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Sfermion decays into singlets and singlinos in the NMSSM

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Abstract

We investigate how the addition of the singlet Higgs field in the NMSSM changes the sfermion branching ratios as compared to the MSSM. We concentrate in particular on the third generation, discussing decays of the heavier stop, sbottom or stau into the lighter mass eigenstate plus a scalar or pseudoscalar singlet Higgs. We also analyse stop, sbottom and stau decays into singlinos. It turns out that the branching ratios of these decays can be large, markedly influencing the sfermion phenomenology in the NMSSM. Moreover, we consider decays of first and second generation sfermions into singlinos.

1 Introduction

The Next-to-Minimal Supersymmetric Standard Model (NMSSM) provides an elegant solution to the μ problem of the MSSM by the addition of a gauge singlet superfield $\hat{S}$ [1]. The superpotential of the Higgs sector then has the form $\lambda\hat{S}(\hat{H}_d \cdot \hat{H}_u) + \frac{1}{3}\kappa\hat{S}^3$. When $\hat{S}$ acquires a vacuum expectation value, this creates an effective μ term, $\mu \equiv \lambda\langle S\rangle$, which is automatically of the right size, *i.e.* of the order of the electroweak scale. In this way, in the NMSSM the electroweak scale originates entirely from the SUSY breaking scale.

The addition of the singlet field leads to a larger particle spectrum than in the MSSM: in addition to the MSSM fields, the NMSSM contains two extra neutral (singlet) Higgs fields –one scalar and one pseudoscalar– as well as an extra neutralino, the singlino. Owing to these extra states, the phenomenology of the NMSSM can be significantly different from the MSSM; not at least because the usual LEP limits do not apply to the singlet and singlino states.

The NMSSM has recently become very popular. Most of the Feynman rules are given in [2]. The NMSSM Higgs phenomenology has been investigated extensively in [3–14] (a model variant without a $\hat{S}^3$ term is discussed in [15]). Detailed studies of the neutralino sector are available in [16–23]. The relic density of (singlino) dark matter has been studied in [24]. The sfermion sector, on the other hand, has so far received very little attention, although here, too, one may observe differences as compared to the MSSM. In this letter we therefore investigate the decays of squarks and sleptons in the framework of the NMSSM and contrast them to the MSSM case. We concentrate on the third generation (stops, sbottoms and staus) where we expect the largest effects, but also consider decays of selectrons, smuons and 1st/2nd generation squarks.

In the MSSM, squarks and sleptons can decay via $\tilde{f}_i \to f\tilde{\chi}^0_k$, $\tilde{f}_i \to f'\tilde{\chi}^\pm_j$ with $i,j = 1,2$ (or $i = L,R$ in case of no mixing) and $k = 1,...,4$. Squarks can also decay into gluinos, $\tilde{q}_i \to q\tilde{g}$, if the gluino is light enough. In addition, sfermions of the third generation ($\tilde{f} = \tilde{t}, \tilde{b}, \tilde{\tau}$) can have the bosonic decay modes $\tilde{f}_i \to \tilde{f}'_j + W^\pm$, $H^\pm$ and $\tilde{f}_2 \to \tilde{f}_1 + Z^0$, h^0, H^0, A^0 [25]. In the NMSSM, we have additional decay modes into singlinos and singlet Higgs states: $\tilde{f}_i \to f\tilde{\chi}^0_n$ with $n = 1, ..., 5$ and $\tilde{f}_2 \to \tilde{f}_1 + A^0_1$, A^0_2, H^0_1, H^0_2, H^0_3. Pure singlets and singlinos couple in general very weakly to the rest of the spectrum. There are hence two potentially interesting cases: a) large mixing of singlet and doublet Higgs states and/or large mixing of singlinos with gauginos-higgsinos and b) (very) light singlet/singlino states. In case b) A^0_1 and H^0_1 are almost pure singlets with masses well below the LEP bound of $m_h \geq 114$ GeV, and the singlino is the LSP. We investigate these cases in this letter and show that they can markedly influence the sfermion phenomenology.

The paper is organized as follows. In Section 2 we explain our notation, the potential and the relevant Feynman rules. In Section 3 we perform a numerical analysis, and in Section 4 we present our conclusions.

2 Notation and couplings

2.1 Potential

We follow the notation of NMHDECAY [26]. The superpotential is then given as [1,2]

$$\mathcal{W} = h_t\, \hat{Q} \cdot \hat{H}_u \hat{T}^c_R + h_b\, \hat{H}_d \cdot \hat{Q} \hat{B}^c_R + h_\tau\, \hat{H}_d \cdot \hat{L}\, \hat{\tau}^c_R - \lambda \hat{S}\, \hat{H}_d \cdot \hat{H}_u + \frac{1}{3}\kappa \hat{S}^3 \tag{1}$$

with the $SU(2)$ doublet superfields

$$\hat{Q} = \begin{pmatrix} \hat{T}_L \\ \hat{B}_L \end{pmatrix}, \quad \hat{L} = \begin{pmatrix} \hat{\nu}_L \\ \hat{\tau}_L \end{pmatrix}, \quad \hat{H}_u = \begin{pmatrix} \hat{H}^+_u \\ \hat{H}^0_u \end{pmatrix}, \quad \hat{H}_d = \begin{pmatrix} \hat{H}^0_d \\ \hat{H}^-_d \end{pmatrix} \tag{2}$$

and the product of two $SU(2)$ doublets

$$\hat{X}_1 \cdot \hat{X}_2 = \hat{X}^1_1 \hat{X}^2_2 - \hat{X}^2_1 \hat{X}^1_2 \,. \tag{3}$$

From Eq. (1) we derive the F-terms

$$F_{T_L} = h_t \widetilde{T}^c_R H^0_u - h_b \widetilde{B}^c_R H^-_d \,, \qquad F_{T^c} = h_t(\widetilde{T}_L H^0_u - \widetilde{B}_L H^+_u)\,, \tag{4}$$

$$F_{B_L} = -h_t \widetilde{T}^c_R H^+_u + h_b \widetilde{B}^c_R H^0_d \,, \qquad F_{B^c} = -h_b(\widetilde{T}_L H^-_d - \widetilde{B}_L H^0_d)\,, \tag{5}$$

$$F_{\tau_L} = h_\tau \tilde{\tau}^c_R H^0_d \,, \qquad F_{\tau^c} = -h_\tau(\tilde{\nu}_L H^-_d - \tilde{\tau}_L H^0_d)\,, \tag{6}$$

$$F_{H^0_u} = h_t \widetilde{T}_L \widetilde{T}^c_R - \lambda S H^0_d \,, \qquad F_{H^+_u} = -h_t \widetilde{B}_L \widetilde{T}^c_R + \lambda S H^-_d \,, \tag{7}$$

$$F_{H^0_d} = h_b \widetilde{B}_L \widetilde{B}^c_R + h_\tau \tilde{\tau}_L \tilde{\tau}^c_R - \lambda S H^0_u \,, \qquad F_{H^-_d} = -h_b \widetilde{T}_L \widetilde{B}^c_R - h_\tau \tilde{\nu}_L \tilde{\tau}^c_R + \lambda S H^+_u \,, \tag{8}$$

yielding the Yukawa part of the scalar potential $\mathcal{V}_F = \sum_i F_i F^*_i$. Note that the F-terms in Eqs. (7) and (8) imply direct interactions between the S field and the sfermions:

$$\begin{aligned} \mathcal{V}_F \;\supseteq\; & -h_t\lambda^* \left(\widetilde{B}_L \widetilde{T}^c_R S^* H^{-*}_d + \widetilde{T}_L \widetilde{T}^c_R S^* H^{0*}_d\right) - h_b\lambda^* \left(\widetilde{B}_L \widetilde{B}^c_R S^* H^{0*}_u + \widetilde{T}_L \widetilde{B}^c_R S^* H^{+*}_u\right) \\ & -h_\tau\lambda^* \left(\tilde{\tau}_L \tilde{\tau}^c_R S^* H^{0*}_u + \tilde{\nu}_L \tilde{\tau}^c_R S^* H^{+*}_u\right) + \text{h.c.}\;. \end{aligned} \tag{9}$$

[1]Note the different signs of the h_b and h_τ terms as compared to Eq. (A.1) of Ref. [26].

[2]The superpotential Eq. (1) posesses a discrete Z_3 symmetry which is spontaneously broken at the electroweak phase transition. This results in cosmologically dangerous domain walls [27]. We implicitly assume a solution [28] to this domain wall problem which does not impact collider phenomenology.

We also need the soft SUSY-breaking potential for the derivation of the couplings, c.f. Eq. (A.4) of [26],

$$\mathcal{V}_{soft} = h_t A_t \widetilde{Q} \cdot H_u \widetilde{T}^c_R + h_b A_b H_d \cdot \widetilde{Q} \widetilde{B}^c_R + h_\tau A_\tau H_d \cdot \widetilde{L} \widetilde{\tau}^c_R - \lambda A_\lambda S\, H_d \cdot H_u + \frac{1}{3} \kappa A_\kappa S^3 \,. \tag{10}$$

2.2 Sfermion–Higgs interaction

In the following, we denote the neutral scalar and pseudoscalar Higgs bosons by H^0_i $(i = 1, 2, 3)$ and A^0_l $(l = 1, 2)$, respectively. The interaction of H^0_i and A^0_l with a pair of sfermions $\tilde{f}_j \tilde{f}^*_k$ $(j, k = 1, 2)$ can be written as:

$$\mathcal{L}_{\tilde{f}\tilde{f}\phi} \;=\; g^S_{ijk}\, H^0_i \tilde{f}_j \tilde{f}^*_k + g^P_{ljk}\, A^0_l \tilde{f}_j \tilde{f}^*_k \,. \tag{11}$$

Apart from D-term contributions, the Higgs–sfermion couplings are proportional to the Yukawa couplings h_f. We therefore write $g^{S,P}_{ijk}$ explicitly for the third generation. For stops, we have

$$\begin{aligned}(g^S_{ijk})^{\tilde{t}} \;=\;& \left(\frac{1}{\sqrt{2}}\, h_t \,(\mu^*_{\mathit{eff}} R^{\tilde{t}\,*}_{j1} R^{\tilde{t}}_{k2} + \mu_{\mathit{eff}} R^{\tilde{t}\,*}_{j2} R^{\tilde{t}}_{k1}) - v_d D_{jk}\right) S_{i2} \\ &- \left(\frac{1}{\sqrt{2}}\, h_t \,(A_t R^{\tilde{t}\,*}_{j1} R^{\tilde{t}}_{k2} + A^*_t R^{\tilde{t}\,*}_{j2} R^{\tilde{t}}_{k1}) - v_u D_{jk} + \sqrt{2} v_u h_t^2 \delta_{jk}\right) S_{i1} \\ &+ \frac{1}{\sqrt{2}}\, v_d h_t \,(\lambda^* R^{\tilde{t}\,*}_{j1} R^{\tilde{t}}_{k2} + \lambda R^{\tilde{t}\,*}_{j2} R^{\tilde{t}}_{k1})\, S_{i3} \,,\end{aligned} \tag{12}$$

$$\begin{aligned}(g^P_{ljk})^{\tilde{t}} \;=\;& -\frac{i}{\sqrt{2}}\, h_t \,(\mu^*_{\mathit{eff}} P_{l2} + A_t P_{l1} + v_d \lambda^* P_{l3}) R^{\tilde{t}\,*}_{j1} R^{\tilde{t}}_{k2} \\ &+ \frac{i}{\sqrt{2}}\, h_t \,(\mu_{\mathit{eff}} P_{l2} + A^*_t P_{l1} + v_d \lambda P_{l3}) R^{\tilde{t}\,*}_{j2} R^{\tilde{t}}_{k1} \,,\end{aligned} \tag{13}$$

where μ_{eff} is the effective μ term:

$$\mu_{\mathit{eff}} \equiv \lambda s \tag{14}$$

with $s = \langle S \rangle$ the vev of the singlet S. (In the presence of an additional generic μ term $\mu \hat{H}_d \hat{H}_u$, $\mu_{\mathit{eff}} \to \mu_{\mathit{eff}} = \lambda s + \mu$.) For sbottoms, we get

$$\begin{aligned}(g^S_{ijk})^{\tilde{b}} \;=\;& -\left(\frac{1}{\sqrt{2}}\, h_b \,(A_b R^{\tilde{b}\,*}_{j1} R^{\tilde{b}}_{k2} + A^*_b R^{\tilde{b}\,*}_{j2} R^{\tilde{b}}_{k1}) + v_d D_{jk} + \sqrt{2} v_d h_b^2 \delta_{jk}\right) S_{i2} \\ &+ \left(\frac{1}{\sqrt{2}}\, h_b \,(\mu^*_{\mathit{eff}} R^{\tilde{b}\,*}_{j1} R^{\tilde{b}}_{k2} + \mu_{\mathit{eff}} R^{\tilde{b}\,*}_{j2} R^{\tilde{b}}_{k1}) + v_u D_{jk}\right) S_{i1} \\ &+ \frac{1}{\sqrt{2}}\, v_u h_b \,(\lambda^* {R^{\tilde{b}}_{j1}}^* R^{\tilde{b}}_{k2} + \lambda R^{\tilde{b}\,*}_{j2} R^{\tilde{b}}_{k1})\, S_{i3} \,,\end{aligned} \tag{15}$$

$$\begin{aligned}(g^P_{ljk})^{\tilde{b}} \;=\;& -\frac{i}{\sqrt{2}}\, h_b \,(A_b P_{l2} + \mu^*_{\mathit{eff}} P_{l1} + v_u \lambda^* P_{l3}) R^{\tilde{b}\,*}_{j1} R^{\tilde{b}}_{k2} \\ &+ \frac{i}{\sqrt{2}}\, h_b \,(A^*_b P_{l2} + \mu_{\mathit{eff}} P_{l1} + v_u \lambda P_{l3}) R^{\tilde{b}\,*}_{j2} R^{\tilde{b}}_{k1} \,,\end{aligned} \tag{16}$$

and analogously for staus with the obvious replacements $h_b \to h_\tau$, $A_b \to A_\tau$ and $R^{\tilde{b}} \to R^{\tilde{\tau}}$.

	T_{3L}	Q_f	Y_L	Y_R
$\tilde{t}$	$1/2$	$2/3$	$1/3$	$4/3$
$\tilde{b}$	$-1/2$	$-1/3$	$1/3$	$-2/3$
$\tilde{\tau}$	$-1/2$	-1	-1	-2

Table 1: Isospin, electric charge and hypercharges of stops, sbottoms and staus.

In Eqs. (16)–(19), S_{ij} and P_{ij} are the Higgs mixing matrices as in [26], and $R^{\tilde{f}}$ are the sfermion mixing matrices diagonalizing the sfermion mass matrices in the notation of [29, 30]:

$$\begin{pmatrix}\tilde{f}_1\\ \tilde{f}_2\end{pmatrix} = R^{\tilde{f}}\begin{pmatrix}\tilde{f}_L\\ \tilde{f}_R\end{pmatrix}, \quad \mathrm{diag}(m^2_{\tilde{f}_1}, m^2_{\tilde{f}_2}) = R^{\tilde{f}}\begin{pmatrix} m^2_{\tilde{f}_L} & a_f^* m_f \\ a_f m_f & m^2_{\tilde{f}_R}\end{pmatrix}(R^{\tilde{f}})^\dagger \tag{17}$$

where $a_t = A_t - \mu^*_{\mathit{eff}}\cot\beta$ and $a_{b,\tau} = A_{b,\tau} - \mu^*_{\mathit{eff}}\tan\beta$. Furthermore, $v_d = \langle H^0_d\rangle$, $v_u = \langle H^0_u\rangle$ and $\tan\beta = v_u/v_d$. The D-terms are given by

$$D^{\tilde{f}}_{jk} \;=\; \frac{1}{\sqrt{2}}\Big((g^2 T_{3L} - Y_L g'^2)R^{\tilde{f}\,*}_{j1}R^{\tilde{f}}_{k1} - Y_R g'^2 R^{\tilde{f}\,*}_{j2}R^{\tilde{f}}_{k2}\Big) \tag{18}$$

where g and g' are the $SU(2)$ and $U(1)$ gauge couplings: $g^2 = 4\sqrt{2}G_F M_W^2$ with G_F the Fermi constant and $g'^2 = g^2\tan^2\theta_W = 4\sqrt{2}G_F(M_Z^2 - M_W^2)$ using the on-shell relation $\sin^2\theta_W = (1 - M_W^2/M_Z^2)$ for the Weinberg angle θ_W; T_{3L} and $Y_{L(R)}$, are the 3rd component of the isospin and the hypercharge of the left (right) sfermion, respectively. We have $Y = 2(Q_f - T_3)$ where Q_f is the electric charge. For completeness, the sfermion quantum numbers are listed in Table 1.

Notice that in the CP-conserving case the pseudoscalars only couple to $\tilde{f}_1\tilde{f}_2$ combinations and hence $g^P_{l11} = g^P_{l22} = 0$ in Eqs. (13) and (16); moreover $g^S_{i12} = g^S_{i21}$ and $g^P_{l12} = -g^P_{l21}$. Notice also that in the CP-violating case, the scalar and pseudoscalar Higgs states will mix to mass eigenstates $h^0_{1\ldots5}$ similar to the MSSM case. The couplings of the charged Higgs bosons to sfermions are the same as in the MSSM.

2.3 Sfermion–neutralino interaction

The sfermion interaction with neutralinos has the same form as in the MSSM,

$$\mathcal{L}_{f\tilde{f}\tilde{\chi}^0} = g\,\bar{f}\,(a^{\tilde{f}}_{in}P_R + b^{\tilde{f}}_{in}P_L)\,\tilde{\chi}^0_n\,\tilde{f}_i + \mathrm{h.c.}\,. \tag{19}$$

The only difference is the addition of the singlino state $\tilde{S}$:

$$\tilde{\chi}^0_n = N_{n1}\tilde{B} + N_{n2}\tilde{W} + N_{n3}\tilde{H}_d + N_{n4}\tilde{H}_u + N_{n5}\tilde{S} \tag{20}$$

with $n = 1...5$, N the matrix diagonalizing the 5×5 neutralino mass matrix in the basis $(\tilde{B}, \tilde{W}, \tilde{H}_d, \tilde{H}_u, \tilde{S})$:

$$\mathcal{M}_N = \begin{pmatrix} M_1 & 0 & -m_Z s_W c_\beta & m_Z s_W s_\beta & 0 \\ 0 & M_2 & m_Z c_W c_\beta & -m_Z c_W s_\beta & 0 \\ -m_Z s_W c_\beta & m_Z c_W c_\beta & 0 & -\mu_{\mathit{eff}} & -l v s_\beta \\ m_Z s_W s_\beta & -m_Z c_W s_\beta & -\mu_{\mathit{eff}} & 0 & -l v c_\beta \\ 0 & 0 & -l v s_\beta & -l v c_\beta & 2\kappa s \end{pmatrix}, \tag{21}$$

with $s_W = \sin\theta_W$, $c_W = \cos\theta_W$, $s_\beta = \sin\beta$, $c_\beta = \cos\beta$, and

$$N^* \mathcal{M}_N N^\dagger = \mathrm{diag}(m_{\tilde\chi^0_1}, m_{\tilde\chi^0_2}, m_{\tilde\chi^0_3}, m_{\tilde\chi^0_4}, m_{\tilde\chi^0_5})\,. \tag{22}$$

We can hence use the couplings $a^{\tilde f}_{in}$ and $b^{\tilde f}_{in}$ as given in [29,30] with the neutralino index running from 1 to 5 instead from 1 to 4. It is worth noting that the couplings between sfermions and singlinos only occur via the neutralino mixing. This is in contrast to the Higgs couplings where additional terms originating from Eq. (9) are present in the sfermion–singlet interaction. Note also that [26] and [24] use the basis $(\tilde B, \tilde W, \tilde H_u, \tilde H_d, \tilde S)$ for the neutralino system. Therefore, the indices 3 and 4 of the neutralino mixing matrix N need to be interchanged when comparing [24, 26] and this paper.

3 Numerical results

We have implemented all 2-body sparticle decays of the NMSSM in the `SPheno` [31] package. The Higgs sector of the NMSSM is calculated with `NMHDECAY` [26] linked to `SPheno`. For our discussion of sfermion decays in the NMSSM, we choose the benchmark scenario 3 of [24] which is characterized by

$$\begin{aligned} \lambda = 0.4, \quad \kappa = 0.028, \quad \tan\beta = 3, \quad \mu_{eff} = \lambda s = 180 \text{ GeV}, \\ A_\lambda = 580 \text{ GeV}, \quad A_\kappa = -60 \text{ GeV}, \quad M_2 = 660 \text{ GeV}, \end{aligned} \tag{23}$$

with $M_1 = (g_1/g_2)^2 M_2 \simeq 0.5 M_2$ by GUT relations as an illustrative example. This leads to a light $\tilde\chi^0_1$ with a mass of 35 GeV which is to 87% a singlino. The $\tilde\chi^0_2$ weights 169 GeV and is dominantly a higgsino. Moreover, we have a light scalar Higgs with a mass of $m_{H_1} = 36$ GeV and a light pseudoscalar with $m_{A_1} = 56$ GeV, both being almost pure singlet states and thus evading the LEP bounds. H^0_2, H^0_3, and A^0_2 are $SU(2)$ doublet fields similar to h^0, H^0 and A^0 in the MSSM. The relic density in this scenario is $\Omega h^2 = 0.1155$ [24].

In this letter, we are interested in the decays $\tilde f_2 \to \tilde f_1 H^0_1$, $\tilde f_2 \to \tilde f_1 A^0_1$, and $\tilde f_{1,2} \to f \tilde\chi^0_1$, with $\tilde f = \tilde t, \tilde b, \tilde\tau$. In order to see the relevance of these decays, we perform a random scan over the parameters of the third generation, $M_{\tilde Q_3}$, $M_{\tilde U_3}$, $M_{\tilde D_3}$, $M_{\tilde L_3}$, $M_{\tilde E_3}$, A_t, A_b, A_τ. The sfermion mass parameters are varied between 100–800 GeV, and the trilinear couplings in their whole possible range allowed by the absence by charge or colour breaking minima. We compute the mass spectrum and the branching ratios at each scan point, accepting only points which pass the experimental bounds from LEP (the bounds from the LEP Higgs searches are fully implemented in `NMHDECAY` [26]). Owing to radiative corrections, the mass of H^0_2 varies between ~ 100 GeV and 117 GeV in the scan. The effect on the other quantities, in particular m_{H_1}, m_{A_1} and Ωh^2, is negligible.

As a sideremark we note that renormalization-group (RG) arguments can be used to set lower bounds on the $SU(2)$ doublet sfermion masses. Requiring, for example, that $m^2_{\tilde f_L}$ remain positive all the way up to the GUT scale implies $m_{\tilde f_L} \gtrsim 0.9 M_2$ for the first and second generation at the weak scale. The corresponding bounds for the third generation are much lower because the Yukawa couplings contribute to the RG running with opposite sign as the $SU(2)$ gauge coupling. We refrain, however, from imposing any such RG-inspired constraint in our analysis for two reasons: firstly because it is our aim to discuss the weak-scale phenomenology in the most general way using just one illustrative benchmark scenario, and secondly because the actual scale of SUSY breaking is unknown and may well be much lower than M_{GUT} (see e.g. the NMSSM variants of the models presented in [32]).

Let us now discuss the sfermion branching ratios. Figure 1 shows scatter plots of the branching ratios of $\tilde{b}_2$ and $\tilde{\tau}_2$ decays into A_1^0, H_1^0 and H_2^0 as function of the heavy sfermion mass, $m_{\tilde{b}_2}$ or $m_{\tilde{\tau}_2}$. As can be seen, decays into the singlet Higgs bosons H_1^0 or A_1^0 can have sizable branching ratios provided the sfermions are relatively light, $m_{\tilde{b}_2,\tilde{\tau}_2} \lesssim 400$ GeV. This feature can easily be understood from the partial width of the decay of a heavier sfermion into a lighter one plus a massless singlet. For sbottoms we have, for instance,

$$\Gamma(\tilde{b}_2 \to \tilde{b}_1 S) = \frac{c}{32\pi} h_b^2 \lambda^2 \left(\frac{v_u}{m_{\tilde{b}_2}}\right)^2 \left[1 - \left(\frac{m_{\tilde{b}_1}}{m_{\tilde{b}_2}}\right)^2\right] m_{\tilde{b}_2} \,, \tag{24}$$

which is suppressed by a factor $(v_u/m_{\tilde{b}_2})^2$ for heavy sbottoms. The factor c is $c = \cos^2 2\theta_{\tilde{b}}$ for the scalar and $c = 1$ for the pseudoscalar singlet. Analogous expressions hold for staus with $b \to \tau$ and for stops with $b \to t$ and $v_u \to v_d$. Decays into the $SU(2)$ doublet Higgs H_2^0 can also have large branching fractions, provided the splitting of the sfermion mass eigenstates is large enough. Here note that for the parameter choice Eq. (23), A_1^0 and H_1^0 decay predominantly into $b\bar{b}$ and may hence only be distinguished by the different $b\bar{b}$ invariant masses. The H_2^0 on the other hand, decays to about 60% into $H_1^0 H_1^0$ and hence into a $4b$ final state. Both signatures, the one from decay into a H_1^0 or A_1^0 leading to $b\bar{b}$ with small invariant mass as well as the $4b$'s from the decay into H_2^0, are distinct from the usual MSSM case.

We next turn to sfermion decays into singlinos. Figure 2 shows scatter plots of the branching ratios of $\tilde{t}_1$ decays into $t\tilde{\chi}_1^0$, $b\tilde{\chi}_1^\pm$, and of $\tilde{b}_1$ decays into $b\tilde{\chi}_{1,2}^0$. As expected, $\tilde{t}_1 \to b\tilde{\chi}_1^\pm$ and $\tilde{b}_1 \to b\tilde{\chi}_2^0$ ($\tilde{\chi}_1^\pm$ and $\tilde{\chi}_2^0$ being mostly doublet higgsinos) are in general the dominant modes if kinematically allowed. Nevertheless, as can be seen from the upper two plots of Fig. 2, decays into the singlino LSP can have sizable branching fractions, even if other decay modes are open. The size of the branching ratio into the singlino is mainly governed by its admixture from the doublet higgsinos, *i.e.* by the size of the λ parameter. Similar features appear also in the stau decays as illustrated in Fig. 3. The pattern of $\tilde{\tau}_1$ is quite similar to that of $\tilde{b}_1$, with $\gtrsim 10\%$ branching ratio into the singlino for $m_{\tilde{\tau}_1} \lesssim 200$–250 GeV, depending on the L/R character of the stau. For the $\tilde{\tau}_2$, the branching ratio of the decay into the singlino is even more important. In fact it can be 10–50% over a large part of the parameter space even if the decay into $\tilde{\chi}_2^0$ is open.

In this context note also the $\tilde{\chi}_1^\pm$ and $\tilde{\chi}_2^0$ will cascade further into the singlino LSP, see *e.g.* [17, 18, 22]. This is quite distinct from the MSSM, where the singlino is absent and all decay chains end in what in our case is the $\tilde{\chi}_2^0$. Obviously, this also affects cascade decays of squarks of the first and second generation (and likewise of gluinos), since in case of a singlino LSP there is one more step in the chain as compared to the MSSM. At the LHC, a singlino LSP can in fact lead to similar signatures as a gravitino or axino LSP. The presence of light scalar or pseudoscalar Higgs bosons in the decay chains may be a way to distinguish the NMSSM from other scenarios. If λ is large enough the decays $\tilde{f}_i \to f\tilde{\chi}_1^0$ with $\tilde{f}_i \neq$NLSP may also be used for discrimination, since the corresponding $\tilde{f}_i$ decays into gravitino or axino would not occur. We will discuss this in more detail in a forthcoming paper.

Last but not least we consider the decays of sfermions of the first two generations. Owing to the small Yukawa couplings, decays into Higgs bosons are negligible in this case. Decays into singlinos can, however, be important. As an example, Fig. 4 shows the branching ratios of $\tilde{e}_{L,R}$ (being the same as those of $\tilde{\mu}_{L,R}$) for the scenario of Eq. (23). For $m_{\tilde{e}} \simeq 180$–370 GeV,

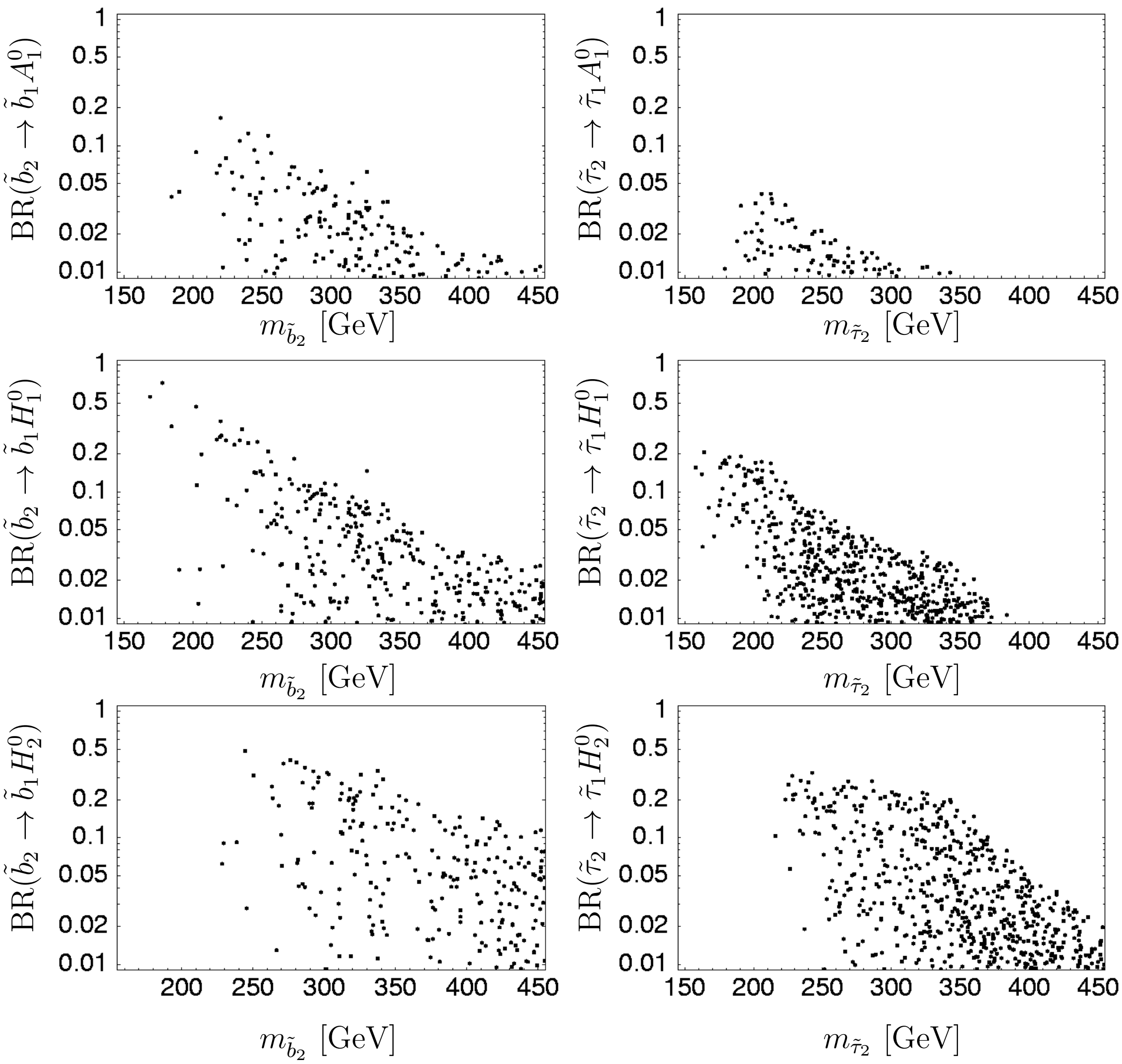

Figure 1: Branching ratios of $\tilde{b}_2$ (left) and $\tilde{\tau}_2$ (right) decays into Higgs bosons A_1^0, H_1^0, H_2^0 for scenario 3 of [24], c.f. Eq. (23), as function of the mass of the decaying particle.

the decays $\tilde{e}_L \to \nu_e \tilde{\chi}_1^-$ and $\tilde{e}_R \to e\tilde{\chi}_2^0$ clearly dominate, giving a 2-step cascade decay into the singlino LSP. Nevertheless even in this case the decays $\tilde{e}_{L,R} \to e\tilde{\chi}_1^0$ have sizable rates, being of the order of 10% for $\tilde{e}_R$. The reason is that the relative importance of the decays into charginos and neutralinos is determined by the gaugino components of the these particles, and $\tilde{\chi}_1^\pm$ and $\tilde{\chi}_{2,3}^0$ are mainly higgsino-like in our scenario. Only when the decay into $\tilde{\chi}_4^0$, which is mainly a bino, gets kinematically allowed the direct decays into the singlino become negligible. Also for the decays of up and down squarks into singlinos we find branching ratios of $\mathcal{O}(1$–$10)\%$. The implications for collider phenomenology will be discussed in detail elsewhere.

4 Conclusions

We have discussed the decays of sfermions in the NMSSM. We have shown that for stops, sbottoms and staus, in addition to the decay modes already present in the MSSM, decays into light singlet Higgs bosons, $\tilde{f}_2 \to \tilde{f}_1 + A_1^0, H_1^0$, as well as decays into a singlino LSP, $\tilde{f}_i \to f\tilde{\chi}_1^0$

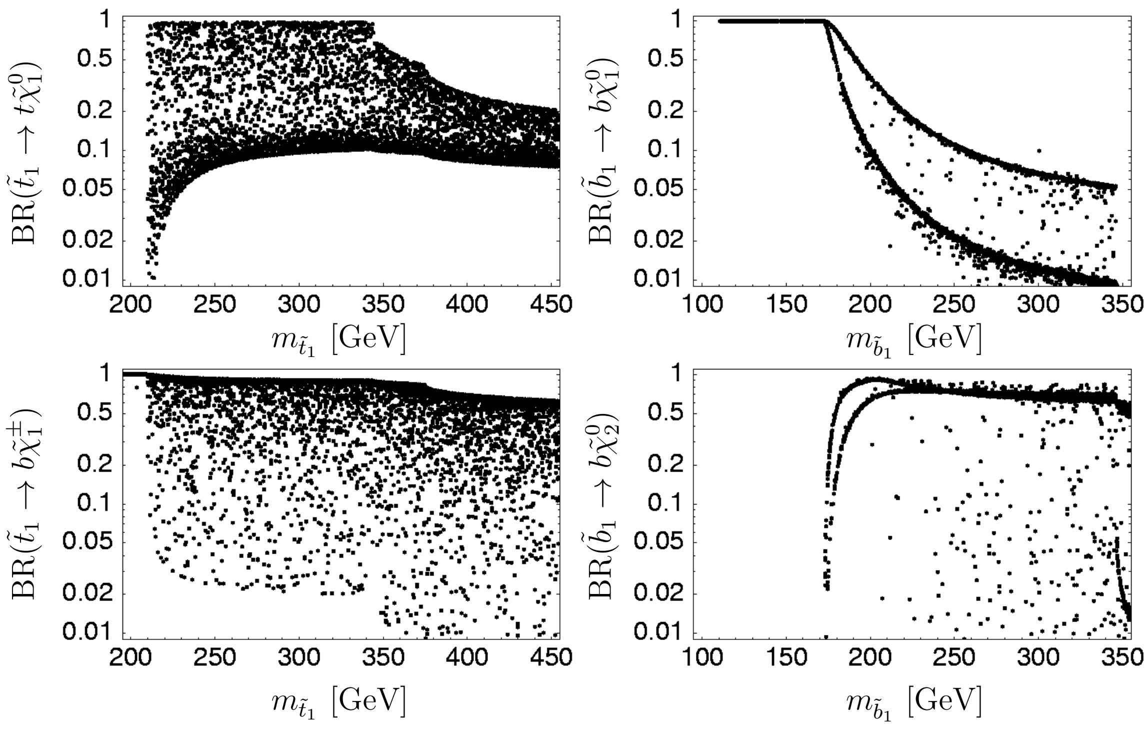

Figure 2: Branching ratios of $\tilde{t}_1$ decays into $t\tilde{\chi}_1^0$ and $b\tilde{\chi}_1^\pm$ (left) and of $\tilde{b}_1$ decays into $b\tilde{\chi}_1^0$ and $b\tilde{\chi}_q^0$ (right) for the parameters of Eq. (23).

$(i = 1, 2)$ with $\tilde{\chi}_1^0 \simeq \tilde{S}$, can be important. This is in particular the case for light sfermions. The presence of these decay modes modifies the signatures of stop, sbottom and stau events as compared to the MSSM. Also for first and second generation sfermions it turned out that the decays into a singlino LSP can be quite important. Even if other decay modes are open, $\tilde{f}_i \to f\tilde{\chi}_1^0$ with $\tilde{\chi}_1^0 \simeq \tilde{S}$ can have $\mathcal{O}(10\%)$ branching ratio. Moreover, decays (of any SUSY particle) into singlinos or singlet Higgs bosons may significantly influence cascade decays of squarks and gluinos at the LHC. In particular in case of a singlino LSP, there is one more possible step in the decay chain than in the MSSM. A singlino LSP in the NMSSM can in fact lead to similar signatures at the LHC as a gravitino or axino LSP in the MSSM. The decays discussed in this letter may help discriminating these scenarios.

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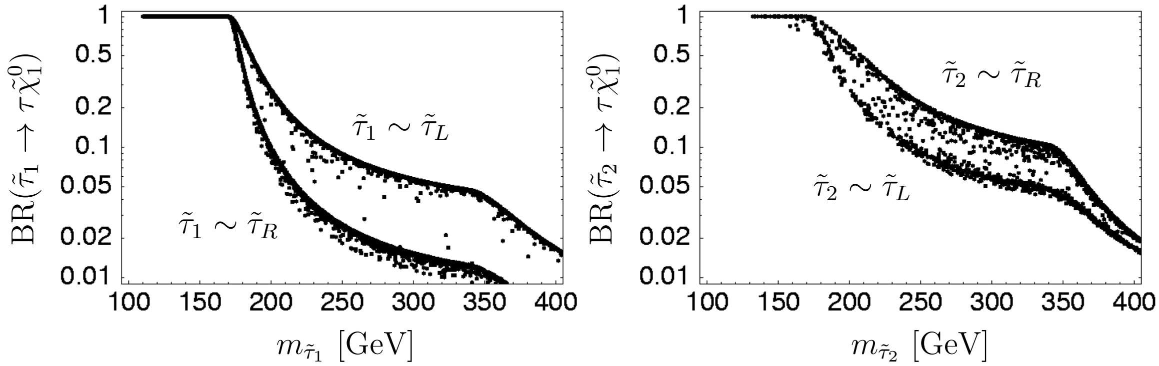

Figure 3: Branching ratios of $\tilde\tau_1 \to \tau\tilde\chi_1^0$ (left) and $\tilde\tau_2 \to \tau\tilde\chi_1^0$ (right) decays for the parameters of Eq. (23).

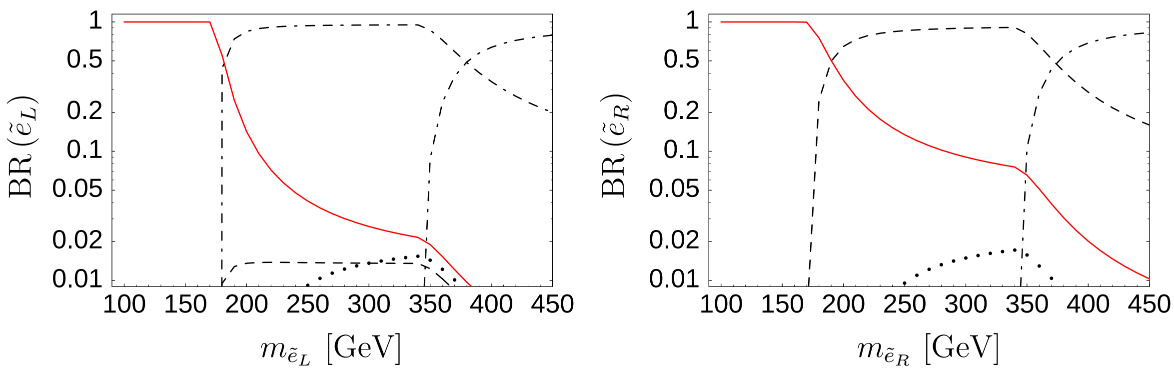

Figure 4: Branching ratios of $\tilde e_L$ (left) and $\tilde e_R$ (right) decays for the parameters of Eq. (23). The full, dashed, dotted, dash-dotted, and dash-dash-dotted lines are for the decays into $e\tilde\chi_1^0$, $e\tilde\chi_2^0$, $e\tilde\chi_3^0$, $e\tilde\chi_4^0$, and $\nu\tilde\chi_1^-$, respectively. The branching ratios for smuons $\tilde\mu_{L,R}$ are the same as for the selectrons.

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