3- and 2-body supersymmetric processes: predictions, challenges and prospectives

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ABSTRACT

The processes subject of this talk are: 1) the $2 \rightarrow 3$ processes contributing, together with the $2 \rightarrow 2$ one, to the hadron-collider production of a charged Higgs boson in association with a $t$-quark; 2) the $2 \rightarrow 2$ and $2 \rightarrow 3$ processes giving rise to the hadron-collider strahlung of a charged slepton from a $t$-quark in $R_p$-violating models in which lepton-number-violating trilinear couplings largely dominate over bilinear ones; 3) 3-body neutralino and chargino decays in similar $R_p$-violating scenarios. The significance of the $R_p$-violating processes is critically assessed against implications from neutrino physics. Contributions to neutrino masses arising at the tree level, at the one- and two-loop levels are reviewed. Comments are also made on the production of sneutrinos via gluon fusion and on the decay of sneutrinos into photons pairs, and on the influence that constraints from neutrino physics have on these processes.

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ABSTRACT

The processes subject of this talk are: 1) the $2 \to 3$ processes contributing, together with the $2 \to 2$ one, to the hadron-collider production of a charged Higgs boson in association with a $t$-quark; 2) the $2 \to 2$ and $2 \to 3$ processes giving rise to the hadron-collider strahlung of a charged slepton from a $t$-quark in $R_p$-violating models in which lepton-number-violating trilinear couplings largely dominate over bilinear ones; 3) 3-body neutralino and chargino decays in similar $R_p$-violating scenarios. The significance of the $R_p$-violating processes is critically assessed against implications from neutrino physics. Contributions to neutrino masses arising at the tree level, at the one- and two-loop levels are reviewed. Comments are also made on the production of sneutrinos via gluon fusion and on the decay of sneutrinos into photons pairs, and on the influence that constraints from neutrino physics have on these processes.

1 Introduction

Now that LEP has been shut down, the upgraded Tevatron, with $\sqrt{s} = 2$ TeV and, possibly, a luminosity of $10 - 30$ fb$^{-1}$ by the end of the experiment \cite{1}, and soon the LHC, with $\sqrt{s} = 14$ TeV and a luminosity of $100$ fb$^{-1}$ per year \cite{2}, will be the next machines to probe physics beyond the Standard Model and to search for new particles. Possible discoveries will then be confirmed and thoroughly studied at hopefully incoming Linear Colliders \cite{3}. The discovery reach of hadron machines, however, may not be very high for particles that carry only weak charges. Typically these particles are produced through the Drell-Yan mechanism or through cascade decays. At the LHC the former mechanism suffers from being quark-initiated instead than gluon-initiated. The latter is, in general, penalized at both machines by the fact that the energy of the decaying particle may be substantially reduced with respect to that of the initial beam. The discovery reach for sleptons, for example, is only $\lesssim 350$ GeV at the LHC \cite{4}. Nevertheless, a sufficient production cross section may be achieved through the enhancement of some large parameter, or if an alternative production mechanism is possible.

In the case of the charged Higgs there are two such large parameters: $y_t$ and $\tan \beta$. The alternative mechanism is the Higgs strahlung, i.e. the production in association with a $t$- and a $b$-quark \cite{5, 6, 7, 8}, or in association with a $t$-quark only \cite{9, 6, 8}. The former
possibility is described by the elementary $2 \rightarrow 3$ processes $gg, qq \rightarrow tH^- b$, the latter, by the $2 \rightarrow 2$ process $gb \rightarrow tH^-$. Both types of elementary processes are discussed in Section 2. Predictions for the production cross sections are also given, in the case in which the subsequent decay of the charged Higgs gives rise to an overall multi-$b$'s topology, as well as in the case in which at most two $b$-quarks are present in the final state. In the first case, the need of minimizing the number of $b$-quarks to be tagged requires the combination of the $2 \rightarrow 3$ and the $2 \rightarrow 2$ cross sections.

The discovery of a charged Higgs boson will imply the existence of a second Higgs doublet and, possibly, of a supersymmetric underlying structure of the world [10]. Charged Higgs bosons give potentially large contributions to flavour-changing-neutral-current processes. Indeed, there are quite strong constraints on their mass coming from these processes, in particular $b \rightarrow s\gamma$ [11]. The constraint is even stronger in a non-supersymmetric 2 Higgs Doublet Model (2HDM) [12], since it is not plagued by the model dependence that it has in supersymmetric models. However, it was shown that, even in this case, this bound can be evaded if the 2HDM is the remnant of a multiHiggs doublet model, with all doublets, except two, very heavy [13]. The cross section for the strahlung production mechanism, as well as the branching ratios for the decays of the resulting charged Higgs are, in such a model, the same as in the 2HDM. Therefore, it is reasonable to keep separated limits from virtual effects and limits from direct production (as indeed it was/is already assumed in direct searches at LEP and at the Tevatron). They are complementary and they possibly reinforce each other. We assume only the current limits on the charged Higgs bosons mass coming from LEP ($\simeq 78.6$ GeV) and the Tevatron ($\sim 160$ GeV) (see for example [14]), although these limits were challenged as model dependent [15].

From charged-Higgs strahlung, the step to slepton strahlung in $R_p$-violating models [16, 17] is quite small. If $R_p$ is violated, or more properly, if the lepton number $L$ is violated, the charged Higgs has all the quantum numbers of a charged slepton, say $\tilde{\tau}$. This can then be singly produced in association with a $t$- and a $b$-quark, or only a $t$-quark, exactly like the charged Higgs. (Note that the single production of charged sleptons is impossible if $R_p$ is conserved.) The corresponding cross section can be obtained generalizing the calculation for the associated production of the charged Higgs. The role of $y_b \cdot \tan\beta$ is plaid in this case by the $L$-violating parameter $\lambda_{333}^\prime$, but there is no counterpart for the parameter $y_t$. The yield of $\tilde{\tau}$'s turns out to be considerable even for modest values of $\lambda_{333}^\prime$.

The same production mechanism holds also for first and second generation sleptons, $\tilde{\mu}$ and $\tilde{e}$. The corresponding cross sections are in these cases proportional to the square of $\lambda_{233}^\prime$ and $\lambda_{133}^\prime$. As in the case of the charged Higgs, the identification of all decay modes for these singly produced sleptons is very important in that it may affect the way the production cross section must be calculated. These issues are discussed in Section 3. One of the slepton decay modes, induced by the same $L$-violating coupling $\lambda_{333}^\prime$ responsible for its production, $b\tilde{t}$, is exactly one of the decay modes of the charged Higgs. For this decay mode the combination of the $2 \rightarrow 3$ and $2 \rightarrow 2$ elementary processes is mandatory. The other decay modes of the charged slepton are into a charged lepton and the lightest neutralino, $\tilde{\ell}^- \rightarrow \ell^- \tilde{\chi}_0$, or into a neutral lepton and the lightest chargino, $\tilde{\ell}^- \rightarrow \nu \ell \tilde{\chi}_1^-$. The relevance for slepton production of other $R_p$-violating couplings involving at least one third
generation index, such as $\lambda'_{ij}$ and $\lambda'_{j3}$ is also discussed. Moreover, it is emphasized that the simultaneous presence of couplings of type $\lambda'_{33}$ and of type $\lambda_{i33}$ may have important consequences for charged Higgs searches.

It is therefore crucial to know whether among the final states obtained after the decays of $\tilde{\chi}_1^0$ and $\tilde{\chi}_1^-$, there are any distinguishable from those obtained from charged Higgs decays. A discussion of the 3-body decays of $\tilde{\chi}_1^0$ and $\tilde{\chi}_1^-$, due to the $R_p$-violating coupling $\lambda'_{i33}$, $\lambda'_{332}$, and $\lambda'_{323}$, is found in Section 4. It follows closely Refs. [18, 19] and it is rather general, i.e. valid for any kinematical region for the lightest neutralino and chargino, and independent of the assumption on gaugino mass unification at some large scale. Some guidelines on the distinguishability of a charged slepton from a charged Higgs boson, both produced via strahlung from a quark line, are given in Section 3.

In the same spirit as in the discussion on charged Higgs boson strahlung, also in the case of slepton strahlung and neutralino/chargino decays, no constraints from virtual effects are included in this analysis. As before, the rational is that these constraints can, in general, be evaded by some (more or less) specific choices of the different parameters contributing to these effects. An exception is made for neutrino physics in Section 5. The smallness of their masses, makes in general neutrino spectra very sensitive to all $L$-number violating couplings. We review in Section 5 the main contributions to neutrino masses in $R_p$-violating models, at the tree level, at the one- and at the two-loop level. The latter ones were recently analyzed in Ref. [20]. In general, imposing that $R_p$-violating couplings are compatible with neutrino physics jeopardizes the visibility of some of the collider signals discussed here. Moreover, the constraints from the two-loop contributions, if those from the tree- and one-loop level are evaded, may affect also searches for single sneutrino $\tilde{\nu}$ production [21, 22] as well as the decay $\tilde{\nu} \rightarrow \gamma\gamma$. We discuss how these constraints can still be evaded if, for example, $R_p$-violating couplings have nontrivial phases. We single out the $R_p$-violating parameters that remain unsuppressed after imposing constraints from neutrino physics and the relevance that they may have not only for slepton searches, but for Higgs searches as well. In any case, it is clear that an independent collider probe of the lepton structure of these couplings may become a precious tool for model building, in that it may help individuating the real mechanism of generation of neutrino masses.

Finally, we comment on the possibility of having very tiny violation of the lepton number $L$. In general, it is assumed that this would affect only the phenomenology of the lightest neutralino and possibly neutrino physics. We conclude our discussion in Section 5 by pointing out that in models in which $\tilde{\chi}_1^0$ and $\tilde{\chi}_1^-$ are nearly degenerate, such small violation of $R_p$ would have dramatic consequences also for the phenomenology of the lightest chargino.

2 Charged Higgs strahlung

Charged Higgs strahlung is induced by the $2 \rightarrow 3$ elementary processes $gg, q\bar{q} \rightarrow tH^-\bar{b}$, which get contributions from the Feynman diagrams in Figs. 1 and 2. The production
cross section for a $p\bar{p}$ ($pp$) collider is obtained by convoluting the hard-scattering cross sections of the quark- and gluon-initiated processes with the quark- and gluon-distribution functions inside $p$ and/or $\bar{p}$:

$$\sigma = \frac{1}{2s} \int_{\tau_{\text{min}}}^{1} \frac{d\tau}{\tau} \int_{x}^{1} \frac{dx}{x} \int d\text{PS}(q_1 + q_2; p_1, p_2, p_3) \times \left\{ \sum_q \left[ q(x, \mu_f) \bar{q}(\tau/x, \mu_f) + q \leftrightarrow \bar{q} \right] |M|_{q\bar{q}}^2 + g(x, \mu_f) g(\tau/x, \mu_f) |M|_{gg}^2 \right\}, \quad (1)$$

where $q_i$ ($p_i$) are the four-momenta of the initial partons (final-state particles); $s$ is the hadron center-of-mass energy squared, whereas the parton center-of-mass energy squared is indicated, as usual, by $\hat{s} = x_1 x_2 s \equiv \tau s$; the functions $q(x, \mu_f)$, $\bar{q}(x, \mu_f)$ and $g(x, \mu_f)$ designate respectively the quark-, antiquark- and gluon-density functions with momentum fraction $x$ at the factorization scale $\mu_f$, and the index $q$ in the sum, runs over the five flavors $u, d, c, s, b$; and $d\text{PS}(q_1 + q_2; p_1, p_2, p_3)$ is an element of phase space of the 3-body final state. The square amplitudes $|M|_{gg}^2$ and $|M|_{q\bar{q}}^2$ can be decomposed as:

$$|M|_{q\bar{q}}^2 = \left( \frac{4G_F}{\sqrt{2}} M_W^2 \right) (4\pi\alpha_S)^2 |V_{tb}|^2 \left( v^2 V_{q\bar{q}}^2 + a^2 A_{q\bar{q}}^2 \right)$$

$$|M|_{gg}^2 = \left( \frac{4G_F}{\sqrt{2}} M_W^2 \right) (4\pi\alpha_S)^2 |V_{tb}|^2 \left( v^2 V_{gg}^2 + a^2 A_{gg}^2 \right), \quad (2)$$

Figure 1: Charged Higgs strahlung from the quark-initiated $2 \rightarrow 3$ elementary process.

Figure 2: Charged Higgs strahlung from the gluon-initiated $2 \rightarrow 3$ elementary process.
where color factors are included in the reduced squared amplitudes $V^{q\bar{q}}$, $A^{q\bar{q}}$ and $V^{gg}$, $A^{gg}$, and where $v$ and $a$ are:

$$v = \frac{1}{2} \left( \frac{m_b}{M_W} \tan \beta + \frac{m_t}{M_W} \cot \beta \right) ; \quad a = \frac{1}{2} \left( \frac{m_b}{M_W} \tan \beta - \frac{m_t}{M_W} \cot \beta \right). \quad (3)$$

In the kinematical region $m_t > m_{H^\pm} + m_b$ the above cross section could be well approximated by the much simpler resonant production cross section, as given by the on-shell $t\bar{t}$ production cross section times the branching fraction for the decay $t \to H^-\bar{b}$. It was shown, however, that in the region $m_t \sim m_{H^\pm} + m_b$, this description fails to account for the correct mechanism of production and decay of the charged Higgs boson [23].

$$v = \frac{1}{2} \left( \frac{m_b}{M_W} \tan \beta + \frac{m_t}{M_W} \cot \beta \right) ; \quad a = \frac{1}{2} \left( \frac{m_b}{M_W} \tan \beta - \frac{m_t}{M_W} \cot \beta \right).$$

Figure 3: Charged Higgs strahlung from the $2 \to 2$ elementary process.

The $2 \to 2$ process $gb \to tH^-$ giving also rise to strahlung of charged Higgs bosons has Feynman diagrams shown in Fig. 3. The hard-scattering cross section is, in this case:

$$\sigma(gb \to H^- t) = \frac{1}{2s} \left( \frac{4 G_F}{\sqrt{2}} M_W^2 \right) (4\pi \alpha_s) |V_{tb}|^2 \left( v^2 V^{gb} + a^2 A^{gb} \right) \quad (4)$$

where the functions $V^{gb}$ and $A^{gb}$ include in this case the reduced squared amplitudes as well as a phase-space integration factor. Note that, since the initial $b$-quark is contained in the proton or antiproton via a gluon, the $2 \to 2$ process is actually of the same order in $\alpha_s$ as the $2 \to 3$ ones.

When the decay of the charged Higgs does not give rise to third-generation quarks, the two production mechanisms are clearly independent. They require the reconstruction of two $b$-quarks, in the case of the $2 \to 3$ processes, or of only one, in the case of the $2 \to 2$ process. This is the case of the channel $H^- \to \tau^- \nu_\tau$, suitable for the discovery of the charged Higgs in the region of large $\tan \beta$, and possibly bound to play a key-role also in the region of intermediate $\tan \beta$ since it is not plagued by QCD background as the $H^- \to b\bar{t}$ mode [24]. The two production mechanisms can be experimentally distinguished and studied separately.

Leading-order predictions for the two production cross sections are given in Fig. 4 for the Tevatron and LHC energies as a function of the charged-Higgs mass for three different values of $\tan \beta$, $\tan \beta = 2$, 10, and 50. The cross sections $\sigma(pp, pp \to tH^-\bar{b}X)$ are shown by the solid lines, those for $\sigma(pp, pp \to tH^-X)$ by dashed lines. The leading-order parton distribution functions CTEQ4L [24] are used, and the renormalization ($\mu_R$) and factorization ($\mu_f$) scales are fixed to the threshold value $m_t + m_{H^\pm}$. A variation of these scales in the interval between $(m_t + m_{H^\pm})/2$ and $2(m_t + m_{H^\pm})$ results in changes up to
\[ \sigma(p\bar{p} \to t(\bar{b})H^- X) \text{ versus } m_{H^-}, \] for \( \sqrt{s} = 2 \text{ and } 14 \text{ TeV}, \) for \( \tan \beta = 2, 10, 50. \) The solid lines indicate the cross sections due to the \( 2 \to 3 \) processes; the dashed lines, those due to the \( 2 \to 2 \) process. Renormalization and factorization scales are fixed as \( \mu_R = \mu_f = m_t + m_{H^-}. \) From Ref. [6].

\( \pm 30\% \) in both cross sections. QCD corrections, therefore, may be important, as was shown in the case of associate production of the neutral Higgs [25], but they are unfortunately not yet available. Part of these corrections is captured by the QCD correction to the \( b \)-quark mass from which these cross sections depend. A study of their variation for different values of the \( b \)-quark mass can be found in Ref. [26].

As expected, the production cross section due to the \( 2 \to 3 \) processes is enhanced in the resonance region \( m_{H^\pm} < m_t - m_b, \) when \( H^\pm \) is obtained as a decay product of one of the two \( t \)-quarks produced on shell. The resonant \( t \)-quark propagator is regularized by the width of the \( t \)-quark, calculated from the SM decay \( t \to bW^\pm \) and from the decay \( t \to H^+b. \) Away from the resonance region, the cross section diminishes rather rapidly and becomes negligible at the Tevatron energy. In this region, the relative size of the two classes of cross section depends on \( \sqrt{s} \) and \( m_{H^\pm}. \) Indeed, at high energies, where the gluon-initiated \( 2 \to 3 \) processes dominate over the quark-initiated ones, the cross section arising from the elementary process \( gb \to tH^- \) is larger than that from \( gg \to tH^-b, \) which is penalized by a 3-body phase space suppression. At the Tevatron center-of-mass energy, the quark-initiated \( 2 \to 3 \) processes still have the dominant role up to intermediate values of \( m_{H^\pm}. \) For the particular choice of scales \( \mu_R \) and \( \mu_f \) made here, the cross-over for the two cross sections is at about \( m_{H^\pm} \sim 265 \text{ GeV}. \) Both classes of cross sections show the typical behaviour as a function of \( \tan \beta, \) with a minimum at around \((m_t/m_b)^{1/2}.\)

When the charged Higgs, decaying, gives rise to at least two \( b \)-quarks, the final state to be identified contains at least three \( b \)’s for the \( 2 \to 2 \) production mechanism, and at least four for the \( 2 \to 3 \) one. This is the case of the decay channel \( H^+ \to t\bar{b}, \) particularly important in the region of small \( \tan \beta, \) when the rate for \( H^- \to \tau^- \nu_\tau \) vanishes. Since tagging more than three \( b \)’s seems quite difficult even at the LHC (see, however, Ref. [27]), the two production mechanisms result into final states that are indistinguishable and a sum of the two cross sections is necessary. As shown in Fig. 5, the kinematical region in
Figure 5: Term of the gluon-initiated $2 \to 3$ process overlapping with a contribution from the $2 \to 2$ one.

which one of the two initial gluons in the $2 \to 3$ processes produces a pair $b\bar{b}$ collinear to the initial $p$ or $\bar{p}$, is, by definition, also kept into account by the $2 \to 2$ process with the initial $b$-quark collinear to the colliding $p$ or $\bar{p}$, in which it is contained via a gluon. A large factor $\alpha_s(\mu_R) \log(\mu_f/m_b)$ ($\mu_f$ is $\mathcal{O}(m_{H\pm})$), is induced in this kinematical region in the cross section from the $2 \to 3$ processes. This same factor is contained in the cross section from the $2 \to 2$ process and it is resummed to all orders, $(\alpha_s(\mu_R) \log(\mu_f/m_b))^n$, when making use of the phenomenological $b$-distribution function. The first order, $n = 1$, must then be subtracted from the sum of the two cross sections [6].

![Figure 5](image)

Figure 6: The cross section $\sigma(p\bar{p} \to tH^-X)$ obtained by adding the contributions from the $2 \to 2$ and $2 \to 3$ processes shown in Fig. 4, and subtracting overlapping terms. From Ref. [6].

Predictions for the appropriately summed inclusive cross section are shown in Fig. 5, both for the Tevatron and the LHC. These cross sections have the same theoretical uncertainty of the individual cross sections of Fig. 4 as well as the same $\tan \beta$ dependence. Once again, the inclusion of QCD corrections is awaited for. (The dependence of these cross sections on the value of $m_b$, which may be most strongly affected by QCD corrections, is shown in Ref. [20].) Similarly, at least in the region of large $\tan \beta$, supersymmetric corrections should be nonnegligible. These have been calculated for the $2 \to 3$ processes, see [28] and references therein. Recently also a calculation of the gluino corrections for the $2 \to 2$ process [28] had appeared, but no complete estimate exists, as yet, for the summed cross sections.
3 Charged slepton strahlung

As the charged Higgs boson, a charged slepton can be radiated from a $t$-quark thanks to the couplings $\lambda_{i3j}'$ in the superpotential

$$W \supset -\lambda_{tmn}' L_l Q_m D_n^c - \frac{1}{2} \lambda_{tmn} L_l L_m E_n^c ,$$

where, for completeness also the operators with couplings $\lambda_{lmn}$ are shown. In particular, for the coupling $\lambda_{333}'$, both the $2 \rightarrow 3$ partonic processes, $q\bar{q}, gg \rightarrow t\bar{b}\tau \bar{\tau}$ and the $2 \rightarrow 2$ process $gb \rightarrow t\tau \bar{\tau}$ contribute to the strahlung of a $\tau$-slepton. The corresponding Feynman diagrams are obtained from those in Figs. 1, 2, and 3 by replacing $H^-$ with $\tau$. The inclusive cross section induced by the $2 \rightarrow 3$ processes is that of Eq. (1) with the square amplitudes $|M|_{gq}^2$ and $|M|_{q\bar{q}}^2$ given by:

$$|M|_{q\bar{q}}^2 = \frac{1}{4} |\lambda_{333}'|^2 (4\pi\alpha S)^2 |V_{tb}|^2 (V_{q\bar{q}}^q + A_{q\bar{q}}^q)$$

$$|M|_{gg}^2 = \frac{1}{4} |\lambda_{333}'|^2 (4\pi\alpha S)^2 |V_{tb}|^2 (V_{gg}^g + A_{gg}^g) ,$$

where the values $v = a = 1/2$ were already substituted in. The reduced square amplitudes $V_{q\bar{q}}^q, A_{q\bar{q}}^q$ and $V_{gg}^g, A_{gg}^g$ coincide with those obtained for the $2 \rightarrow 3$ charged-Higgs-production processes, once the replacement $m_{H^\pm} \rightarrow m_{\tau}$ is made. The partonic cross section for $gb \rightarrow t\tau \bar{\tau}$ can also be obtained from that for charged Higgs production, with obvious changes, i.e. by replacing $4G_F M_W^2/\sqrt{2}$ with $|\lambda_{333}'|^2$ and $m_{H^\pm}$ with $m_{\tau}$.

![Figure 7](image-url)

Figure 7: Same as in Fig. 6, for $\tau$ production, when the coupling $\lambda_{333}'$ is dominant among all other $R_p$-violating couplings. The production cross section is shown for different values of $\lambda_{333}'$. From Ref. [6].

The same coupling $\lambda_{333}'$ inducing the radiation of the charged slepton from a third generation quark line, is also responsible for the decay $\tau \rightarrow b\bar{t}$. In this case, the charged slepton behaves exactly like a charged Higgs decaying into the same final state, and has only the effect of increasing the number of $ttbb$ events from $2 \rightarrow 3$ processes, or of $ttb$
events from the \( 2 \to 2 \) one. In this case, the evaluation of the slepton production cross section requires the combination of the \( 2 \to 3 \) and the \( 2 \to 2 \) cross sections, as outlined in Section 2. Predictions for this cross section are shown in Fig. 6 for the Tevatron and the LHC, where in the kinematical region of a resonant \( t \)-quark, only the two decay modes \( t \to W^+b \) and \( t \to \bar{\tau}^+b \) are considered. It is assumed that the charged Higgs \( H^\pm \) is sufficiently heavy, as to kinematically forbid the decay mode \( t \to H^+b \).

The large cross section obtained in the case of \( \sqrt{s} = 14 \text{ TeV} \) implies that at the LHC, with a luminosity of 100 fb\(^{-1} \) per year, light \( \tilde{\tau} \)'s may be produced in abundance even for couplings as small as 0.01. Charged sleptons are, therefore, potentially more accessible (for not too small values of \( \lambda'_{333} \)) than in models with lepton-number conservation, where their direct production relies on the Drell–Yan process. As it is known, this allows a discovery reach for sleptons only up to \( m_{\tilde{\ell}} \approx 350 \text{ GeV} \) at the LHC. \(^*\) Cascade decays may provide the bulk of slepton production in the lepton-number-conserving models, but such processes are highly model-dependent. They also complement slepton production in the lepton-number-violating case studied here. The potential reach at the Tevatron is limited to large couplings and/or light \( \tilde{\tau} \)'s, \( i.e. \) to \( \tilde{\tau} \) lighter than the \( t \)-quark. In the resonant part of the curves, where the cross section can be approximated by the production cross section of a pair of \( t \)-quarks times the branching ratio for the decay of one of the two \( t \)'s, the direct decay \( t \to \tilde{\tau}^+b \) can significantly constrain the values of \( \lambda'_{333} \). For \( m_{\tilde{\tau}} = 70 \text{ GeV} \), only values up to \( \sim 0.2 \) will remain allowed, if charged slepton searches have a negative outcome at Run II with an integrated luminosity of 10 fb\(^{-1} \).

Besides the decay \( \tilde{\tau} \to \bar{b}b \), two additional 2-body decays, due to gauge interactions, are possible for \( \tilde{\tau} \): \( \tilde{\tau} \to \tilde{\chi}^0\tau^- \) and \( \tilde{\tau} \to \tilde{\chi}^-\nu_\tau \). Whether the \( 2 \to 3 \) processes and the \( 2 \to 2 \) one must be combined also in this case, depends on the final states produced by the decays of charginos and neutralinos. A list of all the final states produced by the couplings \( \lambda'_{333} \) and \( \lambda_{333} \) can be found in Table I, in which also the final states obtained from charged Higgs decays are listed. In first approximation, let us assume that the coupling \( \lambda'_{333} \) is dominant among all \( R_\mu \)-violating couplings. In such a case, the lightest neutralino decays as (see Section 4):

\[
\tilde{\chi}_1^0 \to b\bar{b}\nu_\tau, \quad \tilde{\chi}_1^0 \to t\bar{b}\tau^-, \quad \tilde{\chi}_1^0 \to \bar{b}b\tau^+, \quad (7)
\]

with the second and third decay mode allowed only for a heavy \( \tilde{\chi}_1^0 \), \( i.e. \) for \( m_{\tilde{\chi}_1^0} > m_t + m_b + m_\tau \). The 2-body decays due to tree-level mixing of Higgs–slepton, \( \tilde{\chi}^-\tau^- \), and \( \tilde{\chi}^0-\nu_\tau \) are neglected here. They become competitive with the 3-body decays only when the values of the trilinear couplings are rather small, and therefore no slepton strahlung signal is expected. For a discussion of the decays due to these mixings, see for example Ref. 34. In the same approximation, the lightest chargino decays via gauge interactions as:

\[
\tilde{\chi}_1^- \to W^-\tilde{\chi}_1^0. \quad (8)
\]

The \( R_\mu \)-violating chargino decays,

\[
\tilde{\chi}_1^- \to \bar{b}b\tau^-, \quad \tilde{\chi}_1^- \to b\bar{b}\nu_\tau, \quad \tilde{\chi}_1^- \to t\bar{t}\tau^-; \quad (9)
\]

\(^*\)The analysis of Ref. 32 neglects the decay \( \tilde{\tau} \to \nu_\tau\tilde{\chi}_1^- \), with the subsequent chargino decay \( \tilde{\chi}_1^- \to \bar{b}b\tau^- \), which may be nonnegligible for the values of \( \lambda'_{333} \) considered.
with the second and third decays allowed only if $m_{\tilde{\chi}_1^-} > m_{\nu} + m_b$ and $m_{\tilde{\chi}_1^-} > 2m_t + m_{\tau}$, respectively, are competitive with the decay mediated by gauge interactions only if $\lambda^\prime_333$ is relatively large. The possible final states that can be produced through these decays contain all at least 3 $b$'s. The $2 \rightarrow 3$ and $2 \rightarrow 2$ processes should, then, be combined, if the number of $b$-quarks that can be detected has to be minimized.

$$
\begin{array}{c|c|c}
\hline
H^- / \tilde{\ell}_i^- & \text{Decay mode} & \text{Decay final state} \\
\hline
H^- & b\bar{t} & b(\bar{b}\ell_i^- \nu_i) \\
& \tau^\nu_\tau & \tau^- \nu_\tau \\
& hW^- & (\bar{b}b)(\ell_i^- \nu_i) \\
& \tilde{\chi}^- \tilde{\chi}_0 & W^- \tilde{\chi}_0 \tilde{\chi}_0 \\
\hline
\tilde{\ell}_i^- & \ell_i^- \tilde{\chi}_0 & \ell_i^- (\bar{b}b\nu_i) \\
& & \ell_i^- (\tau^- \tau^\nu_\tau) \\
& & \ell_i^- (\nu_\tau \tau^\nu_\tau \ell_i^-) \\
& & \ell_i^- (\nu_\tau \tau^\nu_\tau \ell_i^-) \\
& \nu_i \tilde{\chi}^- & \nu_i (\bar{b}b\ell_i^-) \\
& & \nu_i (\tau^- \nu_\tau \ell_i^-) \\
& & \nu_i (\nu_\tau \nu_\tau \ell_i^-) \\
& & \nu_i (\nu_\tau \nu_\tau \ell_i^-) \\
\hline
\end{array}
$$

Table 1: Comparison between decay modes and final states obtained from a charged Higgs boson in a $R_p$-conserving scenario and those obtained from a charged slepton of generation $i$ in a $R_p$-violating one, in which the couplings $\lambda^\prime_333$ and $\lambda_i33, (i \neq 3)$, are the dominant one.

In this case, distinguishing a light $\tilde{\tau}$ (with also light neutral gauginos) from a charged Higgs will be difficult. The final states from $\tilde{\tau}$ production and decay are very similar to those obtained from charged Higgs bosons: compare for example $t(\bar{b})b\bar{b}\tau^- \nu_\tau X$ obtained from sleptons, versus the states $t(\bar{b})tbX, t(\bar{b})\tau^- \nu_\tau X$, obtained from the charged Higgs
boson decays into \( \bar{tb} \) and \( \tau^- \nu_\tau \), or versus the state \( t(\bar{b})b\bar{b}\tau^- \nu_\tau X \) obtained, for example, from the decay of the charged Higgs boson into \( hW^- \).

For heavier sleptons and at least one of the two neutral gauginos heavier than the \( t \)-quark, it is possible to obtain the interesting final state \((2t)\bar{b}(2\tau^-)X\). When the two \( t \)'s decay hadronically, the resulting signature is that of two equal-sign leptons, jets and no missing energy. Note that the two equal-sign leptons may be two \( \mu \)'s if the coupling \( \lambda'_{233} \) is the one considered. In this case, the strahlung-produced sleptons are \( \tilde{\mu} \)'s. The production cross sections shown in Fig. 7 apply also to this case, when \( \lambda'_{333} \) is substituted by \( \lambda'_{233} \) and \( m_\tilde{\chi} \) by \( m_\tilde{\mu} \).

The couplings \( \lambda'_{i3j} \) give also rise to strahlung of the slepton \( \tilde{\ell}_i \) in association with a \( t \)- and an \( s \)-quark, when \( j = 2 \), or in association with a \( t \)- and a \( d \)-quark when \( j = 1 \). At the partonic level these production mechanisms are described by the elementary processes \( q\bar{q}, gg \to t\bar{s}\tilde{\ell}_i \), \( gs \to t\tilde{\ell}_i \), and \( q\bar{q}, gg \to t\bar{d}\tilde{\ell}_i \), \( gd \to t\tilde{\ell}_i \). Some representative diagrams for the strahlung associated with a \( t \)- and an \( s \)-quarks are shown in Fig. 8. For these couplings, at most 2 \( b \)'s are present in the final states and they can be easily tagged at the LHC. The \( 2 \to 3 \) and \( 2 \to 2 \) processes can therefore be disentangled and studied separately.

![Figure 8: Strahlung of a charged slepton associated with a \( t \)- and an \( s \)-quark.](image)

Note that the cross sections for the \( 2 \to 3 \) processes \( q\bar{q}, gg \to t\bar{s}\tilde{\ell}_i \), and \( q\bar{q}, gg \to t\bar{d}\tilde{\ell}_i \) are as those obtained for \( q\bar{q}, gg \to t\bar{b}\tilde{\tau} \). The states \( t\bar{s}\tilde{s} \) and \( t\bar{d}\tilde{d} \), obtained when the radiated sleptons decay as \( \tilde{\ell}_i \to s\tilde{t} \) and \( \tilde{\ell}_i \to d\tilde{t} \), cannot be mimicked by charged Higgs bosons production and decays, which are penalized by CKM suppression and by the practically vanishing mass of the \( s \)- and \( d \)-quark.

The production through the \( 2 \to 2 \) processes \( gs \to t\tilde{\ell}_i \) and \( gd \to t\tilde{\ell}_i \) (in particular also \( \tau \) production) is larger than the production through \( gb \to t\ell_i \), because of the larger content of the \( s \)- and \( d \)-quarks in the proton. The possibility of distinguishing which coupling has induced the production of \( \tilde{\ell}_i \), i.e. whether \( \lambda'_{33j} \), \( \lambda'_{332} \), or \( \lambda'_{331} \), strongly depends on how accurately the final states can be studied. The final states induced by the coupling \( \lambda'_{i3j} \) through \( 2 \to 2 \) processes are:

\[
\begin{align*}
\text{t}\bar{t}d_j & \quad (\tilde{\ell}_i \to d_j\tilde{\ell}) \\
\ell_i b\bar{d}_j\nu_i & \quad (\ell_i \to \ell_i^- \chi_1^0 \to \ell_i^- b\bar{d}_j\nu_i) /
\tilde{\ell}_i \to \nu_i\bar{\ell}_i^- \to \nu_i b\bar{d}_j\ell_i^-) \\
2(t\ell_i^-)d_j & \quad (\ell_i \to \ell_i^- \chi_i^0 \to \ell_i^- t\ell_j\ell_i^-) \\
t\bar{t}d_j\nu_i\nu_i & \quad (\tilde{\ell}_i \to \nu_i\bar{\chi}_i^- \to \nu_i t\ell_j\nu_i).
\end{align*}
\]

(Chargino and neutralino decays due to the couplings \( \lambda'_{332} \) are listed in the next Section.) The identification of 3 \( b \)'s versus that of 2 \( b \)'s and one \( s \)-quark, for example, will allow to
distinguish between $\lambda'_{333}$ and $\lambda'_{332}$. Also in this case, the final states obtained when $j = 1$ or $j = 2$ are not likely to be induced by charged Higgs bosons and such signals would be genuine $R_p$-violating ones.

There is another rather interesting possibility that opens up for sleptons obtained from $2 \to 2$ processes, through the couplings $\lambda'_{332}$ and $\lambda'_{331}$. If the sleptons $\tilde{l}_i$ produced in this way decay through the couplings $\lambda_{ij3}$ or $\lambda_{j3i}$, the processes

$$g\tilde{d}, g\tilde{s} \to t\tilde{l}_i \to t\tau^–\nu_j$$

lead to the same final state obtained from charged Higgs bosons, also originating from the $2 \to 2$ process $g\tilde{b} \to tH^-$. Given the larger yield of the s- and the d-quark into the proton, the slepton production cross section may be much larger than that of the charged Higgs and a serious contamination of the charged Higgs signal may be expected even for not too large values of the relevant $R_p$-violating couplings. Such a possibility should be kept in mind when searching for charged Higgs bosons, through the mode $\tau^–\nu_\tau$. This is especially so, when this decay mode is used to extend the region of $\tan \beta$ that can be probed to intermediate values such as $\sim 10$, for which the Yukawa coupling of the $\tau$ is not particularly large.

Finally, the coupling $\lambda'_{ij3}$ give rise to slepton strahlung in association with a $c$- and a $b$-quark or with a $u$- and a $b$-quark, when $j = 2$ and $j = 1$. In this case, the strahlung mechanism must compete with the $2 \to 1$ processes, $b\tilde{c} \to \tilde{l}_i$ and $b\tilde{u} \to \tilde{l}_i$, which lead to final states different from those obtained from the $2 \to 2$ and $2 \to 3$ processes. Among these, the states $c\bar{c}b\bar{b}$ may substantially add up to those obtained from $W$-pair production. In addition, final states such as $(2c)\tilde{b}(2\ell_1)$, with two like-sign leptons, jets and no missing energy can be obtained, in this case, even for light neutralinos, and are distinctive of $R_p$-violating couplings.

4 Neutralino and chargino decays

As already mentioned in Section 3 in $R_p$-violating scenarios with a dominant coupling $\lambda'_{333}$, the lightest neutralino decays as $\tilde{\chi}^0_1 \to b\nu_\tau$, $t\tau^–$, and $t\tau^+$. A detailed analysis of these and other decays was performed in Ref. [18], where no assumption was made as to whether the final state is massive or massless, or whether gaugino masses are unified or not at some large scale. As already mentioned in the previous Section, we neglect here the 2-body decays of the lightest neutralino due to tree-level mixing Higgs–slepton, $\tilde{\chi}^- \to \tau$, and $\tilde{\chi}^0 \to \nu_\tau$. For a discussion of these decays and further references on the topic, we refer the reader to Ref. [31].

The results may be summarized as follows, when the couplings $\lambda'_{333}$, or $\lambda'_{332}$, and $\lambda'_{323}$ are considered. 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 are, e.g., $\tilde{\chi}^0_1 \to b\nu_\tau$, due to a dominant $\lambda'_{333}$ coupling, and $\tilde{\chi}^0_1 \to c\bar{\tau}^–$, $\tilde{\chi}^0_1 \to \bar{s}\nu_\tau$, and their CP-conjugate.

In addition, final states such as $(2c)\tilde{b}(2\ell_1)$, with two like-sign leptons, jets and no missing energy can be obtained, in this case, even for light neutralinos, and are distinctive of $R_p$-violating couplings.
states, induced by $\lambda'_{323}$ and $\lambda'_{332}$. 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 over a wide range of parameter values, see for example Fig. 9. In this and the following Figure, the lightest neutralino decay widths are plotted for values of $\lambda'_{333}$ and $\lambda'_{323}$ equal to 1. Widths for lower values of these couplings are obtained by scaling down those shown in these Figures.

When $\tilde{\chi}_1^0$ is substantially heavy, decays with the $t$-quark in the final state, such as $t\bar{b}\tau^-$, $t\bar{s}\tau^-$, and their CP conjugated states, can become competitive. In general, for moderate/large left-right mixing terms in the sfermions virtually exchanged, and/or of 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 Figs. 4 and 10.

Furthermore, an overall increase in the total decay width of $\tilde{\chi}_1^0$ is observed, if $\tilde{\chi}_1^0$ is mainly a Wino, as it may happen if gaugino mass unification is not imposed, or in anomaly-mediated supersymmetry-breaking scenarios [32], or in grand-unification models with an additional strong hypercolor group [33].

It should be noted here that, if $\tilde{\chi}_1^0$ is lighter than the $t$-quark, the $\tilde{\chi}_1^0$ decays are not suitable to study the lepton structure of the couplings $\lambda'_{i33}$. In this case, only decays with neutrinos in the final states are possible, i.e. $\tilde{\chi}_1^0 \rightarrow b\bar{b}\nu_i$. The modes in which the charged leptons are present, $\tilde{\chi}_1^0 \rightarrow t\bar{b}\ell_i$ and $\tilde{\chi}_1^0 \rightarrow \bar{t}b\ell_i$, are kinematically forbidden. A study of the lepton structure of these couplings could reveal precious information on the origin of neutrino masses (see discussion in Section 5). Unfortunately, independent collider probes are restricted to the case of heavy neutralinos. Couplings such as $\lambda'_{i23}$, which are far less relevant for neutrino physics, on the contrary, allow light final states, containing a charged lepton, $\tilde{\chi}_1^0 \rightarrow c\ell_i$.  

Figure 9: Widths (in GeV) for the four decays $\tilde{\chi}_1^0 \rightarrow b\bar{b}\nu_\tau$, $tb\tau^-$, $b\bar{s}\nu_\tau$, and $c\bar{b}\nu_\tau$, when only the two couplings $\lambda'_{333}$ and $\lambda'_{323}$ are nonvanishing, as function of the sfermion mass $m_s$, equal to $m_{\tilde{Q}} = m_{\tilde{U}_c} = m_{\tilde{E}_c}$, and to (1.4) $m_L$ (with $m_{\tilde{L}} = m_{\tilde{E}_c}$), for $\mu = 1500$ GeV, $m_{\tilde{B}} = 600$ 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. From Ref. [18].
Figure 10: $\Gamma(\tilde{\chi}^0_1 \to b\bar{b}\nu_{\tau})$ and $\Gamma(\tilde{\chi}^0_1 \to t\bar{b}\tau^-)$ (in GeV), respectively solid and long-dashed lines, versus $m^*_s = m_{\tilde{\chi}^0_1} = m_{\tilde{Q}^c} = m_{\tilde{U}^c} = m_{\tilde{D}^c}$, for $\mu = 600$ GeV, $m_{\tilde{L}^c} = m_{\tilde{E}^c} = 600$ GeV, $m_{\tilde{B}} = 600$ GeV, and $m_{\tilde{W}} = 2m_{\tilde{B}}$. The value of $\tan \beta$ is 3 in the left frame, 30 in the right one. All trilinear soft terms are fixed at 350 GeV in the left frame, and at 150 GeV in the right frame. From Ref. [18].

Chargino decays could be ideal to probe the $i$ dependence of the couplings $\lambda'_{i33}$. The final states induced by such couplings,

$$\tilde{\chi}_1^- \to b\bar{b}\ell_i, \quad \tilde{\chi}_1^- \to b\bar{t}\nu_i, \quad \tilde{\chi}_1^- \to t\bar{t}\ell_i,$$

where the last one is due to nonvanishing left-right mixing terms in the sfermion mass matrix. Among these states, the first one, in particular, could easily supply the information that decays of a light neutralino cannot give. These $R_p$-violating 3-body decays, however, must compete with the $R_p$-conserving 2-body decay $\tilde{\chi}_1^- \to W^-\tilde{\chi}^0_1$. In general, for $\tilde{\chi}^0_1$ lighter than $\tilde{\chi}_1^-$, their rates are much smaller than the rate of the dominant mode. For rather large $R_p$-violating couplings, the 3-body decays, although subleading, may be nonnegligible. They may, however, be the dominant ones, when the 2-body decay $\tilde{\chi}_1^- \to W^-\tilde{\chi}^0_1$ is suppressed. This is typically the case of models in which $\tilde{\chi}_1^-$ and $\tilde{\chi}^0_1$ are nearly degenerate. A sufficient condition for this to happen is that $\tilde{\chi}_1^0$ is mainly a Wino or mainly a Higgsino. A detailed analysis of chargino decays in these scenarios will be presented elsewhere [19].

We conclude this Section by listing the chargino decays due to the coupling $\lambda'_{i32}$ and $\lambda'_{i23}$. The coupling $\lambda'_{i23}$ gives rise to the decays:

$$\tilde{\chi}_1^- \to b\bar{s}\ell_i, \quad \tilde{\chi}_1^- \to b\bar{c}\nu_i, \quad \tilde{\chi}_1^- \to t\bar{c}\ell_i, \quad \tilde{\chi}_1^- \to t\bar{s}\nu_i,$$

and all the CP-conjugated ones, whereas the decays obtained through the coupling $\lambda'_{i32}$ are:

$$\tilde{\chi}_1^- \to b\bar{s}\ell_i, \quad \tilde{\chi}_1^- \to t\bar{s}\nu_i,$$

when the left-right mixing of the strange squark is assumed to be negligible.
5 Collider- neutrino-physics interplay

The reasons for dwelling on \( L \)-violating couplings of type \( \lambda'_{ijk} \) with third generation indices, is that they have a particularly important role for neutrino spectra.

\[
\begin{array}{c}
\langle \tilde{\nu}_i \rangle \\
\tilde{\nu}_i \ L_i \\
\end{array}
\]
\[
\begin{array}{c}
\langle \tilde{\nu}'_i \rangle \\
\tilde{\chi}_j^0 \\
\tilde{\nu}'_i \ L_i \\
\end{array}
\]

Figure 11: The tree-level contribution to the Majorana entry in the neutrino mass matrix, induced by sneutrino \( v.e.v \)’s.

It is well known that, additional \( R_p \)-violating interactions, specifically interactions violating \( L \), induce Majorana masses for all neutrinos, without having to introduce new superfields in addition to those present in a Minimal Supersymmetric Standard Model. In generic \( R_p \)-violating models, neutrino may acquire mass at the tree level and at the quantum level \[35\]. The tree-level contributions are mainly governed by \( R_p \)-violating bilinear couplings. The radiative contributions arise at the one- or the two-loop level. At the one loop, they are controlled by trilinear superpotential \( R_p \)-violating couplings and left-right sfermion mixing mass terms. The two-loop contributions are due to trilinear superpotential \( R_p \)-violating couplings, left-right sfermion mixing terms, and to \( R_p \)-violating soft trilinear superpotential terms \[20\].

The tree-level mass is originated by bilinear \( R_p \)-violating terms, in particular by the misalignment of superpotential bilinear terms,

\[
W = \mu_i L_i H_U \quad (i = 1, 2, 3),
\]

and scalar potential bilinear terms,

\[
V = B_i \tilde{L}_i H_U + m_{L_i H_i}^2 \tilde{L}_i H_U \quad (i = 1, 2, 3).
\]

(The symbol \( H_U \) indicates the isospin \( 1/2 \) doublet superfield as well as its scalar component.) This misalignment is expressed by the nonvanishing value of the quantities:

\[
\Delta B_i = B_i - B \tilde{\mu}_i, \quad \Delta m_{L_i H_i}^2 = m_{L_i H_i}^2 - \left( m_{H_D}^2 - m_{L_i}^2 \right) \tilde{\mu}_i,
\]

where \( \tilde{\mu}_i = \mu_i / (\mu^2 + \sum \mu_i^2)^{1/2} \), and \( \mu \) is the usual parameter in the superpotential term \( \mu H_D H_U \). They induce vacuum expectation values (\( v.e.v. \)) for the sneutrinos, \( \nu_i \). Typically,

\[
m_{SUSY}^2 \cdot v_i \sim \Delta B_i < H_U > + \Delta m_{L_i H_i}^2 < H_D >,
\]

where, \( m_{SUSY} \) is, generically, the soft mass for the scalar component of the lepton doublet superfields plus \( D \)-term contributions. In a basis in which the terms \( \mu_i L_i H_U \) in the
superpotential are rotated away, the tree-level contributions to the neutrino mass matrix are given by:

\[ m_{\nu,ii}^\prime \sim \frac{1}{2} g^2 v_i v'_j, \]  

and are obtained by the tree-level diagrams in Fig. [11]. Both SU(2) and U(1) gaugino mediate such a diagram. A suppression of these contributions is granted by rather small v.e.v’s and/or cancellations among the two terms in Eq. (18) (i.e. \( g_1^2/M_1 + g_2^2/M_2 = 0 \), with \( M_1 \) and \( M_2 \), the two gaugino masses).

![Figure 12: A one-loop contribution to the Majorana neutrino mass arising from the \( \Delta L = 1 \) operators of the superpotential in Eq. (5).](image)

The remaining neutrino mass terms are induced by trilinear \( R_p \)-violating terms through loops mediated by quarks-squarks and/or leptons-sleptons. A typical quark-squark loop contribution is shown in Fig. [12]. Such a contribution requires a chirality flip in the sfermion line. Choosing, for simplicity, the basis in which left- and right- components of the down quarks are diagonal, and assuming proportionality of the left-right sfermion matrices to the corresponding Yukawa matrices, i.e. \( (m^2_{dLR})_{ij} = \mathcal{M}_j^d(m_d)\delta_{ij} \) and \( (m^2_{lLR})_{ij} = \mathcal{M}^d_j(m_l)\delta_{ij} \), the contributions to the neutrino mass matrix entries from quark-squark loops read:

\[ m_{\nu,ii}^\prime \sim \frac{3}{8\pi^2} \sum_{k,j} \lambda'_{ikj} \lambda'_{jk}(m_d)_j(m_d)_k \frac{\mathcal{M}_k^d}{m_{d_k}^2}, \]  

where \( m_{d_k} \) is a squark mass eigenvalue, and both parameters \( \mathcal{M}_k^d \) and \( m_{d_k} \) are linked to the mechanism of supersymmetry breaking. The analogous contributions from lepton-slepton loops is obtained by replacing \( \lambda'_{ijk} \) with \( \lambda_{ijk} \) in Eq. (11), and \( m_d \) with \( m_l \). The dominant among these contributions are those mediated by \( b\bar{b} \) and by \( \tau\tilde{\tau} \). Neutrino mass terms \( \ll O(eV) \), imply that \( \lambda_{i33} \) and \( \lambda'_{i33} \) are smaller than \( 10^{-3} \)-\( 10^{-4} \) for supersymmetric masses of 100 GeV. If no special tuning among other parameters occurs, e.g. if no hierarchy such as \( \mathcal{M}_3^d \ll \mathcal{M}_1^d \approx \mathcal{M}_2^d \) is realized, trilinear couplings such as \( \lambda'_{i23} \) or \( \lambda'_{i32} \), and \( \lambda_{32} \) or \( \lambda_{32} \), induce loop contributions suppressed by \( (m_s/m_b) \) and \( (m_\mu/m_\tau) \), with respect to the dominant ones. The suppression would be quadratic in the ratio of quark masses for the one-loop contributions due to the couplings \( \lambda'_{i22} \). Therefore, constraints obtained

\[ \text{The effect of non-diagonal left-right mass matrices on neutrino mass terms was considered in F. Borzumati et al. in Ref. [33].} \]
by imposing that the terms in Eq. (19) are $\lesssim \mathcal{O}(1\text{eV})$, are less powerful for couplings involving one quark index of first or second generation. It is interesting that, in general, collider physics searches act exactly in the opposite way: direct constraints are stronger for trilinear couplings with first and second generations indices, than for those containing third generation indices.

The above formulas show how the predictability of generic $R_p$-violating models is marred by the large number of parameters contributing to the different neutrino mass matrix entries. To cope with this difficulty, two different simplifying assumptions are in general made. They give rise to two different classes of models, which have been widely discussed in the literature.

The first is known as the class of bilinear $R_p$-violating models, in which only bilinear terms exist \[36\]. By rotating away the bilinear terms in the superpotential, the following trilinear couplings are induced:

$$
\lambda'_{ijk} = y_{Djk} \hat{\mu}_i, \quad \lambda_{ijk} = y_{Ljk} \hat{\mu}_i \quad (i \neq j),
$$

where $y_{Djk}$ and $y_{Ljk}$ are, respectively, the down-quark and charged-lepton Yukawa coupling matrices. In general, the largest neutrino mass term comes from the tree-level contribution. The second largest mass terms are due to $b-\tilde{b}$ loops, which give rise to entries in the neutrino mass matrix $m^{\nu,ii'}$ proportional to

$$
\sim (3/8\pi^2)\lambda'_{33}\lambda'_{33}m_b^2M_3^d/m_b^2 = (3/8\pi^2)y_b^2\hat{\mu}_i\hat{\mu}_i'm_b^2M_3^d/m_b^2.
$$

Other large mass terms are induced by $\tau-\tilde{\tau}$ loops:

$$
\sim (1/8\pi^2)\lambda_{33}\lambda_{33}m_{\tau}^2M_3^d/m_{\tau}^2 = (1/8\pi^2)y_{\tau}^2\hat{\mu}_i\hat{\mu}_{i'}m_{\tau}^2M_3^l/m_{\tau}^2,
$$

with $i, i' \neq 3$, due to the fact that the couplings $\lambda_{ijk}$ are antisymmetric in the first two indices. The smallness of neutrino masses constrain the $v.e.v$’s to be small, and in particular the $\mu_i$’s to be small \[1\]. As a consequence, the trilinear couplings are also rather small, and so are the loop contributions to neutrino masses. The lightest neutralino decays mainly as $\tilde{\chi}^0_1 \rightarrow W^- \ell^+$ induced by the mixing $\tilde{\chi}^0_1 - \nu$ proportional to a neutral slepton $v.e.v$. No signal for single-production of charged sleptons is expected in these cases.

Strictly speaking, however, since the $v.e.v$’s are an algebraic sum of two terms, the individual contributions to the $v_i$’s do not have to be small. If a certain amount of tuning of the individual contributions is allowed, these $v.e.v$’s may be sufficiently suppressed. In principle the suppression may be even strong enough to render the tree-level mass contribution of Eq. (18) irrelevant for neutrino physics. Depending on the amount of suppression, the rates for the 3-body decays of the lightest neutralino may be comparable

\[1\] It is often assumed that such a low-energy situation is obtained from a model in which $\Delta B_i, \Delta m^2_{L,H}$ and all the trilinear couplings are vanishing at some fundamental high scale. Nonvanishing $\Delta B_i$ and $\Delta m^2_{L,H}$ are induced by the parameters $\mu_i$ when evolving the model down to low scale, whereas the trilinear terms remain vanishing. As explained above, these result from a rotation of the superfields $L_i$ and $H_D$. It is easy to show that in such a case the neutrino mass matrix generated by the tree-level and the $b-b$-loop contributions has rank 1, and that the second heaviest neutrino is generated by $\tau-\tilde{\tau}$ loops.
to that for the 2-body decay $\tilde{\chi}_1^0 \rightarrow W^- \ell^+$, or even larger. In such a case, the main source for neutrino mass terms are the loop contributions, which must be duly suppressed. Responsible for their suppression may be parameters other than the trilinear couplings, i.e. other than the $\mu_i$’s. For example, nearly vanishing left-right mixing terms in the sfermion mass matrices, may be as effective as very small $\lambda$ and $\lambda'$ parameters. Moreover, it should be noted that the dominant one-loop contributions to $m_{\nu,ii}$, i.e. those mediated by $b\tilde{b}$ and $\tau\tilde{\tau}$ are proportional to the combination of couplings $(\lambda'_{i33})^2$ and $(\lambda_{i33})^2$, respectively, and are therefore sensitive to the phases of these couplings. Thus, it is possible that partial cancellations among the $b\tilde{b}$ loops and the $\tau\tilde{\tau}$ ones occur. Further suppression of the $b\tilde{b}$ loops may come from large sbottom masses. The observation of single production of sleptons would strongly hint at any of these possibilities, i.e. at unexpectedly large sbottom masses (even for not too heavy sleptons), small left-right mixing terms in the sfermion mass matrices, and/or special relations among the couplings of type $\lambda$ and those of type $\lambda'$.

The second class of models that has been considered in the literature is that of models in which only trilinear $R_p$-violating terms exist at some high scale [38]. In such models, bilinear terms are induced by the trilinear couplings $\lambda$ and $\lambda'$, through the renormalization group equation. The misalignment term $\Delta B_i$ and $\Delta m^2_{ijH}$ and the sneutrino $\nu e\tilde{\nu}$'s $\nu_i$ obtained at the electroweak scale are proportional to $\lambda'_{ijj}$ and $\lambda_{ijj}$, in a basis in which down-quark Yukawa couplings are diagonal. Barring accidental cancellations among the different terms inducing the $\nu_i$’s, the smallness of the tree-level contributions to neutrino masses constrain these couplings to be small, if a nonnegligible mixing among left- and right-mass terms for sfermions exists. As a consequence, the leading loop contributions are also suppressed and no significant yield for the strahlung of charged sleptons in association with a $t$- (and possibly a $b$-) quark is expected. As in the previous class of models, this could be observed only if left-right sfermion mixing terms are efficiently suppressed, and/or squark masses are heavy, and/or phases in $R_p$-violation parameters induce cancellations among the different loop contributions. Other one-loop contributions to neutrino masses are due to the trilinear couplings $\lambda_{ijk}$ as well as $\lambda'_{ijk}(i \neq j)$. The constraints on these last couplings are, in general, less limiting and one can expect to observe single production of sleptons, as for example $pp \rightarrow t\tilde{\mu}X$ or $pp \rightarrow t\tilde{\mu}sX$, even without having to advocate the vanishing of other parameters or the mutual cancellation of different loop contributions.

![Figure 13: A two-loop contribution to the Majorana neutrino mass arising from the $\Delta L = 1$ operators in the scalar potential in Eq. (21).](image)

Additional 2-loop contributions to neutrino masses also exists. They are in general
neglected, since assumed to be subleading with respect to the one-loop contributions.
The diagram with exchange of squarks (and sleptons) in the upper loop, shown in Fig. 13,
was already discussed in Ref. [39], in a framework in which hard superpotential trilinear
$R_p$-violating couplings are vanishing, but soft trilinear $R_p$-violating couplings
\[ V = A'_{ijk} \bar{L}_i \bar{Q}_j \bar{D}^c_k + \frac{1}{2} A_{ijk} \bar{L}_i \bar{L}_j \bar{E}^c_k \]  
(21)
are generated by supersymmetry breaking. (See also discussion in Ref. [40].) There are
also other diagrams with exchange of quarks (and leptons) in the upper loop, shown
on the left side of Fig. 14, and with exchange of squarks (and sleptons) but no left-right
mixing in the sfermion lines, as shown on the right side of Fig. 14. If any of the
mechanisms advocated above to suppress the one-loop contribution to neutrino mass
(without suppressing the $\lambda$- and $\lambda'$-couplings) is present, these two-loop diagrams cannot
be neglected any longer [20].

Figure 14: Two-loop contributions to the Majorana neutrino mass involving the same trilinear
superpotential couplings $\lambda'_{ijk}$ present in the one-loop diagram of Fig. 12.

The sneutrino state $\tilde{\nu}$ has a CP-even and a CP-odd component that are degenerate at
the tree level. A splitting in the mass of these two components is induced at the quantum
level [41] by the upper loops in Figs. 13 and 14. Because of this splitting, a nonvanishing
contribution to neutrino mass terms is originated. The contribution coming from the
diagram of Fig. 13 is:
\[ m_{\nu,ii'} \approx \frac{3g_j^2}{(16\pi^2)^2} A'_{i33} A'_{i'33} \left( \frac{m_b M_d^d}{m^2_{\tilde{b}}} \right)^2 \frac{m_{\tilde{\nu}^0}}{m_{\tilde{b}^0}}. \]  
(22)
The symbol $M_d^d$ is that already used in Eq. (19), i.e. $M_d^d m_b$ is the left-right mixing terms
in the sbottom mass matrix. Similarly, the diagrams in Fig. 14 yields the contribution:
\[ m_{\nu,ii'} \approx \frac{3g_j^2}{(16\pi^2)^2} \lambda_{i33} \lambda_{i'33} m_b^2 \frac{m_{\tilde{\nu}^0}}{m_{\tilde{b}^0}}. \]  
(23)
where we have dropped a logarithmic dependence on $m_{\tilde{b}}$. For the overall coefficients in
Eqs. (22) and (23), see Ref. [20].

One way of suppressing the one-loop contribution to neutrino mass terms, and to
evade constraints on the parameters $\lambda'$, is to assume that $M_d^d \approx 0$. In such a case,
obviously, also the contribution from the diagram in Fig. 13 nearly vanishes. The contribution in Eq. (23), however, remain unsuppressed and yield therefore a constraint on the couplings $\lambda_{i33}':$

$$\lambda_{i33}' \lesssim 0.01 \left( \frac{m_\tilde{\nu}}{100 \text{ GeV}} \right) \left( \frac{100 \text{ GeV}}{m_{\tilde{\chi}_0}} \right)^{1/2} ,$$

where the $SU(2)$ gauge coupling was used for the numerical evaluation. This is quite robust since it is not affected by the size of squark masses. Similar constraints are obtained for the couplings $\lambda_{i33}$. 

For such values of $\lambda_{i33}'$, the possibility of observing slepton strahlung at the LHC gets restricted to the kinematical region $m_\tilde{\ell} \leq m_t - m_b$, see the right frame in Fig. 7. Wider reach of values for $m_\tilde{\ell}$, i.e. beyond the $t$-quark mark, are possible if sleptons are produced through the couplings $\lambda_{32}'$ or $\lambda_{31}'$. Both, the one-loop and the two-loop contributions to neutrino mass terms obtained in this case, are suppressed with respect to those of Eq. (23) by the factors $(m_s/m_b)$ and $(m_d/m_b)$. The corresponding constraints on these couplings get weakened by factors $(m_b/m_s)^{1/2}$ and $(m_b/m_d)^{1/2}$. Similarly, also the couplings $\lambda_{ij3}$ and $\lambda_{j3i}$ may remain unsuppressed. As already observed in Section 4, the $2 \to 2$ process induced by $\lambda_{i32}'$ or $\lambda_{i31}'$ give rise to sleptons in association with a $t$-quark only. The subsequent decay of the sleptons through the couplings $\lambda_{ij3}$ or $\lambda_{j3i}$, may give rise to an unsuppressed final state $t\tau^-\nu_j$ indistinguishable from the final state $t\tau^-\nu_\tau$ obtained from a charged Higgs boson. Therefore, through $R_p$-violating couplings, sleptons may contaminate the charged Higgs signal, even when constraints from neutrino physics are imposed.

Last, but not least, it should be kept in mind that even the bounds in Eq. (24) and the corresponding one for the coupling $\lambda_{33}'$ may be evaded. For this, it is sufficient that special relations among the absolute value of $\lambda_{i33}'$ and $\lambda_{i33}$ and the value of their phases induce a partial cancellation between the contribution in Eq. (23) against that from the similar $\tau$ and $\tilde{\tau}$ diagrams. In this case, the possibility of observing slepton-strahlung induced by the coupling $\lambda_{i33}'$ may remain unaffected by neutrino mass constraints.

In turn, it is possible that the one-loop contribution to neutrino masses is, indeed, reduced by the smallness of the $\lambda'$ and $\lambda$ couplings, while the left-right mixing terms in the sfermion mass matrices remain sizable. In such a case, the neutrino mass contribution from the two-loop fully scalar diagram in Fig. 13 induces strong constraints on the $R_p$-violating soft couplings $A'_{i33}$:

$$A'_{i33} M_3^2 \lesssim 100 \text{ GeV}^2 \left( \frac{m_\tilde{\nu}}{100 \text{ GeV}} \right) \left( \frac{100 \text{ GeV}}{m_{\tilde{\chi}_0}} \right)^{1/2} \left( \frac{m_b}{100 \text{ GeV}} \right)^2 ,$$

and similar ones for the couplings $A_{i33}$. Note that this is an order of magnitude estimation. An explicit calculation, in which no approximation is made on the loop function, may attenuate this constraint.

We conclude this Section by observing that constraints such as that in Eq. (24), will also affect the possible production of sneutrinos through gluon fusion [6, 21, 22] and...
their subsequent decays in two photons $\tilde{\nu} \rightarrow \gamma \gamma$. It has been suggested that sneutrinos in $R_p$-violating models may mimic neutral Higgs bosons and spoil their observation and mass determination. They can be produced through the elementary process $b\bar{b} \rightarrow \tilde{\nu}$, as well as strahlung processes, or through gluon fusion, as shown in Fig. [15]. The first two production mechanisms are induced by the couplings $\lambda_{333}'$. These same couplings are also partially responsible for the latter mechanism, as the first two diagrams in Fig. [15] show. The last diagram, on the contrary, depends only on the couplings $A_{333}'$.

Figure 15: Gluon-fusion production of sneutrinos.

For these production mechanisms to be relevant, the generic constraints from neutrino physics must be evaded. Having assumed some accidental cancellation of the tree-level neutrino contributions, vanishing left-right mixing terms in the down-squark mass matrices are enough to leave the couplings $\lambda_{333}'$ unconstrained by the one-loop contribution to neutrino masses. The bound from the two-loop contribution given in Eq. (24), however, may be enough to reduce the $b\bar{b} \rightarrow \tilde{\nu}$ and strahlung production mechanisms to negligible levels. The same is true also for the production through gluon fusion, since the first two diagrams in Fig. [15] would be affected by small couplings $\lambda_{333}'$ and the third by a vanishing left-right mixing in the sbottom mass matrix. In this case, a clean detection of the neutral Higgs boson signal may be hoped for. If on the contrary, the one-loop contribution to neutrino mass are reduced by tiny couplings $\lambda_{333}'$, and the left-right mixing in the sbottom mass matrices are sizable, then single-production of sneutrinos remains possible through the third gluon fusion diagram in Fig. [15]. Parameters $A_{333}'$ of order $\sim 10$ GeV may give smallish cross sections, but perhaps not so small not to affect the neutral Higgs boson searches. Detailed studies on this topic are lacking and are clearly called for.

It is obviously also possible that the one-loop contribution to neutrino masses are subject to other type of suppressions, such as the already mentioned cancellations among the $b - \bar{b}$ and $\tau - \bar{\tau}$ contributions. In such a case, the production of sneutrino at hadron collider may still be relevant. Once again, we can only conclude that direct and indirect searches should be pursued independently and should complement each other.

If neutrino physics constrains the $R_p$-violating trilinear couplings to be below the $10^{-4}$-$10^{-3}$ mark, collider signals such as charged slepton strahlung and sneutrino decays will be completely negligible. Such small couplings would still be relevant for phenomenological searches. They induce a decay of the lightest neutralino, which would be a stable particle for vanishing values of these couplings. It is in general assumed that the effect of such couplings in other sectors would be completely irrelevant. However, in models with nearly-degenerate lightest chargino and lightest neutralino, couplings as small as $10^{-5}$ may affect also the phenomenology of the lightest chargino. For such values of trilinear
couplings, the rate for the main decay mode of the lightest chargino $\tilde{\chi}_1^- \rightarrow W^- \tilde{\chi}_0^1$ may become comparable or even smaller (depending on the actual value of the mass difference $m_{\tilde{\chi}_1^1} - m_{\tilde{\chi}_0^1}$) than that for the $R_p$-violating chargino decay modes. Nearly-degenerate $\tilde{\chi}_1^1$ and $\tilde{\chi}_0^1$ are present in scenarios in which the lightest neutralino is mainly a Wino or a Higgsino. These scenarios are realized in specific model of supersymmetry breaking, such as anomaly-mediated supersymmetry-breaking models \cite{32}, or, for example, in grand-unification models with an additional strong hypercolor group \cite{33}. For studies of the interesting signatures obtained in these cases, the reader is referred to Ref. \cite{19}.

\section{Conclusions and outlook}

In spite of the impressive confirmation that the Standard Model has obtained at LEP and at the Tevatron, the mechanism of electroweak symmetry breaking is still elusive. Searches for Higgs boson(s), which are likely to mediate it, will play a pivotal role in collider physics in the next coming years. Some of the issues that these searches will have to settle are if there exist only one neutral Higgs boson or more, and if there are charged Higgs bosons, i.e. if there are additional Higgs fields that are doublets of $SU(2)$. The discovery of a charged Higgs boson will, undoubtedly, strengthen the case for supersymmetry. In order not to miss such an important element of the Minimal Supersymmetric Standard Model at the LHC, it is important that the cross section for its production is not too small, so that even heavy masses may be reached. As discussed in this review, the charged Higgs strahlung is a suitable production mechanism under this point of view: cross sections at the $fb$-level can be obtained even for charged Higgs masses of $\sim 1$ TeV.

The determination of the charged Higgs mass is also crucial to nail down the parameters responsible for supersymmetry breaking. The cleanliness of this measurement, however, depends strongly on the fact that there are no other particles mimicking the signals that charged Higgs bosons produce. In $R_p$-violating models, charged and neutral sleptons may lead to an incorrect determination of the charged and neutral Higgs boson masses. The trilinear superpotential couplings $\lambda_{i33}$ and $\lambda_{i33}$ can give LHC production cross sections for charged and neutral sleptons very similar to that of the corresponding charged and neutral Higgs bosons.

When the lepton index $i$ is of the third generation, they also give rise to the same final states induced by the decays of Higgs bosons. This is true, however, if these couplings are not much smaller than numbers of $\mathcal{O}(1)$. In the case of charged Higgs bosons and sleptons radiated from a quark line, then, searches of final states such as $(2t)\bar{b}(2\tau^-)X$, obtained for heavy $\tilde{\chi}_1^0$ are necessary to establish whether the charged Higgs signal was contaminated by charged sleptons through sizable violation of the lepton number. Moreover, when the leptonic indices are of first or second generation, the presence of like-sign $e$’s and $\mu$’s in possible final states such as $(2t)\bar{b}(2\mu^-)X$ or $(2t)\bar{b}(2e^-)X$, will unambiguously indicate that charged sleptons have been produced. In addition, also first and second generation quark jets, induced for example by couplings $\lambda_{ij3}$ or $\lambda_{i3j}$, are likely to arise from charged sleptons.
and not from charged Higgs bosons. Interesting final states due to charged Higgs radiated from $t$-quark or $c$-quark lines are, for example, $t\bar{t} s$, $(2t)\bar{s}(2\ell^-_i)X$, and $(2c)\bar{b}(2\ell^-_i)X$, with $\ell^-_i = \tau^-, \mu^-, e^-$. The neutral Higgs bosons at the LHC will be mainly produced through gluon fusion, as well as strahlung off a top quark line. Sneutrino may also be produced in similar ways and with similar values of cross sections thanks to $\mathcal{O}(1)$ couplings $\lambda'_{33}$ [21] as well as electroweak scale massive couplings $A'_{33}$. Also in this case, final states overlapping with those of the neutral Higgs boson may be obtained.

Contaminations between the Higgs bosons and sleptons signals become negligible if $\lambda_{333}$, $\lambda'_{333}$, and $A'_{333}$ are sufficiently small. In general, no strong upper bounds have been obtained so far from direct searches of sparticle production and/or particle/sparticle decays. Neutrino physics, however, may play a crucial role in constraining such couplings. It has been widely discussed in the literature how one-loop contributions to neutrino masses impose that $\lambda'_{333}$ and $\lambda_{333}$ are not larger than $10^{-4}$-$10^{-3}$. Since other unknown parameters enter the one-loop calculation, however, it is simple to evade these constraints. Several possible ways have been listed in this review. In addition, there exist two-loops contributions to neutrino masses that are in general left out of discussion, since considered subleading. Nevertheless, the combinations of various parameters entering in the calculation of the two-loop diagrams are different from those encountered in the calculation of the one-loop diagrams. Therefore, the two-loop diagrams may not be affected by the one-loop constraints. Although suffering a stronger loop-suppression with respect to the one-loop diagrams, they may still give too large contributions to neutrino masses. As it is explicitly shown in this review, they yield constraints on the couplings $\lambda'_{333}$ and $\lambda_{333}$ that are more difficult to evade (although not impossible), once specific choices of parameters are made to cancel the one-loop contributions. It turns out that $\lambda'_{333}$ and $\lambda_{333}$ are constrained to be a few $\%$, whereas the couplings $A'_{333}$ should not exceed the 10 GeV order of magnitude.

Values of $\lambda'_{333}$ at the percent level, can still allow nonnegligible yields of charged sleptons at the LHC, although the production cross section is considerably smaller than that for a charged Higgs boson of equal mass. Their detection, challenging, relies on the search for final states different than those produced by charged Higgs bosons. Sneutrinos may still be produced at the LHC, through gluon fusion and/or strahlung off a quark line. Detailed studies of the impact of the parameters $A'_{333}$ and of the distinguishability of sneutrinos from Higgs bosons are still missing.

Lepton violating couplings $\lambda'_{i3j}$, $\lambda'_{ij3}$, with the quark index $j$ of second and especially first generation, are much less affected by neutrino physics. The same is true for the couplings $\lambda_{ij3}$ and $\lambda_{j3i}$. As it was extensively discussed in Sections 2 and 3, it should be noted that the $2 \to 2$ elementary processes $gs, gd \to t\tilde{\ell}_i$ induced by the couplings $\lambda'_{i3j}$, do not need to be combined with the $2 \to 3$ processes $q\bar{q}, gg \to t\tilde{\ell}_i$ or $q\bar{q}, gg \to t\bar{d}\tilde{\ell}_i$ if the produced sleptons $\tilde{\ell}_i$ decay leptonically. This is exactly the same situation encountered in the case of strahlung-produced charged Higgs bosons that decay into $\tau^-\nu_{\tau}$. Moreover the $2 \to 2$ cross section for slepton-strahlung due to the quark-flavour violating couplings $\lambda'_{i3j}$,
is bound to be larger than that obtained from a coupling $\lambda_{33}$ of equal strength, due to the larger content of the $d$- and the $s$-quark in the proton. Therefore the subsequent decay of such sleptons into $\tau^{-}\nu_{j}$ induced by the couplings $\lambda_{j3}$ and $\lambda_{j3}$ will provide a final state identical to that obtained from the charged Higgs. Sleptons of mass very close to that of the charged Higgs, would therefore affect the search of charged Higgs. A considerable splitting among the mass of these particles may lead to exciting results at the LHC.

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