LSP sneutrino decays into heavy standard model pairs

D. Aristizabal Sierra[a,1], D. Restrepo[b,2], S. Spinner[c,3]

[a] *IFPA, Dep. AGO, Universite de Liege, Bat B5, Sart Tilman B-4000 Liege 1, Belgium.*

[b] *Instituto de Fisica, Universidad de Antioquia, A.A. 1226, Medellin, Colombia.*

[c] *SISSA and INFN, Sezione di Trieste, Via Bonomea 265, 34136 Trieste, Italy.*

Abstract

In bilinear R-parity violation (BRpV), in which the superpotential includes a bilinear term between the lepton doublet and the up-type Higgs superfields, a sneutrino LSP can decay into pairs of heavy standard model states: W's, Z's, tops or Higgs bosons. These finals states can dominate over the traditionally considered bottom pair final state. This would lead to unique and novel supersymmetric signals with each supersymmetric event possibly producing two pairs of these heavy standard model fields. We investigate this possibility and find that the branching ratio into heavier states dominates when the bilinear term is much smaller than the sneutrino vacuum expectation value for a given sneutrino flavor. When BRpV is the only source of neutrino masses these decays can only dominate for one of the sneutrino generations. Relaxing this constraint opens these channels for all three generations.

[1] e-mail address: daristizabal@ulg.ac.be
[2] e-mail address: restrepo@udea.edu.co
[3] e-mail address: sspinner@sissa.it

1 Introduction

As the large hadron collider (LHC) continues to successfully probe the nature of electroweak symmetry breaking (a recent breakthrough being the discovery of a Higgs candidate [1, 2]), a solution to the gauge hierarchy problem, if it exists, has evaded our efforts so far. An elegant candidate for such a solution and one which addresses several other open issues as well, *e.g.* dark matter and gauge coupling unification, is supersymmetry (SUSY). Because of its theoretical appeal, it is important to understand all the guises that SUSY may adopt in order to recognize it if it is produced at the LHC.

An important open issue in SUSY, which has strong ramifications for its LHC phenomenology, as well as its cosmology, is the gauge invariance of lepton and baryon number violating interactions. Aside from introducing many new unknown parameters, these interactions also lead to rapid proton decay. The most common solution to address this issue is the imposing of a discrete symmetry, R-parity, defined as $R_p \equiv (-1)^{3(B-L)+2S}$, see [3] for a review. This forbids all the tree-level lepton and baryon number violating terms and also causes the lightest supersymmetric particle (LSP), which must be neutral, to be stable. This LSP can then play the role of dark matter and its "smoking-gun" signature at colliders is missing energy. However, proton decay requires the elimination of only the lepton number or baryon number violating terms and from a theoretical perspective, R-parity can be *ad hoc*. Furthermore, an open-mindness to possible signals at the LHC should push us to consider alternatives. Finally, when R-parity is violated the stringent constraints on superpartner masses, derived from negative collider searches for missing energy events at Tevatron and LHC [4, 5], can be relaxed[1].

A systematic study of all possible R-parity violating terms and their effects on phenomenology is an arduous task and furthermore, one would like a mechanism for understanding why proton decay is significantly suppressed. A well-motivated solution to both of these issues are models which can predict the fate of R-parity. A natural framework for this endeavor is in the context of $U(1)_{B-L}$ symmetries, see [9, 10, 11] for early examples[2]. While some $B-L$ models predict R-parity conservation [14, 15], the most minimal ones (in terms of particle content) require R-parity violation (RpV) [9, 11, 16, 17, 18]. Even some non-minimal models prefer RpV from considerations of the renormalization group evolution of the soft masses [10, 19]. These and other models of spontaneous R-parity violation, such as [20, 21], have the common feature that they can be described, in an effective field theory way, by bilinear R-parity violation (BRpV): the only R-parity violating terms are the mixings between the lepton doublets and up-type Higgs doublet in the superpotential. This makes BRpV a powerful tool for studying possible signatures of spontaneous RpV. Furthermore, proton decay is highly suppressed and it is important to mention that a gravitino LSP can be a dark matter candidate in such models [21, 22].

Once R-parity is broken the LSP is no longer stable and therefore astrophysical constraints on its nature do not apply [23]. Accordingly, from a purely phenomenological

[1]See ref. [6, 7, 8] for details.

[2]Horizontal symmetries $U(1)_X$ can be also used to construct models where the RpV couplings, arising from effective operators, are intrinsically small [12, 13].

point of view any superpartner can be the LSP, and studies of the different possibilities in BRpV models (and of sleptons and sneutrinos in general models [24, 25]) and their relation to neutrino masses and mixings have been carried out [26, 27].

In this paper we extend upon previous results by considering sneutrino LSP[3] decays in BRpV into heavy standard model (SM) final states: W^+W^-, Z^0Z^0, h^0h^0 and $t\bar{t}$. Such states, to our knowledge, had not been considered before despite the fact that they can dominate sneutrino decays and can yield unique and unanticipated SUSY signals[4]. Specifically, neutrino masses force RpV to be small so its only effect is on the decay of the LSP. Therefore, for a sneutrino LSP, every SUSY event will eventually decay into two sneutrinos which could then decay into one of these heavy states. Evidence for SUSY might then consist of events with two pairs of Ws, Zs, Higgs bosons or tops.

The main goal of this paper is to study the sneutrino decays into these heavy SM final states, which are usually due to the mixing of the sneutrino with the Higgs fields, and to show that they can dominate over the traditionally considered $b\bar{b}$ final state. We find that the latter dominate roughly when the BRpV term, ϵ_i is smaller than the vacuum expectation value (vev) of the sneutrino, v_i, for a given flavor of sneutrino, i. In the case when BRpV is the only source of neutrino masses, this possibility can only hold for one sneutrino flavor, however, when this assumption is relaxed, it can hold true for all three generations. Therefore, if two or more generations of sneutrinos decay via RpV, it might be possible to rule out BRpV as the sole generator of neutrino masses.

The rest of the paper is organized as follows. In section 2 we discuss the generalities of the BRpV model, in particular those related with the neutral scalar sector. In section 3 we derive formulas for BRpV induced mixings. In section 4 we write the relevant couplings for BRpV sneutrino decays, give analytical formulas for the different partial decay widths, analyze the constraints on parameter space enforced by neutrino data and present our results. In section 5 we summarize and present our conclusions.

2 Bilinear R-parity violation

In what follows we will briefly describe the main features of the bilinear R-parity breaking model, in particular those related with the neutral scalar sector. We shall closely follow the notation used in [29] and assume all the parameters to be real[5]. Throughout the text matrices will be denoted in bold-face.

In addition to the MSSM R-parity conserving superpotential (where we have suppressed relevant indices):

$$W_{\text{MSSM}} = \boldsymbol{h^U}\hat{Q}\hat{H}_u\hat{u}^c + \boldsymbol{h^D}\hat{Q}\hat{H}_d\hat{d}^c + \boldsymbol{h^E}\hat{L}\hat{H}_d\hat{e}^c + \mu\hat{H}_u\hat{H}_d \quad (1)$$

[3]A sneutrino NLSP with a gravitino LSP would not change our phenomenological results.

[4]Gauge boson and top quark pair production through a sneutrino resonance in general RpV was studied in [28].

[5]This simplification does not affect our main conclusions.

the bilinear R-parity breaking model also contains the following terms

$$W_{\rm BRpV} = \epsilon_{\alpha\beta}\,\epsilon_i\,\hat{L}_i^\alpha\,\hat{H}_u^\beta\,, \tag{2}$$

$\epsilon_{\alpha\beta}$ is the $SU(2)$ completely antisymmetric tensor, $i = 1, 2, 3$ runs over the SM fermion generations and ϵ_i is the R-parity and lepton number breaking bilinear parameter with units of mass. Consistency then requires a new set of soft SUSY breaking terms in the scalar potential, namely

$$V_{\rm BRpV} = B_i\,\epsilon_i\,\epsilon_{\alpha\beta}\,\tilde{L}_i^\alpha\,H_u^\beta\,. \tag{3}$$

Neglecting soft flavor mixing, the scalar potential relevant for neutral scalars is

$$\begin{aligned} V \supset\; & (m_{H_d}^2+\mu^2)H_d^\dagger H_d + (m_{H_u}^2+\mu^2)H_u^\dagger H_u + m_{L_i}^2\tilde{L}_i^\dagger\tilde{L}_i \\ & + \frac{1}{8}g_Z^2(H_u^\dagger H_u - H_d^\dagger H_d - \tilde{L}_i^\dagger\tilde{L}_i)^2 + |\boldsymbol{\epsilon}|^2 H_u^\dagger H_u + \epsilon_i\,\epsilon_j\,\tilde{L}_i^\dagger\tilde{L}_j \\ & + \Big(-\mu\,\epsilon_i\tilde{L}_i^\dagger H_d - B\mu\epsilon_{\alpha\beta}H_d^\alpha H_u^\beta + B_i\,\epsilon_i\,\epsilon_{\alpha\beta}\,\tilde{L}_i^\alpha\,H_u^\beta + \text{H.c.}\Big)\,, \end{aligned} \tag{4}$$

with $g_Z^2 = g^2 + g'^2$ and $\boldsymbol{\epsilon}^T = (\epsilon_1, \epsilon_2, \epsilon_3)$. Electroweak symmetry is broken once the Higgs and slepton acquire a vev, $\langle H_{d,u}\rangle = v_{d,u}/\sqrt{2}$ and $\langle\tilde{L}_i\rangle = v_i/\sqrt{2}$, with $v = (v_u^2 + v_d^2 + \sum_{i=1,2,3} v_i^2)^{1/2} \simeq 246$ GeV and $M_Z^2 = g_Z^2 v^2/4$. The doublets are parameterized as

$$H_d = \begin{pmatrix} H_d^0 \\ H_d^- \end{pmatrix}\,, \quad H_u = \begin{pmatrix} H_u^+ \\ H_u^0 \end{pmatrix}\,, \quad \tilde{L}_i = \begin{pmatrix} \tilde{L}_i \\ l_i^- \end{pmatrix}\,, \tag{5}$$

with the neutral components given by

$$H_d^0 = \frac{1}{\sqrt{2}}(\sigma_d^0 + i\varphi_d^0 + v_d)\,, \quad H_u^0 = \frac{1}{\sqrt{2}}(\sigma_u^0 + i\varphi_u^0 + v_u)\,, \quad \tilde{L}_i = \frac{1}{\sqrt{2}}(\tilde{\nu}_i^R + i\tilde{\nu}_i^I + v_i)\,. \tag{6}$$

Here we introduce the notation $\tilde{\nu}^{R,I}$ to differentiate the CP-even sneutrinos from the CP-odd.

In the basis $(S^0)^T = (\sigma_d^0, \sigma_u^0, \tilde{\nu}_i^R)^T$ the linear part of the neutral scalar potential can be written as

$$V_{\rm linear}(S_i^0) = \sum_{a=1,\cdots,5} t_a\,S_a^0\,, \tag{7}$$

where the t_a's are the so-called tadpoles. At tree-level these are (see eqs. (4) and (7))

$$\begin{aligned} t_d^{(0)} &= (m_{H_d}^2+\mu^2)v_d + Dv_d - \mu(Bv_u + \boldsymbol{\epsilon}\cdot\boldsymbol{v})\,, \\ t_u^{(0)} &= -B\mu v_d + (m_{H_u}^2+\mu^2)v_u - Dv_u + \sum_{i=1,2,3} B_i\epsilon_i v_i + |\boldsymbol{\epsilon}|^2 v_u\,, \\ t_i^{(0)} &= Dv_i + \epsilon_i(-\mu v_d + B_i v_u + \boldsymbol{\epsilon}\cdot\boldsymbol{v}) + m_{L_i}^2 v_i\,, \end{aligned} \tag{8}$$

where $\boldsymbol{v}^T = (v_1, v_2, v_3)$ and $D = g_Z^2(v_d^2 - v_u^2 + \sum_i v_i^2)/8$. The minimization of the potential, $V_{\rm linear} = 0$, requires the tadpoles to vanish. Thus, the vevs can be determined from the system of equations in (8) by imposing $t_a = 0$ ($a = 1, \ldots, 5$). In particular, considering only leading order BRpV terms, the sneutrino vevs can be written as

$$v_i \simeq \frac{\epsilon_i\, v}{m^2_{\tilde{\nu}_i}}(\mu\, c_\beta - B_i\, s_\beta)\,, \tag{9}$$

where $m^2_{\tilde{\nu}_i} = m^2_{L_i} + M_Z^2(c_\beta^2 - s_\beta^2)/2$ (the tree level sneutrino mass) and $t_\beta \equiv \tan\beta = v_u/v_d$.

3 R-parity violating mixings

Without conserved R-parity, there are no quantum numbers to distinguish the leptons and sleptons from the gauginos and Higgsinos and Higgs bosons respectively. Therefore, in BRpV several mixings between supersymmetric and non supersymmetric particles exist: (*i*) neutralinos mix with neutrinos, (*ii*) charginos mix with charged leptons, (*iii*) Higgs bosons mix with the sneutrinos and (*iv*) charged Higss bosons mix with charged sleptons. These mixings are important for calculating LSP decays, especially mixings of type (*i*) as they allow to fix—via experimental neutrino data—the size of the BRpV parameters. Since mixings of type (*iv*) are not of interest for sneutrino decays, in what follows we will only discuss analytical approximations for mixings of type (*i*)-(*iii*).

3.1 Neutralino-neutrino mixings

In the basis $(\psi^0)^T = (-i\lambda, -i\lambda_3, \tilde{H}_d^0, \tilde{H}_u^0, \nu_i)$ the neutral fermion mass matrix can be written as

$$\mathcal{L}_{\psi^0} = -\frac{1}{2}(\psi^0)^T \boldsymbol{M_N} \psi^0 + \text{H.c.} \tag{10}$$

with

$$\boldsymbol{M_N} = \begin{pmatrix} \boldsymbol{M_{\chi^0}}|_{4\times4} & \boldsymbol{M^T_{\chi\nu}}|_{4\times3} \\ \boldsymbol{M_{\chi\nu}}|_{4\times3} & \boldsymbol{0}|_{3\times3} \end{pmatrix}\,, \tag{11}$$

where $\boldsymbol{M_\chi}$ is the neutralino mass matrix:

$$\boldsymbol{M_{\chi^0}} = \begin{pmatrix} M_1 & 0 & -g'v_d/2 & g'v_u/2 \\ 0 & M_2 & gv_d/2 & -gv_u/2 \\ -g'v_d/2 & gv_d/2 & 0 & -\mu \\ g'v_u/2 & -gv_u/2 & -\mu & 0 \end{pmatrix}\,, \tag{12}$$

and M_1 and M_2 are the soft masses for the bino and wino respectively. $\boldsymbol{M_{\chi\nu}}$ is the 4×3 neutralino-neutrino mixing matrix given by

$$\boldsymbol{M_{\chi\nu}} = \begin{pmatrix} \frac{1}{2}g'v_i & \frac{1}{2}gv_i & 0 & \epsilon_i \end{pmatrix}\,. \tag{13}$$

In the Weyl mass eigenstate basis, defined as[6]

$$F^0 = \boldsymbol{N}\psi^0\,, \tag{14}$$

[6]The Majorana mass eigenstates are defined as $\overline{\chi^0} = (\overline{F_i^0}\; F_i^0)$.

the mass matrix becomes

$$\hat{\boldsymbol{M}}_{\boldsymbol{N}} = \boldsymbol{N}^* \, \boldsymbol{M}_{\boldsymbol{N}} \, \boldsymbol{N}^\dagger \,. \tag{15}$$

Due to the smallness of the BRpV parameters, at order ϵ_i, $\boldsymbol{M}_{\boldsymbol{N}}$ can be block diagonalized by decomposing the diagonalizing matrix $\boldsymbol{N}$ as follows [29]:

$$\boldsymbol{N} = \boldsymbol{\mathcal{N}} \, \boldsymbol{\Xi} \simeq \begin{pmatrix} \boldsymbol{N_C} & \boldsymbol{0} \\ \boldsymbol{0} & \boldsymbol{U_\ell}^\dagger \end{pmatrix} \begin{pmatrix} \mathbb{I} & \boldsymbol{\xi}^T \\ -\boldsymbol{\xi}^* & \mathbb{I} \end{pmatrix} = \begin{pmatrix} \boldsymbol{N_C} & \boldsymbol{N_C}\boldsymbol{\xi}^T \\ -\boldsymbol{U_\ell}^T \, \boldsymbol{\xi} & \boldsymbol{U_\ell}^T \end{pmatrix} . \tag{16}$$

The matrix $\boldsymbol{\Xi}$ block diagonalizes $\boldsymbol{M}_{\boldsymbol{N}}$ to the form $\text{diag}(\boldsymbol{M}_{\boldsymbol{\chi^0}}, \boldsymbol{m}_{\boldsymbol{\nu}}^{\text{eff}})$, where $\boldsymbol{m}_{\boldsymbol{\nu}}^{\text{eff}}$ is the tree-level light neutrino mass matrix. The mixing parameters ξ_{ij} can thus be determined to be

$$\begin{aligned} \xi_{i1} &= \frac{g' M_2 \mu}{2|\boldsymbol{M}_{\boldsymbol{\chi^0}}|}\Lambda_i \,, & \xi_{i2} &= -\frac{g M_1 \mu}{2|\boldsymbol{M}_{\boldsymbol{\chi^0}}|}\Lambda_i \\ \xi_{i3} &= -\frac{\epsilon_i}{\mu} + \frac{(g^2 M_1 + g'^2 M_2) v_u}{4|\boldsymbol{M}_{\boldsymbol{\chi^0}}|}\Lambda_i \,, & \xi_{i4} &= -\frac{(g^2 M_1 + g'^2 M_2) v_u}{4|\boldsymbol{M}_{\boldsymbol{\chi^0}}|}\Lambda_i \,, \end{aligned} \tag{17}$$

where $|\boldsymbol{M}_{\boldsymbol{\chi^0}}| = -(M_1 M_2 \mu^2 - 2 M_1 \mu M_W^2 \, c_\beta s_\beta - 2 M_2 \mu M_W^2 \, c_\beta s_\beta \, t_{\theta_W})$ (with $t_W = \tan\theta_W$, θ_W being the weak mixing angle) and

$$\Lambda_i = \mu v_i + v_d \epsilon_i \,. \tag{18}$$

Finally the block diagonal mixing matrices $\boldsymbol{N_C}$ and $\boldsymbol{U_\ell}$ in (16) diagonlize the neutralino and neutrino effective mass matrix, which reads

$$\left(\boldsymbol{m}_{\boldsymbol{\nu}}^{\text{eff}}\right)_{ij} = \frac{M_1 g^2 + M_2 g'^2}{4|\boldsymbol{M}_{\boldsymbol{\chi^0}}|}\Lambda_i \Lambda_j \,. \tag{19}$$

Since this matrix has two vanishing eigenvalues it can be diagonalized by only two rotation matrices, namely

$$\boldsymbol{U_\ell} = \boldsymbol{U_\ell}(\theta_{23}^{\text{BRpV}}) \, \boldsymbol{U_\ell}(\theta_{13}^{\text{BRpV}}) \,, \tag{20}$$

with

$$\tan^2\theta_{23}^{\text{BRpV}} = \frac{\Lambda_2^2}{\Lambda_3^2} \,, \qquad \tan^2\theta_{13}^{\text{BRpV}} = \frac{\Lambda_1^2}{\Lambda_2^2 + \Lambda_3^2} \,. \tag{21}$$

For sneutrino decays the relevant part of the $\chi - \nu$ mixing turns out to be the $\boldsymbol{U_\ell}^T \, \boldsymbol{\xi}$ block, that from eqs. (17) and (20) can be written as [30]

$$\boldsymbol{U_\ell}^T \, \boldsymbol{\xi} = \begin{pmatrix} 0 & 0 & -\bar{\epsilon}_1/\mu & 0 \\ 0 & 0 & -\bar{\epsilon}_2/\mu & 0 \\ a_1|\boldsymbol{\Lambda}| & a_2|\boldsymbol{\Lambda}| & -\bar{\epsilon}_3/\mu & a_4|\boldsymbol{\Lambda}| \end{pmatrix} \tag{22}$$

where $\bar{\epsilon}_{1,2} = (\boldsymbol{U_\ell}^T)_{(1,2)j}\epsilon_j$ and $\bar{\epsilon}_3 = \tilde{\epsilon}_3 - a_3|\boldsymbol{\Lambda}|\mu$ with $\tilde{\epsilon}_3 = (\boldsymbol{U_\ell}^T)_{3j}\epsilon_j$. The coefficients a_i are given by

$$a_1 = \frac{g' M_2 \mu}{2|\boldsymbol{M}_{\boldsymbol{\chi^0}}|} \,, \; a_2 = -\frac{g M_1 \mu}{2|\boldsymbol{M}_{\boldsymbol{\chi^0}}|} \,, \; a_3 = \frac{M_{\tilde{\gamma}} \; v \; s_\beta}{4|\boldsymbol{M}_{\boldsymbol{\chi^0}}|} \,, \; a_4 = \frac{M_{\tilde{\gamma}} \; v \; c_\beta}{4|\boldsymbol{M}_{\boldsymbol{\chi^0}}|} \,, \tag{23}$$

with $M_{\tilde{\gamma}} = g^2 M_1 + g'^2 M_2$. Taking into account eqs. (20) and (21), explicitly $\bar{\epsilon}_i$ $(i = 1, 2)$ and $\tilde{\epsilon}_3$ are given by

$$\begin{aligned}
\bar{\epsilon}_1 &= \frac{\epsilon_1(\Lambda_2^2 + \Lambda_3^2) - \Lambda_1(\Lambda_2\epsilon_2 + \Lambda_3\epsilon_3)}{|\boldsymbol{\Lambda}|\sqrt{\Lambda_2^2 + \Lambda_3^2}}\,,\\
\bar{\epsilon}_2 &= \frac{\Lambda_3\epsilon_2 - \Lambda_2\epsilon_3}{\sqrt{\Lambda_2^2 + \Lambda_3^2}}\,,\\
\tilde{\epsilon}_3 &= \frac{\boldsymbol{\Lambda}\cdot\boldsymbol{\epsilon}}{|\boldsymbol{\Lambda}|}\,.
\end{aligned} \tag{24}$$

3.2 Chargino-charged lepton mass matrices and mixings

In the bases $(\psi^\pm)^T = (-i\lambda^\pm, H^\pm_{u,d}, e^\pm_{R,L})$ the chargino charged lepton mass matrix is determined by the following Lagrangian

$$-\mathcal{L}_{\psi^\pm} = \frac{1}{2}\Psi^T \begin{pmatrix} \mathbf{0} & \boldsymbol{M_C^T} \\ \boldsymbol{M_C} & \mathbf{0} \end{pmatrix} \Psi + \text{H.c.}\,, \tag{25}$$

where $\Psi^T = (\psi^+, \psi^-)^T$ and in the basis in which the charged lepton mass matrix is diagonal $\boldsymbol{M_C}$ can be written as

$$\boldsymbol{M_C} = \begin{pmatrix} \boldsymbol{M_\chi}|_{2\times2} & \boldsymbol{M_{R\chi}}|_{2\times3} \\ \boldsymbol{M_{L\chi}}|_{3\times2} & \boldsymbol{\hat{M}_\ell}|_{3\times3} \end{pmatrix}\,. \tag{26}$$

The block diagonal matrices correspond to the MSSM chargino and charged lepton mass matrices whereas the off-diagonal mass matrices read

$$\boldsymbol{M_{R\chi}} = \begin{pmatrix} 0 \\ -\frac{1}{\sqrt{2}}h_i^E v_i \end{pmatrix}\,, \qquad \boldsymbol{M_{L\chi}} = \begin{pmatrix} \frac{1}{\sqrt{2}}g v_i & -\epsilon_i \end{pmatrix}\,. \tag{27}$$

Defining the Weyl mass eigenstates as[7]

$$F^- = \boldsymbol{U}\psi^- \quad \text{and} \quad F^+ = \boldsymbol{V}\psi^+\,, \tag{28}$$

the diagonal mass matrix $\boldsymbol{\hat{M}_C}$ is obtained through the biunitary transformation

$$\boldsymbol{U}\;\boldsymbol{M_C}\;\boldsymbol{V}^T\,. \tag{29}$$

Approximate analytical expressions for the mixing matrices $\boldsymbol{U}, \boldsymbol{V}$ have been discussed in [30, 31, 32]. We here—for completeness—describe the method. The off-diagonal block matrix $\boldsymbol{M_{R\chi}}$, being proportional to the charged lepton Yukawa couplings, can be neglected, and due to the smallness of the BRpV parameters the mixing matrices $\boldsymbol{U}, \boldsymbol{V}$ can

[7]The corresponding Dirac eigenstates are defined as $\overline{\chi_i^-} = (\overline{F_i^-}\; F_i^+)$.

be written according to

$$
\begin{aligned}
\boldsymbol{U} &= \boldsymbol{\mathcal{U}}\,\boldsymbol{\Xi_L} \simeq \begin{pmatrix} \boldsymbol{U_L} & \boldsymbol{0} \\ \boldsymbol{0} & \mathbb{I} \end{pmatrix} \begin{pmatrix} \mathbb{I} & \boldsymbol{\xi}_{\boldsymbol{L}}^{\dagger} \\ -\boldsymbol{\xi_L} & \mathbb{I} \end{pmatrix}, \\
\boldsymbol{V} &= \boldsymbol{\mathcal{V}}\,\boldsymbol{\Xi_R} \simeq \begin{pmatrix} \boldsymbol{V_R} & \boldsymbol{0} \\ \boldsymbol{0} & \mathbb{I} \end{pmatrix} \begin{pmatrix} \mathbb{I} & \boldsymbol{\xi}_{\boldsymbol{R}}^{T} \\ -\boldsymbol{\xi}_{\boldsymbol{R}}^{*} & \mathbb{I} \end{pmatrix},
\end{aligned} \tag{30}
$$

where, in first approximation in the BRpV parameters, the matrices $\boldsymbol{\Xi_{L,R}}$ block-diagonalize the mass matrix, namely

$$
\begin{pmatrix} \mathbb{I} & \boldsymbol{\xi}_{\boldsymbol{L}}^{\dagger} \\ -\boldsymbol{\xi_L} & \mathbb{I} \end{pmatrix} \begin{pmatrix} \boldsymbol{M_\chi} & \boldsymbol{0} \\ \boldsymbol{M_{L\chi}} & \hat{\boldsymbol{M}}_{\boldsymbol{\ell}} \end{pmatrix} \begin{pmatrix} \mathbb{I} & \boldsymbol{\xi}_{\boldsymbol{R}}^{T} \\ -\boldsymbol{\xi}_{\boldsymbol{R}}^{*} & \mathbb{I} \end{pmatrix} \simeq \begin{pmatrix} \boldsymbol{M_\chi} & \boldsymbol{0} \\ \boldsymbol{0} & \hat{\boldsymbol{M}}_{\boldsymbol{\ell}} \end{pmatrix}, \tag{31}
$$

and the matrices $\boldsymbol{U_L}, \boldsymbol{V_R}$ diagonalize in turn the chargino mass matrix, with the rotation angles given by

$$
\tan 2\theta_L = -\frac{2\sqrt{2}M_W(Mc_\beta + \mu s_\beta)}{M^2 - \mu^2 - 2M_W^2 c_{2\beta}}, \quad \tan 2\theta_R = -\frac{2\sqrt{2}M_W(Ms_\beta + \mu c_\beta)}{M^2 - \mu^2 - 2M_W^2 c_{2\beta}}. \tag{32}
$$

From equation (31) the matrices $\boldsymbol{\xi_{L,R}}$ are found to be

$$
\boldsymbol{\xi}_{\boldsymbol{R}}^{T} = \boldsymbol{M}_{\boldsymbol{\chi}}^{-1}\,\boldsymbol{\xi}_{\boldsymbol{L}}^{T}\,\hat{\boldsymbol{M}}_{\boldsymbol{\ell}} \quad \text{and} \quad \boldsymbol{\xi}_{\boldsymbol{L}}^{*} = \boldsymbol{M_{L\chi}}\,\boldsymbol{M}_{\boldsymbol{\chi}}^{-1}, \tag{33}
$$

which implies $\boldsymbol{\xi_R}$ is suppressed with respect $\boldsymbol{\xi_L}$ by a factor m_ℓ/m_{susy} and thus can be neglected ($\boldsymbol{\Xi_R} = \mathbb{I}_{5\times5}$). Explicitly $\boldsymbol{\xi_L}$ can be written in terms of the BRpV parameters and the coefficients entering in the chargino mass matrix:

$$
\begin{aligned}
\xi_{L_{i1}} &= \frac{g}{\sqrt{2}}\frac{\Lambda_i}{|\boldsymbol{M_\chi}|} \quad \text{with} \quad |\boldsymbol{M_\chi}| = M\mu - M_W^2 s_{2\beta}, \\
\xi_{L_{i2}} &= -\frac{2M_W^2 s_\beta}{v|\boldsymbol{M_\chi}|\mu}\Lambda_i - \frac{\epsilon_i}{\mu}.
\end{aligned} \tag{34}
$$

3.3 CP-even neutral scalars mass matrices and mixings

In the basis $(S^0)^T = (\sigma_d^0, \sigma_u^0, \tilde{\nu}_i^R)^T$ the mass matrix of the CP-even neutral scalars S^0 is determined by the following quadratic terms

$$
\mathcal{L}_S^0 = \frac{1}{2}(S^0)^T \boldsymbol{M}_{\boldsymbol{S^0}}^2 S^0 = \frac{1}{2}(S^0)^T \begin{pmatrix} \boldsymbol{M}_{\boldsymbol{HH}}^2|_{2\times2} & \boldsymbol{M}_{\boldsymbol{H\tilde{\nu}}}^2|_{2\times3} \\ \boldsymbol{M}_{\boldsymbol{H\tilde{\nu}}}^2|_{3\times2} & \boldsymbol{M}_{\boldsymbol{\tilde{\nu}\tilde{\nu}}}^2|_{3\times3}\,. \end{pmatrix} S^0\,. \tag{35}
$$

In what follows we will discuss approximate analytical formulas for the $\sigma_{u,d}^0 - \tilde{\nu}_i^R$ mixing. The entries of the mass matrix in (35) involve the parameters μ, $\tan\beta$, the soft SUSY breaking coefficients m_{L_i} and B, and the R-parity breaking parameters ϵ_i, v_i and B_i. We use the minimization conditions of the scalar potential ($t_{u,d,i} = 0$) to remove the parameters B_i. In doing so the matrix $\boldsymbol{M}_{\boldsymbol{HH}}^2$ can be written in terms of the CP-odd neutral scalar mass m_{A^0}, the sneutrino masses $m_{\tilde{\nu}_i}$, M_Z and the BRpV parameters ϵ_i and v_i, namely

$$
\begin{aligned}
(\boldsymbol{M^2_{HH}})_{11} &= m^2_{A^0} s^2_\beta + M^2_Z c^2_\beta + \frac{\mu}{v_d} \boldsymbol{\epsilon} \cdot \boldsymbol{v} \\
(\boldsymbol{M^2_{HH}})_{12} &= (\boldsymbol{M^2_{HH}})_{21} = -(m^2_{A^0} + M^2_Z) c_\beta s_\beta \\
(\boldsymbol{M^2_{HH}})_{22} &= m^2_{A^0} c^2_\beta + M^2_Z s^2_\beta + \sum_i \bar{m}^2_{\tilde{\nu}_i} \frac{v_i^2}{v_u^2} - \frac{\mu c^2_\beta}{v_d s^2_\beta} \sum_{i=1,2,3} \epsilon_i^2 v_i^2 + 2 \frac{|\boldsymbol{\epsilon} \cdot \boldsymbol{v}|^2}{v_u^2} + \frac{1}{4} \frac{g_Z^2}{v_u^2} \sum_{\substack{i<j \\ j=1,2,3}} v_i^2 v_j^2 \,.
\end{aligned}
\tag{36}
$$

Where the following relations have been used

$$
m^2_{A^0} = \frac{2B\mu}{s_{2\beta}} \,, \qquad \bar{m}^2_{\tilde{\nu}_i} = m^2_{\tilde{\nu}_i} + \epsilon_i^2 + \frac{1}{8} g_Z^2 v_i^2 \,. \tag{37}
$$

Note that phenomenologically consistency requires the inclusion of the one-loop correction in the $(\boldsymbol{M^2_{HH}})_{22}$ entry, which in the limit of no stop mixing is

$$
(\boldsymbol{M^2_{HH}})_{22}^{\text{1-loop}} = \frac{3 m_t^4}{4\pi^2 v_u^2} \log \left(\frac{m^2_{\tilde{t}_1} \, m^2_{\tilde{t}_2}}{m_t^4} \right) . \tag{38}
$$

See [33] for a more complete expression.

In the absence of BRpV, $\boldsymbol{M^2_{HH}}$ corresponds to the neutral CP-even Higgs mass matrix of the R-parity conserving MSSM, which can be diagonalized via the rotation matrix

$$
\boldsymbol{R_\alpha} = \begin{pmatrix} c_\alpha & s_\alpha \\ -s_\alpha & c_\alpha \end{pmatrix} \quad \text{with} \quad t_{2\alpha} = \frac{m_A^{2(0)} + M_Z^2}{m_A^{2(0)} - M_Z^2} \, t_{2\beta} \,. \tag{39}
$$

The elements of the 3×3 right-lower block sneutrino mass matrix in (35) are given by

$$
\begin{aligned}
(\boldsymbol{M^2_{\tilde{\nu}\tilde{\nu}}})_{(i+2)(i+2)} &= \bar{m}^2_{\tilde{\nu}_i} + \frac{1}{4} v_i^2 \,, \\
(\boldsymbol{M^2_{\tilde{\nu}\tilde{\nu}}})_{(i+2)(j+2)} &= \frac{1}{4} g_Z^2 v_i v_j + \epsilon_i \epsilon_j \qquad (\text{with } i < j \text{ and } i = 1,2).
\end{aligned}
\tag{40}
$$

Finally, for the $\sigma^0_{u,d} - \tilde{\nu}^R_i$ mixing 2×3 block we have

$$
\begin{aligned}
(\boldsymbol{M^2_{H\tilde{\nu}}})_{1(i+2)} &= -\mu \epsilon_i + M_Z^2 \, c_\beta \frac{v_i}{v} \,, \\
(\boldsymbol{M^2_{H\tilde{\nu}}})_{2(i+2)} &= \frac{\mu}{t_\beta} \epsilon_i - (\bar{m}^2_{\tilde{\nu}_i} + M_Z^2 s_\beta^2) \frac{v_i}{v s_\beta} - \frac{1}{2} M_Z^2 \frac{v_i}{v} \sum_{j \neq i} \frac{v_j^2}{v^2} - \frac{\epsilon_i}{v s_\beta} \sum_{i \neq j} v_j \epsilon_j \,.
\end{aligned}
\tag{41}
$$

In the mass eigenstate basis defined as

$$
S'^0 = \boldsymbol{R^{S^0}} S^0 \,, \tag{42}
$$

the Lagrangian in (35) becomes

$$\mathcal{L}_{S^0} = \frac{1}{2}(S'^0)^T \, \hat{\boldsymbol{M}}^{\boldsymbol{2}}_{\boldsymbol{S^0}} \, S'^0 \quad \text{with} \quad \hat{\boldsymbol{M}}^2_{\boldsymbol{S^0}} = \boldsymbol{R}^{\boldsymbol{S^0}} \boldsymbol{M}_{\boldsymbol{S^0}}{}^2 \boldsymbol{R}^{\boldsymbol{S^0}T} , \tag{43}$$

where $\hat{\boldsymbol{M}}^2_{\boldsymbol{S^0}}$ is diagonal. Assuming real parameters $\boldsymbol{R}^{\boldsymbol{S^0}}$ can be parameterized as

$$(\boldsymbol{R}^{\boldsymbol{S^0}})^T = \prod_{\substack{i<j \\ j=1,\cdots,5}} (\boldsymbol{R}_{ij})^T , \tag{44}$$

where the $\boldsymbol{R}_{ij} \equiv \boldsymbol{R}(\theta_{ij})$ are 5×5 rotation matrices.

If the $\sigma^0_{d,u} - \tilde{\nu}^R_i$ mixing is small—as expected due to the smallness of the BRpV parameters required by neutrino data—a perturbative diagonalization of the mass matrix in (35) can be done. By neglecting the BRpV parameters in $\boldsymbol{M}^2_{\boldsymbol{HH}}$ and $\boldsymbol{M}^2_{\tilde{\boldsymbol{\nu}}\tilde{\boldsymbol{\nu}}}$ the rotation matrix reduces to

$$\boldsymbol{R}^{\boldsymbol{S^0}} = \boldsymbol{R_{25}} \, \boldsymbol{R_{24}} \, \boldsymbol{R_{23}} \, \boldsymbol{R_{15}} \, \boldsymbol{R_{14}} \, \boldsymbol{R_{13}} \, \boldsymbol{R_{12}} \, . \tag{45}$$

When acting on $\boldsymbol{M}^2_{\boldsymbol{S^0}}$ the matrix $\boldsymbol{R_{12}}$ diagonalizes the 2×2 block $\boldsymbol{M}^2_{\boldsymbol{HH}}$ according to $\hat{\boldsymbol{M}}^2_{\boldsymbol{HH}} = \text{diag}(m^2_{H^0}, m^2_{h^0})$ (where h^0 and H^0 are the light and heavy CP-even Higgs bosons) and modifies the BRpV mixing matrix $\boldsymbol{M}^2_{\boldsymbol{H}\tilde{\boldsymbol{\nu}}}$:

$$\boldsymbol{M}^2_{\boldsymbol{H}\tilde{\boldsymbol{\nu}}} \to \begin{pmatrix} s_\alpha \, (\boldsymbol{M}^2_{\boldsymbol{H}\tilde{\boldsymbol{\nu}}})_{2(i+2)} + c_\alpha (\boldsymbol{M}^2_{\boldsymbol{H}\tilde{\boldsymbol{\nu}}})_{1(i+2)} \\ c_\alpha \, (\boldsymbol{M}^2_{\boldsymbol{H}\tilde{\boldsymbol{\nu}}})_{2(i+2)} - s_\alpha (\boldsymbol{M}^2_{\boldsymbol{H}\tilde{\boldsymbol{\nu}}})_{1(i+2)} \end{pmatrix} . \tag{46}$$

For the mixing angle we have $\theta_{12} = \alpha$. The matrices $\boldsymbol{R_{1(i+2)}}$ eliminate the first row entries in (46) and up to $\mathcal{O}(\epsilon_i^2)$ leave the block-diagonal matrices $\hat{\boldsymbol{M_{HH}}}^2$ and $\hat{\boldsymbol{M}_{\tilde{\boldsymbol{\nu}}\tilde{\boldsymbol{\nu}}}}^2$ diagonal. The rotation angles are given by

$$\theta_{1(i+2)} \simeq \frac{s_\alpha \, (\boldsymbol{M}^2_{\boldsymbol{H}\tilde{\boldsymbol{\nu}}})_{2(i+2)} + c_\alpha (\boldsymbol{M}^2_{\boldsymbol{H}\tilde{\boldsymbol{\nu}}})_{1(i+2)}}{m^2_{H^0} - m^2_{\tilde{\nu}_i}} . \tag{47}$$

The matrices $\boldsymbol{R_{2(i+2)}}$ instead eliminate the second row elements in (46) leaving again, up to order ϵ_i^2, the block-diagonal matrices diagonal. The rotation angles in this case read

$$\theta_{2(i+2)} \simeq \frac{c_\alpha \, (\boldsymbol{M_{H\tilde{\nu}}}{}^2)_{2(i+2)} - s_\alpha (\boldsymbol{M_{H\tilde{\nu}}}{}^2)_{1(i+2)}}{m^2_{h^0} - m^2_{\tilde{\nu}_i}} . \tag{48}$$

With these results at hand and neglecting terms of $\mathcal{O}(\theta^2)$ the full rotation matrix in (45) can be written as

$$\boldsymbol{R}^{\boldsymbol{S^0}} \simeq \begin{pmatrix} \boldsymbol{R_\alpha} & \boldsymbol{R_{\tilde{\nu}}} \\ \boldsymbol{R_\sigma} & \mathbb{I} \end{pmatrix} , \tag{49}$$

where $\boldsymbol{R_{\tilde{\nu}}}$ and $\boldsymbol{R_\sigma}$ account for the sneutrino and CP-even neutral Higgs components of the mass eigenstates, and read

$$\boldsymbol{R_{\tilde{\nu}}} = \begin{pmatrix} \theta_{1(i+2)} \\ \theta_{2(i+2)} \end{pmatrix} , \qquad \boldsymbol{R_\sigma} = \begin{pmatrix} -c_\alpha \theta_{1(i+2)} + s_\alpha \theta_{2(i+2)} & -s_\alpha \theta_{1(i+2)} - c_\alpha \theta_{2(i+2)} \end{pmatrix} . \tag{50}$$

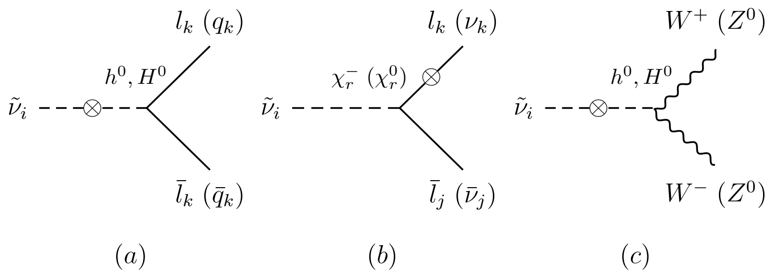


Figure 1: *BRpV induced sneutrino decay modes. In* (*a*) *and* (*c*) *sneutrino decays are induced by sneutrino-Higgs mixing while in figure* (*b*) *by chargino-charged lepton or neutralino-neutrino mixing, depending on whether the final states involves charged lepton or neutrinos. The open circles with a cross inside indicate a BRpV mixing insertion.*

4 Sneutrino decays

With the results of section 3 we are now in a position to discuss approximate formulas for sneutrino decays. From now on we will consider only CP-even sneutrino decays, $\tilde{\nu}^R$, and so will drop the superscript R. Possible tree-level two-body sneutrino final states include fermionic modes $\tilde{\nu}_i \to l_j^+ l_k^-$, $\nu_j\nu_k$ and $q_k\bar{q}_k$; electroweak gauge bosons modes $\tilde{\nu}_i \to W^+W^-$ and Z^0Z^0 and Higgs bosons modes $\tilde{\nu}_i \to h^0h^0$, H^0H^0, A^0A^0 and H^+H^-.

The goal of this paper and section is not a full study of the parameter space and all decays but rather to show that these heavy states can dominate the traditionally considered decay to $b\bar{b}$ and to identify the relevant parameter space for this dominance.

4.1 Relevant Lagrangians

Taking into account our approximate results for the chargino-right-handed lepton mixing, $\xi_R \sim (m_\ell/m_{\rm susy})$, mixing with right-handed leptons is zero ($\xi_R \sim 0$, see eq. (33)), the full Lagrangian for sneutrino decays into charge leptons is reduced to [29]

$$-\mathcal{L}^{(\pm)} = \frac{h_k^E}{\sqrt{2}} (\boldsymbol{R}^{\boldsymbol{S^0}})_{(i+2)1}\, \overline{l_k^-}\ P_L\ l_k^- \tilde{\nu}_i - \frac{h_i^E}{\sqrt{2}} \boldsymbol{U}_{(j+2)2}\, \overline{l_i^-}\ P_L\ l_j^- \tilde{\nu}_i + \text{H.c.}\,, \tag{51}$$

where the mixing matrices are given by eqs. (34) and (49). Due to the smallness of the BRpV induced mixing we take $S'^0_{(i+2)} \to \tilde{\nu}_i$. Note that while the first term in (51) necessarily leads to two same-flavor opposite-sign charged lepton final states the second term can lead to different-flavor signatures. The corresponding Feynman diagrams for these processes are depicted in figure 1((*a*),(*b*)). Being proportional to SM charged lepton Yukawa couplings these decays are dominated by final states involving τ's.

For sneutrino decays into neutrinos (invisible decays $\tilde{\nu}_i \to \sum_{k,j} \nu_k\nu_j$), the relevant interactions are given by

$$-\,\mathcal{L}^{(0)} = \frac{1}{2}\overline{\nu_k}\,(C_{kji} + C_{jki})\,\nu_j\ \tilde{\nu}_i + \text{H.c.}\,, \tag{52}$$

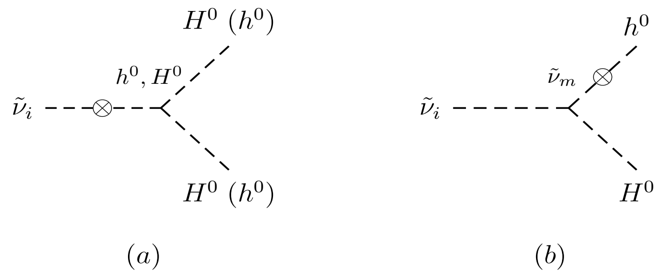


Figure 2: *BRpV induced sneutrino decays involving Higgs final states. The open circles with a cross inside indicate a BRpV mixing insertion.*

where the coupling reads

$$C_{kji} = \bar{N}_{k+4} \sum_{n=1}^{5} s_n \, \boldsymbol{N}_{(j+4)(n+2)} \, (\boldsymbol{R^{S^0}})_{(i+2)n} \, , \tag{53}$$

where $s_n = (1, -1, 1, 1, 1)$ and

$$\bar{N}_{k+4} = g' \, \boldsymbol{N}_{(k+4)1} - g \, \boldsymbol{N}_{(k+4)2} \, . \tag{54}$$

The elements of the neutral fermion mixing matrix correspond to the entries of the matrix given in eq. (16). These interactions induce invisible sneutrino decays as the ones shown in figure 1(b).

In the up and down quark mass eigenstate basis the sneutrino-quark-quark interactions are dictated by

$$\mathcal{L}^{(q)} = -\frac{1}{\sqrt{2}} h_k^U (\boldsymbol{R^{S^0}})_{(i+2)2} \, \bar{u}_k \, u_k \, \tilde{\nu}_i - \frac{1}{\sqrt{2}} h_k^D (\boldsymbol{R^{S^0}})_{(i+2)1} \bar{d}_k \, d_k \, \tilde{\nu}_i \, . \tag{55}$$

As for the charged lepton final states, these decays are controlled by SM quark Yukawa couplings and thus are dominated by $b\bar{b}$ and $t\bar{t}$, the last one if kinematically allowed.

The Lagrangian for gauge boson final states is given by

$$\mathcal{L}^{(V)} = g_V \left(c_\beta (\boldsymbol{R^{S^0}})_{(i+2)1} + s_\beta (\boldsymbol{R^{S^0}})_{(i+2)2} + \frac{v_i}{v} \right) \, V^\mu \, V_\mu \, \tilde{\nu}_i \, , \tag{56}$$

with $V_\mu = W_\mu, Z_\mu$ and $g_{W,Z} = g \, M_W, g \, M_Z / c_W$. For Higgs final states we write the Lagrangian involving h^0 and H^0[8]:

$$\mathcal{L}^{S^0} = -g_{ijk} \, S_j^0 \, S_k^0 \, \tilde{\nu}_i \, , \tag{57}$$

where the fully symmetric coupling $g_{(i+2)jk}$ is given by

$$g_{ijk} = -\frac{1}{4} \, g_Z^2 \sum_{n=1}^{5} u_n \left[(\boldsymbol{R^{S^0}})_{(i+2)n} \, \overline{R}_{jk} + (\boldsymbol{R^{S^0}})_{jn} \, \overline{R}_{(i+2)k} + (\boldsymbol{R^{S^0}})_{kn} \, \overline{R}_{(i+2)j} \right] \, , \tag{58}$$

[8] We do not study $\tilde{\nu}_i \to A^0 \, A^0$ and $\tilde{\nu}_i \to H^+ \, H^-$ decays and so do not present the Lagrangians that govern these interactions.

with $u_n = (u_1, u_2, u_3 \ldots) = (v_d, -v_u, v_1 \ldots)$ and

$$\overline{R}_{jk} = (\boldsymbol{R}^{\boldsymbol{S^0}})_{j1}(\boldsymbol{R}^{\boldsymbol{S^0}})_{k1} - (\boldsymbol{R}^{\boldsymbol{S^0}})_{j2}(\boldsymbol{R}^{\boldsymbol{S^0}})_{k2} + \sum_{n=3}^{5}(\boldsymbol{R}^{\boldsymbol{S^0}})_{jn}(\boldsymbol{R}^{\boldsymbol{S^0}})_{kn} \,. \tag{59}$$

The interactions in (57) for $j = k = 1, 2$ lead to decays of type (a) in figure 2 while for $j = 1, k = 2$ to those shown in figure 2(b).

4.2 Partial decay widths

Fermionic final states are dominated by third generation quark and charged leptons. Due to the structure of the Higgs-sneutrino mixing, $\tau\bar{\tau}$, $b\bar{b}$ and $t\bar{t}$ final states are possible independently of the sneutrino flavor, whereas $\bar{\tau}(e,\mu)$ final states are only sizable for tau sneutrinos. Neglecting the final state masses, the partial decay widths for $\bar{\tau}(e,\mu)$, $\tau\bar{\tau}$ and $b\bar{b}$ decays can be written as

$$\Gamma(\tilde{\nu}_\tau \to \bar{\tau}(e,\mu)) = \frac{m_\tau^2\, G_F}{4\sqrt{2}\pi\, c_\beta^2} m_{\tilde{\nu}_\tau} \boldsymbol{U}^2_{(3,4)2} \,, \tag{60}$$

$$\Gamma(\tilde{\nu}_i \to \tau\bar{\tau}) = \frac{m_\tau^2\, G_F}{4\sqrt{2}\pi\, c_\beta^2} m_{\tilde{\nu}_i} \left[(\boldsymbol{R}^{\boldsymbol{S^0}})_{(i+2)1} - \boldsymbol{U}_{52}\delta_{i\tau}\right]^2 \,, \tag{61}$$

$$\Gamma(\tilde{\nu}_i \to b\bar{b}) = \frac{3m_b^2\, G_F}{4\sqrt{2}\pi\, c_\beta^2} m_{\tilde{\nu}_i} (\boldsymbol{R}^{\boldsymbol{S^0}})^2_{(i+2)1} \,, \tag{62}$$

where G_F is the Fermi constant. For invisible decay modes the partial decay width, summing over lepton flavors, can be written according to

$$\Gamma(\tilde{\nu}_i \to \sum_{k,j} \nu_k \nu_j) = \frac{m_{\tilde{\nu}_i}}{16\pi} \sum_{k,j} (C_{kji} + C_{jki})^2 \,, \tag{63}$$

with C_{kji} given by (53). For $t\bar{t}$ final states the phase space factors are relevant, accordingly the corresponding decay width reads

$$\Gamma(\tilde{\nu}_i \to t\bar{t}) = \frac{3m_t^2\, G_F}{4\sqrt{2}\pi\, s_\beta^2} m_{\tilde{\nu}_i} (\boldsymbol{R}^{\boldsymbol{S^0}})^2_{(i+2)2} \left(1 - 4\frac{m_t^2}{m_{\tilde{\nu}_i}^2}\right)^{3/2} \tag{64}$$

For Gauge boson final states the partial decay width is given by

$$\Gamma(\tilde{\nu}_i \to VV) = \frac{G_F m_{\tilde{\nu}_i}^3}{16\sqrt{2}\,\pi} \delta_V \sqrt{1 - 4\frac{M_V^2}{m_{\tilde{\nu}_i}^2}} \left(1 - 4\frac{M_V^2}{m_{\tilde{\nu}_i}^2} + 12\frac{M_V^4}{m_{\tilde{\nu}_i}^4}\right) \left|\mathcal{A}_i^V\right|^2 \,, \tag{65}$$

where $V = Z, W$, $\delta_{Z,W} = 1, 2$ and the amplitude A_i^V reads

$$\mathcal{A}_i^V = c_\beta (\boldsymbol{R}^{\boldsymbol{S^0}})_{(i+2)1} + s_\beta (\boldsymbol{R}^{\boldsymbol{S^0}})_{(i+2)2} + \frac{v_i}{v} \,. \tag{66}$$

For $V V^*$ final states the partial decay width is given by [34]

$$\frac{d\Gamma(\tilde{\nu}_i \to VV^*)}{dx_1 dx_2} = K_{\tilde{\nu}_i VV} \frac{(1-x_1)(1-x_2) + \kappa_V(2x_1 + 2x_2 - 3 + 2\kappa_V)}{(1-x_1-x_2)^2 + \kappa_V \gamma_V}\,, \tag{67}$$

with

$$\kappa_V = \frac{M_V^2}{m_{\tilde{\nu}_i}^2}\,, \quad K_{\tilde{\nu}_i VV} = \frac{3G_F^2 M_W^4}{16\pi^3} |\mathcal{A}_V^i|^2 m_{\tilde{\nu}_i}\, 3\delta'_V \tag{68}$$

and

$$\delta'_W = 1\,, \quad \delta'_Z = \frac{1}{c_W^4}\left(\frac{7}{12} - \frac{10}{9}\, s_W^2 + \frac{40}{9}\, s_W^4\right)\,. \tag{69}$$

The integration variables $x_{1,2}$ lie in the ranges $x_1 = [1 - x_2 - \kappa_V, 1 - \kappa_V/(1-x_2)]$ and $x_2 = [0, 1-\kappa_V]$. The parameter $\gamma_V = \Gamma_V^2/M_V^2$ (Γ_V being the total decay width of the gauge boson V) allows a smooth transition in the threshold region where the off-shell gauge boson becomes on-shell ($m_{\tilde{\nu}_i} = 2M_V$), the calculation of $\Gamma(\tilde{\nu}_i \to VV^*)$ thus requires numerical integration over the variables $x_{1,2}$ in the transition region. Outside that region, i.e. for $m_{\tilde{\nu}_i} \lesssim 2M_V - \Gamma_V$ it can be written according to [35]

$$\Gamma(\tilde{\nu}_i \to VV^*) = \frac{3\,G_F^2 M_V^4}{16\sqrt{2}\,\pi} m_{\tilde{\nu}_i} \delta'_V R_T\left(\frac{M_V^2}{m_{\tilde{\nu}_i}^2}\right) \left|\mathcal{A}_i^V\right|^2\,, \tag{70}$$

where $R_T(x)$ is a kinematic function given by

$$\begin{aligned} R_T(x) = {} & \frac{3(1-8x+20x^2)}{\sqrt{4x-1}} \arccos\left(\frac{3x-1}{2x^{3/2}}\right) - \frac{1-x}{2x}(2 - 13x + 47x^2) \\ & - \frac{3}{2}(1-6x+4x^2)\log x\,. \end{aligned} \tag{71}$$

Finally for CP-even Higgs bosons final states the partial decay width can be written as

$$\Gamma(\tilde{\nu}_i \to S_j^0 S_k^0) = \frac{g^2_{(i+2)jk}}{16\pi\, m_{\tilde{\nu}_i}^3} \lambda^{1/2}(m_{\tilde{\nu}_i}^2, m_{S_j^0}^2, m_{S_k^0}^2) \qquad (j,k = 1,2)\,, \tag{72}$$

where $\lambda(a,b,c) = (a-b-c)^2 - 4bc$, $S_{1,2}^0 = H^0, h^0$ and the dimensionful coupling $g_{(i+2)jk}$ is given in (58). For the particular case $j = k$ (h^0h^0 and H^0H^0 final states) the width reduces to

$$\Gamma(\tilde{\nu}_i \to S_j^0 S_j^0) = \frac{g_{ijj}^2}{32\pi\, m_{\tilde{\nu}_i}} \sqrt{1 - 4\frac{m_{S_j^0}^2}{m_{\tilde{\nu}_i}^2}}\,. \tag{73}$$

4.3 Constraints on BRpV parameters from neutrino data

Neutrino data plays an important role in BRpV, since BRpV parameters contribute to neutrino masses. Therefore, even if BRpV is not solely responsible for neutrino masses,

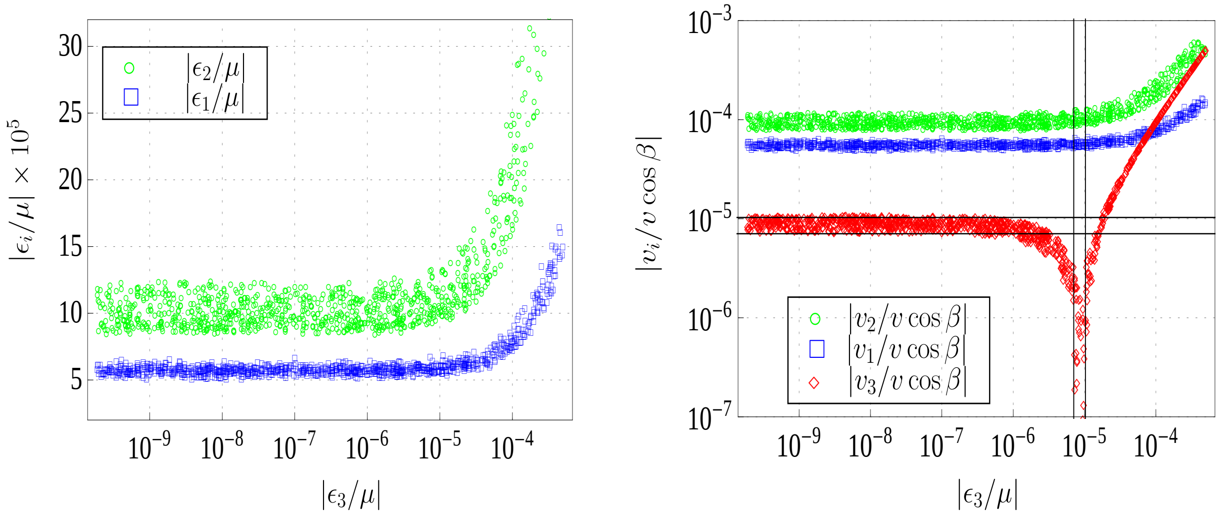


Figure 3: *Numerical ranges for ϵ_i $(i = 1, 2, 3)$ parameters and sneutrino vevs as required for explaining neutrino data for the R-parity conserving parameters discussed in the text. The horizontal (vertical) solid lines in the right-panel plot indicate the values where Λ_3 is dominated by v_3 (ϵ_3), i.e. where $\Lambda_3 \simeq \mu v_3$ $(\Lambda_3 \simeq v\cos\beta\,\epsilon_3)$.*

e.g. [11, 21, 37, 38, 39], one must not saturate neutrino masses via BRpV. Therefore it would be worthwhile to discuss this correlation briefly here. We proceed by assuming BRpV as the sole source of neutrino masses and later consider relaxing this assumption.

At tree-level, BRpV allows for only one massive neutrino, as mentioned earlier. One-loop contributions then leave only one generation massless. Furthermore, for the approximations made in this paper, the tree-level mass is larger than the one-loop mass. The upshot of all this is that only allow so-called normal hierarchy neutrino mass spectrum is allowed here. Recent neutrino data [40, 41, 42] then allows the following neutrino masses:

$$m_1 = 0, \qquad m_2 = \sqrt{\Delta m_{21}^2} = 0.00873 \text{ eV}, \qquad m_3 = \sqrt{\Delta m_{31}^2} = 0.0505 \text{ eV}. \tag{74}$$

In terms of the model parameters

$$\Delta m_{32}^{\text{BRpV}} = \frac{M_{\tilde{\gamma}}}{4|\boldsymbol{M_\chi^0}|}\,|\boldsymbol{\Lambda}|^2\,. \tag{75}$$

The experimental values of the atmospheric and reactor angles yield two additional constraints given by eq. (21). Therefore, the atmospheric sector entirely fixes the Λ_i parameters.

Constraints on the ϵ parameters arise from the solar sector. As long as the one-loop contribution to neutrino masses is smaller than the tree-level one (an assumption we use throughout), $\Delta m_{21}^{\text{BRpV}}$ is accurately determined by $b - \tilde{b}$ one-loop diagrams [30]:

$$\Delta m_{21}^{\text{BRpV}} = \frac{3}{8\pi^2}\sin 2\theta_{\tilde{b}}\frac{m_b^3}{v^2\,c_\beta^2}\Delta B_0^{\tilde{b}_2\tilde{b}_1}\frac{(\bar{\epsilon}_1^2 + \bar{\epsilon}_2^2)}{\mu^2}\,, \tag{76}$$

where $\theta_{\tilde{b}}$ stands for sbottom mixing, $\bar{\epsilon}_{1,2}$ are defined in eq. (24) and $\Delta B_0^{\tilde{b}_2\tilde{b}_1} = B_0(0, m_{\tilde{b}_2}^2, m_b^2) - B_0(0, m_{\tilde{b}_1}^2, m_b^2)$ (with $B_0(0, x, y)$ a scalar Passarino-Veltman function [43]). Note that due to $m_b \ll m_{\tilde{b}_{1,2}}$,

$$\Delta B_0^{\tilde{b}_2\tilde{b}_1} \simeq \log\left(\frac{m_{\tilde{b}_2}^2}{m_{\tilde{b}_1}^2}\right) . \tag{77}$$

Moreover the solar mixing angle can be written as [30]

$$\tan^2\theta_{12}^{\text{BRpV}} = \frac{\bar{\epsilon}_1^2}{\bar{\epsilon}_2^2} . \tag{78}$$

Thus, eqs. (76) and (78) provide two constraints and determine, from eqs. (24), $\epsilon_{1,2}$ ($\epsilon_{1,3}$ or $\epsilon_{2,3}$) as a function of ϵ_3 (ϵ_2 or ϵ_1). Once the Λ's and ϵ's are fixed the sneutrino vevs are automatically fixed as well through

$$v_i = \frac{\Lambda_i - v_d\epsilon_i}{\mu} , \tag{79}$$

see eq. (18). This in turn fixes all the relevant mixings for sneutrino decays ($\chi^0 - \nu$, $\chi^- - \ell_L^-$, $\sigma_{d,u}^0 - \tilde{\nu}_i^R$) once the R-parity conserving supersymmetric parameters are specified. With this knowledge in hand, we start to explore the consequences of the neutrino sector on the BRpV parameters. We note that our results have been verified using SPheno [44][9].

Figure 3 shows typical values for $\epsilon_{1,2}$ and sneutrino vevs as a function of ϵ_3. The plots were obtained by fixing $\theta_{\tilde{b}} = \pi/16$, $\tan\beta = 10$ and

$$(\mu,\ M_1,\ M_2,\ m_{\tilde{b}_2},\ m_{\tilde{b}_1},\ m_A,\ m_{\tilde{\nu}_i}) = (650, 550, 600, 10^3, 700, 10^3, 300)\ \text{GeV},$$

where we have assumed a conservative lower bound on the sneutrino mass of 100 Gev, see [45] for more details. We will use this parameter point throughout the paper. Neutrino observable ($\Delta m_{32}, \Delta m_{21}, \theta_{ij}$) are varied in their 3σ experimental range [40, 41, 42]. We further assume that ϵ_3 is positive. For these assumptions we see from the left-hand side that ϵ_1 and ϵ_2 have a lower bound and from the the right-hande side that this lower bound is larger than v_1 and v_2 respectively. On the other hand, since ϵ_3 is undetermined, v_3 dominates it up to about $|\epsilon_3/\mu| = 10^{-5}$. While allowing negative values for the parameters changes this picture, it leaves one important qualitative property the same: **only in one generation, j, can the sneutrino vev, v_j, be larger than the bilinear mixing term, ϵ_j, when BRpV is the sole contributor to neutrino masses.** This will have important consequences in the next section.

Also of note in the left-panel plot, ϵ_3 obeys only an upper bound determined by the condition $(\boldsymbol{m}_{\boldsymbol{\nu}}^{\text{eff}})^{\text{tree}} > (\boldsymbol{m}_{\boldsymbol{\nu}}^{\text{eff}})^{\text{1-loop}}$, which in terms of the bilinear R-parity violating parameters translates into $|\boldsymbol{\Lambda}| > |\boldsymbol{\epsilon}|^2$. In contrast the $\epsilon_{1,2}$ parameters, due to solar neutrino physics constraints, are forced to lie in a "narrow" range and are such that a region where

[9] We thank Werner Porod for fixing a bug in the decay routines which allowed us to calculate sneutrino decays to gauge bosons, discussed in section 4.4

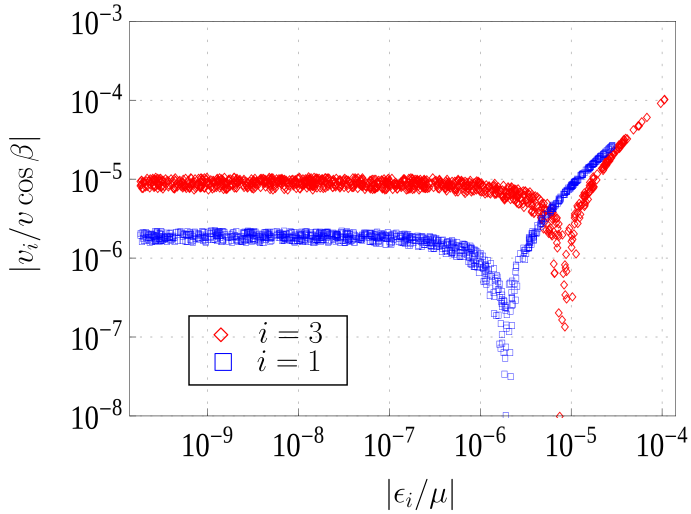


Figure 4: *Values of R-parity breaking parameters for which BRpV is not the sole source of neutrino physics. The parameters are such that* Δm_{32}^{BRpV} *fits the experimental* 3σ *range but* Δm_{21}^{BRpV} *falls below the measured value (see the text for more details).*

$\epsilon_{1,2} \gg \epsilon_3$ exist. Consequently, while $\Lambda_{1,2}$ are mostly determined by $\epsilon_{1,2}$, Λ_3 is controlled by v_3 in the region where $\epsilon_3/\mu \lesssim 10^{-7}$, as demonstrated by the horizontal solid lines in figure 3 (right-panel) which correspond to $\Lambda_3 \simeq \mu v_3$.

The bilinear R-parity breaking parameters selected as described above satisfy neutrino data, and thus lead to BRpV models that can account for neutrino masses and mixings. However, it might be that these parameters are not sufficiently large to account for the neutrino mass scales. In that case their contribution to the atmospheric and solar masses are still determined by eqs. (75) and (76) but are such that, for example, $\Delta m_{31}^{\text{BRpV}} < (\ll)\Delta m_{31}^{\text{Exp}}$ and $\Delta m_{21}^{\text{BRpV}} < (\ll)\Delta m_{21}^{\text{Exp}}$. Figure 4 shows the results for an illustrative case where the BRpV model fits the atmospheric mass scale as well as the atmospheric and reactor angles in their 3σ experimental range, but the contribution to Δm_{21} is subdominant[10]. Note that we plot only $v_{1,3}$, as we have found that $v_2 \sim v_3$. This result as well as $v_1 \ll v_3$ in the region of small $\epsilon_{1,3}$ are due to the constraints arising from fitting θ_{23} and θ_{13}.

In summary, if data is not explained by the BRpV parameters all the ϵ's can be small mainly due to the absence of the solar data constraint, thus implying that, in this case, a region where the three sneutrino vevs are large ($\Lambda_i \simeq \mu\, v_i$) exist.

4.4 LSP sneutrino phenomenology

We are finally ready to address the main aim of this paper: showing that sneutrino LSP decays into heavy SM final states, *i.e.* W^+W^-, ZZ, h^0h^0, and $t\bar{t}$, can dominate over the traditionally considered lighter states. Of the latter, $b\bar{b}$ has been shown to generally have

[10]Models of this kind have been considered in the literature, see e.g. [36, 37].

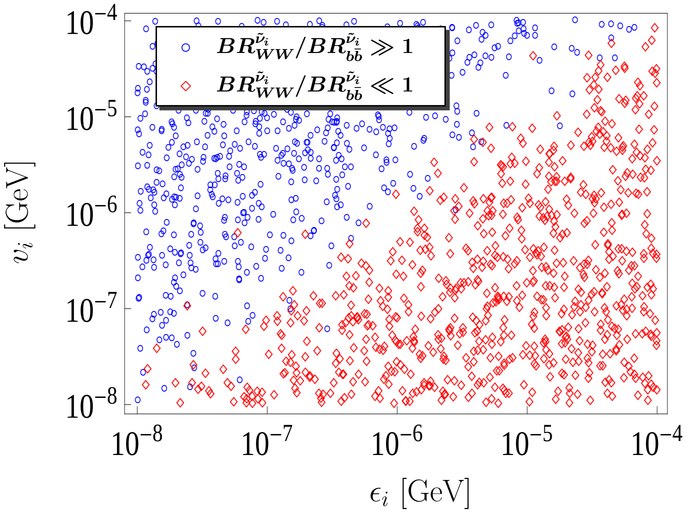


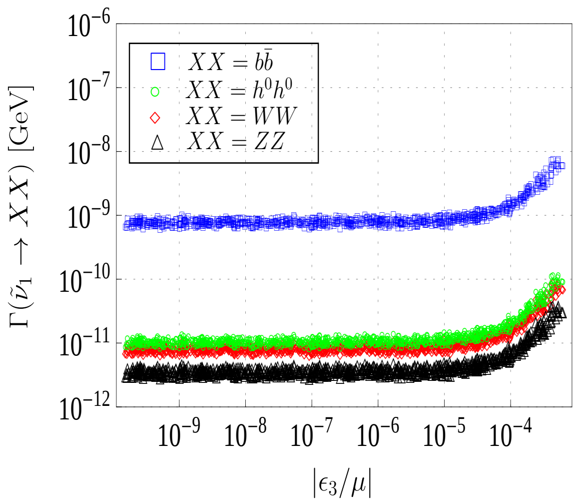
Figure 5: *Results of a scan over SUSY parameters, as shown in the text, comparing the sneutrino decay width into W^+W^- and $b\bar{b}$ final states left plotted in the ϵ_i-v_i plane. It has been assumed that BRpV does not saturate neutrino masses. Blue (red) points indicate where the W^+W^- ($b\bar{b}$) final state dominates. This plot therefore shows an important general result: the W^+W^- dominates only when the sneutrino vev is much larger than its bilinear mixing parameter.*

the largest partial width [25] and we compare all partial widths to it.

We begin with a very general study ignoring all neutrino constrains (SUSY parameters can always be chosen in such a way so that BRpV does not saturate the neutrino masses) and exploring in the ϵ_i-v_i parameter space where W^+W^- final states dominate the $b\bar{b}$ decays of $\tilde{\nu}_i$. Without neutrino constraints, the decay properties are independent of generation. We scan over the following values:

$$\tan\beta = 2 - 50, \quad m_A = 500 - 1000 \text{ GeV}, \quad m_{\tilde{\nu}_i} = 161 - 500 \text{ GeV},$$
$$v_i = 10^{-8} - 10^{-4} \text{ GeV}, \quad \epsilon_i = 10^{-8} - 10^{-4} \text{ GeV},$$

and μ below 1000 GeV but large enough so that the sneutrino is the LSP. We also use $m_h = 125$ GeV as suggested by recent LHC results. The the results are displayed in figure 5 where blue dots indicate the points at which the W^+W^- width is much larger than the $b\bar{b}$ width, while the red points show the opposite. It is striking that the ratio of the partial widths are relatively independent of R-parity conserving SUSY parameters and depend only on the BRpV parameters. From this figure, one can conclude that the W^+W^- partial width dominates the $b\bar{b}$ width only when the sneutrino vev is much larger than the bilinear mixing parameter and, as will be shown below, this condition holds for the other heavy SM final states. Couple this with the main result of the previous section: satisfying neutrino masses solely through BRpV means that $v_i \gg \epsilon_i$ can only be satisfied in one generation, indicates that heavy final states can only dominate the sneutrino BRpV decays for one generation of sneutrinos. Of course, this generation must also be the LSP in order for these decays to be observable.



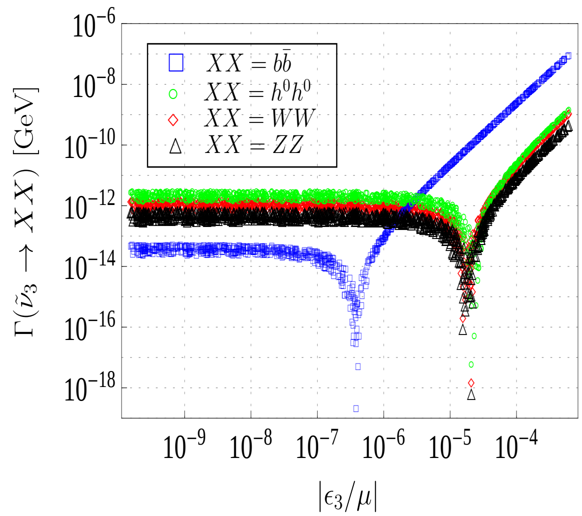


Figure 6: *Partial decay widths for $\tilde{\nu}_1$ (left) and $\tilde{\nu}_3$ (right) to $b\bar{b}$, h^0h^0, WW and ZZ final states as a function of $|\epsilon_3/\mu|$. BRpV parameters have been fixed by using neutrino oscillation data (see the text for more details). Decays of $\tilde{\nu}_2$ are not shown since they are very similar to those of $\tilde{\nu}_1$.*

We continue by considering the specific SUSY point discussed in the last section and examining sneutrino decays into $b\bar{b}$, W^+W^-, ZZ and h^0h^0 assuming that BRpV is solely responsible for neutrino masses (the $t\bar{t}$ channel is suppressed due to the off-shell top for this sneutrino mass). The results are displayed in figure 6 versus ϵ_3, where it can be seen that for $\tilde{\nu}_1$, the dominant decay mode always corresponds to $b\bar{b}$, the h^0h^0 decay branching ratio barely reaches values of $\sim 10^{-2}$ (this holds for $\tilde{\nu}_2$, which was not plotted since its features are similar to $\tilde{\nu}_1$). For $\tilde{\nu}_3$ the situation, however, is quite different. While in the region of large ϵ_3 the $b\bar{b}$ mode dominates, in the region of small ϵ_3, h^0h^0 ZZ and W^+W^- final states become the dominant channels, eventually exceeding the $b\bar{b}$ mode by almost two orders of magnitude.

Figure 6 supports the argument made in reference to figure 5. Reiterating: for a sufficiently heavy sneutrino LSP, even if bilinear R-parity violation accounts for neutrino data, the $b\bar{b}$ channel is not necessarily the dominant decay mode. Therefore, observing a sneutrino decaying dominantly to h^0h^0 (or WW)[11] does not rule out bilinear R-parity violation as responsible for neutrino masses and mixings. However, if the sneutrino mass splittings do not allow the heavier sneutrinos to decay into the LSP, two or all three sneutrino generations will decay via R-parity violating couplings. In this case observing at least two different sneutrino flavors[12] decaying dominantly into heavy SM final states will prove that bilinear R-parity violation does not account for neutrino data, because this would requires $v_i \gg \epsilon_i$ for those generations, which is not possible if neutrino physics

[11] Although h^0h^0 final states have a larger decay branching fraction, depending on the experimental environment, WW modes could be an easier experimental target, as it turns out to be the case at LHC.

[12] In principle for sneutrinos stemming from chargino decays the corresponding sneutrino flavor could be tagged [25].

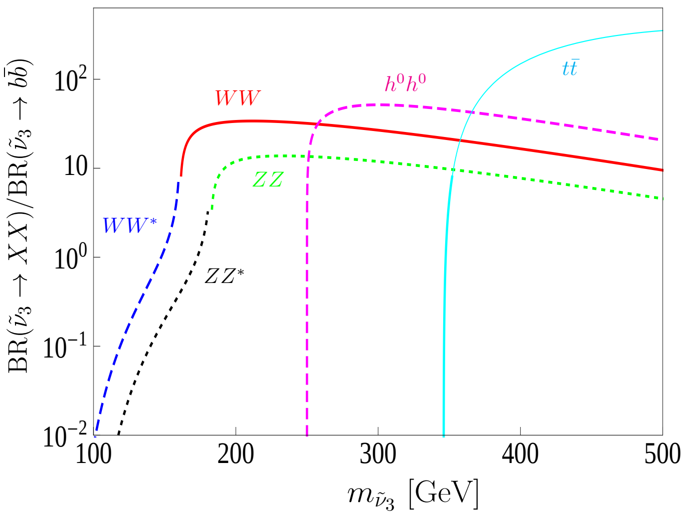


Figure 7: *Decay branching fractions for final states with "strong" kinematical thresholds normalized to $BR(\tilde{\nu}_3 \to b\bar{b})$ as a function of $m_{\tilde{\nu}_3}$. The Branching ratios have been calculated in the small ϵ_3 region ($|\epsilon_3/\mu| = 10^{-9}$) and for bilinear R-parity breaking parameters fixed by neutrino parameters according to their best fit point values [40, 41, 42].*

is determined solely by BRpV, see figure 4.

Focusing further on the heavy SM final states of $\tilde{\nu}_3$, we study its decays as a function of its mass in figure 7 allowing a heavy enough $\tilde{\nu}_3$ so that even the $t\bar{t}$ channel is opened. All the R-parity conserving parameters have been fixed according to the supersymmetric point used in the previous section, the R-parity breaking parameters have been fixed by adjusting neutrino observables to their best fit point values [40, 41, 42] and $\epsilon_3 = \mu \times 10^{-9}$. We have checked the ratios are quite insensitive to changes in $\tan\beta$ in the range [2,30]. Calculations of the WW^* and ZZ^* modes are performed to show where these overtake the $b\bar{b}$ channel, close to the relative thresholds. However, the off-shell calculations for the Higgs and top states have been neglected since in that region the gauge boson states dominate.

Note that this analysis has been done by fixing the bilinear R-parity breaking parameters via neutrino data. If one sticks to BRpV models that do not account for neutrino masses (as in the second case discussed in section 4.3) the results will not drastically change, as what determines the relative importance of the ratios $\mathrm{BR}^{\tilde{\nu}_i}_{XX}/\mathrm{BR}^{\tilde{\nu}_i}_{b\bar{b}}$ is the condition $v_i \gg \epsilon_i$. However, in that case, a quantity that becomes relevant is the sneutrino decay length $L(\tilde{\nu}_i)$ since it can be that the size of the bilinear R-parity breaking parameters leads to a sneutrino decaying out of the detector. Neglecting the Lorentz boost factor and considering only the leading processes (the ones discussed in figure 7), we have found that as long as the R-parity violating parameters fit data, in general $L(\tilde{\nu}_i)$ is well below 1 mm (inline with [27]). We have also checked that if $r_\nu = \Delta m^{\mathrm{BRpV}}_{32,21}/\Delta m^{\mathrm{Exp}}_{32,21} = 10^{-3}$ the decay length is generically below ~ 10 cm. Note that values of r_ν in the range $(10^{-3}, 1)$ might be in conflict with neutrino data when the contributions from the mechanism re-

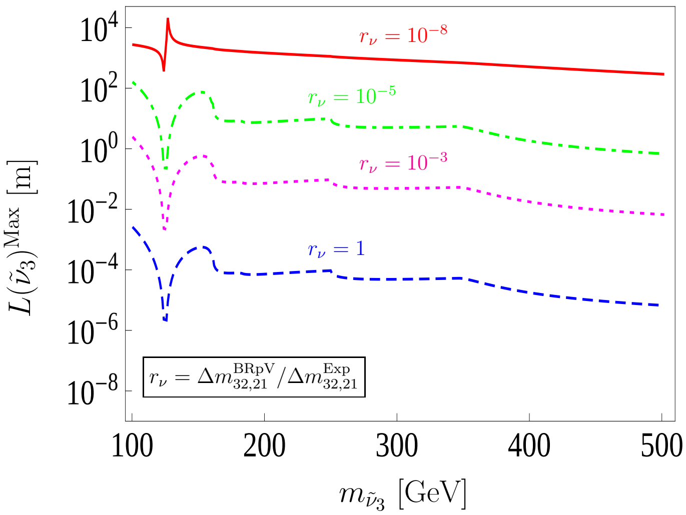


Figure 8: *Estimated upper limit for the sneutrino decay length for different choices of the bilinear R-parity breaking parameters. The BRpV contribution to the atmospheric and solar mass scales is specified by the parameter $r_\nu \equiv \Delta m_{32,21}^{BRpV}/\Delta m_{32,21}^{Exp}$, where $\Delta m_{32,21}^{Exp}$ have been taken according to their best fit point values [40, 41, 42].*

sponsible for neutrino masses are taken into account (for example a standard seesaw). Once the contributions to the atmospheric and solar mass scales fall below $r_\nu = 10^{-3}$ the decay length starts exceeding 1 m for certain sneutrino mass ranges. Figure 8 shows the upper bound on the decay length (we consider contributions only from $b\bar{b}$ and the heavy SM states) for $r_\nu = 1, 10^{-3}, 10^{-5}, 10^{-8}$. The downward spike at $m_{\tilde{\nu}_3} = 125$ GeV is due to the singularity of the neutral scalar mixing at $m_{\tilde{\nu}_i} = m_h$ (see eq. (48)) while the remaining irregularities are a consequence of the different kinematical thresholds (WW, ZZ, h^0h^0 and $t\bar{t}$).

5 Conclusions

BRpV is a good effective theory for models of spontaneous R-parity violation, in which proton decay as well as the number of new parameters is under control compared to explicit R-parity violation. In the presence of BRpV, a sneutrino LSP decays in some way resemble a heavy Higgs and include light SM states, as well as heavy states which have not been considered in the literature before. We have studied these decays into W^+W^-, ZZ, h^0h^0 and $t\bar{t}$ and found that long as $\epsilon_i \ll v_i$, the heavy SM modes dominate. As discussed in section 4.3, for models where the BRpV parameters are fixed by neutrino data, due to the constraints from the solar sector, this is possible only for a single sneutrino flavor. Therefore, if 2 or more sneutrino generations are degenerate enough to decay via BRpV into these heavy SM states one could rule out BRpV as the sole source of neutrino masses.

For models where the bilinear R-parity breaking parameters do not contribute significantly to neutrino masses, the constraint $\epsilon_i \ll v_i$ for all flavors is viable. Accordingly, in these BRpV models large sneutrino LSP branching fractions to W^+W^-, ZZ, h^0h^0 and $t\bar{t}$ are possible, regardless of the sneutrinos mass spectrum. In this case, however, special attention has to be paid to the sneutrino decay length $L(\tilde{\nu}_i)$. We calculated upper limits for $L(\tilde{\nu}_i)$, finding that as long as the BRpV contributions to neutrino masses are not below $\sim 0.1\%$ the decay length is genericaly below ~ 10 cm.

Regardless of these considerations, sneutrino decays into heavy SM final states (W^+W^-, ZZ, h^0h^0 and $t\bar{t}$) are an interesting phenomenological possibility that has not received much attention before. Because BRpV only affects the decays of the LSP, assuming a sneutrino LSP means that SUSY events could cascade decay to two pairs of these heavy SM states, which constitutes a unique and novel SUSY signature.

6 Acknowledgements

We want to thank Martin Hirsch for his comments in the early stage of this project and Werner Porod for SPheno support. SS would also like to thank Pavel Fileviez Perez for comments. DAS is supported by a Belgian FNRS postdoctoral fellowship. DR has been partially supported by Sostenibilidad-UdeA.

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