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Testing the mechanism of R-parity breaking with slepton LSP decays

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Abstract: In supersymmetric models R-parity can be violated through either bilinear or trilinear terms in the superpotential, or both. If charged scalar leptons are the lightest supersymmetric particles, their decay properties can be used to obtain information about the relative importance of these couplings. We show that in some specific scenarios it is even possible to decide whether bilinear or trilinear terms give the dominant contribution to the neutrino mass matrix.

Keywords: Supersymmetry Phenomenology, Solar and Atmospheric Neutrinos, Beyond Standard Model, Neutrino Physics.

Contents

1. Introduction **1**

2. The model **3**

3. Decay widths **5**

4. Neutrino mass matrix **6**

5. Neutrino physics versus collider physics **7**

5.1 Scenario I. $\lambda'_{ijk} \ll \lambda_{ijk}$ 8

5.2 Scenario II. $\lambda_{ijk} \ll \lambda'_{ijk}$ 10

6. Conclusions **12**

A. Determination of λ_{ijk} **13**

1. Introduction

The most general form for the R-parity violating and lepton number violating part of the superpotential is given by

$$W_{\not R_p} = \varepsilon_{ab} \left[\frac{1}{2} \lambda_{ijk} \widehat{L}_i^a \widehat{L}_j^b \widehat{E}_k + \lambda'_{ijk} \widehat{L}_i^a \widehat{Q}_j^b \widehat{D}_k + \epsilon_i \widehat{L}_i^a \widehat{H}_u^b \right]. \tag{1.1}$$

Taking into account the asymmetry of λ_{ijk} in i and j, eq. (1.1) contains nine different λ_{ijk}, 27 λ'_{ijk} and three ϵ_i, for a total of 39 parameters. In principle - as already pointed out in [1] - the number of parameters in eq. (1.1) can be reduced by 3 by a suitable rotation of basis and a number of authors have used this freedom to absorb the bilinear parameters ϵ_i into re-defined trilinear parameters [2].

However, for consistency one has to add three more bilinear terms to V_{soft}, the SUSY breaking potential, [1]

$$V_{\mathrm{soft},\not R_p} = -\varepsilon_{ab} B_i \epsilon_i \widetilde{L}_i^a H_u^b \,. \tag{1.2}$$

It is important to note that there is no rotation which can eliminate the bilinear terms in eq. (1.2) and eq. (1.1) simultaneously [1][2] and thus we are back to 39 independent parameters.

[1]We will not consider the corresponding lepton number violating trilinear soft masses, A_{ijk} and A'_{ijk}, since they will not affect our conclusions.

[2]Unless the parameters B_i and the corrsponding MSSM parameter B are equal (as well as $m^2_{L_i} \equiv m^2_{H_d}$). However, such an equality is unstable under RGE running [3].

With this huge number of parameters it seems hardly possible to make any definite prediction for phenomenology. Could one reduce the number of free parameters in the theory? On theoretical grounds one could argue that a model with only bilinear terms is a self-consistent theory, whereas a model with only trilinear terms is not. The reasoning behind this argument is quite simple: Renormalization group running [3] will generate bilinears whenever trilinears are present, but RGEs which start with bilinears only will never generate trilinears [4]. In addition, models which break R-parity spontaneously [5] at the TeV scale [6] will lead to effective low-energy models where only bilinear terms are present.

On the other hand, experiments are done in the physical mass-eigenstate basis. Before exploring the implications of eq. (1.1), one therefore has to re-diagonalize all mass matrices of the model. The essential point here is that the bilinears introduce mixing among the various states of the model and thus lead to the appearance of "effective" trilinears in the mass basis, for example,

$$\lambda'_{333} \sim \frac{\epsilon_3}{\mu} h_b \,, \tag{1.3}$$

where h_b and μ are the bottom Yukawa coupling and the Higgs mixing parameter. These "effective" trilinears are unavoidable even in models which start with bilinear terms in the superpotential only. The essential point to realize here, however, is that these "effective" trilinears necessarily always follow the hierarchy implied by the Yukawa couplings of the standard model and therefore are *not* new free parameters of the theory in *contrast* to "genuine" trilinears.

All terms in eq. (1.1) violate lepton number by one unit and thus will necessarily contribute to the (Majorana) neutrino mass matrix. Quite a number of articles have investigated the consequences of eq. (1.1) for neutrino physics, some of them considering only the bilinear terms [7, 8] others only the trilinear terms [9] and a few have also entertained the possibility that both kind of terms exist [10].

An obvious question to ask then is: Which of the terms in eq. (1.1) give the dominant contribution to the neutrino masses? And, is there *any* possibility to settle this question *experimentally*? The aim of this paper is to demonstrate that, if charged scalar leptons are the lightest supersymmetric particles (LSPs), there are different broken R parity scenarios where one can probe for the origin of neutrino mass.

How likely is it that charged sleptons are the LSPs? In models of supersymmetry breaking based on minimal supergravity boundary conditions (mSugra) [11], in gauge mediated SUSY breaking [12] and in anomaly mediated SUSY breaking [13, 14] this actually happens in sizeable portions of the parameter space. Usually this possibility is declared "ruled out cosmologically" and simply discarded (except in GMSB, where the charged sleptons are actually the NLSPs). However, with R-parity violated the charged scalars decay, so cosmological limits simply do not apply.[3]

It is exactly the observation that the LSP decays through the same R-parity violating operators, which also govern the entries in the neutrino mass matrix, which forms the basic idea of the current paper. To give a trivial example, in a world were only λ_{ijk}'s

[3] A possible dark matter candidate in this case is the axion.

are non-zero, charged scalars can decay only to leptonic final states, whereas if only λ'_{ijk}'s are non-zero, there will only be hadronic final states. Of course, this example is vastly oversimplied, but here we will argue that, with experimental information from neutrino physics [15], it is possible to fix the size of the various couplings assuming one or the other combination being dominant and check the consequences for the charged slepton decays.

Upper limits on all parameters in eq. (1.1) exist and have been extensively discussed in the literature [2] and [16]–[18]. However, in agreement with [19], we find that limits on the couplings imposed by current neutrino oscillation data are usually much stronger than all other indirect limits.[4] Notable exceptions from this general rule are the limits from double beta decay [18], from non-observation of $\mu \to 3e$, $\mu Ti \to eTi$, and from Δm_B and Δm_K [16, 19]. Here we note only that none of the existing limits is strong enough to require charged scalar leptons decaying with visible decay lengths.

Scalar tau LSPs have been discussed in the literature before. In [20] scalar tau decays within bilinear R-parity violation have been discussed for a one generation model of R-parity violation. In [21] possible similarities between the phenomenology of a charged Higgs and a stau have been discussed within trilinear R-parity violation. In [22] charged scalar LSP decays within bilinear R-parity violation (with 3 generations) have been discussed with special emphasis on their relation to neutrino physics.[5] And in GMSB the scalar tau can be a quite long-lived next-to-lightest supersymmetric particle (NLSP), which from the collider point of view looks like a stable, charged LSP, see for example [12]. Finally charged scalar LSPs have been discussed in an R-parity *conserving* extension of the MSSM in [24]. These authors argue that one should not take into account the cosmological limits on charged LSPs, because alternative cosmologies can provide loopholes to the usual "overclosure" argument [24].

This paper is organized as follows: in the next section 2 we will set up the model definitions. In section 3 the charged scalar decay widths are given. Section 4 summarizes some necessary formulas for the neutrino mass matrix, while section 5 discusses charged scalar lepton decays under different hypotheses both analytically and numerically. We then close with a short summary and outlook.

2. The model

We work in the minimally supersymmetrized version of the standard model (SM) [11], augmented with the R-parity violating terms in eq. (1.1). Superpotential and mass matrices of the MSSM have been given exhaustively in the literature and will not be repeated here. For mass matrices including the bilinear R-parity violating terms see, for example [7].

[4] This claim is true only if the left-right mixing in the squark/slepton sector is not exactly zero. Exactly vanishing left-right mixing can not be excluded at present, but requires fine tuning of parameters. We will not consider this possibility.

[5] After completion of this article, the preprint [23] appeared. The authors study general R-parity violation in mSugra considering also the case of a scalar tau LSP. However, the emphasis of [23] is on the supersymmetry breaking parameters and no attempt is made to explain current neutrino data by the R-parity violating parameters.

JHEP11(2003)005

As will be seen in the next section, the slepton decay widths depend on left-right mixing in the scalar sector. We therefore give a short repetition of $\tilde{\ell}_L$ — $\tilde{\ell}_R$ mixing. The masses and couplings of the $\tilde{\ell}$ follow from the symmetric 2×2 mass matrix which in the basis $(\tilde{\ell}_R, \tilde{\ell}_L)$ reads [11, 25]

$$\mathcal{L}_M^{\tilde{\ell}} = -\left(\tilde{\ell}_R^\dagger, \tilde{\ell}_L^\dagger\right) \begin{pmatrix} M^2_{\tilde{\ell}_{RR}} & M^2_{\tilde{\ell}_{RL}} \\ M^2_{\tilde{\ell}_{LR}} & M^2_{\tilde{\ell}_{LL}} \end{pmatrix} \begin{pmatrix} \tilde{\ell}_R \\ \tilde{\ell}_L \end{pmatrix}, \tag{2.1}$$

where

$$\begin{aligned} M^2_{\tilde{\ell}_{LL}} &= M^2_{\tilde{L}} + m_Z^2 \cos 2\beta \left(\sin^2\Theta_W - \frac{1}{2}\right) + m_\ell^2 \,, \\ M^2_{\tilde{\ell}_{RR}} &= M^2_{\tilde{E}} - m_Z^2 \sin^2\Theta_W \cos 2\beta + m_\ell^2 \,, \\ M^2_{\tilde{\ell}_{RL}} &= M^2_{\tilde{\ell}_{LR}} = m_\ell (A_\ell - \mu \tan\beta) \,. \end{aligned} \tag{2.2}$$

Here $\tan\beta = v_u/v_d$ with $v_d(v_u)$ being the vacuum expectation value of the Higgs field $H_d^0(H_u^0)$, m_ℓ is the mass of the appropriate lepton and Θ_W is the weak mixing angle, μ is the Higgs–Higgsino mass parameter and $M_{\tilde{L}}$, $M_{\tilde{E}}, A_\ell$ are the soft SUSY–breaking parameters of the slepton system [11]. The mass eigenstates $\tilde{\ell}_i$ are $(\tilde{\ell}_1, \tilde{\ell}_2) = (\tilde{\ell}_R, \tilde{\ell}_L)\mathcal{R}^{\tilde{\ell}^T}$ with

$$\mathcal{R}^{\tilde{\ell}} = \begin{pmatrix} \cos\theta_{\tilde{\ell}} & \sin\theta_{\tilde{\ell}} \\ -\sin\theta_{\tilde{\ell}} & \cos\theta_{\tilde{\ell}} \end{pmatrix}, \tag{2.3}$$

with

$$\cos\theta_{\tilde{\ell}} = -\frac{M^2_{\tilde{\ell}_{LR}}}{\sqrt{(M^2_{\tilde{\ell}_{LR}})^2 + (m^2_{\tilde{\ell}_1} - M^2_{\tilde{\ell}_{RR}})^2}}\,, \qquad \sin\theta_{\tilde{\ell}} = \frac{M^2_{\tilde{\ell}_{RR}} - m^2_{\tilde{\ell}_1}}{\sqrt{(M^2_{\tilde{\ell}_{LR}})^2 + (m^2_{\tilde{\ell}_1} - M^2_{\tilde{\ell}_{RR}})^2}} \tag{2.4}$$

while the mass eigenvalues are given by

$$m^2_{\tilde{\ell}_{1,2}} = \frac{1}{2}\left((M^2_{\tilde{\ell}_{LL}} + M^2_{\tilde{\ell}_{RR}}) \mp \sqrt{(M^2_{\tilde{\ell}_{LL}} - M^2_{\tilde{\ell}_{RR}})^2 + 4(M^2_{\tilde{\ell}_{LR}})^2}\right). \tag{2.5}$$

Due to the appearance of the lepton mass in eq. (2.2) one expects the left-right mixing is tiny in the seletron, small in the smuon and potentially sizeable in the stau sectors.

In the R-parity violating version of the MSSM a priori any SUSY particle can be the LSP. However, supplementing the MSSM with mSugra motivated boundary conditions, one usually finds that either the lightest neutralino or one of the charged scalars is the LSP (CSLSP). As a rule of thumb, charged scalars are LSPs if $m_0 \ll M_{1/2}$ and μ large. CSLSPs in these models are usually mainly right sleptons. Furthermore, even though one expects some splitting between $\tilde{\tau}$, $\tilde{\mu}$ and $\tilde{e}$ from RGE running, the latter are not much heavier than the former and so might also dominantly decay through R-parity violating operators even though they are not strictly the LSP.[6]

[6] This is similar to what happens quite generically in GMSB models, where such a scenario is called co-NLSP [12].

JHEP11(2003)005

3. Decay widths

Charged scalar leptons lighter than all other supersymmetric particles will decay through $\not{R}_p$ vertices. Possible final states are either $\ell_i \nu_j$ or $q\bar{q}'$. We will present the decay widths taking into account only the trilinear parameters. For the corresponding formulas for the bilinear terms see [22].

The relevant lagrangian to study these two-body decays of $\tilde{\ell}_1$ is obtained from eq. (1.1). It is given by [2],

$$\begin{aligned} \mathcal{L}_{\rm int} &= \tilde{e}_{j1}\bar{e}_k[(\sin\theta_{\tilde{e}_j}\lambda_{ijk})P_L + (\cos\theta_{\tilde{e}_j}\lambda_{ikj})P_R]\nu_i - \\ &- \sum_j V_{nj}\lambda'_{ijk}\sin\theta_{\tilde{e}_i}\tilde{e}_{i1}\bar{d}_k P_L u_n + \text{h.c.}\,, \end{aligned} \tag{3.1}$$

where $P_{R,L} = (1 \pm \gamma_5)/2$ and V_{nj} is the corresponding element of the CKM matrix. Here we work in a basis where the CKM matrix is solely due to the mixing of the up-type quarks. [$\tilde{e}_{j1}$ denotes the lighter slepton mass eigenstate with family index j]. In the limit $m_{\ell_j} \ll m_{\tilde{\ell}_1}$ the leptonic two–body decay widths of $\tilde{\ell}_1$ read,

$$\Gamma(\tilde{e}_{j1} \to e_k \sum_i \nu_i) = \frac{m_{\tilde{e}_{j1}}}{16\pi}\sum_i[(\sin\theta_{\tilde{e}_j}\lambda_{ijk})^2 + (\cos\theta_{\tilde{e}_j}\lambda_{ikj})^2]. \tag{3.2}$$

For $\tilde{\ell}_1 \simeq \tilde{\ell}_R$, we find for the branching ratios

$$Br^{\tilde{\ell}}_{(1,2,3)} \simeq \frac{1}{2}\left[1 - \frac{(\lambda^2_{23\tilde{\ell}}, \lambda^2_{13\tilde{\ell}}, \lambda^2_{12\tilde{\ell}})}{\sum_{i<j}\lambda^2_{ij\tilde{\ell}}}\right], \tag{3.3}$$

where we have introduced the following shorthand notation: $Br^{\tilde{\ell}}_{(1,2,3)} \equiv Br(\tilde{\ell}_1 \to (e,\mu,\tau) \times \sum \nu_i)$. Corrections to eq. (3.3) are $\propto \theta^2_{\tilde{\ell}} \lll 1$. In the Appendix we give the solutions to eq. (3.2) with respect to the trilinear couplings λ^2_{ijk} by making an expansion in the small parameters $\theta_{\tilde{\ell}}$. In the limit $\theta_{\tilde{\ell}} \to 0$, we derive form eq. (3.3) that

$$Br^{\tilde{\ell}}_i < 0.5\,, \qquad \forall\, i\,. \tag{3.4}$$

It is interesting to note that eq. (3.4) is a *definite prediction* (correct up to order $\theta^2_{\tilde{\ell}}$) *of trilinear R-parity violation* despite the large number of parameters. An experimental result contradicting eq. (3.4) would be a clear sign for bilinear R-parity breaking, if the CSLSPs are mainly right (isosinglet) sleptons.

In order to determine λ^2_{ijk} a measurement of the total width of the sleptons is necessary. Information on the widths could be obtained if a finite decay length is observed, we will return to this point later. However, because the next to leading term in eqs. (A.1)–(A.3) is only of the order $\theta^2_{\tilde{\ell}}$ ratios $\lambda^2_{ijk}/\lambda^2_{i'j'k}$ may be determined without this information from the branching ratios $Br(\tilde{e}_k \to e_l \sum \nu_m)$ alone.

Note also that there are 9 different λ_{ijk} and 9 observables, which implies that in principle it is possible to determine all λ_{ijk} if there is *no* bilinear R-parity breaking. Conversely, adding also the bilinear terms in the superpotential and/or in the soft SUSY breaking

JHEP11(2003)005

potential a full reconstruction of all parameters from the leptonic decays of the CSLSPs is impossible in principle. Differentiating between bilinear and trilinear R-parity breaking therefore necessarily requires the construction of conditions such as eq. (3.4).

For the hadronic two-body decay widths of $\tilde{\ell}_1$ one finds

$$(\Gamma_h^{\tilde{\ell}})_{nk} \equiv \Gamma(\tilde{\ell}_1 \to \bar{u}_n d_k) = N_c \, \beta \, \frac{m_{\tilde{\ell}}}{16\pi} \, \sin^2\theta_{\tilde{\ell}} \left| \sum_j V_{nj} \, {\lambda'}_{\tilde{\ell}jk} \right|^2 , \tag{3.5}$$

where $N_c = 3$ is the number of colours. $\beta = 1$ for $j \neq 3$ and $\beta = (1 - (m_t/m_{\tilde{\ell}})^2)^2$ for $j = 3$. For the expected small mixing in the first generation of sleptons the hadronic width is highly suppressed compared to the leptonic width if $\lambda' \lesssim \lambda$.

4. Neutrino mass matrix

The Majorana mass term for the neutrinos is of the form

$$\mathcal{L}_{\text{mass}} = -\frac{1}{2} \, (\nu_{Li})^T \, C \, (M_\nu)_{ii'} \, \nu_{Li'} + \text{h.c.} , \tag{4.1}$$

where ν_{Li} denote the left-handed weak eigenstates and M_ν is a complex symmetric 3×3 matrix [26].

Contributions to the neutrino mass matrix are induced at tree-level by the bilinear $R\!\!\!/_p$ terms. They can be expressed as [27, 28]

$$m_{\text{eff}} = \frac{M_1 g^2 + M_2 g'^2}{4 \, \det(\mathcal{M}_{\chi^0})} \begin{pmatrix} \Lambda_e^2 & \Lambda_e\Lambda_\mu & \Lambda_e\Lambda_\tau \\ \Lambda_e\Lambda_\mu & \Lambda_\mu^2 & \Lambda_\mu\Lambda_\tau \\ \Lambda_e\Lambda_\tau & \Lambda_\mu\Lambda_\tau & \Lambda_\tau^2 \end{pmatrix} , \tag{4.2}$$

where $\Lambda_i = \mu v_i + v_d \epsilon_i$. Due to the projective nature of the effective neutrino mass matrix m_{eff}, only one neutrino acquires mass at the tree level. Therefore m_{eff} is diagonalized with only two mixing angles which can be expressed in terms of Λ_i:

$$\tan\theta_{13} = -\frac{\Lambda_e}{(\Lambda_\mu^2 + \Lambda_\tau^2)^{\frac{1}{2}}} , \tag{4.3}$$

$$\tan\theta_{23} = -\frac{\Lambda_\mu}{\Lambda_\tau} . \tag{4.4}$$

One-loop contributions to the neutrino mass matrix from bilinear terms have been discussed extensively, see for example [7] and will not be repeated here. At 1-loop level there are also contributions due to trilinear $R\!\!\!/_p$ couplings. The relevant λ-terms are given in eq. (3.1), for λ'_{ijk} type couplings we need:

$$\mathcal{L}_{\lambda'} \supset \lambda'_{ijk} \, [\tilde{d}_{jL} \bar{d}_k P_L \nu_i + \tilde{d}^\dagger_{kR} \bar{\nu}^c_i P_L d_j] + \text{h.c.} . \tag{4.5}$$

The full mass matrix from trilinear terms is then given by $m^{1-loop} = m^\lambda + m^{\lambda'}$, where

$$m_{ii'}^{\lambda(\lambda')} = -\frac{1(N_c)}{32\pi^2} \, \lambda^{(\prime)}_{ijk} \lambda^{(\prime)}_{i'kj} \left[m_k \sin 2\theta_j \ln\left(\frac{m_{2j}^2}{m_{1j}^2}\right) + m_j \sin 2\theta_k \ln\left(\frac{m_{2k}^2}{m_{1k}^2}\right) \right] , \tag{4.6}$$

where m_k is the appropriate fermion mass, θ_j denotes the appropriate sfermion mixing angle and $m^2_{(1,2)j}$ are the corresponding sfermion masses. Note the manifest symmetry of this matrix as well as the presence of logarithmic factors.

We stress that eq. (4.6) is an approximation to the full 1-loop calculation. Numerically we have found that eq. (4.6) is good only up to factors of ~ 2 in extreme cases. In the numerical part of this work we therefore always diagonalize the full 1-loop corrected (7×7) neutrino-neutralino mass matrix.

5. Neutrino physics versus collider physics

In the following we consider some specific scenarios with bilinear and trilinear R-parity breaking terms being present in different combinations and discuss the resulting decay patterns of the charged sleptons. We will calculate the decay lengths

$$L = \frac{\hbar c}{\Gamma}\sqrt{\frac{s}{4m^2} - 1} \tag{5.1}$$

assuming a centre-of-mass energy $\sqrt{s} = 0.8\,\text{TeV}$, as typically could be reached in a future linear collider [29, 30].

In order to reduce the number of parameters, the numerical calculations were performed in a constrained version of the MSSM. We have scanned the parameters in the following ranges: $M_{1/2} \in [0, 1.2]\,\text{TeV}$, $|\mu| \in [0, 2.5]\,\text{TeV}$, $m_0 \in [0, 1.2]\,\text{TeV}$, A_0/m_0 and $B_0/m_0 \in [-3, 3]$ and $\tan\beta \in [2.5, 10]$. All randomly generated points were subsequently tested for consistency with the stationary conditions of the scalar potencial as well as for phenomenological constraints from supersymmetric particle searches [31]. To be conservative we require the charged scalars to be heavier than 100 GeV, although existing limits are somewhat weaker and depend on the flavour of the charged scalar [31]. We then select points in which charged scalars are the LSPs. This latter cut strongly prefers $m_0 \ll M_{1/2}$ and large values of μ. Note again that our charged scalars are mainly right sleptons, as discussed above. Also the low values of $\tan\beta$ in our scan lead to masses for the three different generations of charged sleptons which are rather similar such that all three charged sleptons decay dominantly through R-parity violating decay modes.

The $R\!\!\!/_p$ parameters were chosen in order to fulfill the requirements of the various scenarios to be discussed and as input from neutrino physics we have used, unless noted otherwise, the following ranges, as currently determined from neutrino oscillation experiments [15]:

$$0.3 < \sin^2\theta_{Atm} < 0.7 \tag{5.2}$$

$$1.2 \times 10^{-3}\,\text{eV}^2 < \Delta m^2_{Atm} < 4.8 \times 10^{-3}\,\text{eV}^2 \tag{5.3}$$

$$0.29 < \tan^2\theta_\odot < 0.86 \tag{5.4}$$

$$5.1 \times 10^{-5}\,\text{eV}^2 < \Delta m^2_\odot < 1.9 \times 10^{-4}\,\text{eV}^2 \tag{5.5}$$

$$\sin^2\theta_{13} < 0.05\,. \tag{5.6}$$

JHEP11(2003)005

Before starting the discussion, let us briefly summarize the extreme case in which neutrino oscillation parameters are determined only by bilinear parameters [22]. Full details of this particular situation can be found in [22] and will not be repeated here. We will simply mention the following features which are the most important tools to distinguish the pure bilinear case from more complicated scenarios:

- The hierarchy in Yukawa couplings implies that the slepton decay lengths must be very different, $L(\tilde{\tau}_1) \ll L(\tilde{\mu}_1) \ll L(\tilde{e}_1)$, in particular $L(\tilde{\tau}_1)/L(\tilde{\mu}_1) \sim (m_\mu/m_\tau)^2$
- Ratios of branching ratios must obey the following condition:

$$\frac{Br(\tilde{\tau}_1 \to e \sum \nu_i)/Br(\tilde{\tau}_1 \to \mu \sum \nu_i)}{Br(\tilde{\mu}_1 \to e \sum \nu_i)/Br(\tilde{\mu}_1 \to \tau \sum \nu_i)} \simeq \frac{Br(\tilde{e}_1 \to \tau \sum \nu_i)}{Br(\tilde{e}_1 \to \mu \sum \nu_i)}. \tag{5.7}$$

- Hadronic final states have branching ratios which are much too small to be observable.
- The $\tilde{e}_1$ decays dominantly to the final state $e \sum \nu$ (electrons plus missing (tranvserse) momentum).

Next we will discuss two scenarios with non-zero trilinears.

5.1 Scenario I. $\lambda'_{ijk} \ll \lambda_{ijk}$

Assume the atmospheric mass scale is generated by $m_{\rm eff}$ in eq. (4.2), whereas the solar mass scale is due to m^λ in eq. (4.6). If the couplings λ_{ij1} and λ_{ij2} are not much larger than λ_{ij3} [7] the leading contribution to $m^\lambda_{ii'}$ comes from the $\tilde{\tau}_i - \tau$ loop and the contributions of all other $\tilde{\ell}_i - \ell$ loops are subleading. Then $m^\lambda_{ii'}$ is given approximately as:

$$m^\lambda \approx -\frac{1}{16\pi^2} m_\tau \sin 2\theta_{\tilde{\tau}} \ln\left(\frac{m^2_{\tilde{\tau}_2}}{m^2_{\tilde{\tau}_1}}\right) \begin{pmatrix} \lambda^2_{133} & \lambda_{133}\lambda_{233} & 0 \\ \lambda_{133}\lambda_{233} & \lambda^2_{233} & 0 \\ 0 & 0 & 0 \end{pmatrix}. \tag{5.8}$$

If the parameters in eq. (5.8) are indeed such, that the solar mass squared difference is correctly explained, the scalar tau will decay with a very short decay length, as demonstrated in figure 1. Note that for a possible linear collider currently one expects to be able to measure decay lengths down to 10 μm [32] in an event-by-event analysis. For typical expected luminosities [33] and production cross sections for charged scalars as calculated in [22] the expected number of events is about 10^4 events per year, implying that around 10 events decay with length $\sim$ 10 μm if its mean value is $\sim$ 1.5 μm. In view of the small decay lengths implied by this scenario (figure 1) we think that attempts by experimentalists at improving the measurement of very short decay lengths would be very interesting.

For other λ_{ijk} to provide the solar mass scale, the decay lengths of the corresponding sleptons would have to be even shorter. In this sense, scenario I is a "worst-case" scenario, since only ratios of λ_{ijk} - and a *lower limit* on the absolute values of these parameters - could be measured. However, it is possible to turn this argument around and state that, if a finite decay length for $\tilde{e}_1$ and $\tilde{\mu}_1$ is found λ_{ij1} and λ_{ij2} *can not contribute significantly* to the neutrino mass matrix.

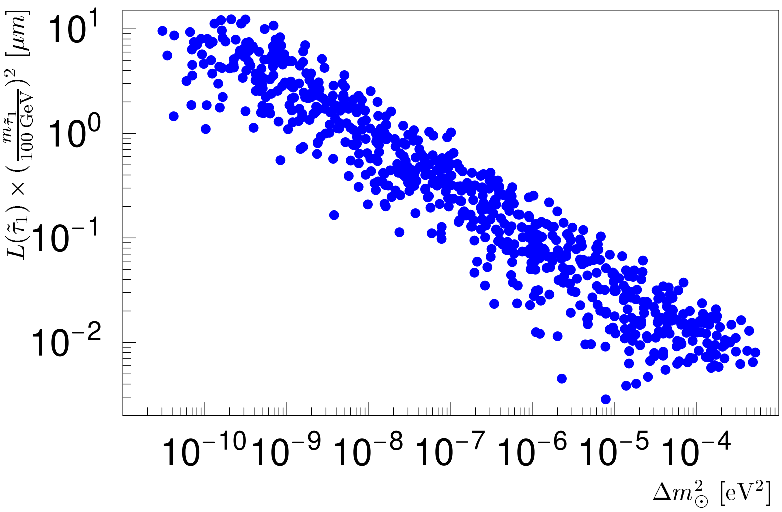

Figure 1: Stau decay length times $(\frac{m_{\tilde{\tau}_1}}{100\,\mathrm{GeV}})^2$ in [μm] versus solar mass squared difference $\Delta m^2_\odot$ [eV^2] in scenario I, defined as λ_{ijk} being responsible for the solar mass scale. If the trilinear loop is responsible for the solar mass, the stau will decay with a length too small to be measured.

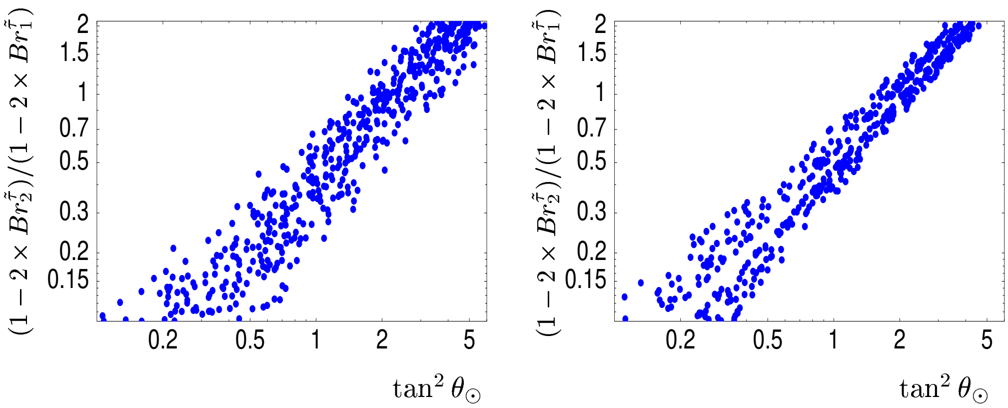

Figure 2: Ratio of branching ratios $(1 - 2 \times Br^{\tilde{\tau}}_2)/(1 - 2 \times Br^{\tilde{\tau}}_1)$ versus the solar angle, $\tan^2\theta_\odot$. The figure to the left is for $\forall\lambda_{ijk}$ non-zero and of similar magnitude, while the figure to the right assumes $\lambda_{123} = 0$.

With a decay length of $\tilde{\tau}_1$ too small to be measured, only a consistency check of this scenario is possible. This check is provided by the measurement of a certain ratio of branching ratios, which is correlated with the solar angle, as demonstrated in figure 2. We want to stress, however, that this measurement can not distinguish trilinear terms from a bilinear-only world, since the latter leads to a very similar prediction for the branching ratios. Instead only disagreement between prediction and measurement would establish that neither scenario I nor bilinear loops are responsible for the solar neutrino mass scale.

[7]Whether this is indeed the case may be checked, in principle, by a comparison of the decay length of the sleptons.

5.2 Scenario II. $\lambda_{ijk} \ll \lambda'_{ijk}$

We now turn to the case where $\lambda_{ijk} \ll \lambda'_{ijk}$. If the λ'_{ijk} do not follow an *inverse* hierarchy, then the main contribution of the trilinears to the neutrino mass, eq. (4.6), is approximately given by

$$m^{\lambda'} \approx -\frac{3}{16\pi^2} m_b \sin 2\theta_{\tilde{b}} \ln\left(\frac{m^2_{\tilde{b}_2}}{m^2_{\tilde{b}_1}}\right) \begin{pmatrix} \lambda'^2_{133} & \lambda'_{133}\lambda'_{233} & \lambda'_{133}\lambda'_{333} \\ \lambda'_{133}\lambda'_{233} & \lambda'^2_{233} & \lambda'_{233}\lambda'_{333} \\ \lambda'_{133}\lambda'_{333} & \lambda'_{233}\lambda'_{333} & \lambda'^2_{333} \end{pmatrix}, \tag{5.9}$$

since the $\tilde{b}_i - b$ loop is largest due to the larger left-right mixing expected in the sbottom sector and $m_d \ll m_s \ll m_b$. Note that $m^{\lambda'}$ has also a projective structure, similar to $m_{\rm eff}$ in eq. (4.2), but with different parameters.

There are two limiting case, which we will study in some detail:

- Scenario IIA: Bilinear R-parity breaking responsible for solar physics, λ'_{i33} for the atmospheric mass scale
- Scenario IIB: Bilinear R-parity breaking responsible for atmospheric physics, λ'_{i33} for the solar mass scale

Consider scenario IIA first. Here, due to the special structure of eq. (5.9), $\tan\theta_{23}$ and $\tan\theta_{13}$ can be approximately expressed as

$$\tan\theta_{23} \approx -\frac{\lambda'_{233}}{\lambda'_{333}}, \tag{5.10}$$

$$\tan\theta_{13} \approx -\frac{\lambda'_{133}}{\sqrt{(\lambda'_{233})^2 + (\lambda'_{333})^2}}. \tag{5.11}$$

The ratio in eq. (5.10) squared can be related to the observable

$$\frac{(\Gamma^{\tilde{\mu}}_h)_{33}}{(\Gamma^{\tilde{\tau}}_h)_{33}} \simeq \left(\frac{\sin\theta_{\tilde{\mu}}\ \lambda'_{233}}{\sin\theta_{\tilde{\tau}}\ \lambda'_{333}}\right)^2 \approx \frac{m^2_\mu}{m^2_\tau}\ \tan^2\theta_{23}\,, \tag{5.12}$$

where $(\Gamma^{\tilde{\ell}}_h)_{33}$ denotes the partial width of $\tilde{\ell}_1$ decaying into the final state $\bar{t}b$. Due to the tiny selectron mixing we expect the number of events $N(\tilde{e}_1 \to \bar{t}b) \ll N(\tilde{\mu}_1 \to \bar{t}b)$, such that the corresponding relation for eq. (5.11) will be impossible to test.

For simplicity let us first consider the case in which only the λ'_{i33} are non-zero. Modifications of the results for other $\lambda'_{ijk} \neq 0$ will be discussed at the end of this section.

In figure 3 we show the atmospheric angle versus $\lambda'_{233}/\lambda'_{333}$ (on the left panel) and $(\Gamma^{\tilde{\mu}}_h)_{33}/(\Gamma^{\tilde{\tau}}_h)_{33}$ (on the right panel) demonstrating to which accuracy one expects eqs. (5.10) and (5.12) to work. Obviously eq. (5.12) can be tested only if $m_{\tilde{\tau}_1}$ is bigger than $m_t + m_b$.

Figure 4 (to the left) then shows the calculated decay lengths for $\tilde{\tau}_1$, $\tilde{\mu}_1$ and $\tilde{e}_1$ in scenario IIA. The top threshold for the $\tilde{\tau}_1$ and $\tilde{\mu}_1$ decays is clearly visible, while it is invisible for the selectron case due to the negligible left-right mixing dictated by the small electron Yukawa coupling. This explains why the $\tilde{e}_1$ decay length does not vary much over the range of masses shown. This is due to the fact that the decay of the

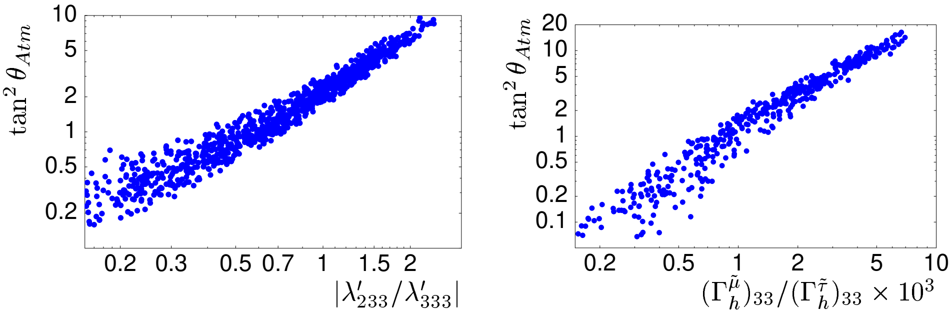

Figure 3: Ratio $\lambda'_{233}/\lambda'_{333}$ (left panel) and $(\Gamma_h^{\tilde{\mu}})_{33}/(\Gamma_h^{\tilde{\tau}})_{33}$ (right panel) versus the atmospheric angle, $\tan^2\theta_{Atm}$, for data points satisfying the criteria defined as scenario IIA. The plot to the right contains only points in which the decay $\bar{t}b$ is kinematically possible.

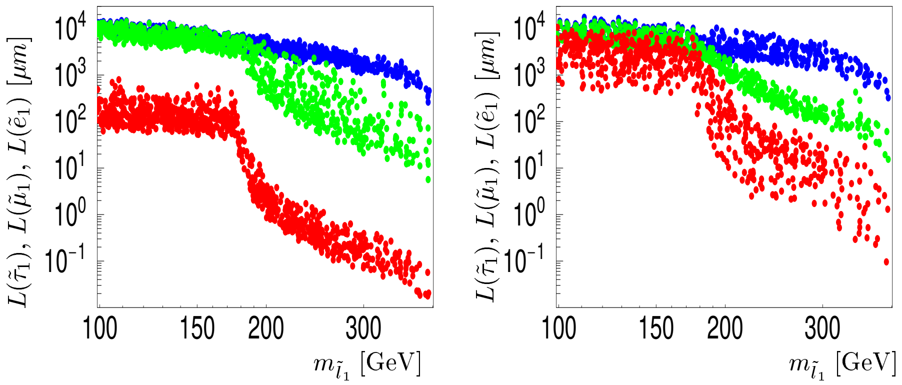

Figure 4: Decay lengths for $\tilde{\tau}_1$, $\tilde{\mu}_1$ and $\tilde{e}_1$, $L(\tilde{\tau}_1)$, $L(\tilde{\mu}_1)$, $L(\tilde{e}_1)$ $[\mu m]$, (bottom to top) versus the mass of the decaying particle, $m_{\tilde{l}_1}$ [GeV]. On the left panel: Scenario IIA, right panel scenario IIB. The plots assume that only λ'_{i33} are different from zero, see text.

$\tilde{e}_1$ is completely dominated by the final state $e\sum\nu$, induced by the bilinear term $|\vec{\Lambda}|$. For the $\tilde{\tau}_1$ below threshold the decay is dominated by the final state $\bar{c}b$, which is induced from λ'_{333} due to the non-zero value of V_{cb}, with a branching ratio into $\tau\sum\nu$ being a few percent. Above threshold the branching into $\bar{t}b$ quickly grows to more than 99 %, with the final state $\bar{c}b$ being suppressed by a factor of V_{cb}^2 and the final state $\tau\sum\nu$ is tiny. For the decay of the $\tilde{\mu}_1$, on the other hand, below threshold $\mu\sum\nu$ dominates, with a branching ratio into $\bar{c}b$ of about $\sim$ (1-20) %, while above threshold $\bar{t}b$ becomes quickly dominant, but (some) percent of the decays still go to $\mu\sum\nu$. Decays $\tilde{e}_{i1} \to e_j\sum\nu$ for $i \neq j$ are always very small, due to the assumed smallness of λ_{ijk}.

From figure 4 one concludes that an experimental test of eq. (5.12) might be possible, but the decay lengths of the $\tilde{\tau}_1$ are so short such that either (i) an improvement in detecting short decay lengths will be necessary or (ii) additional input from the sbottom sector is needed.[8]

[8]Fixing the sbottom mixing angle and masses would allow to *calculate* the absolute size of λ'_{333} needed for scenario IIA, thus fixing the total $\tilde{\tau}_1$ width.

Turning to scenario IIB, in the right panel of figure 4 we plot the calculated decay lengths of the charged scalar leptons as a function of the charged scalar mass. The main difference to scenario IIA is that the $\tilde{\tau}_1$ now decays with a visible length, apart from some exceptional points where $m_{\tilde{\tau}_1}$ is very close to $\sqrt{s}/2$. This can be traced to the fact that smaller values of λ'_{i33} are needed in this scenario since the solar mass scale is smaller than the atmospheric one. Note that the plot assumes that all λ'_{i33} are of the same order of magnitude. While λ'_{333} can not be much larger than the values used in this plot without violating the main assumption of scenario IIB, it could well be smaller than λ'_{133} or λ'_{233}, since non-zero values for the latter are sufficient to create a non-zero solar neutrino mass difference (and mixing). We note in passing that in the (unlikely) case of $\lambda'_{333} = 0$ the $\tilde{\tau}_1$ decay lengths would approach the $\tilde{e}_1$ decay lengths shown, with the dominating final state being $\tau \sum \nu$ (again due to the assumed absence of λ_{ij3}).

Let us finally discuss modifications of the results discussed for scenarios IIA and IIB once we allow for other $\lambda'_{ijk} \neq 0$. Non-zero λ'_{ijk} would, of course, affect the branching ratios discussed above. Since final states $\bar{q}_j q_k$ can not be distinguished experimentally for $j = 1$ and $k = 1, 2$, only the sum of the corresponding $(\lambda'_{ijk})^2$ are measurable. More important for us, however, is that additional final states will shorten the decay lengths of the sleptons compared to the ones shown in figure 4. Numerically we find that in scenario IIA the decay length of the $\tilde{\mu}_1$ is always visible unless λ'_{2jk} for $j, k \neq 3$ are several times *larger* than the values for λ'_{233} needed to explain the atmospheric neutrino mass difference. The situation is different for $\tilde{\tau}_1$. Here, values of λ'_{3jk} for $j, k \neq 3$ smaller than the "correct" λ'_{333} by a factor of a few are needed, otherwise the decay length of the $\tilde{\tau}_1$ becomes to short to be realistically measurable. In scenario IIB similar comments apply, although in this case even larger ratios $\lambda'_{ijk}/\lambda'_{233}$ are allowed for a visible decay length of the $\tilde{\mu}_1$ decays. For the $\tilde{\tau}_1$ in scenario IIB λ'_{3jk} as big as λ'_{333} are possible for decay lengths to be observable. The decay of the $\tilde{e}_1$, on the other hand, never shows any sensitivity to λ'_{1jk}, again due to the tiny left-right mixing in the selectron sector.

6. Conclusions

We have studied the decay properties of a right charged scalar lepton LSP (CSLSP) in a general model of R-parity violation containing both bilinear and trilinear terms. Branching ratios and decay lengths of CSLSPs contain information about ratios and absolute values of R-parity violating couplings which at the same time contribute to the neutrino mass matrix. We have investigated to what extent it is possible to test experimentally whether bilinear or trilinear terms give the dominant contribution to the neutrino masses.

Due to the huge number of parameters characterizing trilinear models the resulting phenomenology can be quite diverse. In fact one can envisage a mixed situation in which bilinear terms are responsible for generating the neutrino masses required to account for current oscillation data, yet the CSLSP decays are governed by trilinears. Thus in general the complexity of the physics resulting from eq. (1.1) is such that the existence of non-zero bilinear terms can not be established.

JHEP11(2003)005

Nevertheless, we have shown that different scenarios for neutrino masses exist which lead to considerably different phenomenology and thus allows for the origin of neutrino mass to be experimentally probed. In some cases we find that bilinears and trilinears can be clearly distinguished. Especially noteworthy is that in trilinear-only models all right CSLSPs should obey $Br(\tilde{\ell}_1 \to (e,\mu,\tau)\sum \nu_i) < 0.5$. This is to be contrasted with the bilinear model predicting [22] $Br(\tilde{e}_1 \to e\sum \nu_i) \simeq 1$.

Decay lengths for the CSLSPs could be measurably large or very small depending mainly on the absolute values of λ_{ijk}. Observing a finite length for $\tilde{e}_1$ and/or $\tilde{\mu}_1$ would establish that the corresponding λ_{ij1} and λ_{ij2} do not contribute significantly to the neutrino mass matrix. If λ_{ijk} are somewhat smaller than λ'_{ijk} hadronic final states will have measurable branching ratios, at least for the $\tilde{\tau}_1$ and depending on $\lambda_{ij2}/\lambda'_{2jk}$ also for the $\tilde{\mu}$. In contrast, note that in the bilinear-only model [22] hadronic final states are never visible.

Finally, since reasonable scenarios for the neutrino mass matrix exist in which the decay length of the $\tilde{\tau}_1$ is just at or below the borderline of what is currently thought of being experimentally accessible [30] we stress that efforts to optimize decay length measurements might be worth undertaking.

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A. Determination of λ_{ijk}

Here we give the solutions for the trilinear couplings λ^2_{ijk} by inverting eq. (3.2) and making an expansion in $\theta_{\tilde{\ell}}$ up to $\mathcal{O}(\theta^2_{\tilde{\ell}})$. This yields

$$
\begin{aligned}
\lambda^2_{121} &= \frac{1}{2}\, C^{-1}\, \{\Gamma^{\tilde{e}}_{\rm tot}(1-2Br^{\tilde{e}}_3) + \theta^2_{\tilde{e}}(\Gamma^{\tilde{e}}_{\rm tot}(Br^{\tilde{e}}_2 - Br^{\tilde{e}}_3) - \Gamma^{\tilde{\mu}}_{\rm tot}Br^{\tilde{\mu}}_1 + \Gamma^{\tilde{\tau}}_{\rm tot}Br^{\tilde{\tau}}_1)\}\,, \\
\lambda^2_{131} &= \frac{1}{2}\, C^{-1}\, \{\Gamma^{\tilde{e}}_{\rm tot}(1-2Br^{\tilde{e}}_2) + \theta^2_{\tilde{e}}(\Gamma^{\tilde{e}}_{\rm tot}(Br^{\tilde{e}}_3 - Br^{\tilde{e}}_2) + \Gamma^{\tilde{\mu}}_{\rm tot}Br^{\tilde{\mu}}_1 - \Gamma^{\tilde{\tau}}_{\rm tot}Br^{\tilde{\tau}}_1)\}\,, \\
\lambda^2_{231} &= \frac{1}{2}\, C^{-1}\, \{\Gamma^{\tilde{e}}_{\rm tot}(1-2Br^{\tilde{e}}_1) + \theta^2_{\tilde{e}}(\Gamma^{\tilde{e}}_{\rm tot}(Br^{\tilde{e}}_3 + Br^{\tilde{e}}_2) - \Gamma^{\tilde{\mu}}_{\rm tot}Br^{\tilde{\mu}}_1 - \Gamma^{\tilde{\tau}}_{\rm tot}Br^{\tilde{\tau}}_1)\}\,, \qquad \text{(A.1)}
\end{aligned}
$$

$$
\begin{aligned}
\lambda^2_{122} &= \frac{1}{2}\, C^{-1}\, \{\Gamma^{\tilde{\mu}}_{\rm tot}(1-2Br^{\tilde{\mu}}_3) + \theta^2_{\tilde{\mu}}(\Gamma^{\tilde{\mu}}_{\rm tot}(Br^{\tilde{\mu}}_2 - Br^{\tilde{\mu}}_3) - \Gamma^{\tilde{e}}_{\rm tot}Br^{\tilde{e}}_2 + \Gamma^{\tilde{\tau}}_{\rm tot}Br^{\tilde{\tau}}_2)\}\,, \\
\lambda^2_{132} &= \frac{1}{2}\, C^{-1}\, \{\Gamma^{\tilde{\mu}}_{\rm tot}(1-2Br^{\tilde{\mu}}_2) + \theta^2_{\tilde{\mu}}(\Gamma^{\tilde{\mu}}_{\rm tot}(Br^{\tilde{\mu}}_2 + Br^{\tilde{\mu}}_3) - \Gamma^{\tilde{e}}_{\rm tot}Br^{\tilde{e}}_2 - \Gamma^{\tilde{\tau}}_{\rm tot}Br^{\tilde{\tau}}_2)\}\,,
\end{aligned}
$$

$$\lambda_{232}^2 = \frac{1}{2}\, C^{-1}\, \{\Gamma^{\tilde\mu}_{\rm tot}(1-2Br^{\tilde\mu}_1) + \theta^2_{\tilde\mu}(\Gamma^{\tilde\mu}_{\rm tot}(Br^{\tilde\mu}_3 - Br^{\tilde\mu}_2) + \Gamma^{\tilde e}_{\rm tot}Br^{\tilde e}_2 - \Gamma^{\tilde\tau}_{\rm tot}Br^{\tilde\tau}_2)\}\,, \quad (A.2)$$

$$\begin{aligned}
\lambda_{123}^2 &= \frac{1}{2}\, C^{-1}\, \{\Gamma^{\tilde\tau}_{\rm tot}(1-2Br^{\tilde\tau}_3) + \theta^2_{\tilde\tau}(\Gamma^{\tilde\tau}_{\rm tot}(Br^{\tilde\tau}_1 + Br^{\tilde\tau}_2) - \Gamma^{\tilde e}_{\rm tot}Br^{\tilde e}_3 - \Gamma^{\tilde\mu}_{\rm tot}Br^{\tilde\mu}_3)\}\,,\\
\lambda_{133}^2 &= \frac{1}{2}\, C^{-1}\, \{\Gamma^{\tilde\tau}_{\rm tot}(1-2Br^{\tilde\tau}_2) + \theta^2_{\tilde\tau}(\Gamma^{\tilde\tau}_{\rm tot}(Br^{\tilde\tau}_1 - Br^{\tilde\tau}_2) - \Gamma^{\tilde e}_{\rm tot}Br^{\tilde e}_3 + \Gamma^{\tilde\mu}_{\rm tot}Br^{\tilde\mu}_3)\}\,,\\
\lambda_{233}^2 &= \frac{1}{2}\, C^{-1}\, \{\Gamma^{\tilde\tau}_{\rm tot}(1-2Br^{\tilde\tau}_1) + \theta^2_{\tilde\tau}(\Gamma^{\tilde\tau}_{\rm tot}(Br^{\tilde\tau}_2 - Br^{\tilde\tau}_1) + \Gamma^{\tilde e}_{\rm tot}Br^{\tilde e}_3 - \Gamma^{\tilde\mu}_{\rm tot}Br^{\tilde\mu}_3)\}\,, \quad (A.3)
\end{aligned}$$

where $\Gamma^{\tilde\ell}_{\rm tot}$ denotes the total leptonic decay width of the appropriate scalar lepton, which is, for the case where $\tilde\ell_1 \simeq \tilde\ell_R$, not much different from the total decay width if $\lambda' \lesssim \lambda$. $C \equiv \frac{m_{\tilde\ell_1}}{16\pi}$, where we have made use of the fact that $m_{\tilde e_1} \simeq m_{\tilde\mu_1} \simeq m_{\tilde\tau_1} (\equiv m_{\tilde\ell_1})$. As can be seen in Eqs. (A.1)-(A.3) the next to leading term is only of the order $\theta^2_{\tilde\ell}$. The measurement of $\Gamma^{\tilde\ell}_{\rm tot}$ and $\theta_{\tilde\ell}$ provides the information whether the expansion to $\mathcal{O}(\theta^2_{\tilde\ell})$ is sufficient.

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