (1)$$

where Q_L are the quark doublets, U_R, D_R are the quark singlets and $H_{u,d}$ are the two Higgs multiplets. For the MSSM the dilaton dominated soft supersymmetry breaking terms are a universal scalar mass m_0 (which coincides with the gravitino mass $m_{3/2}$) and a B_h -parameter generated at some high scale such as M_{GUT} . The (canonically normalised) gaugino masses $m_{1/2}$ and the A -terms are also universal and at leading order related to m_0 by [12]

$$m_{1/2} = \sqrt{3}m_0, \quad A_0 = -\sqrt{3}m_0. \quad (2)$$

Thus before imposing electroweak symmetry breaking there are three independent parameters μ_h, m_0 and B_h . Requiring correct electroweak symmetry breaking and imposing $v \sim 246$ GeV will effectively fix one of these parameters (up to a discrete choice). In fact it is common practice to use the vacuum equations to replace μ_h and B_h by $\tan\beta \equiv v_u/v_d$ and $\text{sign}(\mu_h)$, leaving $m_0, \tan\beta$ and $\text{sign}(\mu_h)$ as free parameters. For the more constrained case where

μ_h is generated by the Giudice–Masiero mechanism [23] the original parameter space is only two-dimensional as the additional constraint $B_h = 2\mu_h m_0$ holds [12]. This latter case necessarily has a light Higgs which has been ruled out already for some time [13, 14]. We will therefore concentrate on the general case with three independent parameters in the next section and indeed find that it is compatible with current experimental constraints.

Before we proceed, let us briefly discuss the corrections to the universal scalar mass and the relations given in Eq. (2). First of all there is the anomaly mediated contribution to the gaugino masses $\delta m_{1/2} = \frac{b}{16\pi^2} m_0$ where b is the one-loop coefficient of the β -function [24,25]. In addition, there are loop corrections to the dilaton Kähler potential which induce further corrections to m_0 and (2) [15,19]. The size of the latter is difficult to estimate and also depends on the specific scenario under consideration. In the following we will assume a setup in which these model-dependent corrections are small and the model independent tree-level relations (2) hold to good approximation.

2.2. Dilaton domination in singlet extensions

The most general extension of the MSSM by a gauge singlet chiral superfield S has a superpotential of the form

$$\mathcal{W}_{\text{GNMSSM}} = \mathcal{W}_{\text{MSSM}} + \lambda S H_u H_d + \xi S + \frac{1}{2} \mu_s S^2 + \frac{1}{3} \kappa S^3, \quad (3)$$

where λ, ξ, μ_s and κ are conventionally chosen as real parameters. As the \mathbb{Z}_3 symmetric version with $\mu_h = \mu_s = \xi = 0$ is usually referred to as *the* NMSSM,² we will denote this more general case as the GNMSSM.

The $\mathcal{W}_{\text{GNMSSM}}$ given in Eq. (3) seems to have a ‘double’ μ problem, as the SM symmetries do not prevent arbitrarily high scales for the dimensionful mass terms. However these terms can be naturally of order the supersymmetry breaking scale if there is an underlying \mathbb{Z}_4^R or \mathbb{Z}_8^R symmetry [27]. In such a case the superpotential is identical to the \mathbb{Z}_3 symmetric NMSSM before supersymmetry breaking with $\mu_h = \mu_s = \xi = 0$. However, after supersymmetry breaking these additional superpotential terms are generated with a size that is set by the supersymmetry breaking scale in the visible sector, i.e. the gravitino mass $m_{3/2}$. Not only is the fine tuning particularly promising in this setup, but also the severe tadpole [28] and domain wall [29] problems of the \mathbb{Z}_3 symmetric NMSSM are avoided. For broken supersymmetry also the discrete R symmetry is broken but the subgroup \mathbb{Z}_2^R , corresponding to the usual matter parity, remains unbroken [27]. As a result the lightest supersymmetric particle (LSP) is stable and a candidate for dark matter.

The general soft supersymmetry breaking terms associated with the Higgs and singlet sectors are

$$\begin{aligned} V_{\text{soft}} = & m_s^2 |s|^2 + m_{h_u}^2 |h_u|^2 + m_{h_d}^2 |h_d|^2 + (B_h h_u h_d + \lambda A_\lambda s h_u h_d \\ & + \frac{1}{3} \kappa A_\kappa s^3 + \frac{1}{2} B_s s^2 + \xi_s s + \text{h.c.}). \end{aligned} \quad (4)$$

As in the MSSM, the soft supersymmetry breaking terms are universal at leading order in the dilaton domination scenario. All soft scalar masses including the soft mass m_s^2 are given by a universal m_0^2 while the gaugino masses and A -terms obey (2) unchanged,

$$\begin{aligned} m_{1/2} = \sqrt{3}m_0, \quad A_0 = A_\lambda = A_\kappa = -\sqrt{3}m_0, \\ m_{h_u} = m_{h_d} = m_s = m_0. \end{aligned} \quad (5)$$

For the B terms we again have two choices. They are either independent or, in the more constrained situation when μ_h and μ_s

² Dilaton domination in this \mathbb{Z}_3 symmetric NMSSM was also considered in [26].

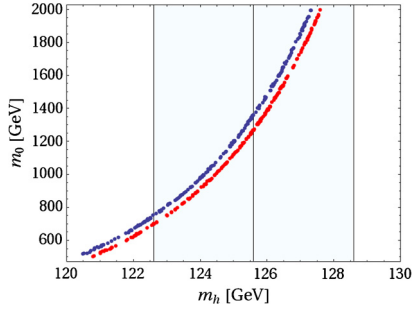


Fig. 1. Higgs mass as a function of m_0 for large $\tan\beta$ and $\mu_h > 0$ (red) and $\mu_h < 0$ (blue) in the MSSM. (For interpretation of the references to colour in this figure legend, the reader is referred to the web version of this article.)

are generated by the Giudice–Masiero mechanism, related by $B_h = 2\mu_h m_0$, $B_s = 2\mu_s m_0$. Thus before electroweak symmetry breaking is imposed the parameter space is either nine- or seven-dimensional with parameters $(m_0, \mu_h, B_h, \lambda, \kappa, \mu_s, B_s, \xi, \xi_s)$ or $(m_0, \mu_h, \lambda, \kappa, \mu_s, \xi, \xi_s)$ respectively. Requiring the correct electroweak symmetry breaking and imposing $v \sim 246$ GeV will effectively fix one of these parameters. In addition the vacuum equations allow us to replace two of the input parameters of the model by $\tan\beta$ and the vacuum expectation value (VEV) of s which we denote as v_s . Moreover, for the case that all variables are taken to be independent, i.e. in the case where no constraint on the B -terms is applied, one of the dimensional parameters can be eliminated by a shift in v_s , which we take to be ξ , i.e. we choose $\xi = 0$. In summary, after electroweak symmetry breaking one ends up with seven independent input parameters in the general case, e.g. $(m_0, \tan\beta, \lambda, \kappa, \mu_s, B_s, v_s)$ and six in the constrained situation, $(m_0, \tan\beta, \mu_h, \lambda, \kappa, \mu_s)$.

3. Phenomenology

3.1. SUSY, Higgs and DM cuts

In this section we briefly describe the cuts that we impose for the numerical analyses of the MSSM and the GNMSSM. For the (first generation of) squarks and gluinos we use a cut of $m_{\tilde{q}} > 1.7$ TeV and $m_{\tilde{g}} > 1.7$ TeV. We further require the chargino and slepton masses to be above 100 GeV. We also require that the lightest supersymmetric particle (LSP) is a neutralino which is a good dark matter candidate and its relic density is within the 5σ PLANCK [30] range of $0.1064 \leq \Omega h^2 \leq 0.1334$. In addition to constraints from the relic density there are constraints from dark matter direct detection searches which limit the cross section of the lightest neutralino to nucleons. We require that the points are consistent with dark matter direct detection constraints, in particular with the latest constraint from LUX [31]. Finally, for the Higgs mass we take the average of the CMS and ATLAS best fit values of 125.7 GeV [32] and 125.5 GeV [33] respectively and allow for a 3 GeV uncertainty, $m_h = 125.6 \pm 3$ GeV.

For our numerical analyses we use SPheno [34,35] created by SARAH [36–39]. This version performs a complete one-loop calculation of all SUSY and Higgs masses [40,41] and includes the dominant two-loop corrections for the scalar Higgs masses [42–45]. The dark matter relic density as well as the direct detection bounds are calculated with MicrOmegas [46–48].

3.2. The MSSM

As explained in Section 2 our independent input parameters are $m_0, \tan\beta$ and $\text{sign}(\mu_h)$. In Fig. 1 we show the resulting dependence of the Higgs mass on m_0 for the case that $\tan\beta$ is large. It

can be seen that there is no problem to achieve a large enough Higgs mass via radiative corrections.³ In fact for $\tan\beta > 10$ there is almost a one-to-one correspondence between m_0 and the resulting Higgs mass for a given $\text{sign}(\mu_h)$. The minimal values of m_0 to achieve a Higgs mass consistent with observation within the theoretical and experimental uncertainties are $m_0 = 695$ GeV ($m_0 = 760$ GeV) for $\mu_h > 0$ ($\mu_h < 0$), well above the value of $m_0 \gtrsim 500$ GeV required by the SUSY cuts. The precise numbers are however rather sensitive to the top Yukawa coupling, which enters the dominant radiative corrections. If we were to take the one sigma upper limit on the top mass rather than its central value, radiative corrections are somewhat increased and the lower bound on m_0 is relaxed to $m_0 = 615$ GeV ($m_0 = 660$ GeV) for $\mu_h > 0$ ($\mu_h < 0$). In the following discussion we will stick to the central value of the top mass. (See Fig. 2.)

Turning to the neutralino sector, we find that the LSP is always an almost pure bino, with mass $m_{\tilde{\chi}_0^1} \simeq 0.8 \cdot m_{3/2}$, somewhat below the gravitino mass. The fact that the gravitino is rather close in mass to the neutralino is potentially dangerous cosmologically. Whether this scenario is viable or not however depends on the cosmological history of our Universe, which we are agnostic about in this letter. Independently, a bino LSP implies that the thermal relic abundance is generically too large because of a rather small annihilation cross section. One therefore has to go to special regions in parameter space where the relic abundance can be small. Generically there is a tension between large m_0 implying a heavy SUSY spectrum (needed in this setup to achieve the correct Higgs mass) and a small enough relic abundance. One possible option would be to have coannihilations with another particle such as the stau, if the mass difference between neutralino and stau is sufficiently small. This can be achieved even for our very restrictive boundary conditions (for positive μ_h), but just about falls short of raising the Higgs mass into the observed window. In fact for the case of positive μ_h the maximal value for m_0 for which the correct relic abundance can be achieved is $m_0 \simeq 675$ GeV, whereas the minimal value to obtain a large enough Higgs mass is $m_0 \simeq 695$ GeV. Another option to deplete the relic abundance would be to have a mediating particle close to resonance, e.g. $m_A \sim 2m_{\tilde{\chi}}$. This can also be realised (for negative μ_h) and turns out to be more promising. In this case we do find viable points within the theoretical uncertainties, with an allowed range of $760 \text{ GeV} \leq m_0 \leq 875 \text{ GeV}$.

Given that the requirement of a relic abundance in accordance with observation is crucial to establishing an upper bound on the superpartner mass scale, let us discuss in some more detail how this comes about. The annihilation cross section which sets the relic abundance is maximised on the pole, where it is proportional to $g_{\tilde{\chi}\tilde{\chi}A}^2/\Gamma_A^2$, with $g_{\tilde{\chi}\tilde{\chi}A}$ the LSP–LSP–Higgs coupling and Γ_A the width of the mediating Higgs scalar. The width Γ_A scales directly with m_A and hence with m_0 , decreasing the annihilation cross section with increasing SUSY breaking scale. Furthermore the coupling $g_{\tilde{\chi}\tilde{\chi}A}$ originates from the Higgs–Higgsino–Gaugino term and is zero for a pure bino.⁴ The coupling is therefore proportional to the (small) Higgsino component of the LSP, which decreases with increasing m_0 (see Fig. 3). The combination of these two effects quickly makes sufficiently efficient annihilation impossible, even on the pole, resulting in this rather stringent upper bound on m_0 .

³ This of course ignores the fine tuning problem.

⁴ One might worry that in the case of an almost pure bino where the tree-level coupling is very suppressed the one-loop coupling (see e.g. [49]) might give a non-negligible correction. We find however that for the case of interest the one-loop coupling gives a correction of $\mathcal{O}(\%)$ and hence should not change the picture appreciably.

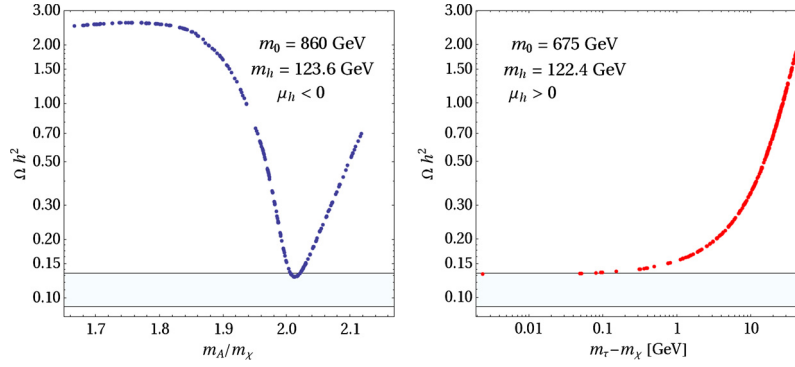


Fig. 2. The MSSM case. Left panel: $\mu < 0$. Right panel: $\mu > 0$. A small enough relic abundance can only be achieved for sufficiently small values of m_0 – we show the critical values for both cases. It can be seen that in order to achieve a sufficiently small relic abundance the process has to be resonant with $m_A \sim 2m_\chi$ for $\mu_h < 0$ or in the stau coannihilation region in the case of $\mu_h > 0$. Only in the case $\mu_h < 0$ do we achieve a Higgs mass consistent within the uncertainties.

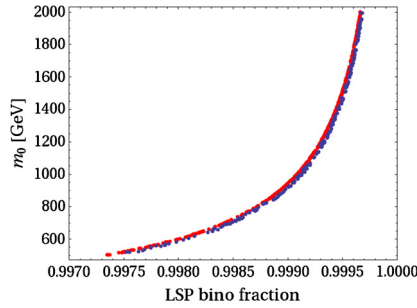


Fig. 3. LSP bino fraction as a function of m_0 for large $\tan\beta$ and $\mu_h > 0$ (red) and $\mu_h < 0$ (blue) in the MSSM. (For interpretation of the references to colour in this figure legend, the reader is referred to the web version of this article.)

This leaves us with a very predictive scenario for dilaton domination in the MSSM, as all superpartner masses scale directly with m_0 . In particular the allowed range of m_0 corresponds to gluino and squark masses of $2750 \text{ GeV} \lesssim m_{\tilde{g}} \lesssim 3150 \text{ GeV}$ and $2450 \text{ GeV} \lesssim m_{\tilde{q}} \lesssim 2800 \text{ GeV}$ respectively, which should be fully testable with the upcoming LHC-14 run [50,51]. While this is a viable scenario given the theoretical and experimental uncertainties, an obvious shortcoming is that it is not possible to raise the Higgs mass to its observed central value. This is different in singlet extensions, which we will discuss next.

3.3. Singlet extensions

The superpotential and soft terms of the general NMSSM have been given in Section 2. Before going to the most general case, let us briefly discuss some restricted cases that might be of interest. To start with, the very well known \mathbb{Z}_3 symmetric NMSSM requires $A_k^2 \gtrsim 9m_s^2$ (see e.g. [1]) in order to have a real VEV for S , which is not possible in the dilaton domination scenario as the relations in (5) show. Another problem is that the region of large λ and small $\tan\beta$ which is preferred due to fine tuning considerations is phenomenologically unacceptable as mixing effects decrease the lightest Higgs mass, overcompensating the additional tree-level contribution. That the \mathbb{Z}_3 symmetric NMSSM does not work in the context of dilaton domination was also found in [26]. In contrast the GNMSSM allows for a natural limit of heavy singlet states, suppressing the mixing in the Higgs sector. In Table 1 we show example points for some cases of interest, including a case in which both linear terms are set to zero, $\xi = \xi_s = 0$, and where the constraints on the B-terms hold, $B_h = 2\mu_h m_0$, $B_s = 2\mu_s m_0$.

To illustrate the general dependence on the different input parameters we performed a scan for the general case in the region of interest corresponding to small $\tan\beta$ and large λ . In Fig. 4 we

show the dependence of the Standard Model like Higgs mass on the input parameters $\tan\beta$, λ and μ_s , which are most relevant. Note that λ and μ_s are the GUT scale parameters. One can clearly see the increased Higgs mass for small $\tan\beta$ and large λ . It can also be seen that the singlet mass term μ_s has to be sufficiently large to avoid too large mixing in the Higgs sector.

For the neutralino sector this implies that the LSP is again a bino-like neutralino, as the singlino is typically rather heavy. To achieve the correct relic abundance one is therefore forced to special regions as in the MSSM case. Given that the stau coannihilation region cannot be accessed for small values of $\tan\beta$ with the given boundary conditions, the only remaining option is resonant annihilation via the Higgs funnel, implying $m_A \sim 2m_\chi$. It turns out that for the GNMSSM, due to the additional tree-level contribution, there is no tension between a large enough Higgs mass and the relic abundance. Nevertheless, the mass scale of the MSSM superpartners is still very constrained, as the annihilation channel ceases to be effective for too large m_0 , as the width as well as the bino fraction increases, similar to the MSSM case.⁵ We find it remarkable that in spite of the many free parameters, the dilaton domination turns out to be rather predictive even in general singlet extensions of the MSSM in the large λ regime. The only possibility to circumvent this conclusion is to go to the decoupling regime of small λ , where the LSP can be singlino-like with the correct relic abundance while all MSSM states can be heavy. We also checked the direct detection cross section for all cases and find that they are significantly below the current bound set by LUX, which is not surprising given that the predominant LSP component is either bino or singlino.

Acknowledgements

We would like to thank Florian Staub for very useful discussions. We would also like to thank Ricardo Marquez for his persistent help and support. J.L. thanks the Weizmann Institute for their kind hospitality during the final stages of this work. This work was supported by the German Science Foundation (DFG) within the Collaborative Research Center (SFB) 676 “Particles, Strings and the Early Universe”, G.I.F. – the German–Israeli Foundation for Scientific Research and Development (GIF I-1-03847.7/2009) and the I-CORE program of the Planning and Budgeting Committee and the Israel Science Foundation (grant number 1937/12).

⁵ For an explicit expression of the LSP–LSP–Higgs coupling in the case of general NMSSM singlet extensions (which trivially reduces to the MSSM case) see e.g. the appendix of [52].

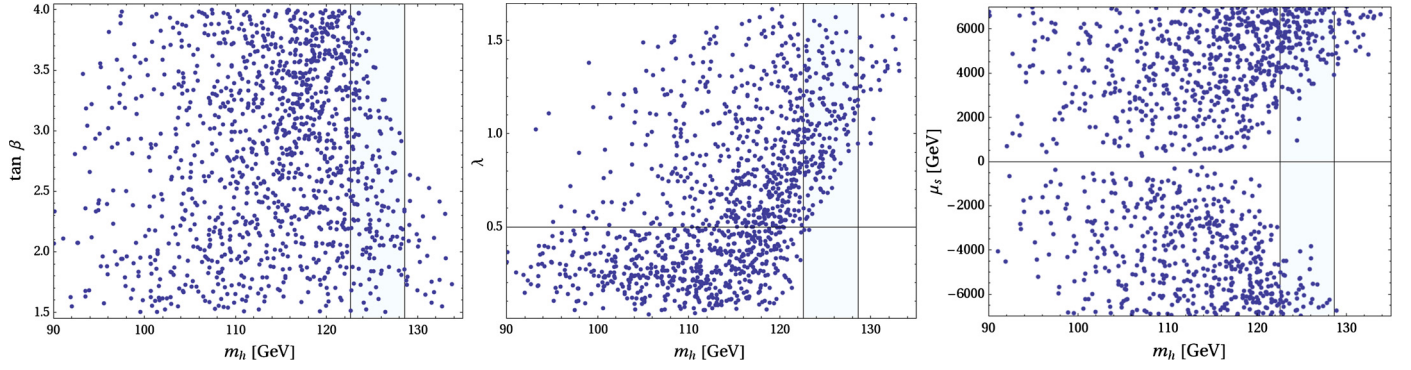


Fig. 4. Dependence of the Standard Model like Higgs mass on $\tan\beta$, λ and μ_s for the GNMSSM. For these figures we have applied the appropriate SUSY cuts, but not the cut on the dark matter relic abundance.

Table 1

Benchmark points for the MSSM and the GNMSSM. All input parameters except $\tan\beta$ and v_s are given at the GUT scale. Values marked with a * are output values at the electroweak scale determined by the electroweak symmetry breaking conditions.

	MSSM	GNMSSM1	GNMSSM2	GNMSSM3	GNMSSM4
m_0 [GeV]	860	500	500	500	1400
$\tan\beta$	36.3	2.7	2.8	2.7	33
μ_h [GeV]	tad	1060*	1000	1350*	2700*
B_h [GeV ²]	tad	$2.7 \cdot 10^{5*}$	$2m_0\mu$	$9.7 \cdot 10^{5*}$	$2.6 \cdot 10^{5*}$
λ	–	1.5	1.5	1.5	0.003
κ	–	–0.5	–0.58	–0.5	0.14
v_s [GeV]	–	500	1330*	–102*	3440
μ_s [GeV]	–	–5000	5400	–5000	–128
B_s [GeV ²]	–	2850^2	$2m_0\mu_s$	$8.2 \cdot 10^6$	$6.3 \cdot 10^5$
ξ [GeV ²]	–	0	$-2.5 \cdot 10^{6*}$	0	0
ξ_s [GeV ³]	–	$-3.7 \cdot 10^{9*}$	$-3.3 \cdot 10^{9*}$	0	$-4.3 \cdot 10^{9*}$
$m_{\tilde{g}}$ [GeV]	3100	1900	1900	1900	4800
m_{squark} [GeV]	2100–2900	1600–1800	1300–1750	1300–1750	3300–4500
m_{slepton} [GeV]	800–1300	600–750	600–750	600–750	1300–2050
$m_{\tilde{\chi}_1^\pm}$ [GeV]	1250	710	710	710	2050
m_{h_1} [GeV]	123.4	125.7	123.6	127.2	125.9
m_{h_2} [GeV]	1385	1600	1550	1580	1140
m_{A_1} [GeV]	1410	785	785	785	1680
$m_{\tilde{\chi}_1^0}$ [GeV]	700	390	390	390	540
$\tilde{\chi}_1^0$ bino part	0.999	0.998	0.998	0.998	$\mathcal{O}(10^{-9})$
$\tilde{\chi}_1^0$ wino part	$\mathcal{O}(10^{-6})$	$\mathcal{O}(10^{-4})$	$\mathcal{O}(10^{-4})$	$\mathcal{O}(10^{-4})$	$\mathcal{O}(10^{-10})$
$\tilde{\chi}_1^0$ higgsino part	0.001	0.002	0.002	0.002	$\mathcal{O}(10^{-8})$
$\tilde{\chi}_1^0$ singlino part	–	$\mathcal{O}(10^{-5})$	$\mathcal{O}(10^{-5})$	$\mathcal{O}(10^{-5})$	1
Ωh^2	0.13	0.11	0.12	0.11	0.11
σ_p [cm ²]	$\sim 10^{-48}$	$\sim 10^{-46}$	$\sim 10^{-46}$	$\sim 10^{-46}$	$\sim 10^{-49}$

References

- [1] U. Ellwanger, C. Hugonie, A.M. Teixeira, Phys. Rep. 496 (2010) 1–77, arXiv:0910.1785.
- [2] M. Maniatis, Int. J. Mod. Phys. A 25 (2010) 3505–3602, arXiv:0906.0777.
- [3] R. Dermisek, J.F. Gunion, Phys. Rev. D 76 (2007) 095006, arXiv:0705.4387.
- [4] G.G. Ross, K. Schmidt-Hoberg, Nucl. Phys. B 862 (2012) 710–719, arXiv:1108.1284.
- [5] G.G. Ross, K. Schmidt-Hoberg, F. Staub, J. High Energy Phys. 1208 (2012) 074, arXiv:1205.1509.
- [6] A. Kaminska, G.G. Ross, K. Schmidt-Hoberg, J. High Energy Phys. 1311 (2013) 209, arXiv:1308.4168.
- [7] A. Kaminska, G.G. Ross, K. Schmidt-Hoberg, F. Staub, arXiv:1401.1816, 2014.
- [8] A. Delgado, C. Kolda, J.P. Olson, A. de la Puente, Phys. Rev. Lett. 105 (2010) 091802, arXiv:1005.1282.
- [9] A. Delgado, C. Kolda, J.P. Olson, A. de la Puente, Phys. Rev. D 82 (2010) 035006, arXiv:1005.4901.
- [10] O. Lebedev, S. Ramos-Sanchez, Phys. Lett. B 684 (2010) 48–51, arXiv:0912.0477.
- [11] M. Cvetič, J. Halverson, P. Langacker, J. High Energy Phys. 1009 (2010) 076, arXiv:1006.3341.
- [12] V.S. Kaplunovsky, J. Louis, Phys. Lett. B 306 (1993) 269–275, arXiv:hep-th/9303040.
- [13] R. Barbieri, J. Louis, M. Moretti, Phys. Lett. B 312 (1993) 451–460, arXiv:hep-ph/9305262.
- [14] A. Brignole, L.E. Ibanez, C. Munoz, Nucl. Phys. B 422 (1994) 125–171, arXiv:hep-ph/9308271.
- [15] J. Louis, Y. Nir, Nucl. Phys. B 447 (1995) 18–34, arXiv:hep-ph/9411429.
- [16] G. Kraniotis, Z. Phys. C 71 (1996) 163–170, arXiv:hep-ph/9507432.
- [17] S. Khalil, A. Masiero, F. Vissani, Phys. Lett. B 375 (1996) 154–162, arXiv:hep-ph/9511284.
- [18] J. Casas, A. Lleyda, C. Munoz, Phys. Lett. B 380 (1996) 59–67, arXiv:hep-ph/9601357.
- [19] J. Casas, Phys. Lett. B 384 (1996) 103–110, arXiv:hep-th/9605180.
- [20] A. Riotto, E. Roulet, Phys. Lett. B 377 (1996) 60–66, arXiv:hep-ph/9512401.
- [21] J.L. Evans, D.E. Morrissey, J.D. Wells, Phys. Rev. D 80 (2009) 095011, arXiv:0812.3874.
- [22] S. Abel, B. Allanach, F. Quevedo, L. Ibanez, M. Klein, J. High Energy Phys. 0012 (2000) 026, arXiv:hep-ph/0005260.
- [23] G. Giudice, A. Masiero, Phys. Lett. B 206 (1988) 480–484.

- [24] L. Randall, R. Sundrum, Nucl. Phys. B 557 (1999) 79–118, arXiv:hep-th/9810155.
- [25] G.F. Giudice, M.A. Luty, H. Murayama, R. Rattazzi, J. High Energy Phys. 9812 (1998) 027, arXiv:hep-ph/9810442.
- [26] G. Kraniotis, Z. Phys. C 68 (1995) 491–494, arXiv:hep-ph/9501393.
- [27] H.M. Lee, et al., Nucl. Phys. B 850 (2011) 1–30, arXiv:1102.3595.
- [28] S.A. Abel, Nucl. Phys. B 480 (1996) 55–72, arXiv:hep-ph/9609323.
- [29] S.A. Abel, S. Sarkar, P.L. White, Nucl. Phys. B 454 (1995) 663–684, arXiv:hep-ph/9506359.
- [30] Planck Collaboration, P. Ade, et al., arXiv:1303.5076, 2013.
- [31] LUX Collaboration, D. Akerib, et al., arXiv:1310.8214, 2013.
- [32] CMS Collaboration, 2013, CMS-PAS-HIG-13-005.
- [33] ATLAS Collaboration, 2013, ATLAS-CONF-2013-014.
- [34] W. Porod, Comput. Phys. Commun. 153 (2003) 275–315, arXiv:hep-ph/0301101.
- [35] W. Porod, F. Staub, Comput. Phys. Commun. 183 (2012) 2458–2469, arXiv:1104.1573.
- [36] F. Staub, arXiv:0806.0538, 2008.
- [37] F. Staub, Comput. Phys. Commun. 181 (2010) 1077–1086, arXiv:0909.2863.
- [38] F. Staub, Comput. Phys. Commun. 182 (2011) 808–833, arXiv:1002.0840.
- [39] F. Staub, Comput. Phys. Commun. 184 (2013) 1792–1809, arXiv:1207.0906.
- [40] D.M. Pierce, J.A. Bagger, K.T. Matchev, R.-j. Zhang, Nucl. Phys. B 491 (1997) 3–67, arXiv:hep-ph/9606211.
- [41] F. Staub, W. Porod, B. Herrmann, J. High Energy Phys. 1010 (2010) 040, arXiv:1007.4049.
- [42] A. Dedes, G. Degrassi, P. Slavich, Nucl. Phys. B 672 (2003) 144–162, arXiv:hep-ph/0305127.
- [43] A. Dedes, P. Slavich, Nucl. Phys. B 657 (2003) 333–354, arXiv:hep-ph/0212132.
- [44] A. Brignole, G. Degrassi, P. Slavich, F. Zwirner, Nucl. Phys. B 643 (2002) 79–92, arXiv:hep-ph/0206101.
- [45] A. Brignole, G. Degrassi, P. Slavich, F. Zwirner, Nucl. Phys. B 631 (2002) 195–218, arXiv:hep-ph/0112177.
- [46] G. Belanger, F. Boudjema, A. Pukhov, A. Semenov, Comput. Phys. Commun. 176 (2007) 367–382, arXiv:hep-ph/0607059.
- [47] G. Belanger, F. Boudjema, A. Pukhov, A. Semenov, Comput. Phys. Commun. 177 (2007) 894–895.
- [48] G. Belanger, F. Boudjema, A. Pukhov, A. Semenov, Nuovo Cimento C 033 (2) (2010) 111–116, arXiv:1005.4133.
- [49] A. Djouadi, M. Drees, P. Fileviez Perez, M. Muhlleitner, Phys. Rev. D 65 (2002) 075016, arXiv:hep-ph/0109283.
- [50] ATLAS Collaboration, arXiv:1307.7292, 2013.
- [51] H. Baer, V. Barger, A. Lessa, X. Tata, arXiv:1306.5343, 2013.
- [52] K. Schmidt-Hoberg, F. Staub, M.W. Winkler, J. High Energy Phys. 1301 (2013) 124, arXiv:1211.2835.