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Lightest-neutralino decays in R_p-violating models with dominant λ' and λ couplings

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Abstract: Decays of the lightest neutralino are studied in R_p-violating models with operators $\lambda' LQD^c$ and λLLE^c involving third-generation matter fields and with dominant λ' and λ couplings. Generalizations to decays of the lightest neutralino induced by subdominant λ' and λ couplings are straightforward. Decays with the top-quark among the particles produced are considered, in addition to those with an almost massless final state. Phenomenological analyses for examples of both classes of decays are presented. No specific assumption on the composition of the lightest neutralino is made, and the formulae listed here can be easily generalized to study decays of heavier neutralinos. It has been recently pointed out that, for a sizable coupling λ'_{333}, tau-sleptons may be copiously produced at the LHC as single supersymmetric particles, in association with top- and bottom-quark pairs. This analysis of neutralino decays is, therefore, a first step towards the reconstruction of the complete final state produced in this case.

Keywords: Supersymmetry Phenomenology, Beyond Standard Model, Supersymmetric Standard Model.

Contents

1. Introduction

Within the Standard Model (SM), electroweak gauge invariance is sufficient to ensure conservation of both the lepton number L and the baryon number B, at least in a perturbative context. This is not the case in supersymmetric theories, wherein it is possible to write terms in the superpotential that are invariant under supersymmetric and SM gauge transformations, but that do not conserve B, nor L. These terms violate a discrete symmetry [1, 2], called R-parity, R_p, which implies a conserved quantum number

$R_p \equiv (-1)^{3B+L+2S}$, where S stands for the spin of the particle. It is clear from this definition that all the SM particles have $R_p = +1$, whereas all superpartners have $R_p = -1$. It is this discrete symmetry that guarantees the stability of the lightest supersymmetric particle (LSP) and thus also supplies a candidate for the cold dark matter. The R_p-violating terms in the superpotential can be written as

$$W_{\not R_p} = \frac{1}{2}\lambda''_{ijk} U_i^c D_j^c D_k^c - \frac{1}{2}\lambda_{ijk} L_i L_j E_k^c - \lambda'_{ijk} L_i Q_j D_k^c + \kappa_i L_i H_2 \,, \tag{1.1}$$

where L_i, Q_i represent the SU(2) lepton and quark doublet superfields, E_i^c, D_i^c and U_i^c, the SU(2) lepton and quark singlet superfields, and i, j, k are generation indices. The symbol H_2 denotes the Higgs superfield that gives rise to the entries in the up-type quark mass matrix, through the Yukawa term $h_U H_2 Q U^c$.

Imposition of R_p invariance gets rid of the potentially dangerous terms in eq. (1.1) that can cause a too fast proton decay. To this aim, however, it is sufficient to forbid the simultaneous presence of the B-violating λ'' terms and the L-violating λ' terms. The existence of B-violating terms that are not negligibly small can wash out baryogenesis [3, 4]. (Baryogenesis may be triggered by leptogenesis through nonperturbative effects. In this case, leptogenesis requires that at least one combination of the L-violating couplings is very small [5, 4].) Furthermore, in studies of unified string theories it was shown [6] that there exist discrete symmetries which treat leptons and quarks differently, as R_p does. In particular, in this context, given the particle content of the minimal supersymmetric Standard Model (MSSM) [7], the lack of rapid proton decay implies two symmetries: R_p and B. Imposing the first eliminates only the dimension-*four* operators that cause proton decay, whereas B conservation removes also the dimension-*five* operators. Thus, the option of B conserved and R_p broken through L violation seems the theoretically interesting one, and it is the one adopted in this paper.

Data from neutrino oscillation experiments show now clear evidence for nonvanishing neutrino masses [8]. R_p-violating models, with their new interaction terms, provide an alternative way to generate these masses at the tree and quantum level [9]–[13] consistent with all informations from the different oscillation experiments, without having to introduce new superfields in addition to those of the MSSM [14].

This discussion makes it clear that studies of R_p-violating models are interesting from a theoretical point of view. A much more pragmatic interest in them is due to the instability of the LSP induced by R_p-violating couplings. If at least one of the R_p-violating couplings is $> 10^{-6}$ or so, then, the LSP decays within the detector [15]. The instability of the LSP changes the phenomenology of sparticle searches at colliders quite drastically (see, among others, ref. [16]), as it changes the type, the number and the energies of the final state particles. If not too small, these couplings may give rise to interesting collider signatures, other than those due to the LSP decay (see for example ref. [17]).

The flip side of considering R_p-violating models, of course, is the existence of additional unknown parameters, as many as 48 in the superpotential. Unfortunately, there is no theoretical indications about their size, as for their R_p-conserving counterparts, the Yukawa couplings. Lacking any theoretical guideline about the magnitude of these couplings, one

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can rely only on phenomenological constraints obtained in various experiments to get a clue to their size. As discussed in section 2, the dimensionful couplings κ_i are severely restricted by tree-level contributions to neutrino masses. Dimensionless trilinear couplings are, in general, also constrained. If one excludes neutrino data, however, the restrictions are less severe for couplings with at least one third-generation index. It is, therefore, not unreasonable to expect that the former are the largest couplings, as the corresponding third generation Yukawa couplings are larger than the first and second generation ones.

Hence, it is an interesting question to ask what role can LHC play in probing these couplings. In ref. [17], it was pointed out that a sizable λ'_{333} can give rise to the process $pp \to t\tilde{\tau}X$, or $pp \to t\bar{b}\tilde{\tau}X$, depending on how many b-quarks in the final state can actually be tagged. (In the following, for the sake of clarity, we shall neglect this complication and list final states with all produced b-quarks.) The production cross section is, for example, $\sim 10fb$ for $m_{\tilde{\tau}}$ up to 1 TeV, if $\lambda'_{333} \sim 0.5$. Such a large value of λ'_{333} may control also the decay of $\tilde{\tau}$-slepton, as well as that of the lightest neutralino and chargino, thus providing very distinctive final states for the signal. To be specific, the slepton $\tilde{\tau}$ can decay into the lightest neutralino, $\tilde{\chi}_1^0$, and, if kinematically allowed, in the lightest chargino, $\tilde{\chi}_1^-$, through gauge interactions:

$$\tilde{\tau} \to \tau\tilde{\chi}_1^0\,, \qquad \tilde{\tau} \to \nu_\tau\tilde{\chi}_1^-\,. \tag{1.2}$$

If heavy enough, it may also decay with a substantial rate into:

$$\tilde{\tau} \to b\bar{t}\,, \tag{1.3}$$

through the R_p-violating interaction $\lambda'_{333}L_3Q_3D_3^c$.[1] In turn, the lightest neutralino obtained from $\tilde{\tau}$, as in eq. (1.2), decays as:

$$\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau\,, \qquad \tilde{\chi}_1^0 \to t\bar{b}\tau\,, \qquad \tilde{\chi}_1^0 \to \bar{t}b\bar{\tau}\,, \tag{1.4}$$

with the second and third decay mode allowed only for a heavy $\tilde{\chi}_1^0$, i.e. for $m_{\tilde{\chi}_1^0} > m_t + m_b + m_\tau$. One should notice that in the case of the decay of eq. (1.3), the final signature

$$pp \to t\bar{b}\,\tilde{\tau}X \to t\bar{b}\,\bar{t}bX\,, \tag{1.5}$$

is also that obtained from a charged Higgs boson[2] produced in association with a t- and b-quark pair through R_p-conserving interactions [17, 22], or from pair-produced charged Higgs bosons. On the contrary, the final states obtained when the $\tilde{\tau}$-slepton decays, for example, as $\tilde{\tau} \to \tau\tilde{\chi}_1^0$, i.e.,

$$pp \to t\bar{b}\,\tilde{\tau}X \to (2t)\,(2\bar{b})\,(2\tau)X\,, \qquad t\bar{t}\,b\bar{b}\,\tau\bar{\tau}X\,, \qquad \text{and} \quad tb\,(2\bar{b})\,\tau\nu_\tau X\,, \tag{1.6}$$

[1] If the slepton $\tilde{\tau}$ is the LSP, and light, as advocated for example in [18], $\tilde{\chi}_1^0$, $\tilde{\chi}_1^-$, and possibly also t, are off-shell in eqs. (1.2) and (1.3).

[2] In R_p-violating models, the $\tilde{\tau}$-sleptons are not necessarily lighter than the charged Higgs bosons, and they may therefore give similar contributions to the above cross section, when their couplings to third-generation quarks are of the same order of magnitude of the b- and t-quark Yukawa couplings. Moreover, their masses are not constrained by low-energy processes such as $b \to s\gamma$ [19], in sharp contrast to the situation for the charged-Higgs boson mass in R_p-conserving models [21]. See also discussion in the next section.

cannot be confused with those originated by a charged Higgs boson, if more than $2b$'s can be tagged. In particular, the first final state in eq. (1.6) is quite distinctive of R_p-violating models since it gives rise to a lepton-number violating final state, with two like-sign τ's. This is an identifying signal of singly-produced τ-sleptons, when the W^+ bosons obtained from the t-quarks decay hadronically.

Two like-sign τ's will be also produced by heavy lightest neutralinos pair-produced at an e^+e^- collider,[3] when λ'_{333} is much larger than all other R_p-violating couplings. Under this assumption, in fact, the final states from a pair of $\tilde{\chi}_1^0$'s are:

$$2(b\bar{b})\,(2\nu_\tau)\,, \tag{1.7}$$

and

$$(2t)\,(2\bar{b})\,(2\tau)\,, \qquad (2\bar{t})\,(2b)\,(2\bar{\tau})\,, \qquad t\bar{t}\,b\bar{b}\,\tau\bar{\tau}\,, \qquad tb\,(2\bar{b})\,\tau\nu_\tau\,, \qquad \bar{t}\bar{b}\,(2b)\bar{\tau}\nu_\tau\,, \tag{1.8}$$

if the lightest neutralino is heavier than the t-quark.

It is clear, then, that the 3-body decays of both, $\tilde{\chi}_1^0$ and $\tilde{\chi}_1^-$, induced by large R_p-violating terms are worth a detailed study. For masses of $\tilde{\chi}_1^0$ and $\tilde{\chi}_1^-$ of interest to LHC and/or Linear Collider studies, decays of $\tilde{\chi}_1^0$ and $\tilde{\chi}_1^-$ into third generation fermions are likely to be allowed kinematically. Further, we expect the third generation sfermions, which are exchanged as virtual particles in these processes, to be lighter than the others. Effects of the mass of the fermions produced in 3-body decays of the charginos and heavier neutralinos were considered in the MSSM [23] and in mSUGRA [24], in the context of R_p-conserving models. These show that 3-body decays into final states containing third generation fermions, such as $b\bar{b}, \tau^+\tau^-$ $(\tau\nu)$, are dominant. In certain regions of the parameter space the contributions due to the Higgs boson exchange are large even for moderate values of $\tan\beta$. It is interesting to find out if these predictions remain valid in the context of R_p-violating models. Moreover, it is also important to know whether the decays of $\tilde{\chi}_1^0$ and $\tilde{\chi}_1^-$ into the heaviest third generation quark, the t-quark, have a competitive rate, in spite of the large kinematical suppression.

The 3-body decays of the lightest neutralino in the presence of R_p-violating couplings have been studied earlier and expressions for the differential decay widths exist in the literature, see refs. [25]–[29]. Correct matrix elements squared, however, were obtained first in ref. [27], up to a typographical error in the sign of the width of the sfermions exchanged as virtual particles in these decays. This was corrected in ref. [29]. No phenomenological analysis is presented in [27], whereas the analysis in [28] concentrates on one of the signatures caused by these decays without addressing the issue of relative decay widths and branching ratios. The phenomenological analysis in [29] deals with the λ'' couplings, although formulae for all R_p-violating couplings are listed. The attitude in ref. [29] is opposite to that taken here: a solution to the problem of a too fast proton decay is obtained by allowing only the B-violating couplings instead than the L-violating ones.

[3] Like-sign dileptons were recognized since long as the typical signal for the decay of a pair of lightest neutralinos into light fermions in R_p-violating models. See for example ref. [16].

The aim of this paper is to investigate the 3-body decays of the lightest neutralino into mainly third generation fermions, including the t-quark, induced by the R_p-violating couplings λ' and λ. This is also a preliminary and nontrivial step towards a complete study of the 3-body decays of the lightest chargino. If kinematically possible, this may decay into the lightest neutralino through gauge interactions, i.e. $\tilde{\chi}_1^- \to \tilde{\chi}_1^0 \bar{f}_u f_d$ and $\tilde{\chi}_1^- \to \tilde{\chi}_1^0 \nu_l l^-$. Such a study is left for further work [30].

This paper is organized as follows. After a discussion in section 2 of the experimental constraints on λ' and λ couplings, we list in section 3 all the relevant interaction terms. In section 4 are given the complete analytical formulae for the 3-body decays of $\tilde{\chi}_1^0$, whether it is the LSP or not, in the case of a dominant λ'_{333} coupling. No assumption is made on the actual composition of $\tilde{\chi}_1^0$ (i.e. if pure Bino ($\tilde{B}$), or mixed Bino-Wino ($\tilde{B}$-$\tilde{W}$), or mixed gaugino and Higgsino ($\tilde{H}$)), and complex values of the left-right mixing parameters in the sfermion mass matrices are assumed. These formulae provide a further check on the calculation of the widths of the 3-body neutralino decays, which has proven to be rather nontrivial. We confirm the analytic results of ref. [29], when taking the limit of real left-right mixing masses in the sfermion mass matrices. Our formulae can be easily generalized to study decays of heavier neutralinos and to situations in which other R_p-violating couplings are present. In particular, we use them to discuss the decays induced by sizable values of λ_{233} and λ'_{233}. Moreover, the generality in the composition of $\tilde{\chi}_1^0$, allows model-independent analyses. In spite of the large number of parameters on which the numerical results depend, we can capture some essential features of the different decay widths by choosing different values of slepton/squark masses, for different gaugino/higgsino contents of the neutralino, and by choosing small and large left-right mixing terms in the sfermion sector. We present and discuss these features in section 6. We end by summarizing our results in section 7. Appendices A and B give our conventions for the neutralino and sfermion mass matrices; appendix C gives the expressions for the R_p-violating 2-body decay widths of squarks and sleptons when λ'_{333} and λ_{233} are nonvanishing, whereas appendix D gives the expressions of terms required to evaluate the matrix element squared for the decay of the lightest neutralino.

2. Limits on R_p-violating couplings

Hints on the size of the R_p-violating couplings in eq. (1.1) may be obtained from various experiments. Two types of experimental constraints are possible. There are "direct" constraints due to effects of these couplings on sparticle production and on the decays of particle/sparticles at colliders. "Indirect" constraints [31], instead, are coming from measurements of SM observables to which supersymmetric particles contribute as virtual particles at the quantum level, or even at the tree level, as in scattering processes. Such observables include, at present, the EDM of fermions [32], the anomalous magnetic moment of the muon [33], neutrino masses, the decay widths of the Z- [34, 35, 36] and of the W-boson [36] the strength of four-fermion interactions, with the subsequent production of lepton pairs at hadron [37, 38] and lepton colliders [39], rare processes such as $\mu \to e\gamma$ [40] and $b \to s\gamma$ [19], the $e - \mu - \tau$ universality [36, 41, 42], etc.

Among these, a particular role is played by neutrino masses, which, being very small, restrict quite strongly the size of the dimensionful couplings κ_i in eq. (1.1) and of the vacuum expectation values (vev's) of the neutral scalar components of the fields L_i, v_i [10]. Strictly speaking, neutrino physics constrains only the parameter $(k_i \cdot v_i)/(k_i^2+\mu^2)^{1/2}(v_i^2+v_d^2)^{1/2}$. If not very small, the vectors $\{v_i\}$ and $\{k_i\}$ must be quite accurately aligned. In the following, we assume, for simplicity, that both vectors have very small magnitudes. (Decays of the lightest neutralino induced by bilinear couplings were studied in [43].)

Trilinear couplings such as λ_{i33} and λ'_{i33}, with $i = 1, 2$, may also be severely bound by the value of the Majorana mass terms for ν_e and ν_μ that they induce at the loop level [16]. The specific value of these bounds depends on the mass of the third-generation sfermion doublet and singlet as well as on the mass term that mixes them. Most of all, it depends also on the particular neutrino spectrum assumed, for example whether ν_e and ν_μ are in the sub-eV region, as the neutrino oscillation experiments may imply. If the heaviest neutrino is also $\lesssim 1\,\mathrm{eV}$, then, a constraint may be deduced also for λ'_{333} (but not for λ'_{323} or λ'_{332}), which induces a quantum contribution to the mass of ν_τ, in addition to that generated at the tree-level by the dimensionful couplings κ_i. However, neutrino masses may be efficiently suppressed not only by small λ and λ' couplings, but also by very tiny left-right mixing terms in the slepton and squark sectors, and/or large values of the slepton and squark masses. Notice that very small left-right sfermion mass terms, as well as large sfermion mass eigenvalues, do change quantitatively the predictions for the widths of neutralino and chargino decays, but do not spoil the possibility of visible signals for these decays. Because of all this, and because the neutrino spectrum is not yet known, we prefer to leave the constraints from neutrino masses aside and consider them as complementary to those that can be obtained from collider searches, in a direct way. Nevertheless, it should be remembered that, irrespectively of constraints from neutrino masses, at least one combination of the L-violating couplings must be very small if leptogenesis has to occur [5, 4].

In general, the "direct" and "indirect" observables listed above constrain most of the R_p-violating couplings, in particular those involving first and second generation indices, to small values, smaller than the gauge couplings [31, 42, 44]. These bounds depend on the values of other supersymmetric parameters and, at the 2σ level, they are severe for rather small soft scalar masses, i.e. $\sim 100\,\mathrm{GeV}$ [42]. The situation is somewhat different for couplings involving third generation particles, which are left practically unrestricted by the low energy processes mentioned above. Constraints on λ'_{3jk} (and λ''_{3jk}) come, for example, indirectly, from radiative corrections to $Z \to b\bar{b}$, i.e. $\lesssim 1$ [34, 35, 44], and to $Z \to l^+l^-$, i.e $\lesssim 0.42$ [36]. Both bounds are derived for sfermion (slepton/squark) masses $\sim 100\,\mathrm{GeV}$. For heavier sfermions, the constraint on λ'_{3jk} (and λ''_{3jk}) are somewhat weaker. Rare processes like $b \to s\gamma$ involve far too many parameters to give clear-cut constraints on any of them [19].

In the future, constraints to the R_p-violating couplings involving third-generation indices may be obtained indirectly from other interesting quantities, such as the t-channel contribution to the $t\bar{t}$ production, t-quark decays into SM fermions [45, 46, 47], and the polarization of the $t/\bar{t}$ quark, which may all get contributions for nonvanishing R_p-violating couplings. In particular, studies of the $t/\bar{t}$ quark polarization at the Tevatron can probe λ'_{31i} (and λ''_{311}), to somewhat lower values [46, 48] than those tested by the Z width. Finally, the effect of a $\tilde{t}$-squark exchange in the t channel on the Drell-Yan $\mu^+\mu^-$ pair production, not very significant at the Tevatron energy [37], may probe λ'_{231}, at the LHC energies, to values < 0.2 even for $\tilde{t}$-squark masses as large as 800 GeV [38].

Direct constraints on R_p-violating couplings that are comparable or larger than the gauge couplings, i.e. $\gtrsim 10^{-1}$–10^{-2}, are obtained from searches of: *i)* unusual decays of the t-quark [47] into still allowed light superpartners, *ii)* resonances which violate L or B, *iii)* signals due to the decays of superpartners, among them the lightest neutralino. The last, in particular, may be the only signal to be probed, for couplings smaller than 10^{-2}. At a hadron collider, λ' couplings can be probed by searching for resonant slepton production [49]; λ' and λ couplings when both production and decay of such resonances take place via R_p violating couplings [50]. At a lepton collider the λ couplings [51] are those to be probed.

Searches at the Tevatron rule out a first generation squark up to 100 GeV, almost irrespective of the size of the λ' couplings [52]. These direct searches put limits on the λ and λ' couplings involving mainly the first two generations, as both the colliding particles as well as the final state particles are in general from the first two generations [53].

At the ep collider HERA a squark resonance can be produced via interactions with λ' couplings. The nonobservation of such a resonance in the e^+p data has ruled out a first generation squark up to 260 GeV for a coupling (to a quark and a lepton) $\lambda' \sim \sqrt{4\pi\alpha_{em}}$ [54]; e^-p data should complement this analysis [55] and strengthen this bound. The analysis for the HERA data has been done in the framework of an unconstrained MSSM, with the additional assumption of a common gaugino mass at a high scale, as well as in the mSUGRA framework.

Searches of s-channel resonance formation at LEP have yielded constraints on the couplings λ_{121}, λ_{131} and λ_{232} in a model independent way [56, 57, 58]. (These analyses include also indirect effects of the t-channel sfermion exchange on the four fermion scattering.) These constraints are considerable for low values of sneutrino masses, viz. $100 < m_{\tilde{\nu}} < 200$ GeV, but rise close to the values of gauge couplings for higher sneutrino masses [56, 57, 58]. In general, for sfermion masses $\gtrsim 100$ GeV, LEP experiments do not give substantial limits on any of the R-parity violating couplings involving more than one third generation index. LEP has also probed R-parity violating couplings, by searching for decays of sparticles, including the decay of the lightest neutralino into light fermions (see for example ref. [16]). These searches have essentially yielded only limits on sparticle masses [59].

Some of these analyses, have been carried out in the framework of mSUGRA, without keeping into account the fact that the effect of relatively large R_p-violating couplings cannot be ignored in the determination of the mass spectra.

In view of all the above considerations, in the following analysis we use the values 0.5, 1 for λ'_{333} and λ'_{233}, whereas we choose a somewhat lower value, i.e. 0.2, for λ_{233}. For one particular direction of parameter space, we show how the decay widths change when the value of the relevant R_p-violating couplings are lowered. In the regions where the sfermions mediating these decays are nonresonant, the decay widths are simply scaled down by the

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overall factors $|\lambda'_{i33}|^2$ or $|\lambda_{233}|^2$. They remain, however, marginally affected in the sfermion resonant regions. This behaviour can be extrapolated to the other directions of parameter space studied in this paper. The present analysis therefore applies also to the case in which the relevant couplings are small, provided they dominate over other R_p-violating couplings.

3. Interaction terms relevant for $\tilde{\chi}_1^0$ decays

The interactions relevant for an analysis of neutralino decays are: *i)* the R_p-violating interactions due to the first two terms of the superpotential in eq. (1.1), and *ii)* the neutralino-fermion-sfermion couplings due to gauge and Yukawa interactions.

In the following, we give explicitly the fermion-fermion-sfermion interaction terms derived from $\lambda_{ijk} L_i L_j E^c_k$ and $\lambda'_{ijk} L_i Q_j D^c_k$. Assuming the standard contraction of SU(2) indices:

$$\lambda'_{ijk} L_i Q_j D^c_k = \lambda'_{ijk} \epsilon_{ab} \left(L_i\right)_a \left(Q_j\right)_b D^c_k \,, \tag{3.1}$$

where $\epsilon_{12} = -\epsilon_{21} = 1$, the R_p-violating interaction terms with couplings λ'_{ijk}, are:

$$\begin{aligned} \mathcal{L} = &+\lambda'_{ijk}\, \tilde{u}_{jL} \left(\overline{d_{kR}}\, l_{iL}\right) + \lambda'_{ijk}\, \tilde{l}_{iL} \left(\overline{d_{kR}}\, u_{jL}\right) + \lambda'_{ijk}\, \tilde{d}^*_{kR} \left(\overline{l^c{}_{iR}}\, u_{jL}\right) + \\ &+ \lambda'^*_{ijk}\, \tilde{u}^*_{jL} \left(\overline{l_{iL}}\, d_{kR}\right) + \lambda'^*_{ijk}\, \tilde{l}^*_{iL} \left(\overline{u_{jL}}\, d_{kR}\right) + \lambda'^*_{ijk}\, \tilde{d}_{kR} \left(\overline{u_{jL}}\, l^c{}_{iR}\right) , \end{aligned} \tag{3.2}$$

when the charged lepton l is involved, and

$$\begin{aligned} \mathcal{L} = &-\lambda'_{ijk}\, \tilde{d}_{jL} \left(\overline{d_{kR}}\, \nu_{iL}\right) - \lambda'_{ijk}\, \tilde{\nu}_{iL} \left(\overline{d_{kR}}\, d_{jL}\right) - \lambda'_{ijk}\, \tilde{d}^*_{kR} \left(\overline{\nu^c{}_{iR}}\, d_{jL}\right) - \\ &- \lambda'^*_{ijk}\, \tilde{d}^*_{jL} \left(\overline{\nu_{iL}}\, d_{kR}\right) - \lambda'^*_{ijk}\, \tilde{\nu}^*_{iL} \left(\overline{d_{jL}}\, d_{kR}\right) - \lambda'^*_{ijk}\, \tilde{d}_{kR} \left(\overline{d_{jL}}\, \nu^c{}_{iR}\right) , \end{aligned} \tag{3.3}$$

when the lepton interacting is neutral. The upperscript ^{c} indicates charge conjugation ($\psi^c = C\bar{\psi}^T$).

The interaction terms with couplings of type λ_{ijk} can be obtained in a similar way:

$$\begin{aligned} \mathcal{L} = &+\lambda_{ijk}\, \tilde{\nu}_{jL} \left(\overline{l_{kR}}\, l_{iL}\right) + \lambda_{ijk}\, \tilde{l}_{iL} \left(\overline{l_{kR}}\, \nu_{jL}\right) + \lambda_{ijk}\, \tilde{l}^*_{kR} \left(\overline{l^c{}_{iR}}\, \nu_{jL}\right) + \\ &+ \lambda^*_{ijk}\, \tilde{\nu}^*_{jL} \left(\overline{l_{iL}}\, l_{kR}\right) + \lambda^*_{ijk}\, \tilde{l}^*_{iL} \left(\overline{\nu_{jL}}\, l_{kR}\right) + \lambda^*_{ijk}\, \tilde{l}_{kR} \left(\overline{\nu_{jL}}\, l^c{}_{iR}\right) , \qquad (i > j) \end{aligned} \tag{3.4}$$

where the relation $\lambda_{ijk} = -\lambda_{jik}$ was used. Notice that there are only 9 independent λ couplings versus the 27 of type λ'.

The neutralino eigenstates are obtained diagonalizing the matrix shown in appendix A. Strictly speaking, in R_p-violating models the neutralino mass matrix is a 7×7 matrix. In the approximation of very small bilinear couplings k_i and *vev*'s v_i of the neutral components of the fields $\tilde{L}_i$, this reduces to the conventional 4×4 matrix typical of R_p-conserving models. Thus, the neutralino-sfermion-fermion interaction terms are:

$$\begin{aligned} L = -\sqrt{2} g \sum_i \Big\{ & L_i^{f_L} \left(\overline{(\tilde{\chi}^0_{iR})} \cdot f_L\right) \tilde{f}^*_L - R_i^{f_L} \left(\overline{(\tilde{\chi}^0_{iL})} \cdot f_R\right) \tilde{f}^*_L + \\ &+ L_i^{f_R} \left(\overline{(\tilde{\chi}^0_{iR})} \cdot f_L\right) \tilde{f}^*_R - R_i^{f_R} \left(\overline{(\tilde{\chi}^0_{iL})} \cdot f_R\right) \tilde{f}^*_R \Big\} + \text{h.c.} \end{aligned} \tag{3.5}$$

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where i is now a neutralino index ($i = 1, 2, 3, 4$) and g is the SU(2) gauge coupling. By using the definition of hypercharge $Q = (Y + T_3)$, the coefficients $L_i^{f_L}$ and $R_i^{f_R}$ can be cast in the compact form:

$$L_i^{f_L} = \eta_i^* \left(T_{3\,f_L} O_{i2} + Y_{f_L} O_{i1} \tan\theta_W\right), \qquad R_i^{f_R} = \eta_i\, Y_{f_R} O_{i1} \tan\theta_W\,, \tag{3.6}$$

whereas the coefficients $L_i^{f_R}$ and $R_i^{f_L}$ are different for fermions of up (f_u) and down type (f_d):

$$\begin{aligned} L_i^{(f_u)_R} &= +\frac{1}{2}\,\eta_i^* \,\frac{m_{f_u}}{M_W \sin\beta}\, O_{i4}\,, \qquad & R_i^{(f_u)_L} &= -\frac{1}{2}\,\eta_i \,\frac{m_{f_u}}{M_W \sin\beta}\, O_{i4}\,, \\ L_i^{(f_d)_R} &= +\frac{1}{2}\,\eta_i^* \,\frac{m_{f_d}}{M_W \cos\beta}\, O_{i3}\,, \qquad & R_i^{(f_d)_L} &= -\frac{1}{2}\,\eta_i \,\frac{m_{f_d}}{M_W \cos\beta}\, O_{i3}\,. \end{aligned} \tag{3.7}$$

Notice that all neutrino couplings are vanishing except for $L_i^{\nu_L}$. The matrix O in the previous equations is the neutralino diagonalization matrix. The phase η is due to the fact that the diagonalization of the neutralino mass matrix may yield complex eigenvalues, see appendix A. The neutralino eigenstates may be rotated in such a way to obtain real and positive eigenvalues, but phase factors η appear in this case in the interaction terms. For real values of the parameters $M_{\tilde{B}}$, $M_{\tilde{W}}$, and μ, the mass eigenvalues are real, but with signs that can be positive or negative. A rotation of the eigenstates that reduces all eigenvalues to be positive, brings in factors η_i, which are now i or 1.

In the approximation of pure B-ino for the lightest neutralino, the coefficients $L_i^{f_R}$ and $R_i^{f_L}$ vanish identically, those in eq. (3.6) become $L_i^{f_L} = \eta_i^* Y_{f_L} \tan\theta_W$, $R_i^{f_R} = \eta_i Y_{f_R} \tan\theta_W$. Equation (3.5) has, then, the simple form:

$$\mathcal{L} = -\sqrt{2}\, g_Y \left\{ \eta_1^* Y_{f_L} \left(\overline{\tilde{\chi}^0_{1R}}\, f_L\right) \tilde{f}_L^* - \eta_1 Y_{f_R} \left(\overline{\tilde{\chi}^0_{1L}}\, f_R\right) \tilde{f}_R^* \right\} + \text{h.c} \tag{3.8}$$

where g_Y is expressed in terms of g as $g_Y = g \tan\theta_W$. In this approximation, unless the gaugino mass M_1 is chosen to be negative, the phase factor η can be altogether dropped.

In the approximation of pure B-ino, the lightest neutralino decays via gauge interactions into a fermion-sfermion pair, $\bar{f}\tilde{f}$. The sfermion $\bar{f}\tilde{f}$ is off-shell when $\tilde{\chi}_1^0$ is the LSP. This sfermion decays predominantly through R_p-violating interactions, if the corresponding couplings are large.

In the following, we start considering the case of a dominant λ'_{333} coupling and discuss later the possibility of having simultaneously two relatively large couplings, for example λ'_{333} and λ'_{233}, or λ'_{333} and λ_{233}.

4. Dominant λ'_{333} coupling

If λ'_{333} is the dominant R_p-violating coupling, the only relevant decays of the lightest neutralino into a fermion-sfermion pair are:

$$\bar{b}\tilde{b}\,, \qquad \bar{\nu}_\tau \tilde{\nu}_\tau\,, \qquad \bar{\tau}\tilde{\tau}\,, \tag{4.1}$$

and their CP-conjugated ones, for a lightest neutralino lighter than the top-quark, or into:

$$\bar{b}\tilde{b}\,, \qquad \bar{\nu}_\tau\tilde{\nu}_\tau\,, \qquad \bar{\tau}\tilde{\tau}\,, \qquad \bar{t}\tilde{t}\,, \tag{4.2}$$

again with their CP-conjugated ones, for a lightest neutralino heavier than the top-quark. In the first case, the two decay modes $\bar{b}\tilde{b}$ and $\bar{\nu}_\tau\tilde{\nu}_\tau$ give rise to the final state $\bar{b}b\nu_\tau$. Since neutrinos can be safely considered as massless at the energies at which neutralino decays will be studied, in the following, the two possible final states $\bar{b}b\nu_\tau$ and $\bar{b}b\bar{\nu}_\tau$ are identified. The third mode in eq. (4.1), $\bar{\tau}\tilde{\tau}$, gives rise to an off-shell top-quark through the decay $\tilde{\tau} \to b\bar{t}$ (see eq. (1.3)) and therefore a multi-body final state. If $\tilde{\chi}_1^0$ is heavier than the top-quark, the modes $\bar{\tau}\tilde{\tau}$ and $\bar{t}\tilde{t}$ give rise to the final state $\bar{t}b\bar{\tau}$, to the conjugate of which, also the mode $\bar{b}\tilde{b}$ may contribute. Therefore, in the above approximation of only one R_p-violating coupling, the 3-body final states arising from the decay of $\tilde{\chi}_1^0$ are an almost massless one, $\bar{b}b\nu_\tau$ and a massive one, $\bar{b}t\tau$. In the following, we give amplitudes and widths for both decays.

4.1 Massless final state — $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$

The decay $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$ is mediated by the exchange of an off-shell/on-shell $\tilde{b}$-squark, as in the diagrams of figure 1, and an off-shell/on-shell neutral slepton $\tilde{\nu}_\tau$, as shown by the diagram in figure 2. In both figures, only the diagrams giving rise to matrix elements proportional to λ'^*_{333} are shown. To the same decay contribute also the "crossed" diagrams, i.e. those with the two external b-quarks interchanged and an ingoing neutrino line. The corresponding matrix elements are proportional to λ'_{333}.

It is convenient to split the matrix elements of all contributions to both decays $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$ and $\tilde{\chi}_1^0 \to t\bar{b}\tau$ as follows:

$$M = \sum_{s,t} a^t_{S,s}\, Q^t_{S,s}\,. \tag{4.3}$$

The Q terms are the results of the contractions between initial and final states of all possible operators obtained from the lagrangian interaction terms, taken in absolute value and without numerical coefficients. Possible relative minus signs, which may be obtained in these contractions according to the Wick's theorem, are included in the corresponding coefficients.

The Q terms contributing to the decay $\tilde{\chi}_1^0(q) \to b(p_1)\nu_\tau(p_2)\bar{b}(p_3)$,

$$\begin{aligned}
Q^x_{S,RR} &= \left(\overline{u}_{\nu_\tau}(p_2)P_R\, v_b(p_3)\right)\left(\overline{u}_b(p_1)P_R\, u_{\tilde{\chi}_1^0}(q)\right)\\
Q^x_{S,RL} &= \left(\overline{u}_{\nu_\tau}(p_2)P_R\, v_b(p_3)\right)\left(\overline{u}_b(p_1)P_L\, u_{\tilde{\chi}_1^0}(q)\right)\\
Q^z_{S,RR} &= \left(\overline{u}_{\nu_\tau}(p_2)P_R\, v_b(p_1)\right)\left(\overline{u}_b(p_3)P_R\, u_{\tilde{\chi}_1^0}(q)\right)\\
Q^z_{S,RL} &= \left(\overline{u}_{\nu_\tau}(p_2)P_R\, v_b(p_1)\right)\left(\overline{u}_b(p_3)P_L\, u_{\tilde{\chi}_1^0}(q)\right)\\
Q^y_{S,RR} &= \left(\overline{u}_b(p_1)P_R\, v_b(p_3)\right)\left(\overline{u}_{\nu_\tau}(p_2)P_R\, u_{\tilde{\chi}_1^0}(q)\right).
\end{aligned} \tag{4.4}$$

have as corresponding coefficients:

$$\begin{aligned}
a^x_{S,RR} &= -\sqrt{2}\,g\lambda'^*_{333}\left((L_1^{b_L})^* D^{\tilde{b}}_{LL}(x) + (L_1^{b_R})^* D^{\tilde{b}}_{RL}(x)\right) \;\to -\sqrt{2}\,g_Y\,\lambda'^*_{333}\,\eta_1^*\,Y_{b_L}\,D^{\tilde{b}}_{LL}(x)\\
a^x_{S,RL} &= +\sqrt{2}\,g\lambda'^*_{333}\left((R_1^{b_R})^* D^{\tilde{b}}_{RL}(x) + (R_1^{b_L})^* D^{\tilde{b}}_{LL}(x)\right) \;\to +\sqrt{2}\,g_Y\,\lambda'^*_{333}\,\eta_1^*\,Y_{b_R}\,D^{\tilde{b}}_{RL}(x)\\
a^z_{S,RR} &= -\sqrt{2}\,g\lambda'^*_{333}\left(R_1^{b_R}(D^{\tilde{b}}_{RR}(z))^* + R_1^{b_L}(D^{\tilde{b}}_{LR}(z))^*\right) \;\to -\sqrt{2}\,g_Y\,\lambda'^*_{333}\,\eta_1\,Y_{b_R}\,(D^{\tilde{b}}_{RR}(z))^*\\
a^z_{S,RL} &= +\sqrt{2}\,g\lambda'^*_{333}\left(L_1^{b_L}(D^{\tilde{b}}_{LR}(z))^* + L_1^{b_R}(D^{\tilde{b}}_{RR}(z))^*\right) \;\to +\sqrt{2}\,g_Y\,\lambda'^*_{333}\,\eta_1\,Y_{b_L}\,(D^{\tilde{b}}_{LR}(z))^*\\
a^y_{S,RR} &= +\sqrt{2}\,g\lambda'^*_{333}\left((L_1^{\nu_\tau L})^* D^{\tilde{\nu}_\tau}_{LL}(y)\right) \qquad\qquad\qquad\quad \to +\sqrt{2}\,g_Y\,\lambda'^*_{333}\,\eta_1^*\,Y_{\nu_\tau L}\,D^{\tilde{\nu}_\tau}_{LL}(y)\,,
\end{aligned} \tag{4.5}$$

where the limiting values in the approximation of lightest neutralino as pure B-ino were given. In both definitions, the indices $\{S, RR\}$ and $\{S, RL\}$ are reminders of the fact that these matrix elements are obtained from a product of two scalar currents, each of them with a chirality projector specified by R/L. The upper index x, y, and z indicates the fraction of the initial momentum flowing through the scalar propagator mediating these diagrams. They are respectively:

$$x = \frac{(p_1 - q)^2}{q^2}\,, \qquad y = \frac{(p_2 - q)^2}{q^2}\,, \qquad z = \frac{(p_3 - q)^2}{q^2}\,. \tag{4.6}$$

In the decomposition of eq. (4.3), therefore, s runs over RR and RL, in the terms induced by the diagrams in figures 1 and 2, and over LL and LR in those induced by the "crossed" diagrams, and t may be x, y, and z. Finally, the coefficients $D^{\tilde{f}}_\sigma(t)$ collect the contributions from the scalar propagators of the two states $\tilde{f}_{1,2}$ weighted by the factors projecting states of definite chirality into mass eigenstates. (No intergenerational mixing among sfermions is assumed here.) The explicit expression of the two mass eigenvalues $m^2_{\tilde{f}_{1,2}}$ and projection factors $\sin\theta_f$ and $\cos\theta_f$ in terms of the parameters in the sfermion mass matrix are given in appendix B. The index σ, which runs over LL, LR, RL, RR, indicates the chirality of the ingoing and outgoing scalar fields at each vertex of the diagrams in figures 1 and 2. For definiteness, for $t = x, y, z$, it is:

$$\begin{aligned}
D^{\tilde{f}}_{LL}(t) &= \sin^2\theta_f D_{\tilde{f}_1}(t) + \cos^2\theta_f D_{\tilde{f}_2}(t)\\
D^{\tilde{f}}_{LR}(t) &= \sin\theta_f\cos\theta_f e^{-i\phi_f}\left(D_{\tilde{f}_1}(t) - D_{\tilde{f}_2}(t)\right)\\
D^{\tilde{f}}_{RL}(t) &= \sin\theta_f\cos\theta_f e^{+i\phi_f}\left(D_{\tilde{f}_1}(t) - D_{\tilde{f}_2}(t)\right)\\
D^{\tilde{f}}_{RR}(t) &= \cos^2\theta_f D_{\tilde{f}_1}(t) + \sin^2\theta_f D_{\tilde{f}_2}(t)\,,
\end{aligned} \tag{4.7}$$

with

$$D_{\tilde{f}_{1,2}}(t) = \frac{1}{m^2_{\tilde{\chi}^0_1}}\left(t - \frac{m^2_{\tilde{f}_{1,2}}}{m^2_{\tilde{\chi}^0_1}}\right)^{-1}\,, \tag{4.8}$$

for off-shell sfermions $\tilde{f}_{1,2}$, or

$$D_{\tilde{f}_i}(t) = \frac{1}{m^2_{\tilde{\chi}^0_1}}\left(t - \frac{m^2_{\tilde{f}_i}}{m^2_{\tilde{\chi}^0_1}} + i\frac{\Gamma_{\tilde{f}_i} m_{\tilde{f}_i}}{m^2_{\tilde{\chi}^0_1}}\right)^{-1}\,, \tag{4.9}$$

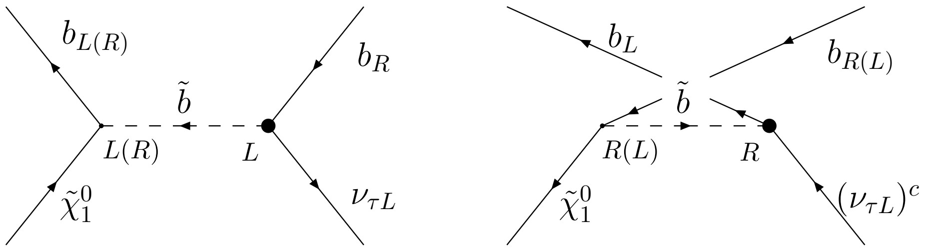

Figure 1: Diagrams contributing to the decay $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$ through the exchange of a virtual $\tilde{b}$-squark. The thick vertex indicates the R_p-violating coupling λ'^*_{333}. The labels L and R at each vertex indicate the chirality of the ingoing and/or outgoing scalar fields. To the same decay contribute also the crossed diagrams, i.e. those with the two external bottom-quarks interchanged and an ingoing neutrino line, and with the thick vertex indicating the R_p-violating coupling λ'_{333}.

in the case in which the sfermion $\tilde{f}_i$ is on shell. The width of possible sfermions lighter than the lightest neutralino $\tilde{\chi}_1^0$ will be discussed in appendix C. Finally, the angle ϕ_f is the argument of the left-right mixing term in the mass matrix of the sfermion $\tilde{f}$, see appendix B.

Similarly, the Q terms for the crossed diagrams are:

$$
\begin{aligned}
Q^x_{S,LL} &\equiv (\overline{u}_{\nu_\tau}(p_2) P_L v_b(p_3)) \left(\overline{u}_b(p_1) P_L\, u_{\tilde{\chi}_1^0}(q)\right) \\
Q^x_{S,LR} &\equiv (\overline{u}_{\nu_\tau}(p_2) P_L v_b(p_3)) \left(\overline{u}_b(p_1) P_R\, u_{\tilde{\chi}_1^0}(q)\right) \\
Q^z_{S,LL} &\equiv (\overline{u}_{\nu_\tau}(p_2) P_L v_b(p_1)) \left(\overline{u}_b(p_3) P_L\, u_{\tilde{\chi}_1^0}(q)\right) \\
Q^z_{S,LR} &\equiv (\overline{u}_{\nu_\tau}(p_2) P_L v_b(p_1)) \left(\overline{u}_b(p_3) P_R\, u_{\tilde{\chi}_1^0}(q)\right) \\
Q^y_{S,LL} &\equiv (\overline{u}_b(p_1) P_L v_b(p_3)) \left(\overline{u}_{\nu_\tau}(p_2) P_L\, u_{\tilde{\chi}_1^0}(q)\right) .
\end{aligned}
\tag{4.10}
$$

They differ from those in eq. (4.4) by an exchange $R \leftrightarrow L$. The related coefficients, with their limiting values in the approximation of lightest neutralino as pure B-ino, are:

$$
\begin{aligned}
a^x_{S,LL} &= +\sqrt{2}\, g \lambda'_{333} \left((R_1^{b_R})^* D^{\tilde{b}}_{RR}(x) + (R_1^{b_L})^* D^{\tilde{b}}_{LR}(x)\right) &&\to +\sqrt{2}\, g_Y \lambda'_{333}\, \eta_1^*\, Y_{b_R} \left(D^{\tilde{b}}_{RR}(x)\right) \\
a^x_{S,LR} &= -\sqrt{2}\, g \lambda'_{333} \left((L_1^{b_L})^* D^{\tilde{b}}_{LR}(x) + (L_1^{b_R})^* D^{\tilde{b}}_{RR}(x)\right) &&\to -\sqrt{2}\, g_Y \lambda'_{333}\, \eta_1^*\, Y_{b_L} \left(D^{\tilde{b}}_{LR}(x)\right) \\
a^z_{S,LL} &= +\sqrt{2}\, g \lambda'_{333} \left((L_1^{b_L})(D^{\tilde{b}}_{LL}(z))^* + (L_1^{b_R})(D^{\tilde{b}}_{RL}(z))^*\right) &&\to +\sqrt{2}\, g_Y \lambda'_{333}\, \eta_1\, Y_{b_L} \left(D^{\tilde{b}}_{LL}(z)\right)^* \\
a^z_{S,LR} &= -\sqrt{2}\, g \lambda'_{333} \left((R_1^{b_R})(D^{\tilde{b}}_{RL}(z))^* + (R_1^{b_L})(D^{\tilde{b}}_{LL}(z))^*\right) &&\to -\sqrt{2}\, g_Y \lambda'_{333}\, \eta_1\, Y_{b_R} \left(D^{\tilde{b}}_{RL}(z)\right)^* \\
a^y_{S,LL} &= +\sqrt{2}\, g \lambda'_{333} \left((L_1^{\nu_{\tau L}})^* D^{\tilde{\nu}_\tau}_{LL}(y)\right) &&\to +\sqrt{2}\, g_Y \lambda'_{333}\, \eta_1^*, Y_{\nu_{\tau L}} \left(D^{\tilde{\nu}_\tau}_{LL}(y)\right) .
\end{aligned}
\tag{4.11}
$$

The results in eqs. (4.5) and (4.11) coincide with those in ref. [29], up to charge conjugations.

The width for the decay mode $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$ is finally obtained after integration of the differential one:

$$
\frac{d\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)}{dx\,dy} = \frac{3\, m_{\tilde{\chi}_1^0}}{512\pi^3}\, |M|^2_{z=1-x-y} \,, \tag{4.12}
$$

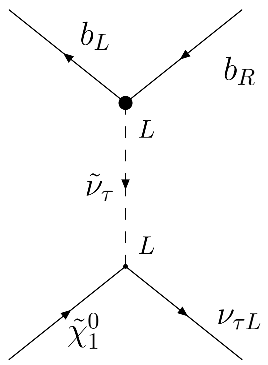

Figure 2: Diagram contributing to the decay $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$, in particular to $\tilde{\chi}_1^0 \to b_L\overline{(b_R)}\,\nu_{\tau L}$, through the exchange of a virtual neutral slepton $\tilde{\nu}_\tau$. The thick vertex indicates the R_p-violating coupling λ'^{*}_{333}. The label L at each vertex indicate the chirality of the ingoing/outgoing scalar field. A similar diagram, with the thick vertex indicating λ'_{333}, contributes to the decay $\tilde{\chi}_1^0 \to b_R\overline{(b_L)}\,\overline{\nu_{\tau L}}$.

given by the standard 3-body phase-space factor multiplied by the square of the matrix elements averaged over the neutralino spin and summed over spin and color of the final fermions. The sum over all spin configurations is included in $|M|^2$, whereas the average over the neutralino spin and the sum over color give an overall factor 3/2 included in the numerical coefficient in eq. (4.12). The square $|M|^2$ can be expressed in terms of the products $\beta^{t,t'}_{s,s'} \equiv Q^t_{S,s} Q^{t'\,\dagger}_{S,s'}$ given explicitly in appendix D. They are evaluated under the assumption that the particle with momentum p_1 has a nonnegligible mass m_1, with $m_1^2 = r m^2_{\tilde{\chi}_1^0}$. The expressions relevant to the case of the decay $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$ are, then, obtained by taking the limit $r \to 0$ in eqs. (D.3), (D.5)–(D.7), and (D.9). For comparisons with ref. [27], one should keep in mind that the formulae listed above combine the partial widths for the decays into $b\bar{b}\nu_\tau$ and into $b\bar{b}\bar{\nu}_\tau$, whereas only the partial width for $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$ is explicitly given in ref. [27]. The formal expression in eq. (4.12) is used also to obtain the partial width of all other decays discussed in this paper with the obvious modifications: *1)* r is nonvanishing in the expression for $|M|^2$ when decays into massive final states are considered; *2)* the overall factor 3 of color must be removed in the case of decays into purely leptonic final states.

The total width is obtained integrating eq. (4.12) over the two variables x and y, with integration bounds given by:

$$r \le y \le 1\,, \qquad 0 \le x \le (1-y)\frac{(y-r)}{y}\,, \tag{4.13}$$

for nonvanishing m_1. Again, the limit $r \to 0$ has to be taken in the case of the decay $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$.

4.2 Massive final state — $\tilde{\chi}_1^0 \to t\bar{b}\tau$

This decay is mediated by the exchange of off-shell $\tilde{t}$- and $\tilde{b}$-squarks, as shown in figure 3, and an off-shell $\tilde{\tau}$-slepton, as in figure 4. In these figures, only the diagrams giving rise to matrix elements proportional to λ'^{*}_{333} are shown. Notice that, since the virtual $\tilde{\tau}$ exchanged in the diagram of figure 4 has both, left- and right-chiralities, there are two independent

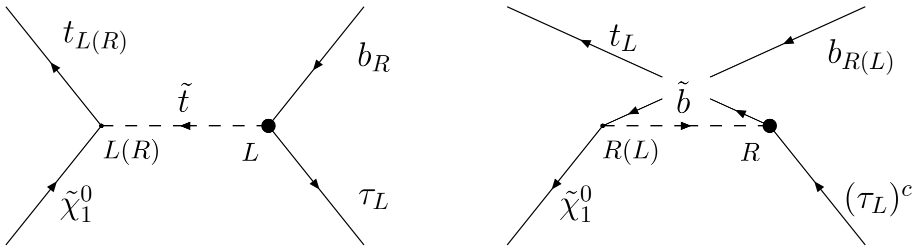

Figure 3: Diagrams contributing to the decay $\tilde{\chi}_1^0 \to t\bar{b}\tau$ through the exchange of a virtual $\tilde{t}$- and $\tilde{b}$-squark. The thick vertex indicates the R_p-violating coupling λ'^{*}_{333}.

contributions coming from this diagrams. Once again, we split the matrix elements for this decay as in eq. (4.3). The Q terms relative to the diagrams in figures 3 and 4 are now *six* instead of *five*:

$$\begin{aligned}
Q^x_{S,RR} &= (\overline{u}_\tau(p_2) P_R\, v_b(p_3)) \left(\overline{u}_t(p_1) P_R\, u_{\tilde{\chi}_1^0}(q)\right) \\
Q^x_{S,RL} &= (\overline{u}_\tau(p_2) P_R\, v_b(p_3)) \left(\overline{u}_t(p_1) P_L\, u_{\tilde{\chi}_1^0}(q)\right) \\
Q^z_{S,RR} &= (\overline{u}_\tau(p_2) P_R\, v_t(p_1)) \left(\overline{u}_b(p_3) P_R\, u_{\tilde{\chi}_1^0}(q)\right) \\
Q^z_{S,RL} &= (\overline{u}_\tau(p_2) P_R\, v_t(p_1)) \left(\overline{u}_b(p_3) P_L\, u_{\tilde{\chi}_1^0}(q)\right) \\
Q^y_{S,RR} &= (\overline{u}_t(p_1) P_R\, v_b(p_3)) \left(\overline{u}_\tau(p_2) P_R\, u_{\tilde{\chi}_1^0}(q)\right) \\
Q^y_{S,RL} &= (\overline{u}_t(p_1) P_R\, v_b(p_3)) \left(\overline{u}_\tau(p_2) P_L\, u_{\tilde{\chi}_1^0}(q)\right) .
\end{aligned} \tag{4.14}$$

For convenience, the same symbols employed in eq. (4.4) are used for these Q terms, in spite of the fact that they are built out of spinors of different particles. The squared products, summed over all possible spin configurations obtained for the two sets of Q's (in eqs. (4.4) and (4.14)) differ only for the value of the parameter $r = m_t^2/m_{\tilde{\chi}_1^0}$, which is nonvanishing in the case of the decay $\tilde{\chi}_1^0 \to t\bar{b}\tau$. The corresponding coefficients are now denoted by the symbols $b^t_{S,s}$ and are:

$$\begin{aligned}
b^x_{S,RR} &= +\sqrt{2}\, g\lambda'^{*}_{333} \left((L^{t_L})^* D^{\tilde{t}}_{LL}(x) + (L^{t_R})^* D^{\tilde{t}}_{RL}(x)\right) && \to +\sqrt{2}\, g_Y \lambda'^{*}_{333}\, \eta_1^*\, Y_{t_L}\, D^{\tilde{t}}_{LL}(x) \\
b^x_{S,RL} &= -\sqrt{2}\, g\lambda'^{*}_{333} \left((R^{t_R})^* D^{\tilde{t}}_{RL}(x) + (R^{t_L})^* D^{\tilde{t}}_{LL}(x)\right) && \to -\sqrt{2}\, g_Y \lambda'^{*}_{333}\, \eta_1^*\, Y_{t_R}\, D^{\tilde{t}}_{RL}(x) \\
b^z_{S,RR} &= +\sqrt{2}\, g\lambda'^{*}_{333} \left((R^{b_R})(D^{\tilde{b}}_{RR}(z))^* + (R^{b_L})(D^{\tilde{b}}_{LR}(z))^*\right) && \to +\sqrt{2}\, g_Y \lambda'^{*}_{333}\, \eta_1\, Y_{b_R} (D^{\tilde{b}}_{RR}(z))^* \\
b^z_{S,RL} &= -\sqrt{2}\, g\lambda'^{*}_{333} \left((L^{b_L})(D^{\tilde{b}}_{LR}(z))^* + (L^{b_R})(D^{\tilde{b}}_{RR}(z))^*\right) && \to -\sqrt{2}\, g_Y \lambda'^{*}_{333}\, \eta_1\, Y_{b_L} (D^{\tilde{b}}_{LR}(z))^* \\
b^y_{S,RR} &= -\sqrt{2}\, g\lambda'^{*}_{333} \left((L^{\tau_L})^* D^{\tilde{\tau}}_{LL}(y) + (L^{\tau_R})^* D^{\tilde{\tau}}_{RL}(y)\right) && \to -\sqrt{2}\, g_Y \lambda'^{*}_{333}\, \eta_1^*\, Y_{\tau_L}\, D^{\tilde{\tau}}_{LL}(y) \\
b^y_{S,RL} &= +\sqrt{2}\, g\lambda'^{*}_{333} \left((R^{\tau_R})^* D^{\tilde{\tau}}_{RL}(y) + (R^{\tau_L})^* D^{\tilde{\tau}}_{LL}(y)\right) && \to +\sqrt{2}\, g_Y \lambda'^{*}_{333}\, \eta_1^*\, Y_{\tau_R}\, D^{\tilde{\tau}}_{RL}(y) .
\end{aligned} \tag{4.15}$$

Notice that no correspondent of the crossed diagrams considered in the massless final state case exist here.

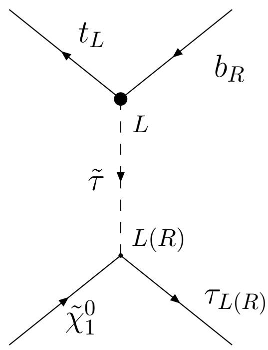

Figure 4: Diagram contributing to the decay $\tilde\chi^0_1 \to t\bar b\tau$ through exchange of a virtual $\tilde\tau$-slepton. The thick vertex indicates the R_p-violating coupling λ'^*_{333}.

4.3 Massive final state — $\tilde\chi^0_1 \to \bar t b \bar\tau$

This decay mode is obtained CP conjugating the previous one $\tilde\chi^0_1 \to t\bar b\tau$, and it has therefore the same width as the previous one. For completeness, however, we give also the matrix elements for this decay. Splitting them again into Q terms and coefficients, we find:

$$\begin{aligned}
Q^x_{S,LL} &= (\overline{u}_\tau(p_2) P_L v_b(p_3)) \left(\overline{u}_t(p_1) P_L\, u_{\tilde\chi^0_1}(q)\right) \\
Q^x_{S,LR} &= (\overline{u}_\tau(p_2) P_L v_b(p_3)) \left(\overline{u}_t(p_1) P_R\, u_{\tilde\chi^0_1}(q)\right) \\
Q^z_{S,LL} &= (\overline{u}_\tau(p_2) P_L v_t(p_1)) \left(\overline{u}_b(p_3) P_L\, u_{\tilde\chi^0_1}(q)\right) \\
Q^z_{S,LR} &= (\overline{u}_\tau(p_2) P_L v_t(p_1)) \left(\overline{u}_b(p_3) P_R\, u_{\tilde\chi^0_1}(q)\right) \\
Q^y_{S,LL} &= (\overline{u}_t(p_1) P_L v_b(p_3)) \left(\overline{u}_\tau(p_2) P_L\, u_{\tilde\chi^0_1}(q)\right) \\
Q^y_{S,LR} &= (\overline{u}_t(p_1) P_L v_b(p_3)) \left(\overline{u}_\tau(p_2) P_R\, u_{\tilde\chi^0_1}(q)\right) .
\end{aligned} \tag{4.16}$$

Notice that these differ form the Q terms in eq. (4.14) for an exchange $R \leftrightarrow L$. The corresponding coefficients,

$$\begin{aligned}
c^x_{S,LL} &= +\sqrt2\, g\lambda'_{333} \left((L^{t_L})(D^{\tilde t}_{LL}(x))^* + (L^{t_R})(D^{\tilde t}_{RL}(x))^*\right) \to +\sqrt2\, g_Y \lambda'_{333}\, \eta_1\, Y_{t_L} \left(D^{\tilde t}_{LL}(x)\right)^* \\
c^x_{S,LR} &= -\sqrt2\, g\lambda'_{333} \left((R^{t_R})(D^{\tilde t}_{RL}(x))^* + (R^{t_L})(D^{\tilde t}_{LL}(x))^*\right) \to -\sqrt2\, g_Y \lambda'_{333}\, \eta_1\, Y_{t_R} \left(D^{\tilde t}_{RL}(x)\right)^* \\
c^z_{S,LL} &= +\sqrt2\, g\lambda'_{333} \left((R^{b_R})^* D^{\tilde b}_{RR}(z) + (R^{b_L})^* D^{\tilde b}_{LR}(z)\right) \quad \to +\sqrt2\, g_Y \lambda'_{333}\, \eta^*_1\, Y_{b_R} \left(D^{\tilde b}_{RR}(z)\right) \\
c^z_{S,LR} &= -\sqrt2\, g\lambda'_{333} \left((L^{b_L})^* D^{\tilde b}_{LR}(z) + (L^{b_R})^* D^{\tilde b}_{RR}(z)\right) \quad \to -\sqrt2\, g_Y \lambda'_{333}\, \eta^*_1\, Y_{b_L} \left(D^{\tilde b}_{LR}(z)\right) \\
c^y_{S,LL} &= -\sqrt2\, g\lambda'_{333} \left((L^{\tau_L})(D^{\tilde\tau}_{LL}(y))^* + (L^{\tau_R})(D^{\tilde\tau}_{RL}(y))^*\right) \to -\sqrt2\, g_Y \lambda'_{333}\, \eta_1\, Y_{\tau_L} \left(D^{\tilde\tau}_{LL}(y)\right)^* \\
c^y_{S,LR} &= +\sqrt2\, g\lambda'_{333} \left((R^{\tau_R})(D^{\tilde\tau}_{RL}(y))^* + (R^{\tau_L})(D^{\tilde\tau}_{LL}(y))^*\right) \to +\sqrt2\, g_Y \lambda'_{333}\, \eta_1\, Y_{\tau_R} \left(D^{\tilde\tau}_{RL}(y)\right)^* ,
\end{aligned} \tag{4.17}$$

are, as expected:

$$c^t_{S,LL} = \left(b^t_{S,RR}\right)^*, \qquad c^t_{S,LR} = \left(b^t_{S,RL}\right)^*, \tag{4.18}$$

where $t = x, y, z$.

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5. Additional large couplings of type λ' and λ

As argued in section 2, there are other couplings of type λ and λ', besides λ'_{333}, that can be large. They originate a variety of final states with possibly interesting experimental signatures. We start discussing some final states due to couplings of type λ', in particular those with two indices of third generation and one of second generation. Generalization to other cases are obvious.

A nonvanishing coupling λ'_{323} gives rise to final states with only light particles. There are two decays into charmed final states: $\tilde{\chi}_1^0 \to c\bar{b}\tau$ and $\tilde{\chi}_1^0 \to \bar{c}b\bar{\tau}$. The diagrams relative to the decay $\tilde{\chi}_1^0 \to c\bar{b}\tau$ can be obtained from those in figures 3 and 4, substituting everywhere $t \to c$. The matrix elements are obtained from those in eqs. (4.14) and (4.15), and in (4.16) and (4.17), respectively. In the evaluation of the matrix elements squared, however, the limit $r \to 0$ has to be taken in appendix D. Decay modes with s-quark in the final state, i.e. $s\bar{b}\nu_\tau$ and $\bar{s}b\nu_\tau$ are also possible. The Feynmman diagrams for $\tilde{\chi}_1^0 \to s\bar{b}\nu_\tau$ due to the coupling λ'^*_{323} are shown explicitly in figures 5 and 6. (Those due to the coupling λ'^*_{332} are obtained from the diagrams in figures 5 and 6 interchanging s with b everywhere.) The matrix elements for the decay mode $\tilde{\chi}_1^0 \to s\bar{b}\nu_\tau$ can be obtained with the replacement $b \to s$ in the first and second line of eq. (4.5) and by taking the momentum p_1 in eq. (4.4) to be the momentum of the s-quark, that is to say, substituting $u_b(p_1)$ with $u_s(p_1)$. Those for the decay mode $\tilde{\chi}_1^0 \to \bar{s}b\nu_\tau$, can be obtained by replacing $b \to s$ in the third and forth line of eq. (4.11) and by taking the momentum p_3 in eq. (4.10) to be the momentum of the s-quark, i.e. substituting $u_b(p_3)$ with $u_s(p_3)$.

As in the case of λ'_{333}, the coupling λ'_{233} induces also decays of the lightest neutralino into a roughly massless final state, $\bar{b}b\nu_\mu$, and into massive ones: $\bar{b}t\mu$ with its conjugate state $b\bar{t}\bar{\mu}$. All formulae derived in section 4 apply also to this case, with the obvious changes: $\nu_\tau \to \nu_\mu$, $\tau \to \mu$, $\tilde{\tau} \to \tilde{\mu}$ and $\lambda'_{333} \to \lambda'_{233}$.

Similarly, a nonvanishing coupling λ'_{332} gives also rise to decay modes with massive particles in the final state: $t\bar{s}\tau$ and $\bar{t}s\bar{\tau}$. The treatment of these decays follows that for $t\bar{b}\tau$ and $\bar{t}b\bar{\tau}$ in sections 4.2 and 4.3. (All needed formulae are obtained from those in these sections with the substitution $b \to s$, $\tilde{b} \to \tilde{s}$, and $\lambda'_{333} \to \lambda'_{332}$.) The massless final states induced by this coupling are $\bar{b}s\nu_\tau$ and $b\bar{s}\nu_\tau$, already discussed in the case of the coupling λ'_{323}. However, λ'_{323} naturally induces a left-handed s-field, whereas λ'_{332} induces a right-handed one. The matrix elements for the decay $\tilde{\chi}_1^0 \to b\bar{s}\nu_\tau$ can be gleaned from the Q terms and coefficients in eqs. (4.4) and (4.5) as follows: substitute $b \to s$ in the third and fourth line of eq. (4.5) and take the momentum p_3 in eq. (4.4) to be the momentum of the s-quark, that is to say, substitute $v_b(p_3)$ with $v_s(p_3)$. The Q terms and coefficients in eqs. (4.10), in turn, can be used to obtain the matrix elements for the decay $\tilde{\chi}_1^0 \to s\bar{b}\nu_\tau$. The replacement recipe is now as follows: substitute $b \to s$ in the first and second line of eq. (4.11) and take the momentum p_1 in eq. (4.10) to be the momentum of the s-quark, i.e. substitute $u_b(p_1)$ with $u_s(p_1)$.

Notice that, among the abovementioned couplings, only λ'_{323} and λ'_{332} remain unconstrained by the requirement that the loop contributions to neutrino masses are not too large, if the left-right mixing terms in the sfermion sector do not play a role in the suppression of these loops.

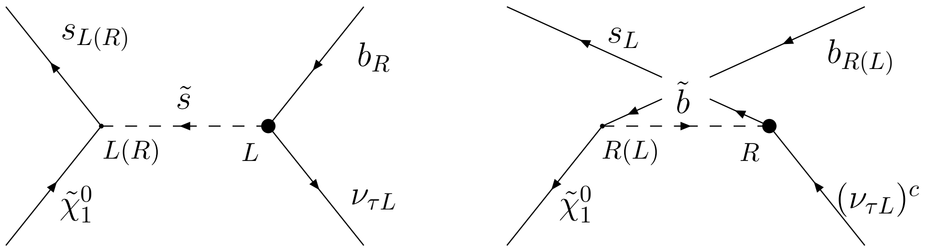

Figure 5: Diagrams contributing to the decay $\tilde{\chi}_1^0 \to s\bar{b}\nu_\tau$ through the exchange of a virtual $\tilde{s}$- and $\tilde{b}$-squark. The thick vertex indicates the R_p-violating coupling λ'^*_{323}.

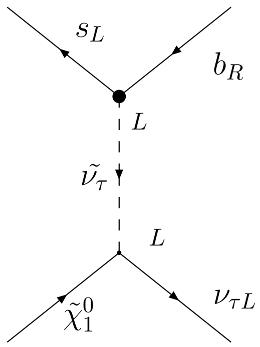

Figure 6: Diagram contributing to the decay $\tilde{\chi}_1^0 \to s\bar{b}\nu_\tau$ through exchange of a virtual $\tilde{\nu_\tau}$-slepton. The thick vertex indicates the R_p-violating coupling λ'^*_{323}.

The typical signatures induced by the coupling λ'_{333} that may be expected at the LHC were already listed in section 1. In a similar way, the couplings λ'_{332} and λ'_{323} give rise to the R_p-conserving final states

$$
\begin{aligned}
pp &\to t\bar{b}\tilde{\mu}X \to t\bar{b}\bar{t}bX\,,\\
pp &\to t\bar{s}\tilde{\tau}X \to t\bar{s}\,\bar{t}sX\,,\\
pp &\to c\bar{b}\tilde{\tau}X \to c\bar{b}\,\bar{c}bX\,,
\end{aligned}
\tag{5.1}
$$

which may also be induced by flavor-conserving and flavor-violating decays of charged-Higgs bosons, or, in the last case, by flavor-violating decays of a pair of W bosons. The couplings λ'_{233}, λ'_{332}, and λ'_{323} also originate R_p-violating final states such as

$$
\begin{aligned}
pp &\to t\bar{b}\tilde{\mu}X \to (2t)\,(2\bar{b})\,(2\mu)X\,,\\
pp &\to t\bar{s}\tilde{\tau}X \to (2t)\,(2\bar{s})\,(2\tau)X\,,\\
pp &\to c\bar{b}\tilde{\tau}X \to (2c)\,(2\bar{b})\,(2\tau)X\,,
\end{aligned}
\tag{5.2}
$$

giving rise to pairs of like-sign dileptons. As already mentioned in section 1, that of like-sign dileptons constitute a distinctive signature of these production and decay mechanisms when the t-quarks involved in these final states decay completely into hadrons. The R_p-violating states $(2t)\,(2\bar{b})\,(2\mu)$, $(2c)\,(2\bar{b})\,(2\tau)$, $(2t)\,(2\bar{s})\,(2\tau)$, and their conjugated ones, are the typical final states obtained from pair-produced lightest neutralinos, (at the LHC or e^+e^- colliders), when couplings λ'_{ijk} with more than two third-generation indices are non-vanishing.

Couplings of type λ are antisymmetric in the first two indices. Therefore, there is only one such coupling with two third-generation indices: λ_{233}. The neutralino decays induced by this coupling give rise to the final states: $\tau\bar{\tau}\nu_\mu$, and $\mu\nu_\tau\bar{\tau}$, together with $\bar{\mu}\nu_\tau\tau$. (Once again, no distinction is made between ν_τ and $\bar{\nu}_\tau$.) The calculations of the corresponding widths match those in section 4, when the following changes are made: $b \to \tau$, $\tilde{b} \to \tilde{\tau}$, $\nu_\tau \to \nu_\mu$, and $\lambda'_{333} \to \lambda_{233}$, in the case of the first decay mode, $\tau\bar{\tau}\nu_\mu$; and $t \to \nu_\tau$, $\tau \to \mu$, $b \to \tau$, and again $\lambda'_{333} \to \lambda_{233}$, in the case of the remaining two decay modes $\mu\nu_\tau\bar{\tau}$ and $\bar{\mu}\nu_\tau\tau$. Typical final states that can be expected at the LHC are:

$$\begin{aligned} pp \to t\bar{b}\tilde{\tau}X &\to t\bar{b}(2\tau)\bar{\tau}\nu_\nu X\,, \\ &\to t\bar{b}(2\tau)\bar{\mu}\nu_\tau X\,; \end{aligned} \tag{5.3}$$

whereas signatures such as $(2\tau)(2\bar{\tau})$ and missing energy, or the typical like-sign dilepton signatures $(2\tau)(2\bar{\mu})$ (and $(2\bar{\tau})(2\mu)$) plus missing energy are obtained from pair-produced lightest neutralinos.

The discussion can be generalized to other couplings of type λ with more than one index different from 3.

6. Numerical results

We are now in a position to discuss the relative size of widths for the final states originated by the decay of the lightest neutralino, when some of the λ' and λ couplings are considerably larger than the others. We concentrate first on the somewhat ideal case in which only one R_p-violating coupling is present, λ'_{333}, and $\tilde{\chi}_1^0$ can only decay into one of the three final states $b\bar{b}\nu_\tau$, $t\bar{b}\tau$, and $\bar{t}b\bar{\tau}$. We shall discuss later other decay modes originated by other couplings of type λ'. Finally we shall show results for a framework in which λ'_{333} and λ_{233} are simultaneously nonvanishing, and dominant among all R_p-violating couplings.

In general, we do not assume the relation $m_{\tilde{W}} \simeq 2m_{\tilde{B}}$ among gaugino masses, which is typical of an mSUGRA scenario, with gaugino mass unification at a high scale. Indeed, gaugino mass unification is not a necessary ingredient of this scenario, only a customary one. Moreover, models with Wino states lighter than the Bino have recently emerged. The relation $m_{\tilde{B}} = km_{\tilde{W}}$, with $k > 1$, is, for example, predicted in models in which the breaking of supersymmetry is transmitted to the visible sector through anomaly meditation (see [60] and references therein). In this case, it is $k \simeq 3$. A more complicated relation between $m_{\tilde{B}}$ and $m_{\tilde{W}}$ is also possible in some grand unification models [61] in which the usual gaugino mass unification in replaced by a more complicated relation among the gaugino masses [62]. In this case, $m_{\tilde{W}}$ and $m_{\tilde{B}}$, together with the gluino mass, $m_{\tilde{g}}$, satisfy a linear constraint, $m_{\tilde{B}} = km_{\tilde{W}} + hm_{\tilde{g}}$, with a wide range of values possible for k.

In the following, only the two specific values $k = 1/2$ and $k = 2$ will be discussed.

6.1 Only $\lambda'_{333} \neq 0$

For a relatively light lightest neutralino, i.e. when $\tilde{\chi}_1^0$ is below the t-quark threshold, the only possible decay mode is $b\bar{b}\nu_\tau$. Above the t-quark threshold, we should distinguish the case in which $\tilde{\chi}_1^0$ is mainly a gaugino with a small contamination from the two Higgsinos,

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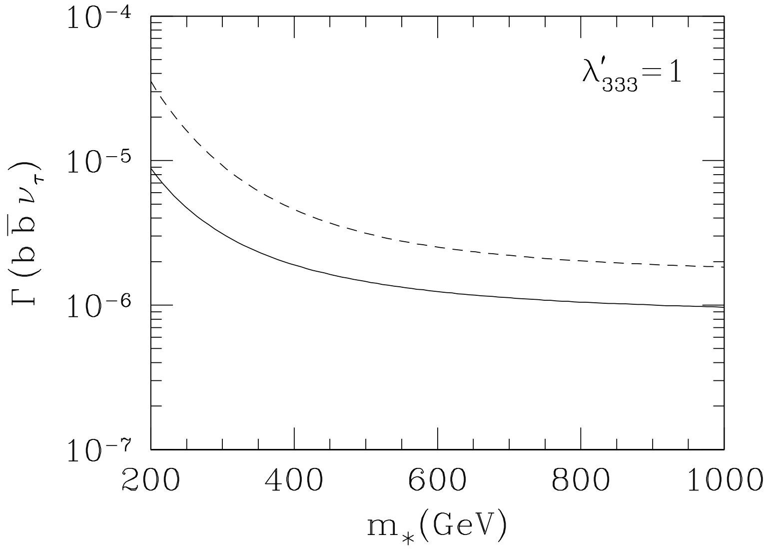

Figure 7: $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$ (in GeV) versus $m_* = m_{\tilde{Q}} = m_{\tilde{U}^c} = m_{\tilde{D}^c}$ for the decay of an almost pure gaugino $\tilde{\chi}_1^0$ for $\mu = 500\,\text{GeV}$, $\tan\beta = 3$, $m_{\tilde{L}} = m_{\tilde{E}^c} = 400\,\text{GeV}$, and gaugino masses $m_{\tilde{B}} = (1/2)m_{\tilde{W}} = 100\,\text{GeV}$ (solid lines), $m_{\tilde{W}} = (1/2)m_{\tilde{B}} = 100\,\text{GeV}$ (dashed lines). The trilinear soft terms were chosen in such a way to cancel the left-right mixing terms in all sfermion mass matrices.

from the case in which $\tilde{\chi}_1^0$ has a substantial Higgsino component. In the former, among the three states in which $\tilde{\chi}_1^0$ can decay, $b\bar{b}\nu_\tau$, $t\bar{b}\tau$, and $\bar{t}b\bar{\tau}$, the state $b\bar{b}\nu_\tau$ has the largest width and therefore the largest branching ratio, when the values of the sfermion masses virtually exchanged in the three decays are not too dissimilar. The branching ratio of the two combined massive modes, however, is, in a wide range of parameter space, of the same order of magnitude of that for the $b\bar{b}\nu_\tau$ mode. This statement is valid irrespective of the value of $\tan\beta$. We observe that the widths for all decay modes are larger if the main gaugino component of $\tilde{\chi}_1^0$ is of Wino type instead than of Bino type.

The two massive modes may become dominant if the gaugino-Higgsino admixture in the lightest neutralino state is substantial and $\tan\beta$ is not too large. This is due to the large t-quark Yukawa coupling. When $\tan\beta$ increases and the b-quark Yukawa coupling becomes comparable to that of the t-quark, the decay mode $b\bar{b}\nu_\tau$ takes over and dominates over the $t\bar{b}\tau$ one.

We show in figure 7 the width $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$, in the case of a light $\tilde{\chi}_1^0$ that is mainly a Bino (solid line), or mainly a Wino (dashed line). In the lower curve, the $\tilde{B}$ mass is fixed to 100 GeV, the $\tilde{W}$ mass is assumed to be twice as large as the $\tilde{B}$ mass, whereas the μ parameter is given the larger value of 500 GeV. The width is shown as a function of the doublet squark mass, $m_{\tilde{Q}}$ assumed to be equal to both, the stop and sbottom singlet masses, $m_{\tilde{U}^c}$ and $m_{\tilde{D}^c}$. The slepton doublet and singlet masses $m_{\tilde{L}} = m_{\tilde{E}^c}$ are fixed to 400 GeV. The soft trilinear terms for squarks and sleptons as well as $\tan\beta$ are chosen in such a way to obtain vanishing left-right entries in the corresponding sfermion mass matrices. Specifically, $\tan\beta = 3$ was used for this figure. This choice, however, does not affect the results shown here: the width $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$ is practically independent of the value of $\tan\beta$, in the approximation of $\tilde{\chi}_1^0$ as a pure gaugino and when the left-right mixing

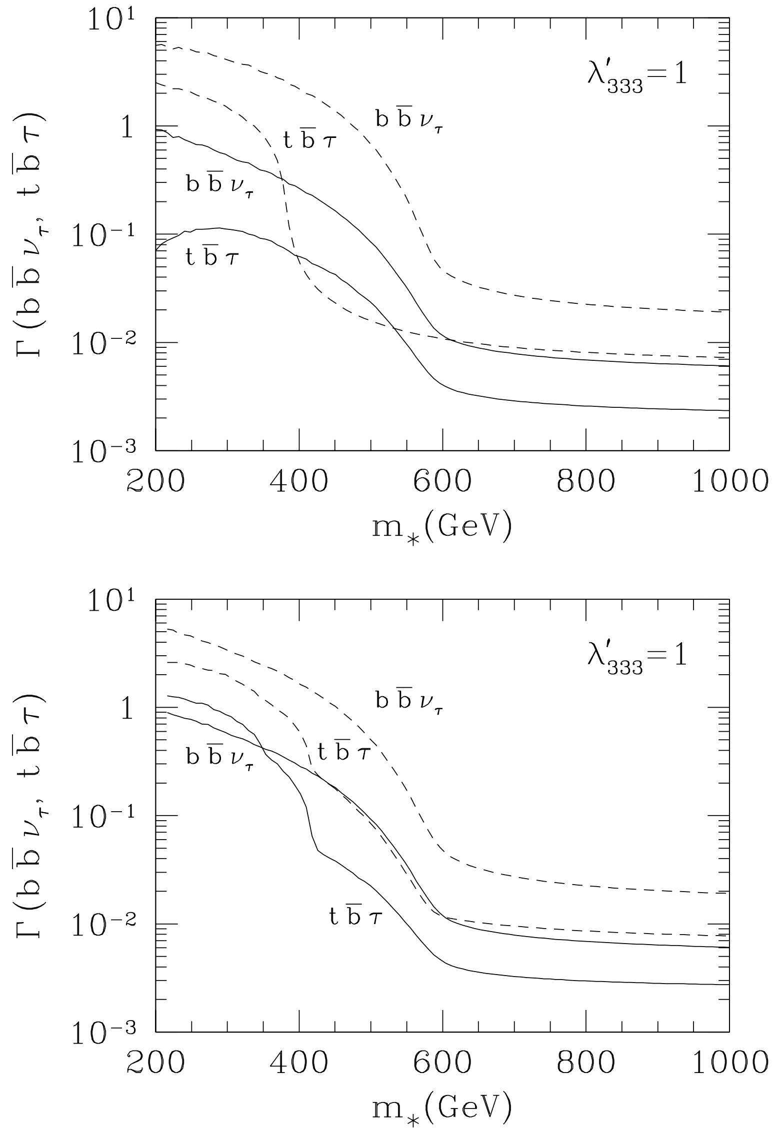

Figure 8: $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$ and $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$ (in GeV) (in GeV) versus $m_* = m_{\tilde{Q}} = m_{\tilde{U}^c} = m_{\tilde{D}^c}$, for $\mu = 1500\,\text{GeV}$, $\tan\beta = 3$, $m_{\tilde{L}} = m_{\tilde{E}^c} = 600\,\text{GeV}$. The gaugino masses are $m_{\tilde{B}} = (1/2)m_{\tilde{W}} = 600\,\text{GeV}$ (solid lines), $m_{\tilde{W}} = (1/2)m_{\tilde{B}} = 600\,\text{GeV}$ (dashed lines). The trilinear soft terms are chosen in such a way to give vanishing left-right entries in all sfermion mass matrices in the upper frame and moderate left-right entries the lower frame.

terms among sfermions are vanishing. In this case, the only dependence on $\tan\beta$ of the widths for the lightest neutralino decays is, in principle, due to the D-term contributions to the sfermion mass eigenvalues. For left- and right-entries in the sfermion mass matrices larger than 200 GeV, this dependence is completely negligible. Thus, for simplicity, in this figure and in the following ones, the D-term contributions to the sfermion mass matrices were ignored. Furthermore, entries in the sfermion mass matrices that are explicitly R_p-violating, are proportional to the bilinear couplings k_i and to the *vev*'s v_i of the fields $\tilde{L}_i$. They were dropped in this analysis, since completely negligible.

The upper curve in this plot (dashed line) shows the width $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$ in the case in which $\tilde{\chi}_1^0$ is mainly a Wino. This approximation was achieved simply inverting the ratio of $m_{\tilde{B}}$ and $m_{\tilde{W}}$, i.e. taking $m_{\tilde{B}} = 2m_{\tilde{W}}$. All other massive parameters have the same values as those chosen in the approximation in which the lightest neutralino is mainly a Bino. Although one diagram less contributes now to $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$ (the right-handed sbottom squarks does not couple to the Wino), the larger value of the couplings fermion-sfermion-Wino accounts for the larger width obtained in this case.

It should be mentioned here that some of the trilinear terms used in this figure (as well as in some of the following figures) to cancel the left-right sfermion mixing terms, may seem, at times, a little too large. However, as it was argued in ref. [12], there are ways to accommodate large trilinear terms in supersymmetric models, provided they do not give rise to negative sfermion mass squared eigenvalues. One of the main constraints to their size remains, the requirement that the radiative contribution to fermion masses induced by such large terms does not exceed the experimentally observed ones.

In figure 8 the almost complete dominance of the widths $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$ and $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$, obtained for a $\tilde{\chi}_1^0$ that is mainly a Wino, over those obtained when $\tilde{\chi}_1^0$ is mainly a Bino, is explicitly shown in the case of a heavy lightest neutralino, i.e. with mass $\simeq$ 600 GeV. Indeed, the lightest of the two gauginos, the Bino in the solid lines ($m_{\tilde{W}} = 2m_{\tilde{B}}$), or the Wino in the dashed lines ($m_{\tilde{B}} = 2m_{\tilde{W}}$), are fixed at 600 GeV, whereas the μ parameter has the large value $\mu = 1500$ GeV. The values $m_{\tilde{L}} = m_{\tilde{E}^c} = 600$ GeV are used for the slepton masses. No left-right mixing in all sfermion mass matrices is assumed in the upper frame of this figure. This is achieved fixing the value of $\tan\beta$ to 3 and choosing consequently the trilinear A couplings. As already mentioned, for a lightest neutralino that is mainly a gaugino, there is practically no $\tan\beta$ dependence in the widths $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$, $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$, and $\Gamma(\tilde{\chi}_1^0 \to \bar{t}b\bar{\tau})$, in the absence of left-right mixing terms among sfermions. In both approximations, that of a mainly Bino and that of a mainly Wino for the lightest neutralino, for similar $\tilde{t}$ and $\tilde{b}$ eigenvalues, the decay mode $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$ dominates over the other two, which are penalized by a large phase-space suppression and by the fact that they are here considerated separately: being $b\bar{b}\nu_\tau$ a self-conjugated state, the decay $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$ collects actually contributions to $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$ and $\tilde{\chi}_1^0 \to \bar{b}b\bar{\nu}_\tau$. The upper curves in the two sets of solid and dashed lines, show $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$ as a function of $m_{\tilde{Q}} = m_{\tilde{U}^c} = m_{\tilde{D}^c}$, and $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$ is shown by the two lower curves. In the dashed curve corresponding to $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$, which is determined by the $\tilde{t}$- and $\tilde{\tau}$-exchange diagrams, (no right-handed $\tilde{b}$ exchange is possible in this case, since the lightest neutralino is mainly a Wino) the resonant region in which the $\tilde{t}$-squark is produced on shell is clearly visible. Between $400 \leq m_{\tilde{Q}} \leq 600$ GeV, the width drops roughly as $1/(m_{\tilde{t}}^2)^2$, whereas the almost flat behavior after 600 GeV is due to the dominance of the $\tilde{\tau}$-exchange diagram. The solid line representing $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$, and corresponding to a mainly Bino lightest neutralino, is determined by all the three diagrams in figure 4. In the resonant region in which $\tilde{b}$-squarks are produced on shell, the diagram with exchange of a right-handed $\tilde{b}$-squark is actually the dominant one. For decreasing values of $m_{\tilde{Q}}$, in the region $m_{\tilde{Q}} < 600$ GeV, the increase of the width, due to the dependence on $(m^2_{\tilde{\chi}_1^0} - m^2_{\tilde{b}})^{-2}$ is damped at the lower end of $m_{\tilde{Q}}$ by the severe phase-space suppression in the branching ratio of the 2-body decay $\tilde{b} \to t\tau$.

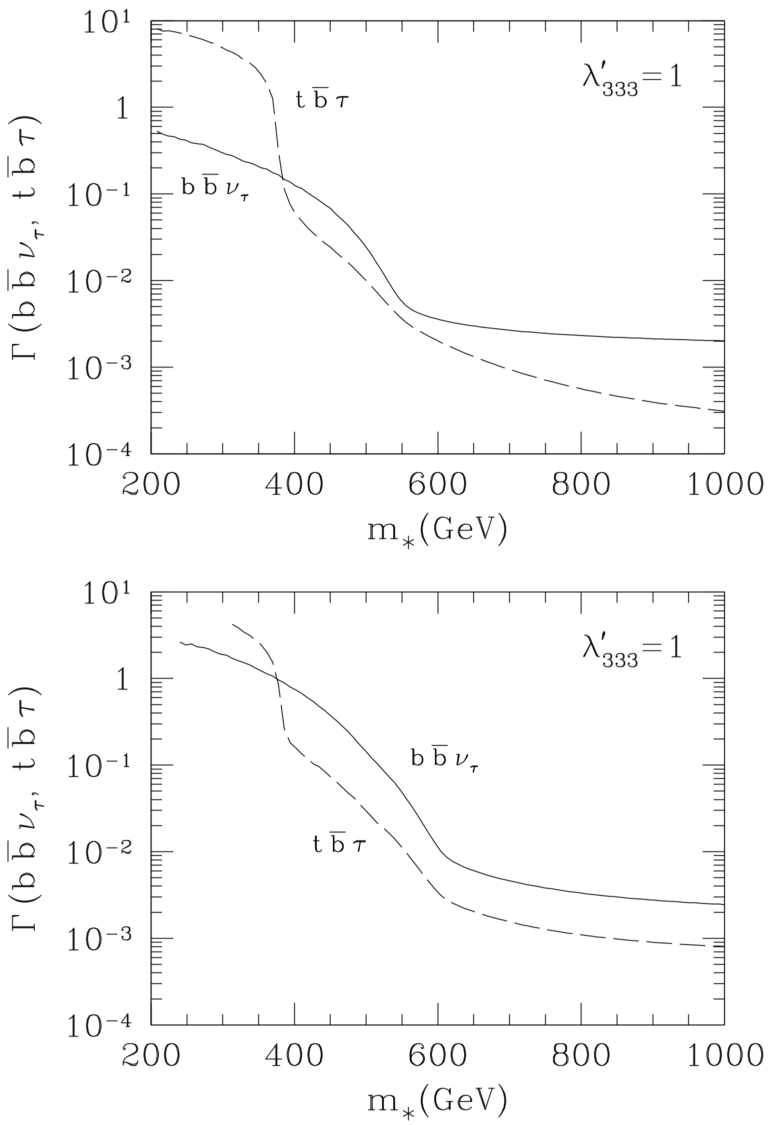

Figure 9: $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$ and $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$ (in GeV), respectively solid and long-dashed lines, versus $m_* = m_{\tilde{Q}} = m_{\tilde{U}^c} = m_{\tilde{D}^c}$, for $\mu = 600\,\mathrm{GeV}$, $m_{\tilde{L}} = m_{\tilde{E}^c} = 600\,\mathrm{GeV}$, $m_{\tilde{B}} = 600\,\mathrm{GeV}$, and $m_{\tilde{W}} = 2m_{\tilde{B}}$. The value of $\tan\beta$ is 3 in the upper frame, 30 in the lower one. All trilinear soft terms are fixed at 350 GeV in the upper frame, and at 150 GeV in the lower frame.

In the lower frame of figure 8, a moderate splitting among the two $\tilde{t}$ and $\tilde{b}$ eigenvalues is allowed, i.e. $A_t - \mu/\tan\beta = 150\,\mathrm{GeV}$, and $(A_b - \mu\tan\beta)m_b \sim (100)^2\,\mathrm{GeV}^2$. Because the left-right mixing terms in the $\tilde{b}$ mass matrix square are nonvanishing, the $\tilde{b}$-mediated diagram contributes now to the width $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$ also when the lightest neutralino is mainly a Wino. This contribution explains the difference in shape of the width $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$ in the region $400 \le m_{\tilde{Q}} \le 600\,\mathrm{GeV}$, with respect to that obtained in the absence of left-right mixing terms. Similarly, the smaller values of the mass of the lightest $\tilde{t}$- and $\tilde{b}$-squarks explain the large enhancement of the width $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$ obtained for $m_{\tilde{Q}} \le 400\,\mathrm{GeV}$ when $\tilde{\chi}_1^0$ is mainly a Bino. Larger values of $\tan\beta$, i.e. larger values of left-right mixing terms in the $\tilde{b}$ mass matrix, would further enhance $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$ with respect to $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$.

If a large mixing Bino-Higgsino is allowed, then, a substantial increase in the width of $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$ is expected for low $\tan\beta$ and an increase of $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$ for large $\tan\beta$. These features are shown explicitly in figure 9, where μ, $m_{\tilde{B}}$ as well as the doublet and singlet slepton masses are all fixed at 600 GeV. The value of $\tan\beta = 3$ and $A_t = A_b = A_\tau = 350$ GeV are used in the upper frame, $\tan\beta = 30$ and $A_t = A_b = A_\tau = 150$ GeV in the lower frame. The curves stop when $m_{\tilde{b}} < m_t$ in the case of the width $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$, and $m_{\tilde{b}} < 70$ GeV (with 70 GeV an average cut that should mimic more complicated and model-dependent experimental lower bounds), in the case of the width $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$.

In conclusion, when only the coupling λ'_{333} is nonnegligible, the lightest neutralino decays, in general, as $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$. A final state with $4b$'s and missing energy is therefore the distinctive signature of two pair-produced neutralinos. Excesses of b-quarks at hadron collider may be obtained from sfermions decaying through gauge interactions into neutralinos and the corresponding fermions. The two massive decay modes $\tilde{\chi}_1^0 \to t\bar{b}\tau$ and $\tilde{\chi}_1^0 \to \bar{t}b\bar{\tau}$ are, in general, subdominant, although there are wide regions of parameter space in which they have rates of the same order of magnitude as that of the decay $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$. The rates for the two decays into t-quarks are, however, the largest when the lightest neutralino has a large Higgsino component and $\tan\beta$ is not too large. This dominance is observed in the region of resonant neutralino decay into an intermediate $\tilde{t}$-squark.

A few more comments are in order here regarding the choice of slepton masses made for the different phenomenological situations illustrated in this section. Fixed values of slepton masses were chosen, smaller than the largest values of squark masses and, in general, bigger than the smallest values. In most of the known supersymmetric models, however, slepton masses are smaller than squark masses and, being linked to the gravitino mass, as the squark masses, they also increase when the squark masses increase. If, for example, the choice $m_{\tilde{L}} = m_{\tilde{E}^c} = (1/2)m_{\tilde{Q}}$ ($m_{\tilde{Q}} = m_{\tilde{U}^c} = m_{\tilde{D}^c}$) is made, then the slepton-exchange diagram becomes more important than it is in the different plots shown in this section. In general the widths are enhanced both in the resonant and nonresonant region. In particular, in the nonresonant region, widths are always decreasing and do not exhibit the typical plateau that indicates now the slepton-exchange contribution to be larger than those with exchange of the lightest squarks. An enhancement is also expected when the slepton-exchange contribution is resonant: i.e. a similar type of enhancement for both $b\bar{b}\nu_\tau$ and $t\bar{b}\tau$ modes at small $\tan\beta$, and possibly a larger enhancement for the $t\bar{b}\tau$ mode at large $\tan\beta$, or more generically, when large left-right mixing terms are present in the charged-slepton mass matrix.

Finally, we show how the widths presented so far change when lower values of the coupling λ'_{333} are considered. We illustrate the dependence of the two widths $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$ and $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$ on this coupling, for the particular direction of parameter space considered in the upper frame of figure 9.

In figure 10, the two widths are shown in solid ($\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$) and dashed ($\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$) lines, for different values of λ'_{333}, i.e. $\lambda'_{333} = 1$, 0.5 and 0.1. In the resonant sfermion region, there is very little dependence on the value of λ'_{333} chosen. Off resonance, the widths are scaled down by the factor $|\lambda'_{333}|^2$. This feature is rather general and holds also in the case of other directions of parameter space.

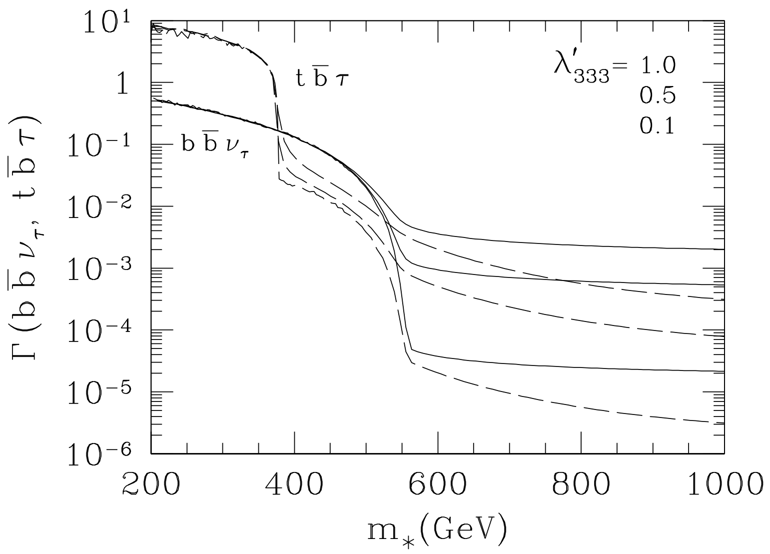

Figure 10: Dependence of $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$ (solid lines) and $\Gamma(\tilde{\chi}_1^0 \to t\bar{b}\tau)$ (long-dashed lines) on the coupling λ'_{333}. All other parameters are chosen as in the upper frame of figure 9. The values of λ'_{333} are 1, 0.5, 0.1 from top to bottom.

6.2 λ'_{333} and $\lambda'_{323} \neq 0$

It is possible that apart from λ'_{333}, some other couplings of type λ' are nonnegligible. As discussed in section 5, couplings such as λ'_{233} and λ'_{332} give rise, as λ'_{333}, to decays of the lightest neutralino into massive and practically massless final states. Qualitatively, the results obtained for λ'_{333} in the previous subsection, therefore, also apply to these couplings.

Different is the situation obtained for the coupling λ'_{323}. As mentioned already, only final states with light particles can be obtained in this case: $\nu_\tau s\bar{b}$, $\nu_\tau \bar{s}b$ and $c\bar{b}\tau$, $\bar{c}b\bar{\tau}$. Those with charged leptons are not phase-space suppressed as the decay modes $t\bar{b}\tau$ and $\bar{t}b\bar{\tau}$ discussed in the previous subsection. The widths of the four decay modes $b\bar{b}\nu_\tau$, $t\bar{b}\tau$, $b\bar{s}\nu_\tau$, and $c\bar{b}\tau$ are shown explicitly in figure 11, when both couplings λ'_{333} and λ'_{323} are equal to 1. They are plotted as functions of the mass of the sfermions virtually exchanged in the decays. Differently than in the previous figures, the common slepton mass ($m_{\tilde{L}} = m_{\tilde{E}^c}$), is now varying together with the common squark mass ($m_{\tilde{Q}} = m_{\tilde{U}^c} = m_{\tilde{D}^c}$) and the ratio $m_{\tilde{L}}/m_{\tilde{Q}}$ is, for simplicity, fixed to the value $\simeq 0.7$. The aim is to visualize, if possible, the dependence on the $\tilde{\tau}$-slepton mass in the resonant region $600 \lesssim m_* \lesssim 860\,\mathrm{GeV}$. In this region, two decay modes are present $\tilde{\tau} \to b\bar{t}$ and $b\bar{c}$, with a nontrivial dependence on the $\tilde{\tau}$-slepton mass. A mainly Bino lightest neutralino is considered (the values $m_{\tilde{B}} = 600\,\mathrm{GeV}$, $m_{\tilde{W}} = 2m_{\tilde{B}}$, and $\mu = 1500\,\mathrm{GeV}$ are used) for this figure. Finally, the choice $\tan\beta = 3$ is made, and all trilinear soft terms are chosen in such a way to have vanishing left-right terms in all sfermion mass matrices.

The distance between the curve relative to $\tilde{\chi}_1^0 \to c\bar{b}\tau$ and $\tilde{\chi}_1^0 \to t\bar{b}\tau$ in figure 11, is essentially explained by the phase-space suppression suffered by the decay mode $\tilde{\chi}_1^0 \to t\bar{b}\tau$. In spite of the more complicated decay possibilities that the $\tilde{b}$-squark has in the resonant region $m_{\tilde{b}} < m_{\tilde{\chi}_1^0}$, the decay $\tilde{\chi}_1^0 \to b\bar{s}\nu_\tau$ has a width which is about 1/2 of that for the

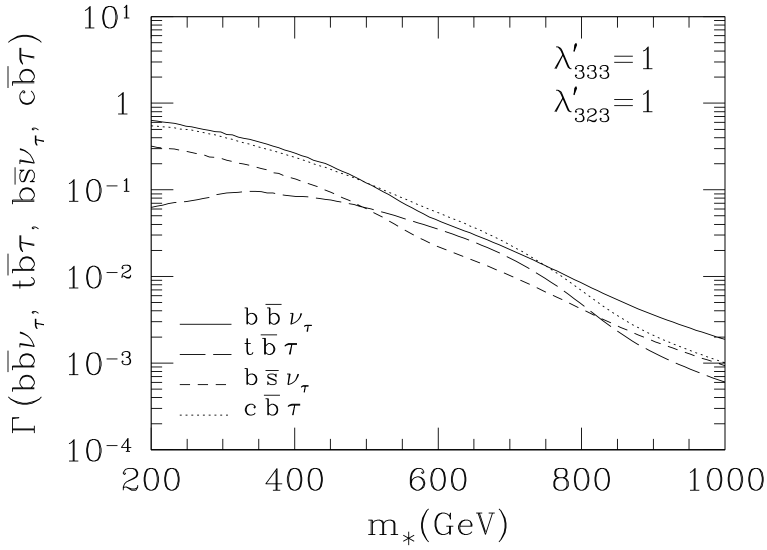

Figure 11: Widths (in GeV) for the four decays $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$, $t\bar{b}\tau$, $b\bar{s}\nu_\tau$, and $c\bar{b}\nu_\tau$, when only the two couplings λ'_{333} and λ'_{323} are nonvanishing, as function of the sfermion mass m_*, equal to $m_{\tilde{Q}} = m_{\tilde{U}^c} = m_{\tilde{D}^c}$, and to $(1.4) \times m_{\tilde{L}}$ (with $m_{\tilde{L}} = m_{\tilde{E}^c}$), for $\mu = 1500\,\mathrm{GeV}$, $m_{\tilde{B}} = 600\,\mathrm{GeV}$, and $m_{\tilde{W}} = 2m_{\tilde{B}}$. The value of $\tan\beta$ is 3 and all trilinear soft terms are chosen in such a way to have vanishing left-right mixing terms in the sfermion mass matrices squared.

decay $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$ throughout the whole range of sfermion masses. Notice that the decay mode $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$ actually collects the final states $b\bar{b}\nu_\tau$ and its conjugated one $\bar{b}b\bar{\nu}_\tau$, which are identical in the approximation of massless neutrinos. Notice further that the width for the channel $\tilde{\chi}_1^0 \to c\bar{b}\tau$ in the resonant slepton region, is about as large as the width for the channel $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$, in spite of the fact that the final state $b\bar{b}\nu_\tau$ includes the state $b\bar{b}\nu_\tau$ and its conjugated one. This can be understood as follows. In this region, the slepton-mediated diagrams are obviously the dominant ones. The lightest neutralino width is then well approximated, say in the case of the decay into $c\bar{b}\tau$, as follows: $\Gamma(\tilde{\chi}_1^0 \to c\bar{b}\tau) \simeq \Gamma(\tilde{\chi}_1^0 \to \tau\tilde{\tau}_1^*) \times BR(\tilde{\tau}_1^* \to c\bar{b})$. In the limit of almost pure Bino for the lightest neutralino, and in the absence of left-right mixing terms in the slepton mass matrices, we have $\Gamma(\tilde{\chi}_1^0 \to \tau\tilde{\tau}_1^*) = \Gamma(\tilde{\chi}_1^0 \to \nu_\tau\tilde{\nu_\tau})$. In this same approximation, and for $\lambda'_{333} = \lambda'_{323}$, we also have $BR(\tilde{\tau}_1^* \to c\bar{b}) \simeq BR(\tilde{\nu}_\tau \to b\bar{b})$. Remember that: $BR(\tilde{\tau}_1^* \to c\bar{b}) = \Gamma_{\tilde{\tau}^*}(\lambda'_{323})/(\Gamma_{\tilde{\tau}^*}(\lambda'_{323}) + \Gamma_{\tilde{\tau}^*}(\lambda'_{333})) \simeq 1/2$ and $BR(\tilde{\nu}_\tau \to b\bar{b}) = \Gamma_{\tilde{\nu}_\tau}(\lambda'_{323})/(\Gamma_{\tilde{\nu}_\tau}(\lambda'_{323}) + \Gamma_{\tilde{\nu}_\tau}(\lambda'_{333})) \simeq 1/2$, see appendix C. This explain the near equality of $\Gamma(\tilde{\chi}_1^0 \to c\bar{b}\tau)$ and $\Gamma(\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau)$, in spite of the fact that the state $b\bar{b}\nu_\tau$ is self-conjugated.

6.3 λ'_{333} and $\lambda_{233} \neq 0$

Let us now assume the somewhat simplified situation in which only one coupling of type λ', λ'_{333}, is nonnegligible together with a coupling of type λ, i.e. λ_{233}. As anticipated in section 5, the lightest neutralino decays induced by λ_{233} are $\tilde{\chi}_1^0 \to \tau\bar{\tau}\nu_\mu$, $\mu\bar{\tau}\nu_\tau$, and $\bar{\mu}\tau\nu_\tau$.

If λ'_{333} and λ_{233} have the same size and $\tilde{\chi}_1^0$ is mainly a Bino, the leptonic decay modes,

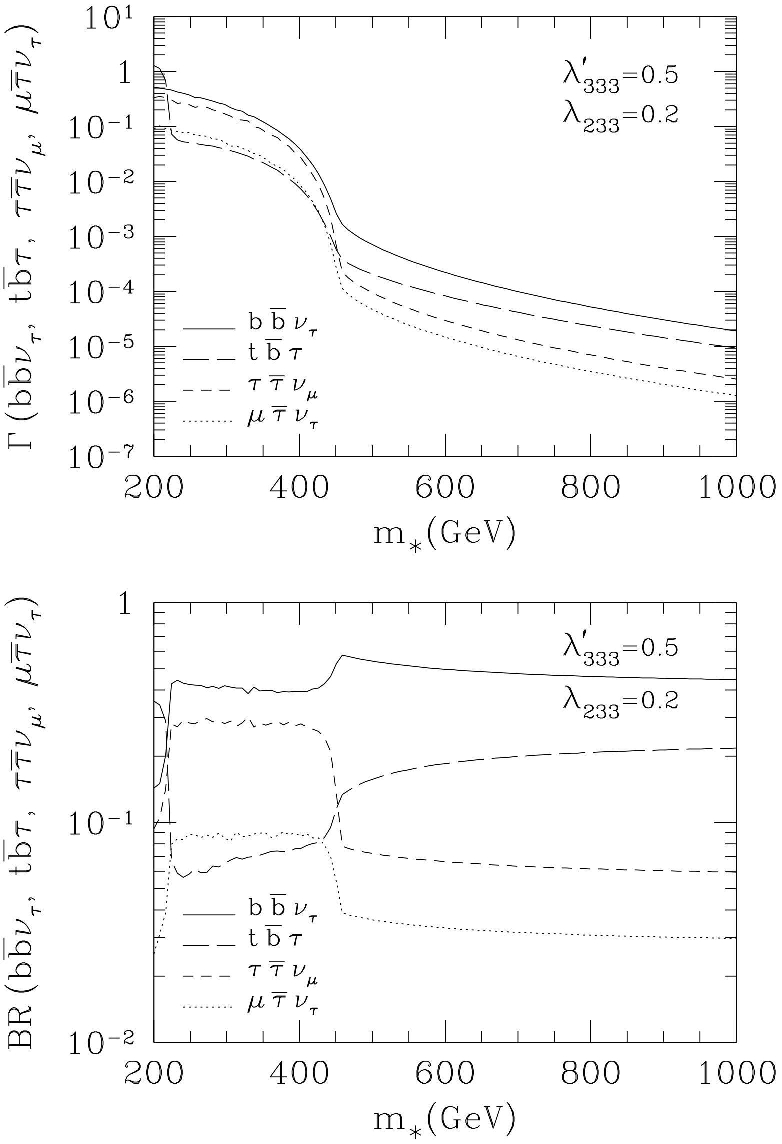

Figure 12: Widths (in GeV) and branching ratios for the four decays $\tilde{\chi}^0_1 \to b\bar{b}\nu_\tau$, $t\bar{b}\tau$, $\tau\bar{\tau}\nu_\mu$, and $\mu\bar{\tau}\nu_\tau$, when only the two couplings λ'_{333} and λ_{233} are nonvanishing, as function of the sfermion mass m_*, equal to $m_{\tilde{Q}} = m_{\tilde{U}^c} = m_{\tilde{D}^c}$, and $m_{\tilde{L}} = m_{\tilde{E}^c}$, for $\mu = 500\,\text{GeV}$, $m_{\tilde{B}} = 500\,\text{GeV}$, and $m_{\tilde{W}} = 2m_{\tilde{B}}$. The value of $\tan\beta$ is 3 and all trilinear soft terms are chosen in such a way to have vanishing left-right mixing terms in the sfermion mass matrices squared.

induced by λ_{233} always dominate over those induced by λ'_{333}. However, if λ_{233} is smaller than λ'_{333}, even only by a factor, say, 2 or 3, then the decay into $b\bar{b}\nu_\tau$, induced by λ'_{333} dominates over the leptonic ones, which, in turn, may have larger widths than the massive decay modes $t\bar{b}\tau$ and $\bar{t}b\bar{\tau}$. This situation is, in particular, realized when the neutralino is a very mixed state with nonnegligible or large Wino and Higgsino components, as shown in figure 12.

In this figure, widths and branching ratios obtained for the four decays, $\tilde{\chi}^0_1 \to b\bar{b}\nu_\tau$,

$t\bar{b}\tau$, $\tau\bar{\tau}\nu_\mu$, and $\mu\bar{\tau}\nu_\tau$, are shown as function of $m_{\tilde{Q}} = m_{\tilde{U^c}} = m_{\tilde{D^c}}$ when $\lambda'_{333} = 0.5$ and $\lambda_{233} = 0.2$. The slepton masses are assumed to be equal to the squark masses: $m_{\tilde{Q}} = m_{\tilde{L}} = m_{\tilde{E^c}}$. The gaugino masses are $m_{\tilde{B}} = (1/2)m_{\tilde{W}} = 500\,\mathrm{GeV}$, and $500\,\mathrm{GeV}$ is also the value assigned to μ. The value of $\tan\beta$ is fixed at 3, and the trilinear A terms are chosen in such a way to have vanishing left-right mixing terms in all sfermion mass matrices. Branching ratios are defined as ratios of the width for the different decay modes over the total width for the lightest neutralino. This, in turn, is given by:

$$\Gamma_{tot}(\tilde{\chi}_1^0) = \Gamma(b\bar{b}\nu_\tau) + 2\Gamma(t\bar{b}\tau) + \Gamma(\tau\bar{\tau}\nu_\mu) + 2\Gamma(\mu\bar{\tau}\nu_\tau)\,. \tag{6.1}$$

As the figure clearly shows, all widths and branching ratios are not too dissimilar, for the choice of λ'_{333} and λ_{233} made here. Notice how the scaling factor of 2 between the curves for $\Gamma(\tilde{\chi}_1^0 \to \tau\bar{\tau}\nu_\mu)$ and $\Gamma(\tilde{\chi}_1^0 \to \mu\bar{\tau}\nu_\tau)$, which would be expected for an almost pure Bino lightest neutralino, is here reduced by the substantial Higgsino and smaller Wino component.

7. Summary and conclusions

The instability of the LSP is, perhaps, the most dramatic feature of R_p-violating models. Whether it is the LSP or not, the lightest neutralino decays through R_p-violating interactions. Thanks to R_p-violating bilinear terms in the superpotential, the lightest neutralino has leptonic components. These leptonic admixtures are, in general, very small, since directly constrained by the smallness of neutrino masses. Among the other three R_p-violating terms in the superpotential, two violate explicitly the lepton number, L, $\lambda' LQD^c$ and λLLE^c, and the third, $\lambda'' U^cD^cD^c$, violates the baryon number, B. Their simultaneous presence induces a too rapid proton decay, which is experimentally excluded. It is, then, in general assumed that one of these two quantum numbers is still a conserved one, whereas R_p-violation is induced by the nonconservation of the second one.

The working assumption in this paper is that B remains the conserved quantum number. It is the violation of L, then, that is responsible for the striking phenomenology of R_p-violating models, and that induces, together with gauge interactions, the decay of the lightest neutralino. L violation may also give rise to unusual signatures for slepton decays directly, through $\tilde{l} \to q'\bar{q}$, or indirectly, through L-violating decays of the lightest neutralino, and/or chargino, and to unusual production mechanisms for sleptons at colliders. Indeed, it was recently pointed out that the coupling λ'_{333} may be responsible for the single production of $\tilde{\tau}$-sleptons at the LHC, through the processes $pp \to t\bar{b}\tilde{\tau}X$, $pp \to t\tilde{\tau}X$, with cross sections comparable to those for the corresponding single-production processes of charged Higgs bosons, $pp \to t\bar{b}H^-X$, $pp \to tH^-X$ [17]. (Single production of charged sleptons and charged Higgs bosons at the Tevatron, although possible, is certainly more difficult to detect than at the LHC.) That the coupling λ'_{333} may be large, or simply the dominant among other R_p-violating couplings, is a possibility that has not yet been challenged by experiment. Perhaps, negative searches for the abovementioned production processes and the combined decays for the $\tilde{\tau}$-sleptons may provide in the future severe restrictions in a region of the supersymmetric parameter space of which λ'_{333}

is one coordinate. For now, however, if such a coupling is the dominant one, it may be the main source of neutralino decays. The distinguishing property of this coupling with respect to other L-violating couplings is that it involves the t-quark in association with a τ-lepton. If heavy enough, the lightest neutralino can therefore give rise to a very massive final state. Such a possibility was never considered before, assuming perhaps that decays of $\tilde{\chi}_1^0$ into massive particles would be handicapped by a severe phase space suppression.

In this paper, neutralino decays are studied without any restrictive assumptions as to whether the final state is massive or massless. The coupling λ'_{333} offers a good ground to study both possibilities and to understand under which conditions the massive decay mode may be sizable with respect to the massless one. It is also representative of other couplings, such as λ'_{233} or λ'_{332}, that may also give rise to final states containing the t-quark. Given the complexity of phenomenological analyses in which many L-violating couplings are simultaneously present, some particular choices are made. First, it is considered the case in which one coupling only, i.e. λ'_{333}, is largely dominating over the others, which can be neglected. Later, it is assumed that two couplings at a time are dominant over the other ones, as for example λ'_{333} and λ'_{323}, or λ'_{333} and λ_{233}, with λ'_{333} the only one of the two capable of inducing a massive final state.

It should be remarked that the importance of the coupling λ'_{333}, and of other couplings of the same type, λ'_{323}, λ'_{332}, etc is also due to fact that, being potentially large, they may pollute the signals at incoming colliders for production and decays of charged Higgs bosons. Through these couplings, a $\tilde{\tau}$-slepton singly produced in association with a t- and a b-quark can easily mimic a charged Higgs boson giving rise to the same final states. It is therefore important to discern all possible consequences that these couplings may have. A detailed knowledge of the decays of $\tilde{\tau}$'s induced directly by these couplings or indirectly, through decays of charginos and neutralinos may be crucial to reach an unambiguous identification of scalar particles at future colliders.

The present analysis has to be considered as a first step of a more complete program of sparticle-decay studies in R_p-violating models [30]. Rather simple and general analytical expressions for neutralino decays into massless or massive final fermions are given. They are easily adaptable to any R_p-violating couplings, and to heavier neutralinos. They include the possibility of both resonant and nonresonant decays, the former one describing the situation in which the lightest neutralino is not the LSP. Some of these 3-body decays were previously discussed in the literature and some were also incorporated in event generators used in previous (LEP) and forthcoming (Tevatron) experimental analyses. For these cases, the present study should provide useful cross checks and guidelines to identify the directions of parameter space in which the R_p-violating signals are more prominent. (We emphasize here again that experimental analyses based on a scan of the parameter space obtained through Renormalization Group Equations from only 5 high-scale parameters, m, μ, M, A_0, B, as in the mSUGRA model, are not valid when trying to obtain limit on not too small R_p-violating couplings.) Analyses of decays into purely-third-generations fermions, including the t-quark are genuinely new. Analytical expressions for these decays are given in a form that can be easily implemented in event generators.[4]

The results may be summarized as follows. Quite generically, if the lightest neutralino is lighter (or only slightly heavier) than the t-quark, it decays predominantly into massless fermions. Depending on the relative size of the R_p-violating couplings, the dominant decays can be, e.g., $\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau$, due to a dominant λ'_{333} coupling, $\tilde{\chi}_1^0 \to c\bar{b}\tau$, $\tilde{\chi}_1^0 \to s\bar{b}\nu_i$ and their CP-conjugate states, induced by λ'_{323}, or $\tilde{\chi}_1^0 \to \mu\bar{\tau}\nu_\tau$, and $\tilde{\chi}_1^0 \to \bar{\mu}\tau\nu_\tau$, due to a sizable or moderate λ_{233} coupling. When two such couplings are simultaneously present and of the same size, the rates for various massless states are of the same order of magnitude for a wide range of parameter values, see for example figure 11. The same holds when λ'_{333} and λ_{233} are dominating over the other couplings, see figure 12.

On the other hand, one can envisage a situation in which the $\tilde{\chi}_1^0$ is substantially heavy, so that its decays into final states containing the t-quark are possible, as for example $t\bar{b}\tau$, $t\bar{s}\tau$, $t\bar{b}\mu$ and their CP conjugated states. Massive decay modes imply the virtual exchange of sfermions of third generations ($\tilde{t}$'s, $\tilde{b}$'s, and $\tilde{\tau}$'s) that can be relatively light if the left-right sfermion mixing is large. In this case, decays with the t-quark in the final state can become competitive. In general, for moderate/large left-right mixing and/or substantial Bino-Higgsino mixing, these decay modes can be large at not too large values of $\tan\beta$. Although in general subdominant, their rate can be comparable to those of massless modes, see figures 9, 11, and 12. When the t-quark does not decay leptonically, these decay modes of a pair of lightest neutralinos, give rise to the like-sign dilepton signal typical of R_p-violating models.

Further, an overall increase in the total decay width of $\tilde{\chi}_1^0$ is observed, if $\tilde{\chi}_1^0$ is of Wino type, as it may happen if gaugino mass unification is not imposed, or in anomaly-mediated supersymmetry-breaking scenarios [60], or in grand-unification models with an additional strong hypercolor group [61, 62].

Given the size of the widths obtained in this analysis, in general, fast decays of the lightest neutralino are expected. If large, the couplings considered here will undoubtedly play an important role in future collider searches. Their impact will be significant also for production mechanisms and decays of squarks and sleptons, as well as chargino decays. Such studies will be the subject of future work [30].

Acknowledgments

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[4]The FORTRAN coded expressions for the amplitudes squared are available upon request.

A. Neutralino mass matrix

In the limit of small k_i in eq. (1.1) and small *vev*'s for the neutral component of $\tilde{L}_i$, the neutralino mass matrix reduces to the ordinary 4×4 matrix of R_p-conserving models. Thus, on the basis $(\tilde{B}, \tilde{W}_3, \tilde{H}_d, \tilde{H}_u,)$ the neutralino mass matrix has the form:

$$M_{neutr} = \begin{pmatrix} M_{\tilde{B}} & 0 & -m_Z s_W c_\beta & m_Z s_W s_\beta \\ 0 & M_{\tilde{W}} & m_Z c_W c_\beta & -m_Z c_W s_\beta \\ -m_Z s_W c_\beta & m_Z c_W c_\beta & 0 & -\mu \\ m_Z s_W s_\beta & -m_Z c_W s_\beta & -\mu & 0 \end{pmatrix}. \tag{A.1}$$

Notice that his matrix is symmetric, but in general not hermitian, if $M_{\tilde{B}}$, $M_{\tilde{W}}$ and μ have phases. Therefore, the mass eigenvalues are, in general, complex. This inconvenience can be avoided by a rotation of the corresponding eigenstates. Thus, mass eigenstates with real and positive eigenvalues are defined as:

$$\tilde{\chi}_i^0 = \eta_i \sum_\alpha O_{i\alpha} \psi_\alpha \,, \tag{A.2}$$

where ψ_a are the states $\psi_\alpha = \tilde{B}, \tilde{W}_3, \tilde{H}_d, \tilde{H}_u$ and η_i are phase factors. If $M_{\tilde{B}}$, $M_{\tilde{W}}$ and μ are real parameters, the matrix M_{neutr} is orthogonal and the mass eigenvalues are real, with positive and/or negative sign. If needed, positive mass eigenvalues are obtained still through a rotation, as before, and the factors η_i are $\pm i$ or ± 1.

B. Scalar superpartner mass and mixing

The 2×2 squark or slepton mass squared matrix,

$$\mathcal{M}_f^2 = \begin{pmatrix} m^2_{f\,RR} & (m^2_{f\,LR}) \\ (m^2_{f\,LR})^* & m^2_{f\,LL} \end{pmatrix}, \tag{B.1}$$

written here in the basis $\{\widetilde{f}_L, \widetilde{f}_R\}$, is hermitean $\mathcal{M}_f^2 = (\mathcal{M}_f^2)^\dagger$, with real eigenvalues

$$m^2_{\widetilde{f}_{1,2}} = \frac{1}{2}\left\{ (m^2_{f\,LL} + m^2_{f\,RR}) \mp \sqrt{\left(m^2_{f\,LL} - m^2_{f\,RR}\right)^2 + 4|m^2_{f\,LR}|^2} \right\}. \tag{B.2}$$

The eigenvectors $\widetilde{f}_1$ and $\widetilde{f}_2$ corresponding to the eigenvalues in (B.2) are obtained from $\widetilde{f}_L$, $\widetilde{f}_R$ through a unitary transformation :

$$\begin{pmatrix} \widetilde{f}_1 \\ \widetilde{f}_2 \end{pmatrix} = U_f \begin{pmatrix} \widetilde{f}_R \\ \widetilde{f}_L \end{pmatrix} \equiv \begin{pmatrix} \cos\theta_f & \sin\theta_f e^{i\phi_f} \\ -\sin\theta_f e^{-i\phi_f} & \cos\theta_f \end{pmatrix} \begin{pmatrix} \widetilde{f}_R \\ \widetilde{f}_L \end{pmatrix}, \tag{B.3}$$

where $\phi_f = \mathrm{Arg}(m^2_{LR})$. The mixing angle θ_f is defined, up to a two-fold ambiguity, by the relations

$$\sin 2\theta_f = -\frac{2m^2_{f\,LR}}{m^2_{\widetilde{f}_2} - m^2_{\widetilde{f}_1}}\,; \qquad \cos 2\theta_f = \frac{m^2_{f\,LL} - m^2_{f\,RR}}{m^2_{\widetilde{f}_2} - m^2_{\widetilde{f}_1}}\,. \tag{B.4}$$

Notice that the entries in the sfermion mass matrix in eq. (B.1) contain in general terms proportional to R_p-violating couplings. However, in the limit of small k_i in eq. (1.1) and small *vev*'s for the neutral component of $\tilde{L}_i$, $m^2_{f\,RR}$, $m^2_{f\,LL}$, and $(m^2_{f\,LR})$ reduce to the usual entries present in R_p-conserving models.

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C. Sfermion widths

If sufficiently light, the sfermions mediating the neutralino decays discussed in sections 4 and 5 may be on shell. The form of the propagators to use in this case is given in eq. (4.9).

C.1 Only $\lambda'_{333} \neq 0$

We consider first the case in which only the coupling λ'_{333} is nonvanishing. In this approximations, there is only one possible decay mode for the lightest $\tilde{\tau}$ eigenstate $\tilde{\tau}_1$, i.e. that induced by the coupling λ'_{333}: $\tilde{\tau}_1 \to \bar{t}b$. (Decay modes mediated by gauge interactions, such as $\tilde{\tau}_1 \to \tau\tilde{\chi}_1^0$ and $\tilde{\tau}_1 \to \nu_\tau\tilde{\chi}_1^-$, with off-shell $\tilde{\chi}_1^0$ and $\tilde{\chi}_1^-$ would give rise to subdominant 4-body decays.) Thus, it is:

$$\Gamma_{\tilde{\tau}_1}(\lambda'_{333}) = \Gamma(\tilde{\tau}_1 \to \bar{t}b) = \frac{3}{16\pi}|\lambda'_{333}|^2 m_{\tilde{\tau}_1} \sin^2\theta_\tau \left(1 - \frac{m_b^2}{m_{\tilde{\tau}_1}^2} - \frac{m_t^2}{m_{\tilde{\tau}_1}^2}\right) K^{1/2}\left(1, \frac{m_b}{m_{\tilde{\tau}_1}}, \frac{m_t}{m_{\tilde{\tau}_1}}\right), \tag{C.1}$$

where K is the Källen function $K(x,y,z) = ((x^2-y^2-z^2)^2 - 4y^2z^2)$.

Similarly, the width for the decay mode of a third generation sneutrino, $\tilde{\nu}_\tau \to \bar{b}b$, is given by:

$$\Gamma_{\tilde{\nu}_\tau}(\lambda'_{333}) = \Gamma\left(\tilde{\nu}_\tau \to \overline{b_L}\, b_R\right) + \Gamma\left(\tilde{\nu}_\tau \to \overline{b_R}\, b_L\right) = 2\,\Gamma\left(\tilde{\nu}_\tau \to \overline{b_L}\, b_R\right), \tag{C.2}$$

with

$$\Gamma(\tilde{\nu}_\tau \to \overline{b_L}\, b_R) = \frac{3}{16\pi}|\lambda'_{333}|^2\, m_{\tilde{\nu}_\tau}\left(1 - 2\frac{m_b^2}{m_{\tilde{\nu}_\tau}^2}\right) K^{1/2}\left(1, \frac{m_b}{m_{\tilde{\nu}_\tau}}, \frac{m_b}{m_{\tilde{\nu}_\tau}}\right). \tag{C.3}$$

The width for the sbottom squark is obtained as sum of the widths for the two decay modes $\tilde{b}_1 \to b\nu_\tau$ and $\tilde{b}_1 \to t\bar{\tau}$, with the second one present only in the case $m_{\tilde{\chi}_1^0} > m_{\tilde{b}_1} > m_t$:

$$\Gamma_{\tilde{b}_1}(\lambda'_{333}) = \Gamma\left(\tilde{b}_1 \to b\nu_\tau\right) + \Gamma\left(\tilde{b}_1 \to t\bar{\tau}\right). \tag{C.4}$$

The two partial widths are respectively:

$$\begin{aligned}
\Gamma\left(\tilde{b}_1 \to b\nu_\tau\right) &= \frac{1}{16\pi}|\lambda'_{333}|^2\, m_{\tilde{b}_1}\left(1 - \frac{m_b^2}{m_{\tilde{b}_1}^2}\right) K^{1/2}\left(1, \frac{m_b}{m_{\tilde{b}_1}}, 0\right), \\
\Gamma\left(\tilde{b}_1 \to t\bar{\tau}\right) &= \frac{1}{16\pi}|\lambda'_{333}|^2\, m_{\tilde{b}_1}\cos^2\theta_b\left(1 - \frac{m_t^2}{m_{\tilde{b}_1}^2} - \frac{m_\tau^2}{m_{\tilde{b}_1}^2}\right) K^{1/2}\left(1, \frac{m_t}{m_{\tilde{b}_1}}, \frac{m_\tau}{m_{\tilde{b}_1}}\right),
\end{aligned} \tag{C.5}$$

where the first is the width for the two decays $\tilde{b}_1 \to b_L\nu_\tau$ and $\tilde{b}_1 \to b_R\bar{n}u_\tau$.

Finally, the width for the lightest $\tilde{t}$ eigenstate, $\tilde{t}_1$:

$$\Gamma_{\tilde{t}_1}(\lambda'_{333}) = \Gamma(\tilde{t}_1 \to \bar{\tau}b) = \frac{1}{16\pi}|\lambda'_{333}|^2\, m_{\tilde{t}_1}\sin^2\theta_t\left(1 - \frac{m_b^2}{m_{\tilde{t}_1}^2} - \frac{m_\tau^2}{m_{\tilde{t}_1}^2}\right) K^{1/2}\left(1, \frac{m_b}{m_{\tilde{t}_1}}, \frac{m_\tau}{m_{\tilde{t}_1}}\right), \tag{C.6}$$

is the width of the decay $\tilde{t}_1 \to \bar{\tau}b$.

Except for the sneutrino case, for which there is only one eigenstate per generation, the widths given above refer to the lightest $\tilde{\tau}$, $\tilde{b}$, and $\tilde{t}$ eigenstates. In the less likely case that any of the heaviest of these sfermion eigenstates, i.e. $\tilde{\tau}_2$, $\tilde{b}_2$, and $\tilde{t}_2$ is also lighter than $\tilde{\chi}_1^0$, the relevant widths can be obtained by interchanging $\cos\theta_f$ with $\sin\theta_f$ in the above formulas.

C.2 λ'_{333} and $\lambda'_{323} \neq 0$

Generalization to the case in which other λ' couplings are nonvanishing, such as λ'_{233}, λ'_{323}, λ'_{332}, etc., are straightforward. In particular, when two such couplings are present, for example λ'_{333} and λ'_{323}, the $\tilde{\tau}$-slepton may decay also into light particles, as $\tilde{\tau} \to \bar{c}b$. The corresponding with is:

$$\Gamma_{\tilde{\tau}_1}(\lambda'_{323}) = \Gamma(\tilde{\tau}_1 \to \bar{c}b) = \frac{3}{16\pi}|\lambda'_{323}|^2\, m_{\tilde{\tau}_1} \sin^2\theta_\tau \left(1 - \frac{m_b^2}{m_{\tilde{\tau}_1}^2} - \frac{m_c^2}{m_{\tilde{\tau}_1}^2}\right) K^{1/2}\left(1, \frac{m_b}{m_{\tilde{\tau}_1}}, \frac{m_c}{m_{\tilde{\tau}_1}}\right), \tag{C.7}$$

and the total with is given by the sum of $\Gamma_{\tilde{\tau}_1}(\lambda'_{323})$ and $\Gamma_{\tilde{\tau}_1}(\lambda'_{333})$.

The sneutrino $\tilde{\nu}_\tau$ can now decay as $\tilde{\nu}_\tau \to \overline{s}b$ and $\tilde{\nu}_\tau \to \overline{b}s$, with the b-quark always right-handed. These decays have identical partial widths,

$$\Gamma(\tilde{\nu}_\tau \to \overline{s}\, b) = \Gamma(\tilde{\nu}_\tau \to \overline{b}\, s) = \frac{3}{16\pi}|\lambda'_{323}|^2 m_{\tilde{\nu}_\tau} \left(1 - \frac{m_s^2}{m_{\tilde{\nu}_\tau}^2} - \frac{m_b^2}{m_{\tilde{\nu}_\tau}^2}\right) K^{1/2}\left(1, \frac{m_s}{m_{\tilde{\nu}_\tau}}, \frac{m_b}{m_{\tilde{\nu}_\tau}}\right). \tag{C.8}$$

The width due to the coupling λ'_{323}, is then given by

$$\Gamma_{\tilde{\nu}_\tau}(\lambda'_{323}) = \Gamma\left(\tilde{\nu}_\tau \to \overline{s}\, b\right) + \Gamma\left(\tilde{\nu}_\tau \to \overline{b}\, s\right) = 2\,\Gamma\left(\tilde{\nu}_\tau \to \overline{s}\, b\right). \tag{C.9}$$

The total width, when both couplings λ'_{333} and λ'_{323} are present, is obtained by summing $\Gamma_{\tilde{\nu}_\tau}(\lambda'_{333})$ and $\Gamma_{\tilde{\nu}_\tau}(\lambda'_{323})$.

Because of the coupling λ'_{323}, the $\tilde{b}$-squark can now decay into light particles, as $\tilde{b}_1 \to c\bar{\tau}$, and $\tilde{b}_1 \to s\nu_\tau$. The corresponding widths are:

$$\Gamma(\tilde{b}_1 \to c\bar{\tau}) = \frac{1}{16\pi}|\lambda'_{323}|^2\, m_{\tilde{b}_1} \cos^2\theta_b \left(1 - \frac{m_c^2}{m_{\tilde{b}_1}^2} - \frac{m_\tau^2}{m_{\tilde{b}_1}^2}\right) K^{1/2}\left(1, \frac{m_c}{m_{\tilde{b}_1}}, \frac{m_\tau}{m_{\tilde{b}_1}}\right) \tag{C.10}$$

$$\Gamma(\tilde{b}_1 \to s\nu_\tau) = \frac{1}{16\pi}|\lambda'_{323}|^2\, m_{\tilde{b}_1} \cos^2\theta_b \left(1 - \frac{m_s^2}{m_{\tilde{b}_1}^2}\right) K^{1/2}\left(1, \frac{m_s}{m_{\tilde{b}_1}}, 0\right). \tag{C.11}$$

The total width $\Gamma_{\tilde{b}_1}$ is given by the sum of $\Gamma_{\tilde{b}_1}(\lambda'_{333})$ and $\Gamma_{\tilde{b}_1}(\lambda'_{323})$, where

$$\Gamma_{\tilde{b}_1}(\lambda'_{323}) = \Gamma\left(\tilde{b}_1 \to s\nu_\tau\right) + \Gamma\left(\tilde{b}_1 \to c\bar{\tau}\right). \tag{C.12}$$

Finally, the $\tilde{c}$- and $\tilde{s}$-squarks decay as $\tilde{c}_1 \to \bar{\tau}b$ and $\tilde{s}_1 \to b\nu_\tau$, respectively, with widths:

$$\Gamma_{\tilde{c}_1}(\lambda'_{323}) = \Gamma(\tilde{c}_1 \to \bar{\tau}b) = \frac{1}{16\pi}|\lambda'_{323}|^2\, m_{\tilde{c}_1} \sin^2\theta_c \left(1 - \frac{m_b^2}{m_{\tilde{c}_1}^2} - \frac{m_\tau^2}{m_{\tilde{c}_1}^2}\right) K^{1/2}\left(1, \frac{m_b}{m_{\tilde{c}_1}}, \frac{m_\tau}{m_{\tilde{c}_1}}\right), \tag{C.13}$$

$$\Gamma_{\tilde{s}_1}(\lambda'_{323}) = \Gamma(\tilde{s}_1 \to b\nu_\tau) = \frac{1}{16\pi}|\lambda'_{323}|^2\, m_{\tilde{s}_1} \sin^2\theta_s \left(1 - \frac{m_b^2}{m_{\tilde{s}_1}^2}\right) K^{1/2}\left(1, \frac{m_b}{m_{\tilde{s}_1}}, 0\right) \tag{C.14}$$

C.3 λ'_{333} and $\lambda_{233} \neq 0$

If besides λ'_{333}, also λ_{233} is nonvanishing, then, $\tilde{\tau}_1$ can also decay into $\tau\nu_\mu$ and $\mu\nu_\tau$, with corresponding widths:

$$\begin{aligned}
\Gamma(\tilde{\tau}_1 \to \tau\nu_\mu) &= \frac{1}{16\pi}|\lambda_{233}|^2\, m_{\tilde{\tau}_1}\left(1-\frac{m_\tau^2}{m_{\tilde{\tau}_1}^2}\right)K^{1/2}\left(1,\frac{m_\tau}{m_{\tilde{\tau}_1}},0\right), \\
\Gamma(\tilde{\tau}_1 \to \bar{\nu}_\tau\mu) &= \frac{1}{16\pi}|\lambda_{233}|^2\, m_{\tilde{\tau}_1}\sin^2\theta_\tau\left(1-\frac{m_\mu^2}{m_{\tilde{\tau}_1}^2}\right)K^{1/2}\left(1,\frac{m_\mu}{m_{\tilde{\tau}_1}},0\right),
\end{aligned} \tag{C.15}$$

where the first is the width for $\tilde{\tau}_1 \to \tau_L\nu_\mu$ and $\tilde{\tau}_1 \to \tau_R n\bar{u}_\mu$. After defining:

$$\Gamma_{\tilde{\tau}_1}(\lambda_{233}) = \Gamma\left(\tilde{\tau}_1 \to \tau\nu_\mu\right) + \Gamma\left(\tilde{\tau}_1 \to \tau\nu_\tau\right). \tag{C.16}$$

the total width for τ_1, $\Gamma_{\tilde{\tau}_1}$, is obtained summing up $\Gamma_{\tilde{\tau}_1}(\lambda_{233})$ to the width $\Gamma_{\tilde{t}_1}(\lambda'_{333})$ in eq. (C.1).

Similarly, $\tilde{\nu}_\tau$ can decay into $\bar{\mu}\tau$ and $\bar{\tau}\mu$:

$$\Gamma(\tilde{\nu}_\tau \to \bar{\mu}\tau) = \Gamma(\tilde{\nu}_\tau \to \bar{\tau}\mu) = \frac{1}{16\pi}|\lambda'_{333}|^2\, m_{\tilde{\nu}_\tau}\left(1-\frac{m_\tau^2}{m_{\tilde{\nu}_\tau}^2}-\frac{m_\mu^2}{m_{\tilde{\nu}_\tau}^2}\right)K^{1/2}\left(1,\frac{m_\tau}{m_{\tilde{\nu}_\tau}},\frac{m_\mu}{m_{\tilde{\nu}_\tau}}\right), \tag{C.17}$$

and

$$\Gamma_{\tilde{\nu}_\tau}(\lambda_{233}) = \Gamma(\tilde{\nu}_\tau \to \bar{\mu}\tau) + \Gamma(\tilde{\nu}_\tau \to \bar{\tau}\mu). \tag{C.18}$$

Also in this case, the total width $\Gamma_{\tilde{\nu}_\tau}$ is obtained summing up $\Gamma_{\tilde{\nu}_\tau}(\lambda_{233})$ and $\Gamma_{\tilde{\nu}_\tau}(\lambda'_{333})$.

Finally, for $\lambda_{233} \neq 0$, $\tilde{\mu}$ and $\tilde{\nu}_\mu$ may also be produced on shell in the decays $\tilde{\chi}_1^0 \to \mu\bar{\tau}\nu_\tau$ and $\tilde{\chi}_1^0 \to \nu_\mu\bar{\tau}\tau$. The width of the $\tilde{\mu}$ slepton is then:

$$\Gamma_{\tilde{\mu}_1}(\lambda_{233}) = \Gamma(\tilde{\mu} \to \tau\nu_\tau) = \frac{1}{16\pi}|\lambda'_{333}|^2\, m_{\tilde{\mu}}\sin^2\theta_\mu\left(1-\frac{m_\tau^2}{m_{\tilde{\mu}}^2}\right)K^{1/2}\left(1,\frac{m_\tau}{m_{\tilde{\mu}}},0\right), \tag{C.19}$$

that for the sneutrino $\tilde{\nu}_\mu$ is

$$\Gamma_{\tilde{\nu}_\mu}(\lambda_{233}) = \Gamma\left(\tilde{\nu}_\mu \to \overline{\tau_L}\,\tau_R\right) + \Gamma\left(\tilde{\nu}_\mu \to \overline{\tau_R}\,\tau_L\right) = 2\,\Gamma\left(\tilde{\nu}_\mu \to \overline{\tau_L}\,\tau_R\right), \tag{C.20}$$

with

$$\Gamma(\tilde{\nu}_\mu \to \overline{\tau_L}\,\tau_R) = \frac{1}{16\pi}|\lambda'_{333}|^2\, m_{\tilde{\nu}_\tau}\left(1-2\frac{m_\tau^2}{m_{\tilde{\nu}_\tau}^2}\right)K^{1/2}\left(1,\frac{m_\tau}{m_{\tilde{\nu}_\tau}},\frac{m_\tau}{m_{\tilde{\nu}_\tau}}\right). \tag{C.21}$$

D. Spin sum of matrix elements squared

We list here the products of all possible Q terms relevant for the decays $\tilde{\chi}_1^0 \to b\bar{b}\nu$ and $\tilde{\chi}_1^0 \to t\bar{b}\tau$ summed over all spin configurations. These products are evaluated under the assumption that the particle with momentum p_1 has a mass m_1, such that the ratio

$$r = \frac{m_1^2}{m_{\tilde{\chi}_1^0}^2} \tag{D.1}$$

is nonnegligible, as in the decay $\tilde{\chi}_1^0 \to t\bar{b}\tau$. When dealing with decays into a massless final state, as $\tilde{\chi}_1^0 \to b\bar{b}\nu$, it is sufficient to take the limit $r \to 0$ in the following expressions. For simplicity, the new symbol

$$\beta^{t,t'}_{s,s'} = \sum Q^t_{S,s} Q^{t'\,\dagger}_{S,s} \,. \tag{D.2}$$

is introduced, with $t, t' = x, y, z$ and where s, s' run over RR, RL, LR, and LL.

$$\begin{aligned}
\beta^{x,x}_{RR,RR} &= m^4_{\tilde{\chi}_1^0}\,(1+r-x)x \\
\beta^{y,y}_{RR,RR} &= m^4_{\tilde{\chi}_1^0}\,(1-y)(y-r) \\
\beta^{z,z}_{RR,RR} &= m^4_{\tilde{\chi}_1^0}\,(1-z)(z-r) \\
\beta^{x,x}_{RL,RL} &= m^4_{\tilde{\chi}_1^0}\,(1+r-x)x \\
\beta^{y,y}_{RL,RL} &= m^4_{\tilde{\chi}_1^0}\,(1-y)(y-r) \\
\beta^{z,z}_{RL,RL} &= m^4_{\tilde{\chi}_1^0}\,(1-z)(z-r)\,,
\end{aligned} \tag{D.3}$$

where x, y, and z are defined in eq. (4.6). The same results are obtained when the chirality indices in the above terms are exchanged, i.e.

$$\begin{aligned}
\beta^{t,t}_{LL,LL} &= \beta^{t,t}_{RR,RR} \\
\beta^{t,t}_{LR,LR} &= \beta^{t,t}_{RL,RL}\,.
\end{aligned} \tag{D.4}$$

Mixed products such as:

$$\begin{aligned}
\beta^{x,y}_{RR,RR} &= \frac{1}{2} m^4_{\tilde{\chi}_1^0}\,[(1+r-x)x + (1-y)(y-r) - (1-z)(z-r)] \\
\beta^{x,z}_{RR,RR} &= \frac{1}{2} m^4_{\tilde{\chi}_1^0}\,[(1+r-x)x - (1-y)(y-r) + (1-z)(z-r)] \\
\beta^{y,z}_{RR,RR} &= \frac{1}{2} m^4_{\tilde{\chi}_1^0}\,[(1+r-x)x - (1-y)(y-r) - (1-z)(z-r)] \\
\beta^{x,y}_{RL,RL} &= 0 \\
\beta^{x,z}_{RL,RL} &= 0 \\
\beta^{y,z}_{RL,RL} &= 0\,,
\end{aligned} \tag{D.5}$$

are invariant when chirality indices are exchanged, $R \leftrightarrow L$, as they are the products

$$\begin{aligned}
\beta^{x,x}_{RR,RL} &= 2m^4_{\tilde{\chi}_1^0}\,\sqrt{r}\,x \\
\beta^{y,y}_{RR,RL} &= 0 \\
\beta^{z,z}_{RR,RL} &= 0\,.
\end{aligned} \tag{D.6}$$

In addition, those in eq. (D.5) are symmetric under the exchange $t \leftrightarrow t'$. Finally, it is:

$$\begin{aligned}
\beta^{x,y}_{RR,RL} &= 0 \\
\beta^{x,z}_{RR,RL} &= 0 \\
\beta^{y,x}_{RR,RL} &= m^4_{\tilde{\chi}_1^0}\,\sqrt{r}\,x \\
\beta^{y,z}_{RR,RL} &= 0 \\
\beta^{z,x}_{RR,RL} &= m^4_{\tilde{\chi}_1^0}\,\sqrt{r}\,x \\
\beta^{z,y}_{RR,RL} &= 0
\end{aligned} \tag{D.7}$$

where these products are invariant under the exchange $R \leftrightarrow L$ and are symmetric under the simultaneous exchange $t \leftrightarrow t'$ and $s \leftrightarrow s'$,

$$\beta^{t,t'}_{s,s'} = \beta^{t',t}_{s',s}\,. \tag{D.8}$$

Mixed products of terms with s in the subset of indices RR, RL, and s' in LL, LR, or vice versa, vanish identically, except for the two terms

$$\beta^{y,z}_{LR,RL} = \beta^{y,z}_{RL,LR} = m^4_{\tilde{\chi}^0_1}\sqrt{r}\,x\,. \tag{D.9}$$

References

[1] P. Fayet, *Spontaneously broken supersymmetric theories of weak, electromagnetic and strong interactions*, *Phys. Lett.* **B 69** (1977) 489.

[2] G.R. Farrar and P. Fayet, *Phenomenology of the production, decay and detection of new hadronic states associated with supersymmetry*, *Phys. Lett.* **B 76** (1978) 575.

[3] B.A. Campbell, S. Davidson, J.R. Ellis and K.A. Olive, *Cosmological baryon asymmetry constraints on extensions of the standard model*, *Phys. Lett.* **B 256** (1991) 457;
W. Fischler, G.F. Giudice, R.G. Leigh and S. Paban, *Constraints on the baryogenesis scale from neutrino masses*, *Phys. Lett.* **B 258** (1991) 45.

[4] E. Ma, M. Raidal and U. Sarkar, *Neutrino masses in supersymmetry: R-parity and leptogenesis*, *Phys. Lett.* **B 460** (1999) 359 [hep-ph/9901406];
U. Sarkar, *Models of neutrino masses and baryogenesis*, *Pramana* **54** (2000) 101 [hep-ph/9906335].

[5] H.K. Dreiner and G.G. Ross, *Sphaleron erasure of primordial baryogenesis*, *Nucl. Phys.* **B 410** (1993) 188 [hep-ph/9207221].

[6] L.E. Ibáñez and G.G. Ross, *Discrete gauge symmetries and the origin of baryon and lepton number conservation in supersymmetric versions of the standard model*, *Nucl. Phys.* **B 368** (1992) 3.

[7] See for example: H.E. Haber and G.L. Kane, *The search for supersymmetry: probing physics beyond the standard model*, *Phys. Rept.* **117** (1985) 75;
K. Hikasa, Lecture Notes *Supersymmetric standard model for collider physicists*, (1996) (unpublished);
MSSM Working Group collaboration, *The minimal supersymmetric standard model: group summary report*, hep-ph/9901246;
SUGRA Working Group collaboration, *Report of the SUGRA working group for run II of the Tevatron*, hep-ph/0003154;
N. Polonsky, *Supersymmetry: structure and phenomena. extensions of the standard model*, *Lect. Notes Phys.* **M68** (2001) 1 [hep-ph/0108236].

[8] Particle Data Group collaboration, D.E. Groom et al., *Review of particle physics*, *Eur. Phys. J.* **C 15** (2000) 1.

[9] L.J. Hall and M. Suzuki, *Explicit R parity breaking in supersymmetric models*, *Nucl. Phys.* **B 231** (1984) 419.

JHEP07(2002)037

[10] A.S. Joshipura and M. Nowakowski, *'Just so' oscillations in supersymmetric standard model*, *Phys. Rev.* **D 51** (1995) 2421 [hep-ph/9408224];
T. Banks, Y. Grossman, E. Nardi and Y. Nir, *Supersymmetry without R-parity and without lepton number*, *Phys. Rev.* **D 52** (1995) 5319 [hep-ph/9505248].

[11] There exists a large body of work on the possibility of having neutrino masses induced by R_p-violating interactions. Only some representative papers are listed here and in refs. [13] and [14]: R. Hempfling, *Neutrino masses and mixing angles in SUSY-GUT theories with explicit R-parity breaking*, *Nucl. Phys.* **B 478** (1996) 3 [hep-ph/9511288];
F.M. Borzumati, Y. Grossman, E. Nardi and Y. Nir, *Neutrino masses and mixing in supersymmetric models without R parity*, *Phys. Lett.* **B 384** (1996) 123 [hep-ph/9606251];
H.-P. Nilles and N. Polonsky, *Supersymmetric neutrino masses, R-symmetries and the generalized mu problem*, *Nucl. Phys.* **B 484** (1997) 33 [hep-ph/9606388];
M.A. Diaz, J.C. Romao and J.W.F. Valle, *Minimal supergravity with R-parity breaking*, *Nucl. Phys.* **B 524** (1998) 23 [hep-ph/9706315].

[12] F.M. Borzumati, G.R. Farrar, N. Polonsky and S. Thomas, *Soft Yukawa couplings in supersymmetric theories*, *Nucl. Phys.* **B 555** (1999) 53 [hep-ph/9902443]; *Precision measurements at the Higgs resonance: a probe of radiative fermion masses*, hep-ph/9712428; and *Fermion masses without Yukawa couplings*, hep-ph/9805314.

[13] E.J. Chun and S.K. Kang, *One-loop corrected neutrino masses and mixing in supersymmetric standard model without R-parity*, *Phys. Rev.* **D 61** (2000) 075012 [hep-ph/9909429];
Y. Grossman and H.E. Haber, *Basis-independent analysis of the sneutrino sector in R-parity violating supersymmetry*, *Phys. Rev.* **D 63** (2001) 075011 [hep-ph/0005276];
S. Davidson and M. Losada, *Neutrino masses in the R(p) violating MSSM*, *J. High Energy Phys.* **05** (2000) 021 [hep-ph/0005080]; and *Basis independent neutrino masses in the R(p) violating MSSM*, *Phys. Rev.* **D 65** (2002) 075025 [hep-ph/0010325];

[14] M. Drees, S. Pakvasa, X. Tata and T. ter Veldhuis, *A supersymmetric resolution of solar and atmospheric neutrino puzzles*, *Phys. Rev.* **D 57** (1998) 5335 [hep-ph/9712392];
B. Mukhopadhyaya, S. Roy and F. Vissani, *Correlation between neutrino oscillations and collider signals of supersymmetry in an R-parity violating model*, *Phys. Lett.* **B 443** (1998) 191 [hep-ph/9808265];
E.J. Chun, S.K. Kang, C.W. Kim and U.W. Lee, *Supersymmetric neutrino masses and mixing with R-parity violation*, *Nucl. Phys.* **B 544** (1999) 89 [hep-ph/9807327];
R. Adhikari and G. Omanovic, *LSND, solar and atmospheric neutrino oscillation experiments and R-parity violating supersymmetry*, *Phys. Rev.* **D 59** (1999) 073003;
O.C.W. Kong, *Neutrino oscillations and flavor structure of supersymmetry without R-parity*, *Mod. Phys. Lett.* **A 14** (1999) 903 [hep-ph/9808304];
S. Rakshit, G. Bhattacharyya and A. Raychaudhuri, *R-parity-violating trilinear couplings and recent neutrino data*, *Phys. Rev.* **D 59** (1999) 091701 [hep-ph/9811500];
S.Y. Choi, E.J. Chun, S.K. Kang and J.S. Lee, *Neutrino oscillations and R-parity violating collider signals*, *Phys. Rev.* **D 60** (1999) 075002 [hep-ph/9903465];
D.E. Kaplan and A.E. Nelson, *Solar and atmospheric neutrino oscillations from bilinear R-parity violation*, *J. High Energy Phys.* **01** (2000) 033 [hep-ph/9901254];
A.S. Joshipura and S.K. Vempati, *Sneutrino vacuum expectation values and neutrino anomalies through trilinear R-parity violation*, *Phys. Rev.* **D 60** (1999) 111303 [hep-ph/9903435];

G. Bhattacharyya, H.V. Klapdor-Kleingrothaus and H. Pas, *Neutrino mass and magnetic moment in supersymmetry without R-parity in the light of recent data*, *Phys. Lett.* **B 463** (1999) 77 [hep-ph/9907432];
A. Abada and M. Losada, *Constraints on both bilinear and trilinear R-parity violating couplings from neutrino laboratories and astrophysics data*, *Phys. Lett.* **B 492** (2000) 310 [hep-ph/0007041];
O. Haug, J. D. Vergados, A. Faessler, and S. Kovalenko, *Neutrino oscillations and consequences for coupling constants λ'_{i33} of the R(p)-minimal supersymmetric standard model*, *Prog. Part. Nucl. Phys.* **44** (2000) 467;
F. Takayama and M. Yamaguchi, *Pattern of neutrino oscillations in supersymmetry with bilinear R-parity violation*, *Phys. Lett.* **B 476** (2000) 116 [hep-ph/9910320].

[15] S. Dawson, *R parity breaking in supersymmetric theories*, *Nucl. Phys.* **B 261** (1985) 297.

[16] R.M. Godbole, P. Roy and X. Tata, *Tau signals of R-parity breaking at LEP-200*, *Nucl. Phys.* **B 401** (1993) 67 [hep-ph/9209251].

[17] F. Borzumati, J.-L. Kneur and N. Polonsky, *Higgs- and slepton-strahlung at hadron colliders*, *Phys. Rev.* **D 60** (1999) 115011 [hep-ph/9905443].

[18] A.G. Akeroyd, C. Liu and J. Song, *Stau LSP and comparison with H^{+-} phenomenology*, *Phys. Rev.* **D 65** (2002) 015008 [hep-ph/0107218].

[19] B. de Carlos and P.L. White, *R-parity violation and quark flavour violation*, *Phys. Rev.* **D 55** (1997) 4222 [hep-ph/9609443].

[20] E.J. Chun, K. Hwang and J.S. Lee, *CP asymmetries in radiative B decays with R-parity violation*, *Phys. Rev.* **D 62** (2000) 076006 [hep-ph/0005013];
T. Besmer and A. Steffen, *R-parity violation and the decay $b \to s\gamma$*, *Phys. Rev.* **D 63** (2001) 055007 [hep-ph/0004067];
D. Chakraverty and D. Choudhury, *$b \to s\gamma$ confronts B-violating scalar couplings: R-parity violating supersymmetry or diquarks*, *Phys. Rev.* **D 63** (2001) 075009 [hep-ph/0008165].

[21] S. Bertolini, F. Borzumati, A. Masiero and G. Ridolfi, *Effects of supergravity induced electroweak breaking on rare B decays and mixings*, *Nucl. Phys.* **B 353** (1991) 591;
F.M. Borzumati, *The decay $b \to s\gamma$ in the MSSM revisited*, *Z. Physik* **C 63** (1994) 291 [hep-ph/9310212];
F.M. Borzumati, *On the minimal messenger model*, hep-ph/9702307;
T. Goto and Y. Okada, *Charged Higgs mass bound from the $b \to s\gamma$ process in the minimal supergravity model*, *Prog. Theor. Phys.* **94** (1995) 407 [hep-ph/9412225];
F.M. Borzumati, M. Olechowski and S. Pokorski, *Constraints on the minimal SUSY* SO*(10) model from cosmology and the $b \to s\gamma$ decay*, *Phys. Lett.* **B 349** (1995) 311 [hep-ph/9412379];
M. Ciuchini, G. Degrassi, P. Gambino and G.F. Giudice, *Next-to-leading QCD corrections to $B \to X_s\gamma$: standard model and two-higgs doublet model*, *Nucl. Phys.* **B 527** (1998) 21 [hep-ph/9710335]; and *Next-to-leading QCD corrections to $B \to X_s\gamma$ in supersymmetry*, *Nucl. Phys.* **B 534** (1998) 3 [hep-ph/9806308];
F.M. Borzumati and C. Greub, *2HDMs predictions for $\bar{B} \to X_s\gamma$ in NLO QCD*, *Phys. Rev.* **D 58** (1998) 074004 [hep-ph/9802391]; *Two Higgs doublet model predictions for $\bar{B} \to X_s\gamma$ in NLO QCD. (Addendum)*, *Phys. Rev.* **D 59** (1999) 057501 [hep-ph/9809438]; and *Lessons from $\bar{B} \to X_s\gamma$ in two Higgs doublet models*, hep-ph/9810240.

[22] See for example: A. Djouadi et al., *The Higgs working group: summary report*, hep-ph/0002258;

JHEP07(2002)037

M. Carena et al., *Report of the Tevatron Higgs working group*, hep-ph/0010338.

[23] A. Bartl, W. Majerotto and W. Porod, *Large Higgs boson exchange contribution in three-body neutralino decays*, *Phys. Lett.* **B 465** (1999) 187 [hep-ph/9907377];
A. Djouadi, Y. Mambrini and M. Muhlleitner, *Chargino and neutralino decays revisited*, *Eur. Phys. J.* **C 20** (2001) 563 [hep-ph/0104115].

[24] H. Baer, C.-h. Chen, M. Drees, F. Paige and X. Tata, *Supersymmetry reach of Tevatron upgrades: the large* $\tan\beta$ *case*, *Phys. Rev.* **D 58** (1998) 075008 [hep-ph/9802441].

[25] J. Butterworth and H.K. Dreiner, *R-parity violation at HERA*, *Nucl. Phys.* **B 397** (1993) 3 [hep-ph/9211204].

[26] H.K. Dreiner and P. Morawitz, *Signals for supersymmetry at HERA*, *Nucl. Phys.* **B 428** (1994) 31 [hep-ph/9405253].

[27] E.A. Baltz and P. Gondolo, *Neutralino decay rates with explicit R-parity violation*, *Phys. Rev.* **D 57** (1998) 2969 [hep-ph/9709445].

[28] D.K. Ghosh, R.M. Godbole and S. Raychaudhuri, *Signals for R-parity-violating supersymmetry at a 500 GeV* e^+e^- *collider*, hep-ph/9904233.

[29] H.K. Dreiner, P. Richardson and M.H. Seymour, *Parton-shower simulations of R-parity violating supersymmetric models*, *J. High Energy Phys.* **04** (2000) 008 [hep-ph/9912407].

[30] F. Borzumati et al., in preparation.

[31] V.D. Barger, G.F. Giudice and T. Han, *Some new aspects of supersymmetry R parity violating interactions*, *Phys. Rev.* **D 40** (1989) 2987.

[32] R.M. Godbole, S. Pakvasa, S.D. Rindani and X. Tata, *Fermion dipole moments in supersymmetric models with explicitly broken R-parity*, *Phys. Rev.* **D 61** (2000) 113003 [hep-ph/9912315];
S.A. Abel, A. Dedes and H.K. Dreiner, *Dipole moments of the electron, neutrino and neutron in the MSSM without R-parity symmetry*, *J. High Energy Phys.* **05** (2000) 013 [hep-ph/9912429].

[33] J.E. Kim, B. Kyae and H.M. Lee, *Effective supersymmetric theory and* $(g-2)_\mu$ *with R parity violation*, *Phys. Lett.* **B 520** (2001) 298 [hep-ph/0103054].

[34] G. Bhattacharyya, J.R. Ellis and K. Sridhar, *Bounds on the masses and couplings of leptoquarks from leptonic partial widths of the z*, *Phys. Lett.* **B 336** (1994) 100 [hep-ph/9406354].

[35] O. Lebedev, W. Loinaz and T. Takeuchi, *Constraints on R-parity violating couplings from LEP/SLD hadronic observables*, *Phys. Rev.* **D 62** (2000) 015003 [hep-ph/9911479].

[36] O. Lebedev, W. Loinaz and T. Takeuchi, *Constraints on R-parity violating couplings from lepton universality*, *Phys. Rev.* **D 61** (2000) 115005 [hep-ph/9910435].

[37] J.L. Hewett and T.G. Rizzo, *Much ado about leptoquarks: a comprehensive analysis*, *Phys. Rev.* **D 56** (1997) 5709 [hep-ph/9703337].

[38] R.M. Godbole and D. Choudhury, *Probing* R_p*-violating couplings through indirect effects on Drell Yan production at the LHC*, IIsc-CTS/9/00. in S. Abdullin et al., *The SUSY working group: summary report*, hep-ph/0005142.

JHEP07(2002)037

[39] See for example: J. Kalinowski, R. Ruckl, H. Spiesberger and P.M. Zerwas, *Supersymmetry with R-parity breaking: contact interactions and resonance formation in leptonic processes at LEP2*, *Phys. Lett.* **B 406** (1997) 314 [hep-ph/9703436];
T.G. Rizzo, *Distinguishing indirect signatures arising from new physics at the NLC*, hep-ph/9907344.

[40] B. de Carlos and P.L. White, *R-parity violation effects through soft supersymmetry breaking terms and the renormalisation group*, *Phys. Rev.* **D 54** (1996) 3427 [hep-ph/9602381];
M. Chaichian and K. Huitu, *Constraints on R-parity violating interactions from $\mu \to e\gamma$*, *Phys. Lett.* **B 384** (1996) 157 [hep-ph/9603412];
M. Frank, *Bilinear R-parity violation parameters from $\mu \to e\gamma$*, *Phys. Rev.* **D 62** (2000) 015006.

[41] F. Ledroit and G. Sajot, Report # GRD-S-008 ISN, Grenoble,1998, also in http://qcd.th.u-psud.fr/GDR_SUSY/GDR_SUSY_PUBLIC/entete_note_publique

[42] B.C. Allanach, A. Dedes and H.K. Dreiner, *Bounds on R-parity violating couplings at the weak scale and at the SUSY scale*, *Phys. Rev.* **D 60** (1999) 075014 [hep-ph/9906209].

[43] B. Mukhopadhyaya, S. Roy and F. Vissani, *Correlation between neutrino oscillations and collider signals of supersymmetry in an R-parity violating model*, *Phys. Lett.* **B 443** (1998) 191 [hep-ph/9808265];
E.J. Chun and J.S. Lee, *Implication of Super-Kamiokande data on R-parity violation*, *Phys. Rev.* **D 60** (1999) 075006 [hep-ph/9811201];
S.Y. Choi, E.J. Chun, S.K. Kang and J.S. Lee, *Neutrino oscillations and R-parity violating collider signals*, *Phys. Rev.* **D 60** (1999) 075002 [hep-ph/9903465];
A. Datta, B. Mukhopadhyaya and F. Vissani, *Tevatron signatures of an R-parity violating supersymmetric theory*, *Phys. Lett.* **B 492** (2000) 324 [hep-ph/9910296];
W. Porod, M. Hirsch, J. Romao and J.W.F. Valle, *Testing neutrino mixing at future collider experiments*, *Phys. Rev.* **D 63** (2001) 115004 [hep-ph/0011248].

[44] G. Bhattacharyya, *A brief review of R-parity-violating couplings*, hep-ph/9709395;
R. Barbier et al., *Report of the group on the R-parity violation*, hep-ph/9810232.

[45] D.K. Ghosh, S. Raychaudhuri and K. Sridhar, *Bounds from $t\bar{t}$ production on R-parity violating models of supersymmetry*, *Phys. Lett.* **B 396** (1997) 177 [hep-ph/9608352].

[46] K.-i. Hikasa, J.M. Yang and B.-L. Young, *R-parity violation and top quark polarization at the Fermilab Tevatron collider*, *Phys. Rev.* **D 60** (1999) 114041 [hep-ph/9908231].

[47] T. Han and M.B. Magro, *Top quark decay via R-parity violating interactions at the Tevatron*, *Phys. Lett.* **B 476** (2000) 79 [hep-ph/9911442];
K.J. Abraham, K. Whisnant, J.M. Yang and B.-L. Young, *Probing R-violating top quark decays at hadron colliders*, *Phys. Rev.* **D 63** (2001) 034011 [hep-ph/0007280]; and *Probing R-violating top quark decays at hadron colliders*, *Phys. Rev.* **D 63** (2001) 034011 [hep-ph/0007280];
K.J. Abraham, K. Whisnant, J.-M. Yang and B.-L. Young, *Probing R-violating top quark decays at the NLC*, *Phys. Lett.* **B 514** (2001) 72 [hep-ph/0104184];
A. Belyaev, J.R. Ellis and S. Lola, *Three-body supersymmetric top decays*, *Phys. Lett.* **B 484** (2000) 79 [hep-ph/0002220].

[48] D. Choudhury, R.M. Godbole, S.D. Rindani and P. Poulose, *R-parity violation and top quark physics at hadronic colliders*, in preparation.

JHEP07(2002)037

[49] G. Moreau, E. Perez and G. Polesello, *Resonant sneutrino production in supersymmetry with R-parity violation at the LHC*, *Nucl. Phys.* **B 604** (2001) 3 [hep-ph/0003012];
H.K. Dreiner, P. Richardson and M.H. Seymour, *Resonant slepton production in hadron hadron collisions*, *Phys. Rev.* **D 63** (2001) 055008 [hep-ph/0007228].

[50] B. Allanach et al., *Searching for R-parity violation at Run-II of the Tevatron*, hep-ph/9906224;
J.L. Hewett and T.G. Rizzo, *Resonant slepton production at hadron colliders in R-parity violating models*, hep-ph/9809525.

[51] J. Erler, J.L. Feng and N. Polonsky, *A wide scalar neutrino resonance and b anti-b production at lep*, *Phys. Rev. Lett.* **78** (1997) 3063 [hep-ph/9612397].

[52] D0 collaboration, *Search for scalar leptoquark pairs decaying to electrons and jets in $\bar{p}p$ collisions*, *Phys. Rev. Lett.* **79** (1997) 4321 [hep-ex/9707033];
CDF collaboration, *Search for first generation leptoquark pair production in $p\bar{p}$ collisions at $\sqrt{s} = 1.8$ TeV*, *Phys. Rev. Lett.* **79** (1997) 4327 [hep-ex/9708017];
D0 collaboration, *Search for first generation scalar leptoquark pairs in $p\bar{p}$ collisions at $\sqrt{s} = 1.8$ TeV*, *Phys. Rev. Lett.* **80** (1998) 2051 [hep-ex/9710032].

[53] D0 collaboration, *Search for R-parity violating supersymmetry in the dielectron channel*, *Phys. Rev. Lett.* **83** (1999) 4476 [hep-ex/9907019]; *Search for second generation leptoquark pairs in $\bar{p}p$ collisions at $\sqrt{s} = 1.8$ TeV*, *Phys. Rev. Lett.* **84** (2000) 2088 [hep-ex/9910040]; and *Search for R-parity violation in multilepton final states in $p\bar{p}$ collisions at $\sqrt{s} = 1.8$ TeV*, *Phys. Rev.* **D 62** (2000) 071701 [hep-ex/0005034].

[54] H1 collaboration, *A search for squarks of R_p-violating SUSY at HERA*, *Z. Physik* **C 71** (1996) 211 [hep-ex/9604006]; *Searches at HERA for squarks in R-parity violating supersymmetry*, *Eur. Phys. J.* **C 20** (2001) 639 [hep-ex/0102050].

[55] H1 collaboration, *A search for leptoquark bosons in e^-p collisions at HERA*, *Phys. Lett.* **B 523** (2001) 234 [hep-ex/0107038].

[56] L3 collaboration, *Search for R-parity breaking sneutrino exchange at LEP*, *Phys. Lett.* **B 414** (1997) 373.

[57] OPAL collaboration, *Tests of the standard model and constraints on new physics from measurements of fermion pair production at 189 GeV at LEP*, *Eur. Phys. J.* **C 13** (2000) 553 [hep-ex/9908008].

[58] ALEPH collaboration, *Search for R parity violating decays of supersymmetric particles in e^+e^- collisions at centre-of-mass energies from 189 GeV to 202 GeV*, *Eur. Phys. J.* **C 19** (2001) 415 [hep-ex/0011008].

[59] For some of the recent analyses, see for example: OPAL collaboration, *Searches for R-parity violating decays of gauginos at 183 GeV at LEP*, *Eur. Phys. J.* **C 11** (1999) 619 [hep-ex/9901037]; and *Search for R-parity violating decays of scalar fermions at LEP*, *Eur. Phys. J.* **C 12** (2000) 1 [hep-ex/9904015];
ALEPH collaboration, *Search for R-parity violating decays of supersymmetric particles in e^+e^- collisions at centre-of-mass energies near 183 GeV*, *Eur. Phys. J.* **C 13** (2000) 29;
DELPHI collaboration, *Search for SUSY with R-parity violating $LL\bar{E}$ couplings at $\sqrt{s} = 189$ GeV*, *Phys. Lett.* **B 487** (2000) 36 [hep-ex/0103006]; and *Search for R-parity violation with a u-bar d-bar d-bar coupling at $\sqrt{s} = 189$ GeV*, *Phys. Lett.* **B 500** (2001) 22 [hep-ex/0103015];

JHEP07(2002)037

L3 collaboration, *Search for R-parity violating decays of supersymmetric particles in e^+e^- collisions at $\sqrt{s} = 189\,GeV$*, Eur. Phys. J. **C 19** (2001) 397 [hep-ex/0011087].

[60] T. Gherghetta, G.F. Giudice and J.D. Wells, *Phenomenological consequences of supersymmetry with anomaly-induced masses*, Nucl. Phys. **B 559** (1999) 27 [hep-ph/9904378];
J.L. Feng and T. Moroi, *Supernatural supersymmetry: phenomenological implications of anomaly-mediated supersymmetry breaking*, Phys. Rev. **D 61** (2000) 095004 [hep-ph/9907319].

[61] T. Yanagida, *Naturally light Higgs doublets in the supersymmetric grand unified theories with dynamical symmetry breaking*, Phys. Lett. **B 344** (1995) 211 [hep-ph/9409329];
T. Hotta, K.I. Izawa and T. Yanagida, *Dynamical models for light Higgs doublets in supersymmetric grand unified theories*, Phys. Rev. **D 53** (1996) 3913 [hep-ph/9509201]; and *Nonabelian duality and Higgs multiplets in supersymmetric grand unified theories*, Prog. Theor. Phys. **95** (1996) 949 [hep-ph/9601320].

[62] N. Arkani-Hamed, H.-C. Cheng and T. Moroi, *Non-unified gaugino masses in supersymmetric missing partner models with hypercolor*, Phys. Lett. **B 387** (1996) 529 [hep-ph/9607463].

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