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Decays of a bino-like particle in the low-mass regime

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Abstract

We study the phenomenology associated with a light bino-like neutralino with mass $\lesssim m_\tau$ in the context of the R-parity violating Minimal Supersymmetric Standard Model. This is a well-motivated example of scenarios producing potentially light and long-lived exotic particles, which might be testable in far-detector experiments, such as the FASER experiment at the Large Hadron Collider. A quantitative assessment of the discovery potential or the extraction of limits run through a detailed understanding of the interactions of the light exotic fermion with Standard Model matter, in particular, the hadronic sector. Here, we propose a systematic analysis of the decays of such a particle and proceed to a model-independent derivation of the low-energy effects, so that this formalism may be transposed to other UV-completions or even stand as an independent effective field theory. We then stress the diversity of the possible phenomenology and more specifically discuss the features associated with the R-parity violating supersymmetric framework, for example neutron-antineutron oscillations.

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1 Introduction

While baryon and lepton numbers are classical symmetries of the Standard Model (SM) of particle physics, the observed baryon-antibaryon asymmetry of the Universe [1, 2] provides clear evidence of their violation [3]. In fact, attempts to unify the interactions of Yang-Mills type [4, 5] generically break such symmetries in an explicit fashion, suggesting that their apparent conservation results from the accidental composition of the low-energy spectrum rather than from deep conceptual roots. On the other hand, matter stability — see *e.g.* the reviews

in Refs. [6–8] — strongly limits the magnitude that baryon- and lepton-number violation can reach, pushing the corresponding dynamics to very high energy scales. This situation leaves open the question of whether baryon- and / or lepton-number violation can be detected in low-energy processes.

In supersymmetric (SUSY) extensions of the SM [9, 10], such as the Minimal SUSY SM (MSSM), conservation of baryon- and lepton-number usually appears as a consequence of the requirement of R-parity conservation [11]. Yet, this property does not extend to non-renormalizable operators, as expected when taking the low-energy limit of *e.g.* a grand-unified ultraviolet (UV) completion. (See instead the discrete symmetry proton hexality [12–14] to cover also the non-renormalizable operators.) On the other hand, through the enforced stability of the lightest R-odd particle, R-parity has a deep impact on the phenomenology of the model, both in view of cosmological (thermal relics of the stable exotic particle) and collider observables (missing energy in SUSY decay chains). While this path has been copiously investigated in the course of the past forty years, its benefits with respect to matter stability may not outweigh the rich alternative, that of R-parity violation (RpV) [15–18]. In this latter case, explicit baryon- and / or lepton-number violating terms enter the superpotential, even though matter stability might still be protected by either of these discrete symmetries. For completeness, we remind here the form of the renormalizable RpV terms completing the superpotential of the MSSM [19]:

$$W_{\text{RpV}} = \mu_i \hat{H}_u \cdot \hat{L}_i + \frac{1}{2}\lambda_{ijk}\hat{L}_i \cdot \hat{L}_j(\hat{E}^c)_k + \lambda'_{ijk}\hat{L}_i \cdot \hat{Q}_{j\alpha}(\hat{D}^c)^\alpha_k + \frac{1}{2}\lambda''_{ijk}\varepsilon_{\alpha\beta\gamma}(\hat{U}^c)^\alpha_i(\hat{D}^c)^\beta_j(\hat{D}^c)^\gamma_k, \quad (1)$$

with $\hat{Q}$, $\hat{U}^c$, $\hat{D}^c$, $\hat{L}$, $\hat{E}^c$ denoting the usual matter chiral superfields [20] and appearing in three generations while $\hat{H}_d$ and $\hat{H}_u$ represent the Higgs supermultiplets. Latin subscripts refer to flavor and Greek superscripts to color. The $SU(2)_L$ invariant product is indicated with '$\cdot$'. Without loss of generality, λ_{ijk} is antisymmetric under exchange of its two first indices, and λ''_{ijk}, under exchange of its two last.

A remarkable scenario of the RpV MSSM [20] involves a bino-dominated neutralino with mass in or below the GeV range [21, 22]. Indeed, as long as the bino mass is not correlated with other gaugino masses, such a light particle escapes constraints from laboratory experiments [23], the main limit being set by the bound on the invisible Z-decay [21], to which the bino would only contribute through subdominant higgsino components. At the LHC, searches with a missing energy signature [24] are in general inapplicable to the RpV case [17, 18, 25]. Similarly, constraints from astrophysical [22, 26, 27] or cosmological [28–30] sources would only apply to a stable particle (or long-lived on cosmological scales). Search proposals for a light unstable or long-lived neutralino in heavy meson, Z or tau decays have been suggested in *e.g.* Refs. [21, 31–39]. Nevertheless, it is clear that a preliminary understanding of the phenomenology of the light fermion, especially its interplay with hadronic matter, is needed in any attempt at quantitatively interpreting the outcome of such searches. Common approximations, such as focusing on individual channels or trivializing mixing effects may be helpful for a qualitative comprehension of the phenomena at stake, but remain insufficient for actual tests and simulations.

In this paper, we specialize to the case of a light bino-dominated neutralino with mass below that of the τ lepton and investigate its decays in an unprejudiced fashion: given the considered field content, we build the most general low-energy effective field theory (EFT) mediating neutralino decays up to dimension 6 terms. While we explicitly fix the short-distance physics through matching conditions with the RpV MSSM, the low-energy analysis should remain valid for a wider range of models predicting a light unstable exotic fermion, such as a right-handed neutrino, see for example Refs. [34, 40, 41]. We also resum QCD UV logarithms at leading order through the renormalization group equations (RGEs). In order to derive the decay widths,

in particular in the hadronic and semi-leptonic channels, we call upon available 'dictionaries' between partonic and hadronic physics. This systematic study definitely improves on earlier estimates and we collect our predictions in a Mathematica package, which is available from the authors upon request.

In the following section, we construct and analyze the low-energy EFT that mediates the decays of the light neutralino, listing all possible operators of dimension 6, matching them to the RpV MSSM and computing the associated RGEs. In section 3, we exploit conversion recipes between the partonic and hadronic languages, in order to derive realistic decay widths of the exotic fermion in terms of the Wilson coefficients of the EFT. Section 4 offers a brief discussion of possible production modes. In section 5, we demonstrate a few numerical applications before a short discussion and conclusion in Section 6. In eleven appendices we present some technical details.

2 Low-energy Effective Field Theory for the Decays of a Light Neutralino

We consider a neutralino state with mass $\lesssim m_\tau$. At such energies, the active partonic degrees of freedom consist of the light quarks (in two-component left-handed spinor notation [44]) $q_i \in \{u,d,s\}$ and antiquarks $q_i^c \in \{u^c,d^c,s^c\}$, the light charged leptons $e_i \in \{e,\mu\}$ and antileptons $e_i^c \in \{e^c,\mu^c\}$, the neutrinos $\nu_i \in \{\nu_e,\nu_\mu,\nu_\tau\}$, the photon A_μ, the gluons G_μ^a and the light neutralino ψ. We assume that there exists no additional light 'exotic' degree of freedom (*e.g.* axinos [42], gravitinos [43]). For a left-handed spinor f, we associate the right-handed conjugate spinor $\bar{f}$. For compactness, we will also employ the four-component spinor notation, with generically $F \equiv (f,\bar{f^c})$, and $Q_i \equiv (q_i,\bar{q}_i^c)^T$, and $E_i \equiv (e_i,\bar{e}_i^c)^T$ representing the Dirac spinors for the quarks and charged leptons, $N_i \equiv (\nu_i,\bar{\nu}_i)^T$ and $\Psi \equiv (\psi,\bar{\psi})^T$ representing the Majorana spinors for the neutrinos and light neutralino, while $P_{L,R}$ denote the chiral projectors. The relevant symmetries are the gauged electromagnetism and chromodynamics $U(1)_{\rm em} \times SU(3)_c$. The light neutralino state is (by definition) a color singlet; limits on light charged $SU(2)_L$ partners, as well as on invisible / exotic Z decays, additionally restrict realistic choices to a dominantly electroweak singlet state. With the MSSM field content, ψ is thus necessarily bino-like [21,23].

The most general renormalizable Lagrangian density at low energy reads:

$$\mathcal{L}_{\rm ren.} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} - \frac{1}{4}G^{a\,\mu\nu}G^a_{\mu\nu} + \sum_{F=Q_i,E_i}\bar{F}\left[i\not{D} - m_F\right]F + \sum_{F=N_i,\Psi}\frac{1}{2}\bar{F}\left[i\not{\partial} - m_f\right]F - m_{\psi\nu_i}\bar{\Psi}N_i\,, \quad (2)$$

where $F_{\mu\nu}$ and $G^a_{\mu\nu}$ represent the field-strength tensors associated with the photon and gluon fields, D_μ is the covariant derivative, and m_f is the mass of the fermion species F. The mass term $m_{\psi\nu_i}$ induces mixing between the neutralino and neutrino fields. It can be absorbed within the definition of ψ, through diagonalization of the mass-matrix. The ν_i are hence no longer electroweak gauge eigenstates, but light mass eigenstates of the neutralino-neutrino sector. So we set $m_{\psi\nu_i} \stackrel{!}{=} 0$ [45,46]. We also systematically neglect the neutrino masses $m_{\nu_i} \ll m_\psi, m_e$ below. Without interactions at the renormalizable level, the light neutralino cannot decay via this Lagrangian. The short-distance effects governing this phenomenon must be introduced through operators of higher dimension.

2.1 Dimension 6 operators for neutralino decay

Given the field content and the symmetries, one can build the following dimension (5 and) 6 operators (together with their hermitian conjugates) involving one neutralino and thus pos-

sibly mediating its decays. (We use here the more compact four-component notation. The lengthier but for computations more convenient two-component expressions are given in Appendix A.)

- Electromagnetic dipoles ($i = 1,2,3$):

$$\mathcal{E}_i \equiv \frac{e}{32\pi^2}(\bar{\Psi}\Sigma^{\mu\nu}P_L N_i)F_{\mu\nu}\,. \tag{3}$$

- Leptonic operators (the chirality indices J,K may take the values L,R; the indices i,j,k correspond to the fermion generation):

$$\begin{aligned}
&\widetilde{\mathcal{N}}_{ijk} \equiv (\bar{\Psi}P_L N_i)(\bar{N}_j^c P_L N_k)\,,\\
&\quad (i,j,k) \in \{(1,2,2),(1,3,3),(1,2,3),(2,1,1),(2,3,3),(2,1,3),(3,1,1),(3,2,2)\}\,,\\
&\mathcal{N}_{ijk} \equiv (\bar{\Psi}\gamma^\mu P_L N_i)(\bar{N}_j\gamma_\mu P_L N_k)\,, \quad (i \leqslant k)\,,\\
&\mathcal{S}_{ijk}^{veK} \equiv (\bar{\Psi}P_L N_i)(\bar{E}_j P_K E_k)\,, \qquad \mathcal{V}_{ijk}^{veK} \equiv (\bar{\Psi}\gamma^\mu P_L N_i)(\bar{E}_j\gamma_\mu P_K E_k)\,,\\
&\mathcal{T}_{ijk}^{ve} \equiv \frac{1}{4}(\bar{\Psi}\Sigma^{\mu\nu}P_L N_i)(\bar{E}_j\Sigma_{\mu\nu}P_L E_k)\,.
\end{aligned} \tag{4}$$

- Semi-leptonic operators ($q = u,d$ separately for the operators involving a neutrino):

$$\begin{aligned}
&\mathcal{S}_{ijk}^{eqJK} \equiv (\bar{\Psi}P_J E_i)(\bar{D}_j P_K U_k)\,, && \mathcal{S}_{ijk}^{vqK} \equiv (\bar{\Psi}P_L N_i)(\bar{Q}_j P_K Q_k)\,,\\
&\mathcal{V}_{ijk}^{eqJK} \equiv (\bar{\Psi}\gamma^\mu P_J E_i)(\bar{D}_j\gamma_\mu P_K U_k)\,, && \mathcal{V}_{ijk}^{vqK} \equiv (\bar{\Psi}\gamma^\mu P_L N_i)(\bar{Q}_j\gamma_\mu P_K Q_k)\,,\\
&\mathcal{T}_{ijk}^{eqJ} \equiv \frac{1}{4}(\bar{\Psi}\Sigma^{\mu\nu}P_J E_i)(\bar{D}_j\Sigma_{\mu\nu}P_J U_k)\,, && \mathcal{T}_{ijk}^{vq} \equiv \frac{1}{4}(\bar{\Psi}\Sigma^{\mu\nu}P_L N_i)(\bar{Q}_j\Sigma_{\mu\nu}P_L Q_k)\,.
\end{aligned} \tag{5}$$

- Hadronic operators (greek indices refer to color; F^c is the charged-conjugate of the fermion F in four-component notation):

$$\begin{aligned}
&\mathcal{H}_{ijk}^{JK} \equiv \varepsilon_{\alpha\beta\gamma}(\bar{\Psi}P_J U_i^{c\,\alpha})(\bar{D}_j^{\,\beta}P_K D_k^{c\,\gamma})\,, \quad (j<k)\,,\\
&\widetilde{\mathcal{H}}_{ijk}^{JK} \equiv \varepsilon_{\alpha\beta\gamma}(\bar{\Psi}P_J D_j^{c\,\beta})(\bar{U}_i^{\,\alpha}P_K D_k^{c\,\gamma})\,.
\end{aligned} \tag{6}$$

  The identity $\widetilde{\mathcal{H}}_{ijk}^{LL} - \widetilde{\mathcal{H}}_{ikj}^{LL} + \mathcal{H}_{ijk}^{LL} = 0$ (and similarly for $LL \to RR$) implies some redundancy (which does not matter as long as the corresponding contributions are not double-counted).

All other combinations of the light fields of dimension 6 can be written in terms of the previous ones, through the two-component spinor (Fierz) identities [44] applied to the expressions in Appendix A, *e.g.*:

$$\begin{aligned}
&(f_1 f_2)(f_3 f_4) + (f_1 f_3)(f_2 f_4) + (f_1 f_4)(f_2 f_3) = 0\,,\\
&(\bar{f}_1\bar{\sigma}^\mu f_2)(\bar{f}_3\bar{\sigma}_\mu f_4) = 2(\bar{f}_1\bar{f}_3)(f_2 f_4)\,,\\
&(\bar{f}_1\bar{\sigma}^\mu f_2)(f_3\sigma_\mu\bar{f}_4) = -2(\bar{f}_1\bar{f}_4)(f_2 f_3)\,,\\
&(f_1\sigma^{\mu\nu}f_2)(f_3\sigma_{\mu\nu}f_4) = -2(f_1 f_4)(f_2 f_3) - (f_1 f_2)(f_3 f_4) = 2(f_1 f_3)(f_2 f_4) + (f_1 f_2)(f_3 f_4)\,.
\end{aligned} \tag{7}$$

Forgetting about neutralino-neutrino mixing, the electromagnetic dipole, leptonic and semi-leptonic operators violate lepton number, by three units for $\widetilde{\mathcal{N}}$ and one unit for the others. The hadronic operators violate baryon number by one unit. We define the Wilson coefficients $C[\Omega]$ associated with any of these operators Ω and add the resulting terms to the Lagrangian density of Eq. (2), hence obtaining our effective field theory with non-renormalizable operators:

$$\mathcal{L}_{\text{dim 6}}^{\text{EFT}} = \mathcal{L}_{\text{ren.}} + \sum_{\Omega} C[\Omega]\,\Omega + h.c. \tag{8}$$

2.2 Matching to the RpV MSSM

For almost any model leading to the considered spectrum in the GeV range, the low-energy dynamics associated with the decays of the light exotic fermion field ψ can be described at leading order by the above EFT.[1] In this paper, however, we choose to focus on a specific type of UV completion, that associated with the RpV MSSM. Then, the short-distance effects contributing to the non-renormalizable operators depend on their mediation by heavy SUSY particles, typically sfermions, or electroweak gauge bosons. We denote the scale corresponding to this UV dynamics as $\mu_0 \sim M_{\rm SUSY}$ and assume that it is not very far from the electroweak scale, *i.e.* in the TeV range. This scale is the high-energy threshold at which the EFT should be matched to the RpV MSSM through the identification of scattering amplitudes (although further intermediate thresholds can be considered). Below, we restrict ourselves to a matching at leading order.

Two types of topologies are involved in mediating the amplitudes that contribute to the lepton- or baryon-number-violating operators leading to neutralino decay. The vector topology employs the electroweak gauge bosons W and Z as mediators and relies exclusively on lepton-electrowikino mixing to generate lepton-number violating amplitudes. On the other hand, the scalar topology proceeds through the exchange of squarks or Higgs-slepton admixtures. Mediators of Higgs-type would convey lepton-number violation through lepton-electrowikino mixing, similarly to the vector topology, and / or through Higgs-slepton mixing. In the case of mediators of sfermion-type, lepton- or baryon-number violation emerges through the trilinear RpV couplings.

For commodity, we introduce the following notation for couplings between a vector V_μ or a scalar S and two fermions F_i, F_j:

$$\mathcal{L}_{\rm MSSM} \ni V_\mu \bar{F}_i \gamma^\mu g_{L,R}^{Vf_i^c f_j} P_{L,R} F_j + S \bar{F}_i g_{L,R}^{Sf_i^c f_j} P_{L,R} F_j \,. \tag{9}$$

Employing the usual algebra, one can derive the following identities:

$$g_{L,R}^{Vf_i^c f_j} = \left(g_{L,R}^{V^* f_j^c f_i}\right)^* = -g_{R,L}^{Vf_j f_i^c} = -\left(g_{R,L}^{V^* f_i f_j^c}\right)^* , \tag{10}$$

$$g_{L,R}^{Sf_i^c f_j} = \left(g_{R,L}^{S^* f_j^c f_i}\right)^* = g_{L,R}^{Sf_j f_i^c} = \left(g_{R,L}^{S^* f_i f_j^c}\right)^* . \tag{11}$$

These rules (for uncolored fields) suffer an exception in the case of non-trivial color products, as those generated by the λ'' couplings, see also Ref. [47]. We list the relevant couplings in Appendix B. Here, we stress [20] that the neutrinos ν_i and the light neutralino ψ are just special (light) occurrences of the neutralino mass eigenstates χ_k^0, while the charged leptons e_i and e_i^c also correspond more generally to light chargino mass eigenstates $\chi_k^\mp$. Similarly, Higgs-slepton mixing leads to a collection of real neutral (pseudo-)scalars S_p^0, built out of the neutral Higgs and sneutrino components, as well as to charged Higgs-slepton admixtures $S_p^\pm$. Squarks are denoted as $\tilde{Q}_p \in \{\tilde{U}_p, \tilde{D}_p\}$ and in general involve left-right as well as generation mixing. The mixing among fields of the RpV MSSM is shortly described in Appendix C.

Let us first consider the decay $\psi \to \bar{\nu}_i \nu_j \bar{\nu}_k$ ($i \leqslant k$): the amplitude in the EFT reads (where $\breve{\Omega}$ represents the spinor structure resulting from the operator Ω):

$$\mathcal{A}^{\rm EFT}[\psi \to \bar{\nu}_i \nu_j \bar{\nu}_k] = i(1+\delta_{ik})\, C[\mathcal{N}_{ijk}](\mu_0)\, \breve{\mathcal{N}}_{ijk} \,. \tag{12}$$

[1] The exception is if specific counting rules suppress the dimension 6 operators with respect to higher-order ones.

The corresponding amplitude in the RpV MSSM (neglecting Yukawa couplings) is mediated either by a Z-boson or by a neutral scalar, and reads (with implicit summation):

$$\mathcal{A}^{\text{MSSM}}[\psi \to \bar{\nu}_i \nu_j \bar{\nu}_k] = i \left\{ \frac{1}{2m^2_{S^0_p}} g_R^{S^0_p \nu^c_j \psi} g_L^{S^0_p \nu_i \nu_k} - \frac{1}{M_Z^2} \left[g_L^{Z\psi\nu_i} g_L^{Z\nu^c_j \nu_k} + g_L^{Z\psi\nu_k} g_L^{Z\nu^c_j \nu_i} \right] \right\} \breve{\mathcal{N}}_{ijk} \,. \quad (13)$$

Identifying both amplitudes, we deduce:

$$C[\mathcal{N}_{ijk}](\mu_0) \stackrel{!}{=} \frac{1}{1+\delta_{ik}} \left\{ \frac{1}{2m^2_{S^0_p}} g_R^{S^0_p \nu^c_j \psi} g_L^{S^0_p \nu_i \nu_k} - \frac{1}{M_Z^2} \left[g_L^{Z\psi\nu_i} g_L^{Z\nu^c_j \nu_k} + g_L^{Z\psi\nu_k} g_L^{Z\nu^c_j \nu_i} \right] \right\} . \quad (14)$$

Performing the same operation for the decay $\psi \to \bar{\nu}_i \bar{\nu}_j \bar{\nu}_k$ (for the same combinations of flavor indices as those of the operators $\widetilde{\mathcal{N}}_{ijk}$ in Eq. (4), and in particular $i \neq j,k$, $j \leqslant k$), we can determine:

$$C[\widetilde{\mathcal{N}}_{ijk}](\mu_0) \stackrel{!}{=} \frac{1}{(1+\delta_{jk}) m^2_{S^0_p}} \left[g_L^{S^0_p \psi \nu_i} g_L^{S^0_p \nu_j \nu_k} - g_L^{S^0_p \psi \nu_k} g_L^{S^0_p \nu_j \nu_i} \right] . \quad (15)$$

Given that the corresponding operators violate lepton-number by three units, they require a high degree of mixing and may thus be seen as typically suppressed.

Turning to the decay $\psi \to \bar{\nu}_i e_j \bar{e}_k$, where both vector and scalar mediators contribute, we can determine the remaining Wilson coefficients of leptonic type:

$$\begin{aligned}
C[\mathcal{S}^{veL}_{ijk}](\mu_0) &\stackrel{!}{=} \frac{1}{m^2_{S^0_p}} g_L^{S^0_p \psi \nu_i} g_L^{S^0_p e^c_j e_k} - \frac{1}{2m^2_{S^\pm_p}} \left[g_L^{S^+_p \psi e_k} g_L^{S^-_p e^c_j \nu_i} + g_L^{S^-_p \psi e^c_j} g_L^{S^+_p e_k \nu_i} \right] , \\
C[\mathcal{S}^{veR}_{ijk}](\mu_0) &\stackrel{!}{=} \frac{2}{M_W^2} g_L^{W\psi e_k} g_L^{We^c_j \nu_i} + \frac{1}{m^2_{S^0_p}} g_L^{S^0_p \psi \nu_i} g_R^{S^0_p e^c_j e_k} , \\
C[\mathcal{V}^{veL}_{ijk}](\mu_0) &\stackrel{!}{=} -\frac{1}{M_Z^2} g_L^{Z\psi\nu_i} g_L^{Ze^c_j e_k} - \frac{1}{M_W^2} g_L^{W\psi e_k} g_L^{We^c_j \nu_i} + \frac{1}{2m^2_{S^\pm_p}} g_R^{S^-_p \psi e^c_j} g_L^{S^+_p e_k \nu_i} , \\
C[\mathcal{V}^{veR}_{ijk}](\mu_0) &\stackrel{!}{=} -\frac{1}{M_Z^2} g_L^{Z\psi\nu_i} g_R^{Ze^c_j e_k} - \frac{1}{2m^2_{S^\pm_p}} g_R^{S^+_p \psi e_k} g_L^{S^-_p e^c_j \nu_i} , \\
C[\mathcal{T}^{ve}_{ijk}](\mu_0) &\stackrel{!}{=} \frac{1}{2m^2_{S^\pm_p}} \left[g_L^{S^-_p \psi e^c_j} g_L^{S^+_p e_k \nu_i} - g_L^{S^+_p \psi e_k} g_L^{S^-_p e^c_j \nu_i} \right] .
\end{aligned} \quad (16)$$

The decay amplitudes $\psi \to \bar{\nu}_i q_j \bar{q}_k$ ($q = u,d$) and $\psi \to \bar{e}_i u_j \bar{d}_k$ allow the identification of the Wilson coefficients of the semi-leptonic operators:

$$\begin{aligned}
C[\mathcal{S}^{vqL}_{ijk}](\mu_0) &\stackrel{!}{=} \frac{1}{m^2_{S^0_p}} g_L^{S^0_p \psi \nu_i} g_L^{S^0_p q^c_j q_k} - \frac{1}{2m^2_{\tilde{Q}_p}} \left[g_L^{\tilde{Q}^*_p \psi q_k} g_L^{\tilde{Q}_p q^c_j \nu_i} + g_L^{\tilde{Q}_p \psi q^c_j} g_L^{\tilde{Q}^*_p q_k \nu_i} \right] , \\
C[\mathcal{S}^{vqR}_{ijk}](\mu_0) &\stackrel{!}{=} \frac{1}{m^2_{S^0_p}} g_L^{S^0_p \psi \nu_i} g_R^{S^0_p q^c_j q_k} , \\
C[\mathcal{V}^{vqL}_{ijk}](\mu_0) &\stackrel{!}{=} -\frac{1}{M_Z^2} g_L^{Z\psi\nu_i} g_L^{Zq^c_j q_k} + \frac{1}{2m^2_{\tilde{Q}_p}} g_R^{\tilde{Q}_p \psi q^c_j} g_L^{\tilde{Q}^*_p q_k \nu_i} , \\
C[\mathcal{V}^{vqR}_{ijk}](\mu_0) &\stackrel{!}{=} -\frac{1}{M_Z^2} g_L^{Z\psi\nu_i} g_R^{Zq^c_j q_k} - \frac{1}{2m^2_{\tilde{Q}_p}} g_R^{\tilde{Q}^*_p \psi q_k} g_L^{\tilde{Q}_p q^c_j \nu_i} , \\
C[\mathcal{T}^{vq}_{ijk}](\mu_0) &\stackrel{!}{=} \frac{1}{2m^2_{\tilde{Q}_p}} \left[g_L^{\tilde{Q}_p \psi q^c_j} g_L^{\tilde{Q}^*_p q_k \nu_i} - g_L^{\tilde{Q}^*_p \psi q_k} g_L^{\tilde{Q}_p q^c_j \nu_i} \right] ,
\end{aligned} \quad (17)$$

$$
\begin{aligned}
C[\mathcal{S}^{eq\,LL}_{ijk}](\mu_0) &\overset{!}{=} \frac{1}{m^2_{S^\pm_p}} g_L^{S^+_p \psi e_i} g_L^{S^-_p d^c_j u_k} - \frac{1}{2m^2_{\tilde{U}_p}} g_L^{\tilde{U}^*_p \psi u_k} g_L^{\tilde{U}_p d^c_j e_i} - \frac{1}{2m^2_{\tilde{D}_p}} g_L^{\tilde{D}_p \psi d^c_j} g_L^{\tilde{D}^*_p u_k e_i}\,,\\
C[\mathcal{S}^{eq\,LR}_{ijk}](\mu_0) &\overset{!}{=} \frac{1}{m^2_{S^\pm_p}} g_L^{S^+_p \psi e_i} g_R^{S^-_p d^c_j u_k}\,,\\
C[\mathcal{V}^{eq\,LL}_{ijk}](\mu_0) &\overset{!}{=} -\frac{1}{M_W^2} g_L^{W \psi e_i} g_L^{W d^c_j u_k} + \frac{1}{2m^2_{\tilde{D}_p}} g_R^{\tilde{D}_p \psi d^c_j} g_L^{\tilde{D}^*_p u_k e_i}\,,\\
C[\mathcal{V}^{eq\,LR}_{ijk}](\mu_0) &\overset{!}{=} -\frac{1}{M_W^2} g_L^{W \psi e_i} g_R^{W d^c_j u_k} - \frac{1}{2m^2_{\tilde{U}_p}} g_R^{\tilde{U}^*_p \psi u_k} g_L^{\tilde{U}_p d^c_j e_i}\,,\\
C[\mathcal{T}^{eq\,L}_{ijk}](\mu_0) &\overset{!}{=} -\frac{1}{2m^2_{\tilde{U}_p}} g_L^{\tilde{U}^*_p \psi u_k} g_L^{\tilde{U}_p d^c_j e_i} + \frac{1}{2m^2_{\tilde{D}_p}} g_L^{\tilde{D}_p \psi d^c_j} g_L^{\tilde{D}^*_p u_k e_i}\,,\\
C[\mathcal{S}^{eq\,RL}_{ijk}](\mu_0) &\overset{!}{=} \frac{1}{m^2_{S^\pm_p}} g_R^{S^+_p \psi e_i} g_L^{S^-_p d^c_j u_k}\,,\\
C[\mathcal{S}^{eq\,RR}_{ijk}](\mu_0) &\overset{!}{=} \frac{1}{m^2_{S^\pm_p}} g_R^{S^+_p \psi e_i} g_R^{S^-_p d^c_j u_k} - \frac{1}{2m^2_{\tilde{U}_p}} g_R^{\tilde{U}^*_p \psi u_k} g_R^{\tilde{U}_p d^c_j e_i} - \frac{1}{2m^2_{\tilde{D}_p}} g_R^{\tilde{D}_p \psi d^c_j} g_R^{\tilde{D}^*_p u_k e_i}\,,\\
C[\mathcal{V}^{eq\,RL}_{ijk}](\mu_0) &\overset{!}{=} -\frac{1}{M_W^2} g_R^{W \psi e_i} g_L^{W d^c_j u_k} - \frac{1}{2m^2_{\tilde{U}_p}} g_L^{\tilde{U}^*_p \psi u_k} g_R^{\tilde{U}_p d^c_j e_i}\,,\\
C[\mathcal{V}^{eq\,RR}_{ijk}](\mu_0) &\overset{!}{=} -\frac{1}{M_W^2} g_R^{W \psi e_i} g_R^{W d^c_j u_k} + \frac{1}{2m^2_{\tilde{D}_p}} g_L^{\tilde{D}_p \psi d^c_j} g_R^{\tilde{D}^*_p u_k e_i}\,,\\
C[\mathcal{T}^{eq\,R}_{ijk}](\mu_0) &\overset{!}{=} \frac{1}{2m^2_{\tilde{D}_p}} g_R^{\tilde{D}_p \psi d^c_j} g_R^{\tilde{D}^*_p u_k e_i} - \frac{1}{2m^2_{\tilde{U}_p}} g_R^{\tilde{U}^*_p \psi u_k} g_R^{\tilde{U}_p d^c_j e_i}\,.
\end{aligned}
$$

We already derived the Wilson coefficients of the hadronic operators in Ref. [48]. For completeness, with $j \leqslant k$ (by assumption):

$$
\begin{aligned}
C[\mathcal{H}^{LL}_{ijk}](\mu_0) &\overset{!}{=} \frac{1}{m^2_{\tilde{U}_p}} g_L^{\tilde{U}_p \psi u^c_i} g_L^{\tilde{U}^*_p d^c_j d^c_k}\,, & C[\mathcal{H}^{RL}_{ijk}](\mu_0) &\overset{!}{=} \frac{1}{m^2_{\tilde{U}_p}} g_R^{\tilde{U}_p \psi u^c_i} g_L^{\tilde{U}^*_p d^c_j d^c_k}\,,\\
C[\widetilde{\mathcal{H}}^{LL}_{ijk}](\mu_0) &\overset{!}{=} \frac{1}{m^2_{\tilde{D}_p}} g_L^{\tilde{D}_p \psi d^c_j} g_L^{\tilde{D}^*_p u^c_i d^c_k}\,, & C[\widetilde{\mathcal{H}}^{RL}_{ijk}](\mu_0) &\overset{!}{=} \frac{1}{m^2_{\tilde{D}_p}} g_R^{\tilde{D}_p \psi d^c_j} g_L^{\tilde{D}^*_p u^c_i d^c_k}\,,\\
C[\widetilde{\mathcal{H}}^{LL}_{ikj}](\mu_0) &\overset{!}{=} \frac{1}{m^2_{\tilde{D}_p}} g_L^{\tilde{D}_p \psi d^c_k} g_L^{\tilde{D}^*_p u^c_i d^c_j}\,, & C[\widetilde{\mathcal{H}}^{RL}_{ikj}](\mu_0) &\overset{!}{=} \frac{1}{m^2_{\tilde{D}_p}} g_R^{\tilde{D}_p \psi d^c_k} g_L^{\tilde{D}^*_p u^c_i d^c_j}\,.
\end{aligned}
\tag{18}
$$

As a consequence of the redundancy $\widetilde{\mathcal{H}}^{LL}_{ikj} = \widetilde{\mathcal{H}}^{LL}_{ijk} + \mathcal{H}^{LL}_{ijk}$ that we mentioned earlier, $C[\widetilde{\mathcal{H}}^{LL}_{ikj}](\mu_0)$ could be distributed onto the other Wilson coefficients. In addition, $g_R^{\tilde{U}^*_p d^c_j d^c_k} \overset{!}{=} 0 \overset{!}{=} g_R^{\tilde{D}^*_p u^c_i d^c_k}$ in the RpV MSSM implies that the remaining effective couplings vanish at tree-level order: $C[\mathcal{H}^{LR}_{ijk}](\mu_0) = 0 = C[\mathcal{H}^{RR}_{ijk}](\mu_0) = C[\widetilde{\mathcal{H}}^{LR}_{ijk}](\mu_0) = C[\widetilde{\mathcal{H}}^{RR}_{ijk}](\mu_0)$.

The neutralino decay into antineutrino and photon, $\psi \to \gamma \bar{\nu}_i$, is a one-loop process in the RpV MSSM and it might be misleading to include it in this analysis while other channels are considered at tree-level. Nevertheless, it was calculated under various approximations in *e.g.* Refs. [45, 49, 50] and can also be viewed as a radiative neutralino decay [51]. Let us parametrize the decay amplitude as follows:

$$
\mathcal{A}^{\text{MSSM}}[\psi(p_\psi) \to \gamma(p_\gamma)\bar{\nu}_i(p_{\nu_i})] \equiv -\frac{e}{16\pi^2}(p_\gamma)_\mu \varepsilon^*(p_\gamma)_\nu \big(\bar{\nu}_\psi(p_\psi)\Sigma^{\mu\nu} P_L \nu_{\nu_i}(p_{\nu_i})\big) C^{\text{MSSM}}[\mathcal{E}_i]\,. \tag{19}
$$

We compute $C^{\text{MSSM}}[\mathcal{E}_i]$, *cf.* Eq. (3), in the Feynman gauge. At one-loop order, it involves vertex diagrams of W-bosons (or charged a Goldstone) and chargino/leptons ($\chi^\pm_j$), as well as loops of scalar fields (S) with an associated fermion (f). The $Z-\gamma$ mixing diagrams ensure

UV-finiteness with the triangle diagrams. Combining all contributions,

$$C^{\rm MSSM}[\mathcal{E}_i] = \frac{1}{M_W^2}\Big[m_\psi\big(g_L^{W^+\psi\chi_j^-}\, g_L^{W^-\chi_j^+\nu} - g_L^{W^-\psi\chi_j^+}\, g_L^{W^+\chi_j^-\nu}\big)\mathcal{I}_L^V\big(\tfrac{m_\psi^2}{M_W^2}, \tfrac{m^2_{\chi_j^\pm}}{M_W^2}\big) + m_{\chi_j}\big(g_R^{W^+\psi\chi_j^-}\, g_L^{W^-\chi_j^+\nu} - g_R^{W^-\psi\chi_j^+}\, g_L^{W^+\chi_j^-\nu}\big)\mathcal{I}_R^V\big(\tfrac{m_\psi^2}{M_W^2}, \tfrac{m^2_{\chi_j^\pm}}{M_W^2}\big)\Big] + \frac{Q_f}{m_S^2}\Big[m_f\big(g_L^{S^*\psi f}\, g_L^{Sf^c\nu} - g_L^{S\psi f^c}\, g_L^{S^* f\nu}\big)\mathcal{I}_L^S\big(\tfrac{m_\psi^2}{m_S^2}, \tfrac{m_f^2}{m_S^2}\big) + m_\psi\big(g_R^{S^*\psi f}\, g_L^{Sf^c\nu} - g_R^{S\psi f^c}\, g_L^{S^* f\nu}\big)\mathcal{I}_R^S\big(\tfrac{m_\psi^2}{m_S^2}, \tfrac{m_f^2}{m_S^2}\big)\Big]. \tag{20}$$

The implicit sum over $\chi_j^\pm$ runs over the five corresponding mass eigenstates. For the scalar diagrams, contained in the second set of brackets, the implicit sum runs over all (S, f, Q_f)'s in the following list: $\big\{(S_m^-, \chi_j^-, -1), (\tilde{U}_m, u, \frac{2}{3}), (\tilde{D}_m, d, -\frac{1}{3})\big\}$, including color factors for the quark/squark contributions. The loop integrals $\mathcal{I}_{L,R}^{V,S}$ are provided in Appendix D.

In the EFT, the amplitude corresponding to the $\psi \to \gamma\, \bar{\nu}_i$ transition receives contributions from $\mathcal{E}_i$ at tree-level, but also (at competitive order) from one-loop diagrams involving the operators $\mathcal{T}_{ijj}^{\nu f\,(L)}$, $f = e, q$, $j = 1, 2, 3$. In particular, it is necessary to decouple step by step the heavy fermions (t, b) that are integrated out at low-energy but not at the scale μ_0. The amplitude reads:

$$\mathcal{A}^{\rm EFT}[\psi(p_\psi) \to \gamma(p_\gamma)\bar{\nu}_i(p_{\nu_i})] = -\frac{e}{16\pi^2}(p_\gamma)_\mu \varepsilon^*(p_\gamma)_\nu \big(\bar{\nu}_\psi(p_\psi)\Sigma^{\mu\nu}P_L\nu_{\nu_i}(p_{\nu_i})\big) \times \left\{C[\mathcal{E}_i] + \delta C[\mathcal{E}_i] + \sum_{f_j} 2Q_{f_j} N_c^{f_j} m_{f_j} C[\mathcal{T}_{ijj}^{\nu f\,(L)}]\left(\frac{1}{2-\frac{D}{2}} - \gamma_E + \ln\frac{4\pi\mu_0^2}{m_{f_j}^2}\right)\right\}(\mu_0), \tag{21}$$

where D represents the dimension used for the regularization of the loop integrals, γ_E is the Euler-Mascheroni constant, μ_0 is the renormalization scale (as already implied), Q_{f_j} is the electric charge of the fermion species f_j, $N_c^{f_j} = 1$ or 3 for color singlets and triplets, respectively. $\delta C[\mathcal{E}_i]$ denotes the counterterm associated with $\mathcal{E}_i$, which is needed for the renormalization of the amplitude. Therefore, the matching condition connecting $C[\mathcal{E}_i](\mu_0)$ and $C^{\rm MSSM}[\mathcal{E}_i]$ is not completely trivial, as the $\mathcal{T}_{ijj}^{\nu f\,(L)}$ continue to contribute below μ_0: the renormalization of the EFT will clarify the meaning of this decomposition for the bookkeeping of logarithmic corrections.

Under the assumption of subdominant mixing and negligible Yukawa couplings, it is possible to cast the contributions to the Wilson coefficients in a much simpler form, where trilinear RpV couplings dominate the phenomenology associated with the light bino-like neutralino: see Appendix E. Comparison with Refs. [33, 39] for the contributions to semi-leptonic operators revealed probable issues in these references with effects of tensor type.

2.3 Renormalization group evolution under QCD

While the short-distance effects that generate the contributions to the operators mediating the neutralino decay are characterized by the high-energy scale μ_0, the physical transition itself (the neutralino decay) takes place at a low energy determined by m_ψ. As perturbative QCD fails at low-energy scales, we will employ an infrared scale $\mu \gtrsim m_c$. Given the large hierarchy between μ_0 and μ, large QCD logarithms of UV type are expected to develop and it thus appears meaningful to resum them through the renormalization group evolution within the EFT. To this end, we renormalize all the parameters of the EFT in the $\overline{\rm MS}$ scheme (following

the traditional procedure, described in *e.g.* Ref. [52] in the context of B decays), thus deriving the Callan-Symanzik equations associated with the Wilson coefficients at leading logarithmic order in QCD.

For commodity, we introduce the strong coupling constant $\alpha_s \equiv \frac{g_s^2}{4\pi}$, the coefficient of the quadratic Casimir operator $C_2(3) \equiv \frac{4}{3}$, the leading contribution to the QCD beta function $\beta_0[n_f] \equiv \frac{11N_c - 2n_f}{3}$, with $N_c \equiv 3$ the number of colors, and n_f the number of active quark flavors. We also employ the following notation for the ratios of the α_s's at various scales: $\eta_\mu \equiv \frac{\alpha_s(\mu)}{\alpha_s(m_b)}$, $\eta_b \equiv \frac{\alpha_s(m_b)}{\alpha_s(m_t)}$, and $\eta_t \equiv \frac{\alpha_s(m_t)}{\alpha_s(\mu_0)}$. The quark masses $m_{t,b}$ appearing in these equations correspond to decoupling thresholds and should be understood as the $\overline{\text{MS}}$ masses $m_{t,b}(m_{t,b})$. The RGE's associated with the hadronic operators were previously derived in Ref. [53] and lead to the following evolution (for $m_c \lesssim \mu < m_b$):

$$C[\mathcal{H}^{JK}_{ijk}](\mu) = \eta_\mu^{2/\beta_0[4]}\, \eta_b^{2/\beta_0[5]}\, \eta_t^{2/\beta_0[6]}\, C[\mathcal{H}^{JK}_{ijk}](\mu_0)\,, \tag{22}$$

and similarly for the operators of $\widetilde{\mathcal{H}}$-type. To appreciate the order of magnitude of this correction, the choice ($\mu \overset{!}{=} 2\,\text{GeV}$, $\mu_0 \overset{!}{=} 1\,\text{TeV}$) returns an enhancement factor of about 1.4 at low energy.

The semi-leptonic operators ($\ell = \nu, e$) receive the following scaling with energy:

$$\begin{aligned}
C[\mathcal{S}^{\ell qJK}_{ijk}](\mu) &= \eta_\mu^{3C_2(3)/\beta_0[4]}\, \eta_b^{3C_2(3)/\beta_0[5]}\, \eta_t^{3C_2(3)/\beta_0[6]}\, C[\mathcal{S}^{\ell qJK}_{ijk}](\mu_0)\,,\\
C[\mathcal{V}^{\ell qJK}_{ijk}](\mu) &= C[\mathcal{V}^{\ell qJK}_{ijk}](\mu_0)\,,\\
C[\mathcal{T}^{\ell qJ}_{ijk}](\mu) &= \eta_\mu^{-C_2(3)/\beta_0[4]}\, \eta_b^{-C_2(3)/\beta_0[5]}\, \eta_t^{-C_2(3)/\beta_0[6]}\, C[\mathcal{T}^{\ell qJ}_{ijk}](\mu_0)\,.
\end{aligned} \tag{23}$$

We here recover the usual scaling of scalar semi-leptonic operators, similar to a quark mass, while the vector operators have a vanishing anomalous dimension. However, we remind the reader that these simple scaling conditions narrowly depend on the choice of operator basis.

For instance, replacing $\mathcal{S}^{\ell qJ}_{ijk}$ and $\mathcal{T}^{\ell qJ}_{ijk}$ by $\hat{\mathcal{S}}^{\ell qJK}_{ijk} \equiv \mathcal{S}^{\ell qJK}_{ijk}$ and $\widetilde{\mathcal{S}}^{\ell qJ}_{ijk} \equiv (\bar{\Psi} P_J Q_k)(\bar{Q}_j P_J E_i) = -\frac{1}{2}\mathcal{T}^{\ell qJ}_{ijk} - \frac{1}{2}\mathcal{S}^{\ell qJJ}_{ijk}$ would lead to a hybrid running for $C[\hat{\mathcal{S}}^{\ell qJJ}_{ijk}]$ (even though $\hat{\mathcal{S}}^{\ell qJJ}_{ijk} = \mathcal{S}^{\ell qJJ}_{ijk}$). The choice ($\mu \overset{!}{=} 2\,\text{GeV}$, $\mu_0 \overset{!}{=} 1\,\text{TeV}$) implies an enhancement factor of about 1.9 in the case of scalar coefficients, while tensor coefficients are suppressed by about 0.8, in this example.

The leptonic operators do not receive QCD corrections at leading order. In the case of the electric dipole operator, the corrections from the $\mathcal{T}^{\nu f}_{ijj}$ operators — see Eq. (21) — require the renormalization of $\mathcal{E}_i$ and imply running conditions for the Wilson coefficient:

$$\begin{aligned}
C[\mathcal{E}_i](\mu) &= C[\mathcal{E}_i](\mu_0) - \sum_{e_j} 2Q_{e_j} m_{e_j} C[\mathcal{T}^{\nu e}_{ijj}](\mu_0)\, \ln \frac{\text{Max}[\mu^2, m_{e_j}^2]}{\mu_0^2}\\
&- \sum_{q_j} \frac{8\pi N_c Q_{q_j}}{\beta_0[6] - 2C_2[3]} \left[\eta_t^{2C_2(3)/\beta_0[6]-1} - 1\right] \left(\frac{m_{q_j}}{\alpha_s} C[\mathcal{T}^{\nu q\,L}_{ijj}]\right)(\mu_0)\\
&- \sum_{q_j \neq t} \frac{8\pi N_c Q_{q_j}}{\beta_0[5] - 2C_2[3]} \eta_t^{2C_2(3)/\beta_0[6]-1} \left[\eta_b^{2C_2(3)/\beta_0[5]-1} - 1\right] \left(\frac{m_{q_j}}{\alpha_s} C[\mathcal{T}^{\nu q\,L}_{ijj}]\right)(\mu_0)\\
&- \sum_{q_j \neq t,b} \frac{8\pi N_c Q_{q_j}}{\beta_0[4] - 2C_2[3]} \eta_t^{2C_2(3)/\beta_0[6]-1} \eta_b^{2C_2(3)/\beta_0[5]-1} \left[\eta_\mu^{2C_2(3)/\beta_0[4]-1} - 1\right] \left(\frac{m_{q_j}}{\alpha_s} C[\mathcal{T}^{\nu q\,L}_{ijj}]\right)(\mu_0)\,.
\end{aligned} \tag{24}$$

The significance of these corrections is made more explicit after expressing $C[\mathcal{E}_i](\mu_0)$ in terms of the corresponding coefficient in the RpV MSSM — see Eq. (21):

$$\begin{aligned}
C[\mathcal{E}_i](\mu) = C^{\text{MSSM}}[\mathcal{E}_i] &- \sum_{e_j} 2Q_{e_j} m_{e_j} C[\mathcal{T}^{\nu e}_{ijj}](\mu_0) \ln \frac{\text{Max}[\mu^2, m^2_{e_j}]}{m^2_{e_j}} \\
&- \sum_{q_j} 8\pi N_c Q_{q_j} \left[\frac{\eta_t^{2C_2(3)/\beta_0[6]-1}-1}{\beta_0[6]-2C_2[3]} + \frac{\alpha_s(\mu_0)}{4\pi} \ln \frac{\mu_0^2}{m_t^2} \right] \left(\frac{m_{q_j}}{\alpha_s} C[\mathcal{T}^{\nu q\, L}_{ijj}] \right)(\mu_0) \\
&- \sum_{q_j \neq t} 8\pi N_c Q_{q_j} \left[\eta_t^{2C_2(3)/\beta_0[6]-1} \frac{\eta_b^{2C_2(3)/\beta_0[5]-1}-1}{\beta_0[5]-2C_2[3]} + \frac{\alpha_s(\mu_0)}{4\pi} \ln \frac{m_t^2}{m_b^2} \right] \left(\frac{m_{q_j}}{\alpha_s} C[\mathcal{T}^{\nu q\, L}_{ijj}] \right)(\mu_0) \\
&- \sum_{q_j \neq t,b} 8\pi N_c Q_{q_j} \left[\eta_t^{2C_2(3)/\beta_0[6]-1} \eta_b^{2C_2(3)/\beta_0[5]-1} \frac{\eta_\mu^{2C_2(3)/\beta_0[4]-1}-1}{\beta_0[4]-2C_2[3]} + \frac{\alpha_s(\mu_0)}{4\pi} \ln \frac{m_b^2}{m_{q_j}^2} \right] \left(\frac{m_{q_j}}{\alpha_s} C[\mathcal{T}^{\nu q\, L}_{ijj}] \right)(\mu_0) .
\end{aligned} \tag{25}$$

It is thus evident that the scale evolution re-inserts in the Wilson coefficient those pieces that were subtracted at the level of the matching for being captured by the semi-leptonic operators. However, quark loops are now included in a form resumming leading QCD logarithms between μ_0 and μ. Finally, we stress that the decay amplitude at the scale μ adds a further factor $2N_c^{f_j} Q_{f_j} \left(m_{f_j} C[\mathcal{T}^{\nu f}_{ijj}] \right)(\mu) \ln \frac{\mu^2}{m^2_{f_j}}$, for all active fermion fields f_j at the scale μ, to the decay amplitude at this scale.

This step completes the determination of the Wilson coefficients at low energy, hence the definition of the EFT for phenomenological use.

3 Decay Widths of the Neutralino

After determining the Wilson coefficients at high-energy through matching, and running them down to the infrared scale, we have obtained a working partonic theory describing the low-energy phenomenology of the light neutralino. In order to make meaningful predictions of the decay widths, it is then necessary to translate the partonic operators into interactions involving the actual physical degrees of freedom at these scales, *i.e.* hadrons instead of quarks.

3.1 Radiative and leptonic decays

The decay widths in the radiative and leptonic channels can be derived in a straightforward fashion. In this paper, we systematically neglect interaction effects between final states, which, strictly speaking, is only valid in situations with enough kinetic energy.

The two-body radiative decay is kinematically open at neutralino masses above the neutrino masses, *i.e.* essentially in all the considered range for m_ψ. However, the choice $\mu \stackrel{!}{=} m_\psi$ for the renormalization scale would be problematic for $m_\psi \lesssim m_c$, as the perturbative description of the strong coupling loses significance. Maintaining $\mu \gtrsim m_c$ means that the resummation of QCD logarithms is not comprehensive for light quarks. Yet, given that the contributions of light fermions to the decay amplitude is suppressed in proportion to their mass, light quark contributions are non-essential. The decay width then reads (for $\mu \gtrsim m_c$, $i = 1,2,3$):

$$\Gamma[\psi \to \gamma \bar{\nu}_i] = \frac{\alpha m_\psi^3}{4096\pi^4} \left| C[\mathcal{E}_i](\mu) + \sum_{f_j \neq t,b,\ldots} 2Q_{f_j} N_c^{f_j} m_{f_j} \ln \frac{\mu^2}{m^2_{f_j}} C[\mathcal{T}^{\nu f}_{ijj}](\mu) \right|^2 = \Gamma[\psi \to \gamma \nu_i] . \tag{26}$$

The decays into three neutrinos are also accessible down to very low m_ψ. As we restricted ourselves to resumming QCD logarithms, the Wilson coefficients of leptonic operators do not

run, and the choice of μ does not matter here. For the seven independent three-neutrino channels $\nu_i\nu_j\nu_k$ ($i\neq j=k$ and $(i,j,k)=(1,2,3)$):

$$\Gamma[\psi\to\bar\nu_i\bar\nu_j\bar\nu_j]=\frac{m_\psi^5}{3072\pi^3}\left|C[\widetilde{\mathcal{N}}_{ijj}](\mu)\right|^2=\Gamma[\psi\to\nu_i\nu_j\nu_k]\,, \tag{27}$$

$$\begin{aligned}\Gamma[\psi\to\bar\nu_1\bar\nu_2\bar\nu_3]&=\frac{m_\psi^5}{6144\pi^3}\Big\{\left|C[\widetilde{\mathcal{N}}_{123}](\mu)\right|^2+\left|C[\widetilde{\mathcal{N}}_{213}](\mu)\right|^2-\mathrm{Re}\big[C[\widetilde{\mathcal{N}}_{123}](\mu)^*C[\widetilde{\mathcal{N}}_{213}](\mu)\big]\Big\}\\&=\Gamma[\psi\to\nu_1\nu_2\nu_3]\,.\end{aligned}$$

Similarly, the neutralino decays into two antineutrinos ($i\leqslant k$) and one neutrino (j) read:

$$\Gamma[\psi\to\bar\nu_i\nu_j\bar\nu_k]=\frac{1+\delta_{ik}}{1536\pi^3}m_\psi^5\left|C[\mathcal{N}_{ijk}](\mu)\right|^2=\Gamma[\psi\to\nu_i\bar\nu_j\nu_k]\,, \tag{28}$$

corresponding to 18 channels (+ CP-conjugates).

The decays into an antineutrino and a pair of charged leptons is computed to be:

$$\Gamma[\psi\to\bar\nu_i e_j\bar e_k]=\frac{m_\psi^5}{256\pi^3}\sum_{\Omega,\Omega'}C[\Omega]C[\Omega']^*I_{\Omega,\Omega'}\big[\tfrac{m_{e_j}}{m_\psi},\tfrac{m_{e_k}}{m_\psi}\big]=\Gamma[\psi\to\nu_i\bar e_j e_k]\,, \tag{29}$$

where the sums over Ω and Ω' run over the list of leptonic νe operators (see Eq. (4)) and the kinematic integrals $I_{\Omega,\Omega'}$ can be explicitly calculated: see Appendix F.

3.2 Hadronic decays

In the presence of baryon-number violation, the hadronic operators of Eq. (6) generate interactions between ψ and the baryons that are difficult to characterize on a quantitative basis. We will restrict ourselves to the chiral picture to shed some light at the qualitative level, considering only the baryon and meson octets of lowest energy. We refer the reader to Refs. [54, 55] for a more detailed derivation of the formalism.

We introduce the light meson and baryon octets of the chiral approach in $3\times\bar3$ space of the vectorial $(u,d,s)^T$ $SU(3)$ symmetry:

$$M\equiv\begin{bmatrix}\frac{\pi^0}{\sqrt2}+\frac{\eta_8^0}{\sqrt6} & \pi^+ & K^+\\ \pi^- & -\frac{\pi^0}{\sqrt2}+\frac{\eta_8^0}{\sqrt6} & K^0\\ K^- & \bar K^0 & -\sqrt{\frac23}\eta_8^0\end{bmatrix}\sim\begin{bmatrix}u^\alpha\bar u^\alpha & u^\alpha\bar d^\alpha & u^\alpha\bar s^\alpha\\ d^\alpha\bar u^\alpha & d^\alpha\bar d^\alpha & d^\alpha\bar s^\alpha\\ s^\alpha\bar u^\alpha & s^\alpha\bar d^\alpha & s^\alpha\bar s^\alpha\end{bmatrix}, \tag{30}$$

$$B\equiv\begin{bmatrix}\frac{\Sigma^0}{\sqrt2}+\frac{\Lambda^0}{\sqrt6} & \Sigma^+ & p^+\\ \Sigma^- & -\frac{\Sigma^0}{\sqrt2}+\frac{\Lambda^0}{\sqrt6} & n^0\\ \Xi^- & \Xi^0 & -\sqrt{\frac23}\Lambda^0\end{bmatrix}\sim\varepsilon_{\alpha\beta\gamma}\begin{bmatrix}u^\alpha(d^\beta s^\gamma) & u^\alpha(s^\beta u^\gamma) & u^\alpha(u^\beta d^\gamma)\\ d^\alpha(d^\beta s^\gamma) & d^\alpha(s^\beta u^\gamma) & d^\alpha(u^\beta d^\gamma)\\ s^\alpha(d^\beta s^\gamma) & s^\alpha(s^\beta u^\gamma) & s^\alpha(u^\beta d^\gamma)\end{bmatrix}.$$

The properties of the hadronic operators $\mathcal{H}$ and $\widetilde{\mathcal{H}}$ under the $SU(3)_L\times SU(3)_R$ symmetry (as well as parity) may then be exploited to convert them into effective terms of the chiral Lagrangian depending on only two low-energy constants $\tilde\alpha$ [for $(3_L,\bar3_R)$ or $(\bar3_L,3_R)$ terms] and

$\tilde{\beta}$ [for $(8_L,1_R)$ or $(1_L,8_R)$ terms]:

$$\begin{aligned}\mathcal{L}_{\psi}^{\chi\mathrm{PT}} \ni &-\tilde{\alpha}\,\bar{\Psi}P_R\mathrm{Tr}\Big\{\Big(C[\mathcal{H}^{LR}_{112}]^*E_{11}-C[\widetilde{\mathcal{H}}^{LR}_{111}]^*E_{32}+C[\widetilde{\mathcal{H}}^{LR}_{112}]^*E_{22}-C[\widetilde{\mathcal{H}}^{LR}_{121}]^*E_{33}+C[\widetilde{\mathcal{H}}^{LR}_{122}]^*E_{23}\Big)\xi^\dagger B\xi^\dagger\Big\}\\ &-\tilde{\alpha}\,\bar{\Psi}P_L\mathrm{Tr}\Big\{\Big(C[\mathcal{H}^{RL}_{112}]^*E_{11}-C[\widetilde{\mathcal{H}}^{RL}_{111}]^*E_{32}+C[\widetilde{\mathcal{H}}^{RL}_{112}]^*E_{22}-C[\widetilde{\mathcal{H}}^{RL}_{121}]^*E_{33}+C[\widetilde{\mathcal{H}}^{RL}_{122}]^*E_{23}\Big)\xi B\xi\Big\}\\ &-\tilde{\beta}\,\bar{\Psi}P_R\mathrm{Tr}\Big\{\Big(C[\mathcal{H}^{LL}_{112}]^*E_{11}-C[\widetilde{\mathcal{H}}^{LL}_{111}]^*E_{32}+C[\widetilde{\mathcal{H}}^{LL}_{112}]^*E_{22}-C[\widetilde{\mathcal{H}}^{LL}_{121}]^*E_{33}+C[\widetilde{\mathcal{H}}^{LL}_{122}]^*E_{23}\Big)\xi^\dagger B\xi\Big\}\\ &-\tilde{\beta}\,\bar{\Psi}P_L\mathrm{Tr}\Big\{\Big(C[\mathcal{H}^{RR}_{112}]^*E_{11}-C[\widetilde{\mathcal{H}}^{RR}_{111}]^*E_{32}+C[\widetilde{\mathcal{H}}^{RR}_{112}]^*E_{22}-C[\widetilde{\mathcal{H}}^{RR}_{121}]^*E_{33}+C[\widetilde{\mathcal{H}}^{RR}_{122}]^*E_{23}\Big)\xi B\xi^\dagger\Big\}\\ &+\mathrm{h.c.}\end{aligned} \tag{31}$$

Here $\xi\equiv\exp\left[\frac{i}{f_\chi}M\right]$, and $f_\chi\simeq 131\,\mathrm{MeV}$ is the chiral-symmetry breaking scale (pion decay constant). The E_{ij}'s represent the elements of the canonical basis of 3×3 real matrices (*i.e.* the entries of E_{ij} are 0 except for the one in the i^{th} line and j^{th} column, which is 1). The low-energy constants $\tilde{\alpha}\approx-0.0144\,\mathrm{GeV}^3$ and $\tilde{\beta}\approx0.0144\,\mathrm{GeV}^3$ have been determined on the lattice [56], but could also be identified from individual hadronic matrix elements of proton decays into mesons, leading to a wider range of variation.

Eq. (31) highlights two physical phenomena. First, it generates mass-mixing terms between the light neutralino and the neutral baryons n^0, Σ^0, Λ^0 and Ξ^0. This effect is analogous to that appearing in the case of a light pseudoscalar mixing with the mesons — see *e.g.* Refs. [57, 58]. A sterile ψ close in mass to a neutral baryon could then decay into lighter baryons via its subleading baryonic component, leading to a spectrum of resonances when varying m_ψ. However, the hadronic decays of hyperons (Σ^0, Λ^0 and Ξ^0; see *e.g.* Ref. [59]) are highly suppressed, because they are mediated by the weak interaction and typically involve an $s\to d$ transition. As a result, the impact of the baryon-neutralino mixing on the neutralino decays remains extremely narrow, as long as we focus on the lightest baryonic octet. That we do not attempt to describe the interactions with heavier and broader baryons, such as the nucleon states at $\sim1.35\,\mathrm{GeV}$, also restricts the validity of our analysis of baryon-number violating decays to $m_\psi\lesssim1.3\,\mathrm{GeV}$. As a complementary feature, the mixing between neutralino and hadrons implies baryon-antibaryon oscillations for the baryon-dominated states, which, in view of the experimental limits on $n^0-\bar{n}^0$ oscillations, constrains the corresponding mixing to comparatively weak values. We discuss this feature in Appendix G. Furthermore, if lepton-number is violated, the subleading ψ component of the baryon-dominated state could mediate baryon- (and lepton)-number violating decays, *e.g.* of the neutron.

The second effect in Eq. (31) is the explicit emergence of ψ-baryon-meson couplings, which allow for neutralino decays into lighter baryon+meson pairs, even though the mixing effect remains negligible. The tree level calculation is presented in Appendix H. Expressing the effective couplings between the neutralino ψ, the mesons $M_i\in\{\pi^0,\pi^+,\pi^-,K^0,\bar{K}^0,K^+,K^-,\eta_8^0\}$ and the baryons $B_j\in\{\Sigma^0,\Sigma^+,\Sigma^-,n^0,\Xi^0,p^+,\Xi^-,\Lambda^0\}$ as $ig_{L,R}^{\psi M_iB_j}M_i\bar{\Psi}P_{L,R}B_j$, we can derive the decay width for the process $\psi\to\bar{M}_i\bar{B}_j$, if kinematically open:

$$\begin{aligned}\Gamma[\psi\to\bar{M}_i\bar{B}_j] &= \frac{m_\psi}{32\pi}\left\{\left(1+\frac{m_{B_j}^2-m_{M_i}^2}{m_\psi^2}\right)\left[\left|g_L^{\psi M_iB_j}\right|^2+\left|g_R^{\psi M_iB_j}\right|^2\right]+4\frac{m_{B_j}}{m_\psi}\mathrm{Re}\left[\left(g_L^{\psi M_iB_j}\right)^*g_R^{\psi M_iB_j}\right]\right\}\\ &\quad\times\left(1-2\frac{m_{B_j}^2+m_{M_i}^2}{m_\psi^2}+\frac{(m_{B_j}^2-m_{M_i}^2)^2}{m_\psi^4}\right)^{1/2}\\ &= \Gamma[\psi\to M_iB_j]\,.\end{aligned} \tag{32}$$

We stress that this description would need to be extended and include the impact of additional (and broad) baryonic states, such as the heavy nucleons at 1.35 GeV or even the Δ's at ~1.2 GeV, as soon as these become kinematically accessible.

Further decays, such as the radiative $\psi \to \gamma B_j$ (for $B_j \in \{\Sigma^0, n^0, \Xi^0, \Lambda^0\}$) could also be assessed in the chiral description. Finally, for $m_\psi < m_p$, hadronic decays of the neutralino are kinematically forbidden, as no hadronic fermion can be reached below the proton mass: it is then the baryons which may decay into ψ and a meson [48].

3.3 Semi-leptonic decays

The semi-leptonic operators produce a collection of neutralino decay channels into mesons and a lepton. The simplest final states, already considered in *e.g.* Ref. [33], consist of two bodies, one lepton (ℓ_i) and one meson M. The decay amplitude may be schematically written as:

$$\begin{aligned}\mathcal{A}[\psi \to \bar{\ell}_i M] = {} & i\,C[\mathcal{S}^{\ell q JK}_{ijk}]\,\langle M(p_M)|\,\bar{Q}_j P_K Q_k\,|0\rangle\;\bar{\nu}_\psi(p_\psi) P_J \nu_{\ell_i}(p_{\ell_i}) \\ & + i\,C[\mathcal{V}^{\ell q JK}_{ijk}]\,\langle M(p_M)|\,\bar{Q}_j \gamma^\mu P_K Q_k\,|0\rangle\;\bar{\nu}_\psi(p_\psi)\gamma_\mu P_J \nu_{\ell_i}(p_{\ell_i}) \\ & + i\,C[\mathcal{T}^{\ell q J}_{ijk}]\,\langle M(p_M)|\,\bar{Q}_j \tfrac{\Sigma^{\mu\nu}}{2} P_J Q_k\,|0\rangle\;\bar{\nu}_\psi(p_\psi)\tfrac{\Sigma_{\mu\nu}}{2} P_J \nu_{\ell_i}(p_{\ell_i})\,,\end{aligned} \tag{33}$$

where the scalar, vector and tensor form factors of the meson M depend on its nature. At this level, the interactions of the final-state lepton with the meson have been neglected.

Focusing first on the pseudoscalar mesons of Eq. (30), to which we may add the $\eta_1^0 \sim \frac{1}{\sqrt{3}}(u^\alpha \bar{u}^\alpha + d^\alpha \bar{d}^\alpha + s^\alpha \bar{s}^\alpha)$ and account for $\eta - \eta'$ mixing, the pseudoscalar form factors in the chiral limit are all connected to the parameter B_0 which relates meson masses to quark mass parameters, *e.g.* $B_0 \approx \frac{f_\chi}{\sqrt{2}m_s}(m^2_{K^\pm} + m^2_{K^0} - m^2_{\pi^0}) \sim 0.5\,\text{GeV}^2$:

$$\langle M(p_M)|\,\bar{Q}_j P_R Q_k\,|0\rangle = -\frac{iB_0}{2\sqrt{2}}\mathcal{R}_M[q_j\bar{q}_k] = -\langle M(p_M)|\,\bar{Q}_j P_L Q_k\,|0\rangle\,. \tag{34}$$

Here $\mathcal{R}_M$ represents the overlap of the $q_j\bar{q}_k$ pair with the valence quarks of $M(p_M)$; for instance $\mathcal{R}_{\pi^0}[u\bar{u}] = \frac{1}{\sqrt{2}} = -\mathcal{R}_{\pi^0}[d\bar{d}]$ and all other $\mathcal{R}_{\pi^0}[q_j\bar{q}_k]$ vanish. The axial form factors then relate to the above through the Dirac equation:

$$\langle M(p_M)|\,\bar{Q}_j \gamma^\mu P_R Q_k\,|0\rangle = -\frac{i\,B_0}{2\sqrt{2}}\frac{m_{q_j} + m_{q_k}}{m_M^2}\,\mathcal{R}_M[q_j\bar{q}_k]\,p_M^\mu = -\langle M(p_M)|\,\bar{Q}_j \gamma^\mu P_L Q_k\,|0\rangle\,. \tag{35}$$

In turn, $(m_{q_j} + m_{q_k})B_0$ can be written in terms of f_χ and meson masses. There is no antisymmetric tensor associated with scalars so that $\langle M(p_M)|\,\bar{Q}_j \frac{\Sigma^{\mu\nu}}{2} P_{L,R} Q_k\,|0\rangle = 0$. In order to account for effects beyond the chiral approximation, it is possible to re-parametrize the pseudoscalar and axial form-factors in terms of the decay constants of the meson $f_M^{q_j\bar{q}_k}$, as already suggested in Ref. [33]: $\langle M(p_M)|\,\bar{Q}_j \gamma^\mu \gamma^5 Q_k\,|0\rangle = -i\,f_M^{q_j\bar{q}_k}\,p_M^\mu$. We finally arrive at the decay widths (with implicit sum over the quark indices):

$$\begin{aligned}\Gamma[\psi \to \bar{\ell}_i M] = {} & \frac{m_\psi}{32\pi}\Big\{\big(1 + \tfrac{m^2_{\ell_i} - m^2_M}{m^2_\psi}\big)\big(|g_L^{\psi\bar{M}\ell_i}|^2 + |g_R^{\psi\bar{M}\ell_i}|^2\big) + 4\tfrac{m_{\ell_i}}{m_\psi}\text{Re}\big[(g_L^{\psi\bar{M}\ell_i})^* g_R^{\psi\bar{M}\ell_i}\big]\Big\} \\ & \times\left[1 - 2\tfrac{m^2_M + m^2_{\ell_i}}{m^2_\psi} + \tfrac{(m^2_M - m^2_{\ell_i})^2}{m^4_\psi}\right]^{1/2},\end{aligned} \tag{36}$$

$$\begin{aligned}g_L^{\psi\bar{M}\ell_i} \equiv {} & \tfrac{f_M^{q_j\bar{q}_k}}{2}\Big\{\tfrac{B_0}{\sqrt{2}f_\chi}\Big(\tfrac{\sqrt{2}f_\chi m_M^2}{B_0(m_{q_j} + m_{q_k})}\Big)\big(C[\mathcal{S}^{\ell q\,LR}_{ijk}] - C[\mathcal{S}^{\ell q\,LL}_{ijk}]\big) \\ & - \Big[m_\psi\big(C[\mathcal{V}^{\ell q\,LR}_{ijk}] - C[\mathcal{V}^{\ell q\,LL}_{ijk}]\big) - m_{\ell_i}\big(C[\mathcal{V}^{\ell q RR}_{ijk}] - C[\mathcal{V}^{\ell q\,LR}_{ijk}]\big)\Big]\Big\}\,,\end{aligned}$$

$$\begin{aligned}g_R^{\psi\bar{M}\ell_i} \equiv {} & \tfrac{f_M^{q_j\bar{q}_k}}{2}\Big\{\tfrac{B_0}{\sqrt{2}f_\chi}\Big(\tfrac{\sqrt{2}f_\chi m_M^2}{B_0(m_{q_j} + m_{q_k})}\Big)\big(C[\mathcal{S}^{\ell q RR}_{ijk}] - C[\mathcal{S}^{\ell q RL}_{ijk}]\big) \\ & - \Big[m_\psi\big(C[\mathcal{V}^{\ell q RR}_{ijk}] - C[\mathcal{V}^{\ell q RL}_{ijk}]\big) - m_{\ell_i}\big(C[\mathcal{V}^{\ell q\,LR}_{ijk}] - C[\mathcal{V}^{\ell q\,LL}_{ijk}]\big)\Big]\Big\}\,.\end{aligned}$$

We consider the 19 such channels involving light leptons (plus their 19 CP-conjugate).

Turning to the light vector mesons $M^*_\mu \sim (\bar{q}_j \gamma_\mu q_k)$,

$$M^*_\mu \equiv \begin{bmatrix} \frac{\rho^0_\mu}{\sqrt{2}} + \frac{\phi^0_{8\mu}}{\sqrt{6}} + \frac{\phi^0_{1\mu}}{\sqrt{3}} & \rho^+_\mu & K^{*+}_\mu \\ \rho^-_\mu & -\frac{\rho^0_\mu}{\sqrt{2}} + \frac{\phi^0_{8\mu}}{\sqrt{6}} + \frac{\phi^0_{1\mu}}{\sqrt{3}} & K^{*0}_\mu \\ K^{*-}_\mu & \bar{K}^{*0}_\mu & -\sqrt{\frac{2}{3}}\phi^0_{8\mu} + \frac{\phi^0_{1\mu}}{\sqrt{3}} \end{bmatrix}, \tag{37}$$

and after mixing ϕ^0_8 and ϕ^0_1 to account for the ϕ and ω, one can define the vector and tensor form-factors:

$$\begin{aligned} \langle M^*(p_M)|\bar{Q}_j \gamma^\mu Q_k|0\rangle &= f^V_{M^*}\, m_{M^*}\, \mathcal{R}_{M^*}[q_j\bar{q}_k]\, \epsilon^{*\mu}_{M^*}(p_M), \\ \langle M^*(p_M)|\bar{Q}_j \Sigma^{\mu\nu} Q_k|0\rangle &= -i f^T_{M^*}\, \mathcal{R}_{M^*}[q_j\bar{q}_k]\left[p^\mu_M \epsilon^{*\nu}_{M^*}(p_M) - p^\nu_M \epsilon^{*\mu}_{M^*}(p_M)\right]. \end{aligned} \tag{38}$$

Then, we may parametrize the decay amplitude as:

$$\mathcal{A}[\psi \to M^* \bar{\ell}_i] = i\bar{v}_\psi(p_\psi)\left\{ g^{M^*\psi\ell_i}_{V\,L,R} \not{\epsilon}^*_{M^*}(p_{M^*}) P_{L,R} + 2\, p_\psi \cdot \varepsilon^*_{M^*}(p_{M^*})\, g^{M^*\psi\ell_i}_{S\,L,R} P_{L,R} \right\} v_{\ell_i}(p_{\ell_i}), \tag{39}$$

$$\begin{aligned} g^{M^*\psi\ell_i}_{V\,J} &\equiv \tfrac{1}{2}\mathcal{R}_{M^*}[q_j\bar{q}_k]\Big\{ m_{M^*} f^V_{M^*}\left(C[\mathcal{V}^{\ell q JL}_{ijk}] + C[\mathcal{V}^{\ell q JR}_{ijk}]\right) \\ &\qquad - f^T_{M^*}\left(m_\psi C[\mathcal{T}^{\ell q J}_{ijk}] + m_{\ell_i} C[\mathcal{T}^{\ell q J+1}_{ijk}]\right)\Big\}, \\ g^{M^*\psi\ell_i}_{S\,J} &\equiv -\tfrac{f^T_{M^*}}{2}\mathcal{R}_{M^*}[q_j\bar{q}_k] C[\mathcal{T}^{\ell q J}_{ijk}], \end{aligned}$$

leading to the decay width:

$$\begin{aligned} \Gamma[\psi \to M^*\bar{\ell}_i] &= \frac{m^3_\psi}{32\pi m^2_{M^*}}\left[1 - 2\frac{m^2_{M^*}+m^2_{\ell_i}}{m^2_\psi} + \frac{(m^2_{M^*}-m^2_{\ell_i})^2}{m^4_\psi}\right]^{1/2} \\ &\times \Bigg\{ \left(|g^{M^*\psi\ell_i}_{V\,L}|^2 + |g^{M^*\psi\ell_i}_{V\,R}|^2\right)\left[1 + \frac{m^2_{M^*}-2m^2_{\ell_i}}{m^2_\psi} - \frac{2m^4_{M^*} - m^2_{\ell_i}(m^2_{M^*}+m^2_{\ell_i})}{m^4_\psi}\right] \\ &\quad - 12\frac{m^2_{M^*} m_{\ell_i}}{m^3_\psi}\mathrm{Re}\left[g^{M^*\psi\ell_i *}_{V\,L} g^{M^*\psi\ell_i}_{V\,R}\right] + m^2_\psi\left[1 - 2\frac{m^2_{M^*}+m^2_{\ell_i}}{m^2_\psi} + \frac{(m^2_{M^*}-m^2_{\ell_i})^2}{m^4_\psi}\right] \\ &\quad \times \left[\left(|g^{M^*\psi\ell_i}_{S\,L}|^2 + |g^{M^*\psi\ell_i}_{S\,R}|^2\right)\left(1 + \frac{m^2_{\ell_i}-m^2_{M^*}}{m^2_\psi}\right) + 4\frac{m_{\ell_i}}{m_\psi}\mathrm{Re}\left[g^{M^*\psi\ell_i *}_{S\,L} g^{M^*\psi\ell_i}_{S\,R}\right]\right] \\ &\quad - 2m_\psi\left[1 - 2\frac{m^2_{M^*}+m^2_{\ell_i}}{m^2_\psi} + \frac{(m^2_{M^*}-m^2_{\ell_i})^2}{m^4_\psi}\right]\mathrm{Re}\left[g^{M^*\psi\ell_i *}_{V\,L} g^{M^*\psi\ell_i}_{S\,L} + g^{M^*\psi\ell_i *}_{V\,R} g^{M^*\psi\ell_i}_{S\,R}\right] \\ &\quad - 2m_{\ell_i}\left[1 - 2\frac{m^2_{\ell_i}}{m^2_\psi} + \frac{(m^2_{M^*}-m^2_{\ell_i})^2}{m^4_\psi}\right]\mathrm{Re}\left[g^{M^*\psi\ell_i *}_{V\,L} g^{M^*\psi\ell_i}_{S\,R} + g^{M^*\psi\ell_i *}_{V\,R} g^{M^*\psi\ell_i}_{S\,L}\right]\Bigg\}. \end{aligned} \tag{40}$$

However, the production of vector mesons can be seen as a partial contribution to the more general final states with two (or three) pseudoscalar mesons, see for example Ref. [60] in the R-parity violating context. For this reason, it is preferable to dismiss the two-body decay formalism for vector mesons and directly exploit the more complete form-factors for meson pairs, *i.e.* three-body semi-leptonic decay channels, as we perform below. Yet, in the case of the ω, mainly decaying into three pseudoscalar mesons, the two-body formalism remains sufficient at low energy.

The properties of scalar and tensor mesons being ill-known (and broad), an application of the two-body formalism with such particles in the final state would provide limited insight.

Instead, it is again more meaningful to regard these states as mediators in the three-body decay widths.

Decays into one lepton and two pseudoscalar mesons are relevant as soon as enough phase-space is available. In view of the CP-conservation observed by the strong interaction, we need not consider pseudoscalar and axial-vector mediations and the decay amplitude schematically reads (still neglecting the interactions between leptons and mesons in the final state):

$$\begin{aligned}\mathcal{A}[\psi\to\bar{\ell}_i M_1 M_2] = &\ \tfrac{i}{2}\,C[\mathcal{S}^{\ell q JK}_{ijk}]\,\langle M_1(p_1)M_2(p_2)|\,\bar{Q}_j Q_k\,|0\rangle\ \bar{v}_\psi(p_\psi)P_J v_{\ell_i}(p_{\ell_i}) \\ &+\tfrac{i}{2}\,C[\mathcal{V}^{\ell q JK}_{ijk}]\,\langle M_1(p_1)M_2(p_2)|\,\bar{Q}_j\gamma^\mu Q_k\,|0\rangle\ \bar{v}_\psi(p_\psi)\gamma_\mu P_J v_{\ell_i}(p_{\ell_i}) \\ &+\tfrac{i}{2}\,C[\mathcal{T}^{\ell q J}_{ijk}]\,\langle M_1(p_1)M_2(p_2)|\,\bar{Q}_j\Sigma^{\mu\nu} Q_k\,|0\rangle\ \bar{v}_\psi(p_\psi)\tfrac{\Sigma_{\mu\nu}}{2}P_J v_{\ell_i}(p_{\ell_i})\,.\end{aligned} \tag{41}$$

We then introduce the scalar, vector and tensor form-factors (with $s\equiv(p_1+p_2)^2$):

$$\begin{aligned}&\langle M_1(p_1)M_2(p_2)|\,\bar{Q}_j Q_k\,|0\rangle \equiv F_{S\,jk}^{M_1M_2}(s)\,, \\ &\langle M_1(p_1)M_2(p_2)|\,\bar{Q}_j\gamma^\mu Q_k\,|0\rangle \equiv (p_1+p_2)^\mu\tfrac{m_j-m_k}{s}F_{S\,jk}^{M_1M_2}(s) \\ &\qquad\qquad + \left[p_1^\mu - p_2^\mu - \tfrac{m^2_{M_1}-m^2_{M_2}}{s}(p_1+p_2)^\mu\right]F_{V\,jk}^{M_1M_2}(s)\,, \\ &\langle M_1(p_1)M_2(p_2)|\,\bar{Q}_j\Sigma^{\mu\nu} Q_k\,|0\rangle \equiv i(p_1^\mu p_2^\nu - p_2^\mu p_1^\nu)F_{T\,jk}^{M_1M_2}(s)\,.\end{aligned} \tag{42}$$

Here $F_{S\,jk}^{M_2M_1}(s) = F_{S\,jk}^{M_1M_2}(s)$, while $F_{V\,jk}^{M_2M_1}(s) = -F_{V\,jk}^{M_1M_2}(s)$ and $F_{T\,jk}^{M_2M_1}(s) = -F_{T\,jk}^{M_1M_2}(s)$. The two-meson form factors automatically include contributions from scalar, vector or tensor resonances. They can be computed from chiral perturbation theory in the low-energy range: we refer the reader to Ref. [61] for a recent and comparatively exhaustive derivation. We note in passing that radiative corrections (including resonance effects) significantly modify the leading order contribution in chiral perturbation theory, so that no simple approximation is viable in this regime.

However, theoretical approaches to the two-meson form factors fail for energies somewhat above 1 GeV because they rely on the two-body approximation for intermediate states in the scattering, this condition not being reliable any longer at higher energies. Phenomenological parametrizations exploiting experimental data from tau decays or e^+e^- scattering can be employed essentially for vector form factors — see *e.g.* Ref. [62]. The case of scalar form-factors, determinant for the RpV MSSM, is far more difficult to systematically address without trusting a specific model [63, 64]. Here, we employ the form factors provided by Ref. [61], which can be used up to an invariant mass of 1.2 GeV for the di-meson system. Beyond this energy, we simply extrapolate the form factors with a rational function $A/(s-M^2)$, mimicking the mediation by the principal resonance: this is a coarse approximation for this higher-energy regime. For the vector and tensor currents, it amounts to neglecting the impact of secondary resonances, which we believe is a reasonable approximation. The situation for scalar currents is more complicated as the various form factors at 1.2 GeV have not yet reached a 'tail' behavior. As the scalar operators contain the main contributions in the RpV MSSM, this means that the description of semi-leptonic decays at masses $m_\psi \gtrsim 1.2$ GeV remains purely qualitative. In addition, it is likely that four-body decay channels become relevant already at energies below the tau mass, so that a partonic description of the inclusive hadronic width at $m_\psi \approx m_\tau$ may become preferable to a channel-by-channel approach.

Once the form-factors for the two-meson channels are numerically defined, we can derive the corresponding 48 (+ 48 CP-conjugate) decay widths for the neutralino:

$$\Gamma[\psi\to\bar{\ell}_i M_1 M_2] = \frac{m_\psi^3 B_0^2}{512(1+\delta_{M_1M_2})\pi^3 f_\chi^2}\sum_{\Omega,\Omega'} C[\Omega]C[\Omega']^*\,G^{\bar{\ell}_i M_1M_2}[\Omega,\Omega']\,, \tag{43}$$

where the sum extends over all relevant pairs of semi-leptonic operators (Ω, Ω'). The kinematic integrals $G^{\bar{\ell}_i M_1 M_2}[\Omega, \Omega']$ are defined in Appendix I and *a priori* depend on all the masses of the process.

Given the considered range of neutralino mass, we will assume that the available semi-leptonic final states are dominated by the one- and two-meson channels, without need of addressing the three-body channels (beyond the resonant contribution resulting from single meson production). At higher masses $\sim m_\tau$, the partonic description — see Appendix J — should become applicable.

4 About the Production of the Light Neutralino

While we focus on the decays of the light exotic fermion, which can be conveniently classified through the listing of the low energy operators, the phenomenology of such a particle also crucially depends on its production modes. The latter can be extremely diverse and model-dependent since they may occur at any scale, depending on the beam energy of the collider. If the light bino-like particle is sufficiently long-lived, it may even possess thermal relics, allowing for tests in Dark-Matter detection experiments. Here, we restrict ourselves to a short discussion of typical production channels at colliders.

4.1 R-parity conserving production modes

Contrary to its decays, the dominant production channels of the light exotic fermion may not necessarily violate the $\mathbb{Z}_2$ symmetry distinguishing it from the SM matter (R-parity in the MSSM). Thus, ψ may appear in cascade decays of heavy resonances produced in high-energy collisions. It then becomes necessary to assess the abundance of these $\mathbb{Z}_2$-conserving cascades with respect to the competition of a direct $\mathbb{Z}_2$-violating decay of the heavy resonances into SM matter.

If we turn to the direct production of the light fermion from SM sources (*i.e.* not involving further new physics states), then $\mathbb{Z}_2$ conservation implies a production in pairs. The most obvious 'portals' with SM matter then involve the Z- or Higgs bosons via renormalizable operators of the form $Z_\mu \bar{\psi}\bar{\sigma}^\mu\psi$ and $h^0\psi\psi$. Limits on invisible (or exotic) Z and Higgs decays place strict limits on such production mechanisms. In the case of the RpV MSSM, the couplings of ψ to the Z-boson or the SM-like Higgs would involve the subleading higgsino components of this state (while the dominant bino component remains inert), hence providing a straightforward suppression mechanism. However, precision experiments, *e.g.* at a linear collider or a Z-factory, might allow to detect small exotic decay widths of the neutral SM bosons. This possibility was examined in *e.g.* Refs. [35, 36].

At low energy, pair production would proceed through non-renormalizable operators of the form (restricting ourselves to dimension 6 and omitting the hermitian conjugates; $f = e, u, d$; $F = (f, \bar{f}^c)$; $j, k = 1, 2, 3$; $J = L, R$):

$$\dot{\mathcal{S}}^{f\,J}_{jk} \equiv (\bar{\Psi} P_L \Psi)(\bar{F}_j P_J F_k)\,, \qquad \dot{\mathcal{V}}^{f\,J}_{jk} \equiv (\bar{\Psi}\gamma^\mu P_L \Psi)(\bar{F}_j \gamma_\mu P_J F_k)\,, \tag{44}$$

satisfying the $U(1)_{\rm em} \times SU(3)_c$ gauge symmetry (and possibly additional discrete symmetries depending on the model) — the tensor operator $\dot{\mathcal{T}}^f_{jk} \equiv \frac{1}{4}(\bar{\Psi}\Sigma^{\mu\nu} P_L \Psi)(\bar{F}_j \Sigma_{\mu\nu} P_L F_k)$ identically vanishes. In particular, such operators involving quarks open up the possibility of ψ production in the decays (or collisions) of hadrons (of various spin). Let us illustrate this feature with an example: we consider the low-energy operators $(\bar{\Psi} P_L \Psi)(\bar{B} P_{L,R} S)$. Then, ψ pairs can be produced in a direct B_s^0 decay (calculable via the decay constant $\langle 0|\, \bar{B}\gamma_5 S\, |B_s^0\rangle$), or in semi-leptonic B decays, such as $B \to K\psi\psi$ (calculable with the $B \to K$ form factor) and more gener-

ally $B \to X_s \psi\psi$ (calculable using the heavy quark expansion). These individual decay modes might be testable at a B factory.[2] Similarly, the $(\bar{\Psi}\gamma^\mu P_L \Psi)(\bar{C}\gamma^\mu P_{L,R} C)$ or $(\bar{\Psi}\gamma^\mu P_L \Psi)(\bar{B}\gamma^\mu P_{L,R} B)$ operators would open up the possibility of production from vector quarkonia, such as the J/ψ or the Υ. For meson decays to two light neutralinos see Ref. [32]. The situation is more complicated if the hadronic initial state produced in the experiment is not clearly identified (because many hadrons can potentially decay via the low-energy operator that couples them to the ψ pair), so that ψ production in a b-jet, for instance, would have to be modeled together with the hadronization string for a quantitative estimate. Leptonic four-fermion operators $(\dot{\mathcal{S}}, \dot{\mathcal{V}}, \dot{\mathcal{T}})^{e\,(J)}_{jk}$ may also be exploited to search for ψ's in lepton-antilepton decays or, in the presence of lepton-flavor violation, in leptonic tau (or muon) decays.

Contributions to the operators of Eq. (44) depend on the UV-completion. In the RpV MSSM at tree level, typical 's-channel' mediators would be the Z- and Higgs bosons. In view of the limits on invisible / exotic decays of the Z or the SM-like Higgs, these two mediators are probably not competitive. Heavy Higgs doublets would be legitimate mediators (at least in couplings to third-generation fermions) provided their coupling to the bino-dominated state is not exceedingly suppressed. Sfermion exchange provides another type of 't-channel' contribution. In fact, the matching conditions for the $(\dot{\mathcal{S}}, \dot{\mathcal{V}}, \dot{\mathcal{T}})^{f\,(J)}_{jk}$ operators in the RpV MSSM at tree-level are virtually identical to those of the $(\mathcal{S}, \mathcal{V}, \mathcal{T})^{vfJ(K)}_{ijk}$ operators ($f = e, u, d$) provided in Eqs.(16,17), up to the substitution of ν_i by ψ and the generalization of the indices j and k, so as to include τ, c or b. At the loop level, additional contributions from *e.g.* SUSY box diagrams are also expected. The evolution under the renormalization group is also unchanged, see Eq. (23), although the infrared scale for production should be adjusted to the corresponding center-of-mass energy.

4.2 R-parity violating production modes

Violation of the $\mathbb{Z}_2$ symmetry allows for single production of particles in the 'exotic' sector. Again, ψ might then occur as the end product of the $\mathbb{Z}_2$-conserving cascade decay of a singly-produced heavy resonance, provided the competition of $\mathbb{Z}_2$-violating decays does not overshadow such modes. In the RpV MSSM, a heavy Higgs state might also exploit its subleading slepton component to decay into ψ and a lepton.

Direct production of ψ from SM particles can proceed through low-energy operators of the type displayed in Eqs.(3-6), but extended to heavier fermions. The matching conditions at tree-level in the RpV MSSM are again a simple generalization of Eqs.(16-18). Lepton decays would then appear as a first potential lepton-number violating source of the light exotic fermion. For instance, the decay widths associated with $\tau \to \psi\, \nu_i\, \bar{\nu}_j$, $\tau \to \psi \ell_i \bar{\ell}_j$, $\tau \to \psi M_i$, or $\tau \to \psi M_i M_j$ can be computed similarly to the leptonic and semi-leptonic decays of the light neutralino [see Eqs.(27,28,29,36,40,43)]. For a very light fermion, a partonic approach may efficiently describe the inclusive τ decays into ψ and light hadrons, since the large phase space suggests a possible saturation of the hadronic width. Individual channels have been considered in Ref. [39]. The limited phase space in muon or electron decays restricts the possibility of ψ production to leptonic channels.

Operators of semi-leptonic and hadronic types also suggest single production of the exotic fermion in hadron decays. Exchanging the fermions in the initial and final states in Eq. (32) opens the path to a lepton-number conserving disintegration of baryons and, in particular, of nucleons [48]: this is a very constrained production channel, but constantly investigated by experiments. Now turning to operators including a bottom (generated in the RpV MSSM via *e.g.* λ''_{113}) or charm quark (via *e.g.* λ''_{212}), one can expect single production from beauty-

[2] See also Ref. [32] for the related decay $K \to \pi\psi\psi$.

or charm-flavored hadrons. The weight of a few inclusive channels, such as $B, D \to X\psi$, may be estimated from the partonic picture. On the other hand, exclusive channels, such as $B^+, D^+ \to p^+\psi$ or $\Lambda^0_{b,c} \to \pi^0\psi$, are much more difficult to assess since they would require the knowledge of baryon-number and flavor violating form factors ($B^+ \to p^+$ or $\Lambda^0_b \to \pi^0$ in our example). As in the case of pair production, the ψ content of a b- (or c-) jet due to the presence of baryon-number violating operators would probably require detailed modeling. The situation is comparable when considering semi-leptonic operators, even though it should be possible to theoretically describe a few individual exclusive channels testable at B-factories: the corresponding form-factors (*e.g.* $B \to K$) are no longer baryon-number violating, hence more easily accessible. A few such processes have been exploited in *e.g.* Refs. [33–35].

To summarize this short discussion, the potential production channels of a light neutralino-like particle are even more diverse than its decay modes. Many transitions may be difficult to describe accurately with the usual techniques, as they would call upon unknown form-factors or require a dedicated hadronization model. The relative magnitude of pair and single production channels should also be systematically compared, as these sources are largely independent from one another. Given that our focus in this paper are the decays of the neutralino-like particle, we will leave this discussion of the production modes at this qualitative level. A few formulae for the production of a light neutralino in the decays of SM particles are provided for reference in Appendix K.

5 Numerical Applications

The formalism described in the previous sections is collected within a Mathematica package, which allows for the computation of all the considered decay widths of a light neutralino for the most general MSSM input. Numerical comparison of a few individual semi-leptonic widths with a private code, derived from Ref. [33] and subsequently employed in simulations of far-detector sensitivities, returned discrepancies by a factor 10 (underestimating the pseudoscalar meson channel) to 100 (overestimating the vector meson channel). These differences were traced back to the impact of QCD running — see Eq. (23) —, the issues in matching raised at the very end of Section 2.2 and numerical bugs. We here demonstrate a few possible applications of these calculations at the numerical level.

5.1 Baryon-number violation

We first assume that lepton number is conserved. Neutralino decays are then driven by baryon-number violating couplings. We stress again that hadronic decays are kinematically forbidden for a fermion with mass below the proton mass, and in fact, as long as we neglect the radiative hadronic decay $\psi \to n^0\gamma$, below the threshold for nucleon+pion production, *i.e.* ~ 1.075 GeV. Then, our description for such channels is at best of qualitative value, due first to the limited validity of the chiral description, especially in the higher mass range, and also to the restriction to the fundamental baryon octet, beyond which further hadronic states become available almost from the opening of the kinematic window for hadronic decays.

If flavor mixing is negligible in the sfermion sector, only one RpV parameter is liable to mediate ψ decays at tree-level: $\lambda''_{112} = -\lambda''_{121}$. Decay widths will then scale like $|\lambda''_{112}|^2$. Yet, such a coupling cannot lead to decays into nucleons and pions (the kinematically available channels at low mass) without a source of $s \to d$ transition, and the kinematic threshold is pushed further up to the opening of $\Lambda^0 + \pi^0$ production (at ~ 1.25 GeV). Nevertheless, the weak interaction is always available as a source of flavor violation and can mediate neutralino decays into nucleons and pions at low mass via the mixings of ψ with the hyperons: the recourse to

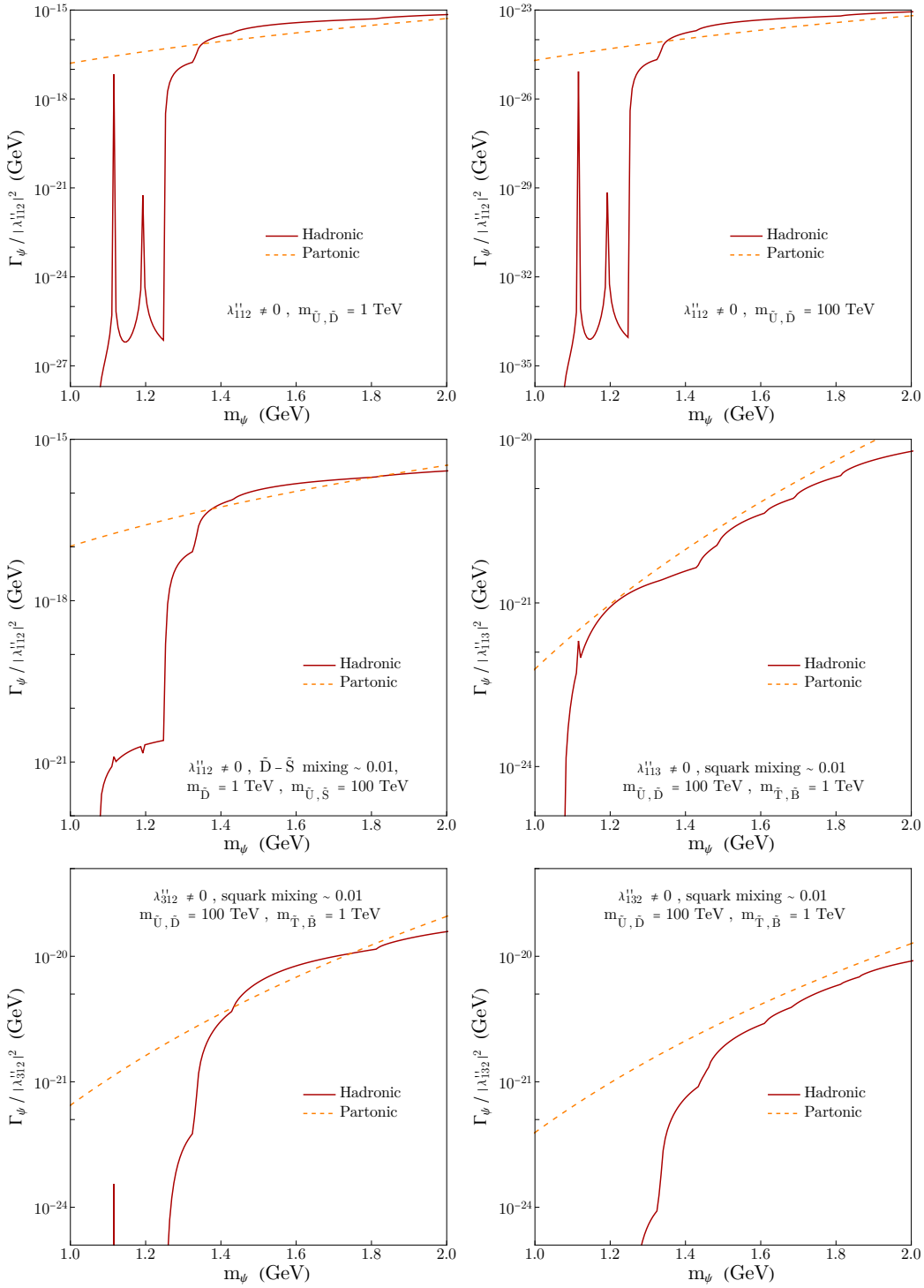

Figure 1: Hadronic decay width of the neutralino in various scenarios. *First row*: No squark mixing; the neutralino decays are mediated by the (right-handed) squarks of first and second generation with mass at 1 TeV (Left) and 100 TeV (Right). *Second row, left*: Mixing of magnitude 10^{-2} between right-handed light down-type (1 TeV) and heavy strange-type squarks (100 TeV). *Second (right) and third row*: Mixing of 10^{-2} between the heavy squarks of first generations (100 TeV) and the light squarks of third generation (1 TeV); the dominant coupling of UDD-type is chosen differently in the three cases: λ''_{113}, λ''_{312} and λ''_{132}. The inclusive hadronic width, normalized to the dominant coupling of UDD-type squared is displayed in red. The partonic approximation is shown in dashed orange.

the weak Lagrangian considerably reduces the widths, however.[3] The total neutralino width in such a setup is presented in the first row of Fig. 1 for squark masses of 1 TeV (left) and 100 TeV (right). The solid red line represents the actual hadronic width, while the dashed orange curve corresponds to the partonic three-body approximation (hence ignores hadronic thresholds). The scaling of the Wilson coefficients as the inverse of the mass-squared of the squark mediator amounts to a 10^{-8} suppression at the level of the decay width in the scenario with the heavier sfermions — QCD logarithms also produce a prefactor of $\mathcal{O}(1)$. In the mass range accessible only to the nucleon+pion final states, the decay width is significantly suppressed (as compared to the partonic width) and clearly dominated by the resonant behavior caused by the ψ mixing with Λ^0 and Σ^0. Above ~ 1.25 GeV, the hadronic neutralino decays can proceed without recourse to the weak interaction and promptly reach a magnitude comparable with the partonic width.[4] The latter may actually offer a quantitatively more reliable description of the (un-charmed) inclusive hadronic decay width for $m_\psi \to 2$ GeV.

The presence of squark mixing may substantially alter the picture presented above, in particular with respect to the relative importance of nucleon+pion channels as opposed to the strange-violating ones. In the second row of Fig. 1, on the left, we allow for a mixing of magnitude 10^{-2} between a light right-handed sdown at 1 TeV and the heavy sstrange at 100 TeV. Then the nucleon+pion decays only receive a mixing squared suppression (10^{-4}) with respect to the partonic width, instead of the $\sim 10^{-9}$ suppression observed in the scenarios without squark mixing of the first row. Squark mixing also allows for other couplings of UDD-type, beyond λ''_{112}, to contribute to the neutralino decays: such effects will always be suppressed in proportion to the needed mixing to couple squark mediators to the light quarks. The last three plots of Fig. 1 illustrate this setup with various choices of dominant UDD-coupling (λ''_{113}, λ''_{312} and λ''_{132}) and a mixing of magnitude 10^{-2} between heavy squarks of first and second generations (100 TeV) and light squarks of third generation (1 TeV). Again, the relative importance of the decays into nucleons and pions depends on the possibility to generate a contribution to the $\widetilde{\mathcal{H}}^{JK}_{111}$ operators via squark mixing. In this respect, the scenario with dominant λ''_{113} results in no suppression of the strange-number conserving channels with respect to the strange-violating ones, since all require the same degree of mixing of the squarks of third generation with those of first and second generations (this mixing was made symmetrical for both generations in our example). On the other hand, strange-number conserving decays in the scenarios with dominant λ''_{312} or λ''_{132} require one additional degree of flavor mixing as compared to the strange-violating ones, explaining the suppression of the nucleon+lepton channels in the corresponding plots.

We note the comparative agreement of the hadronic and partonic widths in the higher mass range (once enough decay channels are kinematically accessible) for all scenarios of Fig. 1. This situation may be largely coincidental since the chiral description is not expected to be predictive (or even exhaustive) in this regime, while the partonic description is insensitive to hadronic thresholds, hence only valid at sufficiently high mass. However, this agreement allows for a comparatively smooth transition between the two models at $m_\psi \approx 1.5$ GeV. Finally, we stress that the hadronic description of the baryon-number violating scenario should also be put into perspective with its consequences for the phenomenology of baryons. Correspondingly, constraints on nucleon decays [48] apply for $m_\psi \lesssim m_p$ while limits on neutron-antineutron oscillations — see Appendix G — restrict the acceptable magnitude of contributions to the $\widetilde{\mathcal{H}}^{JK}_{111}$ operators (those relevant for strange-number conserving neutralino decays).

[3] From a high-energy model building point of view one might expect the coupling λ''_{112} not to appear alone, but those couplings related by CKM mixing to also be present in the theory, correspondingly suppressed. See for example [17, 65].

[4] For a related comparison of partonic decay widths to those incorporating hadronic thresholds in the case of sterile neutrinos see Ref. [66].

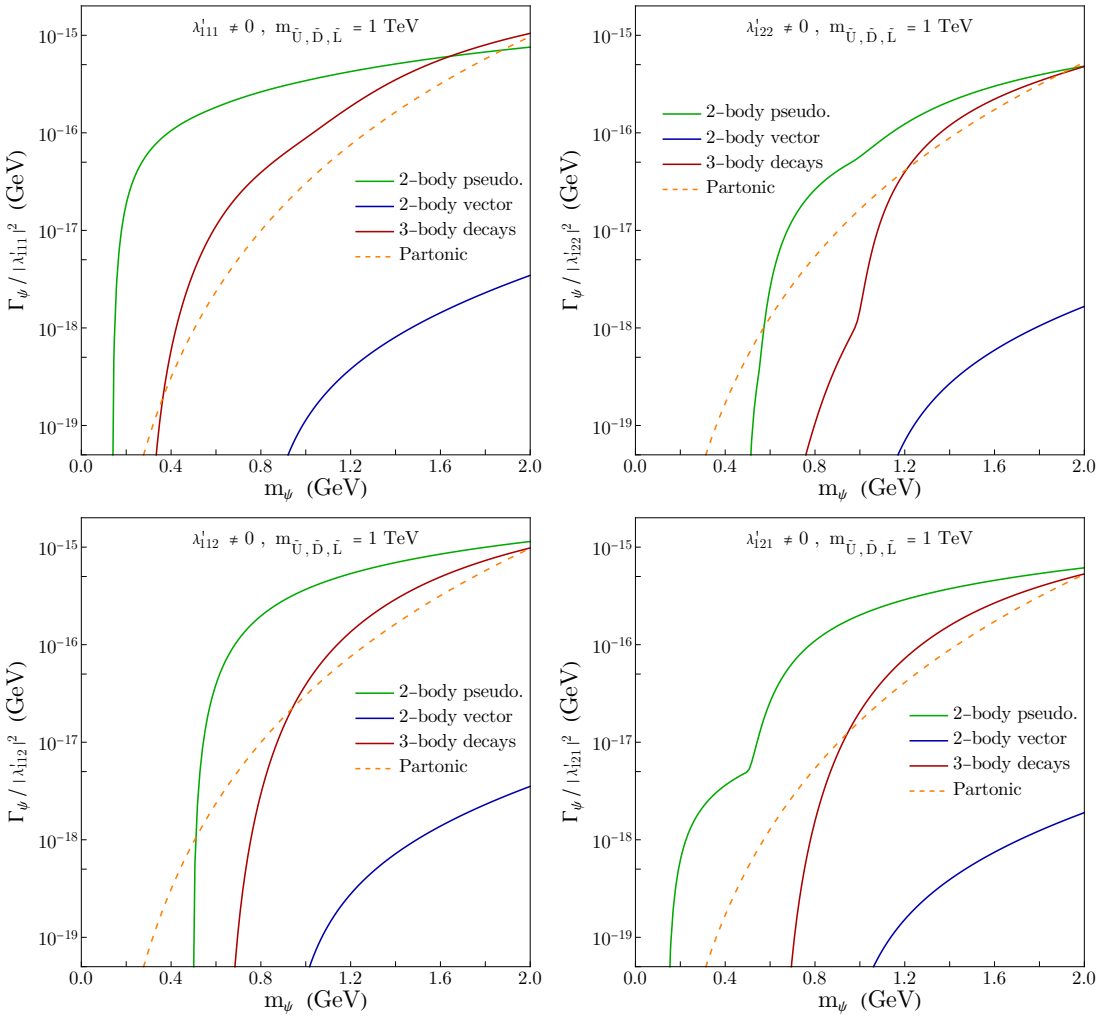

Figure 2: Semi-Leptonic decay width of the neutralino in various scenarios without sfermion mixing. Both slepton and squark masses are set to 1 TeV. *First row*: quark-flavor conserving scenarios; lepton-number violation is mediated by the λ'_{111} coupling (left) or the λ'_{122} coupling (right). *Second row*: the lepton-number violating coupling also violates the strange-number; it is chosen as λ'_{112} in the left plot and as λ'_{121} in the right one. The green, blue and red curves respectively represent the two-body decay width into one lepton and a pseudoscalar meson, the two-body decay width into one lepton and a vector meson and the three-body decay width into one lepton and two pseudoscalar mesons.

5.2 LQ$\bar{\text{D}}$ couplings and semi-leptonic decays

We now turn to the lepton-number violating case and consider the decay channels driven by the couplings of LQD-type. In the absence of RpV bilinear terms, the semi-leptonic decay amplitudes are all expected to scale linearly with a λ' coupling and receive a quadratic suppression with the mass of the mediating sfermions.

In Fig. 2, we investigate the neutralino decays in the mass-range $[0, 2]$ GeV in several scenarios without sfermion mixing. Correspondingly, the squark (and slepton) mediators always belong to the first or second generation. In the upper left-hand quadrant, we consider the decays proceeding through $\lambda'_{111} \neq 0$. Wilson coefficients of scalar-type dominate the considered

regime, resulting in suppressed vector-meson production (mostly driven by tensor operators here; shown in blue) as compared to the channels with pseudoscalar mesons. Unsurprisingly, the two-body decays with one pseudoscalar meson in the final state (shown in green) primarily involve (neutral or charged) pions, and secondarily, η (η'): these are the channels coupling to the $\bar{U}D$ and $\bar{D}D$ pseudoscalar sources opened by λ'_{111}. In addition, despite the kinematic suppression, three-body decays (shown in red) are quite important and eventually dominate the semi-leptonic width at high mass. At $m_\psi \approx 1.2\,\text{GeV}$ (the upper limit for fully reliable form factors), they already amount to about 50% of the magnitude of the two-body widths: the radiation of an extra pion indeed costs little phase space at such energies and the relevant channels thus involve two pions or pion+η in their final state, which also couple to the $\bar{U}D$ and $\bar{D}D$ scalar sources.

In the upper right-hand quadrant of Fig. 2, we turn to the case of a dominant λ'_{122}, hence switching on the $\bar{S}S$ scalar and pseudoscalar sources (charm production is kinematically forbidden). As a result, the two-body decays of the neutralino are dominated by η or η' production. The two-pion channels dominate the three-body decay widths: while $K-\bar{K}$ is allowed, these channels only open in the higher range of mass. In addition, the Cabibbo Kobayashi Maskawa (CKM) mixing between first and second generation ($V_{us}^{\text{CKM}} \approx 0.225$) opens a hardly suppressed $\bar{U}S$ source which controls the channels with charged leptons in the final state and leads to substantial kaon (and $K+\pi$) production. In any case, this scenario is characterized by a higher threshold for neutralino decays, since favored final states are all more massive than in the scenario with dominant λ'_{111}.

In the second row of Fig. 2, we consider LQD couplings that violate the strange number: λ'_{112} (left) and λ'_{121} (right). The resulting $\bar{U}S$, $\bar{D}S$ or $\bar{S}D$ (pseudo)scalar sources trigger kaon production in semi-leptonic neutralino decays. Three-body channels are again competitive in the higher mass window, involving the radiation of one extra pion. A noteworthy difference between the two scenarios emerges from CKM mixing, as the V_{us}^{CKM} allows for pion+electron final states, already accessible in the lower mass range, in the case where the second-generation index is carried by the $SU(2)_L$ doublet, but not if it is carried by the $SU(2)_L$ singlet.

In all the scenarios of Fig. 2, we have displayed the inclusive partonic width as a dashed orange curve. Such a description is obviously unreliable in the lower mass-regime, as it ignores thresholds and tends to underestimate the actual semi-leptonic decay width by orders of magnitude (due to the three-body phase-space suppression). Yet, as the semi-leptonic width becomes increasingly controlled by the production of meson pairs for $m_\psi \gtrsim 1.5\,\text{GeV}$, the partonic approach also starts providing an inclusive width of comparable magnitude. In fact, the partonic description of uncharmed semi-leptonic decays may be quantitatively more meaningful at $m_\psi \approx 2\,\text{GeV}$, where the mesonic form factors are largely speculative.

In Fig. 3, we study the impact of the sfermion sector on the neutralino decays. We restrict ourselves to masses $m_\psi \leqslant 1.2$, where the two-meson form factors are fully reliable. In the plots on the left-hand side, we linger on scenarios with strange-number violation and no squark mixing, but consider very massive squarks of first and second generations at 100 TeV, instead of 1 TeV in the previous set of plots — sleptons continue to be comparatively light, at 1 TeV. The mediation of squarks in low-energy processes is thus inhibited. However, in the upper plot, we observe that the suppression of squark mediation hardly affects the magnitude of neutralino decay channels driven by (pseudo)scalar $\bar{Q}Q$ currents — a factor ~ 0.5 as compared to the corresponding scenario in Fig. 2. Indeed, these operators also receive contributions from slepton mediators, which remain effective independently from the squark sector. On the other hand, the tensor (and vector) currents receive only contributions from squark mediators, so that the widths involving vector mesons in the final state are even more suppressed than in the previous scenarios. In the lower plot, we vary the lepton index of the LQD coupling, setting it to λ'_{321}: as tau production is kinematically inaccessible in the decays of a light neutralino,

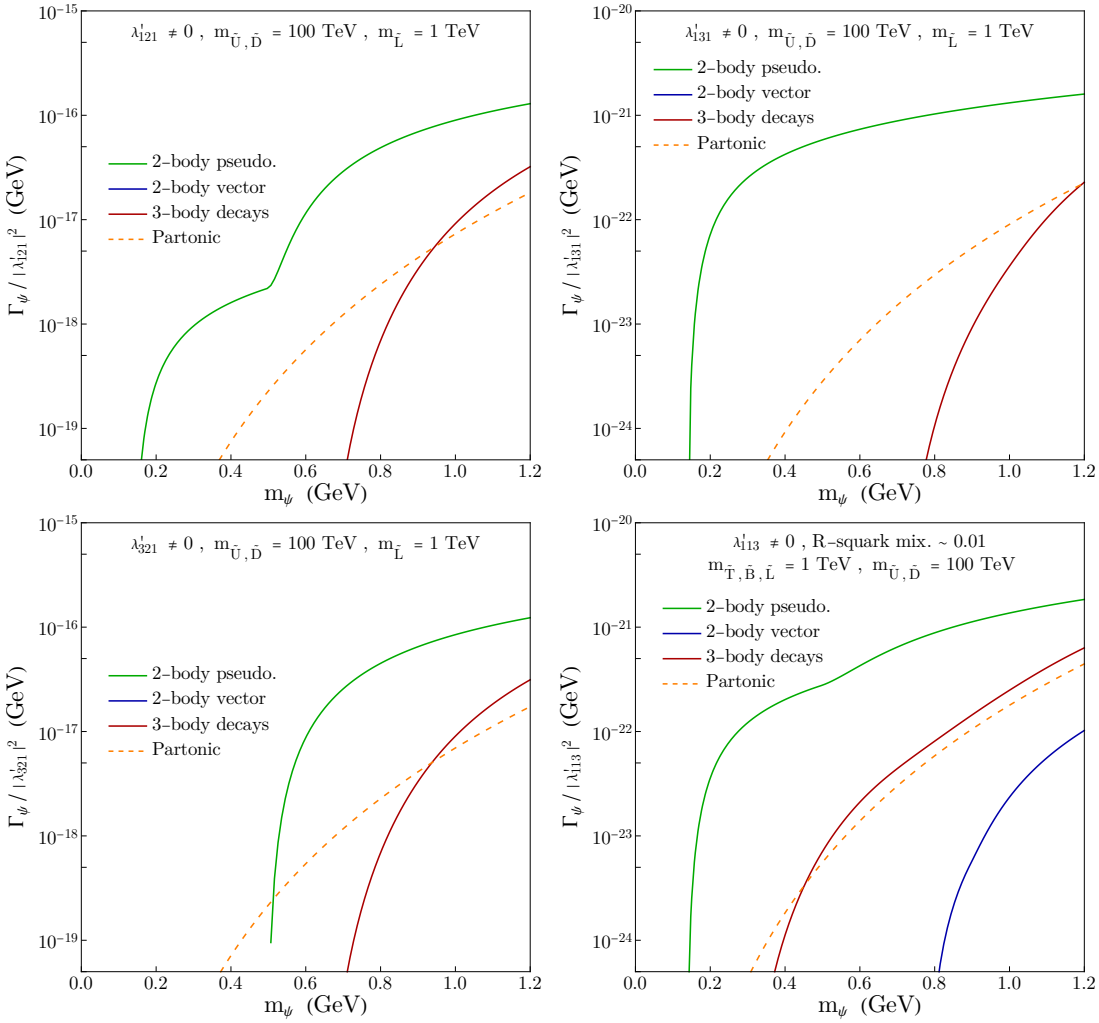

Figure 3: Sfermion effects in semi-Leptonic decay width of the neutralino in scenarios. The squark masses of the first two generations are set to 100 TeV, while the third generation squarks and the sleptons remain at about 1 TeV. *Left*: LQD couplings violating the strange-number are considered, λ'_{121} (up) and λ'_{321} (down) in the presence of heavy squarks; *Right*: LQD couplings involving the third generation squarks, λ'_{131} or λ'_{113}, are considered without (up) or in the presence of flavor-violating squark mixing (down). The color conventions are similar to those of Fig. 2.

only decay channels involving a neutrino in the final state are allowed, which contrasts with the scenario of the upper plot.

On the right-hand side of Fig. 3, we consider mixing effects in the quark or squark sectors. We already stressed, when discussing Fig. 2, that λ'_{121} opened electron+pion decays through the CKM mixing, but not λ'_{112}. The same applies to λ'_{131}, considered in the upper plot, with an additional suppression of $|V_{ub}^{\rm CKM}/V_{us}^{\rm CKM}|^2 \approx 3 \cdot 10^{-4}$. On the other hand, while strange-number violating decays are allowed with λ'_{121} (see plot in the upper left-hand corner), this is no longer possible with third generation quarks, so that the only allowed decay channels in the scenario with dominant λ'_{131} are indeed these CKM-rotated ones involving electrons and pions. In the lower plot, we consider a dominant λ'_{113}: this coupling cannot funnel the decays of a light neutralino except in the presence of squark mixing involving the right-handed sbottom. Therefore, we introduce a mixing of $\mathcal{O}(10^{-2})$ between down-type squarks of the two first

generations and of the third one, resulting in neutralino decays into leptons and pions, η or kaons. As the squark mediation is crucial in the neutralino widths employing squark mixing, we observe that the decays into vector mesons recover their relative, though subleading, importance with respect to the channels involving a single pseudoscalar meson. Finally, we stress that similar effects are triggered by slepton mixing and that, in the scenario with dominant λ'_{321} (in the lower left-end corner of Fig. 3), decays involving electrons or muons in the final state could be restored in this fashion.

To summarize this discussion, we have observed that the semi-leptonic decays of the neutralino driven by LQD-couplings are controlled by scalar and pseudoscalar contributions. In particular, the production of a single meson dominates the low-mass window, with meson pair production becoming competitive as soon as pion radiation costs little relative energy, at $m_\psi \approx 1\,\text{GeV}$. A partonic description may become reliable around the tau mass. Furthermore, it would be relevant to place flavor-violating effects in neutralino decays into the perspective of known constraints on flavor transitions.

5.3 Lepton-number violating bilinear couplings

The interplay of trilinear couplings of LLE-type with leptonic neutralino decays offers limited novelty. Instead, we here consider the decays triggered by RpV bilinear terms. We observe that such terms are generically present in the Lagrangian and generated by the renormalization group as soon as lepton number is violated. There exist two categories of bilinear couplings in the RpV MSSM:

1. the mass terms μ_i, $i = 1,2,3$, introduced in the superpotential in particular produce lepton-electrowikino mixing, opening the possibility of breaking lepton number directly at the level of the external legs of the operators mediating neutralino decays;

2. the mass-squared terms B_i, $i = 1,2,3$, introduced in the soft SUSY-breaking Lagrangian result in admixtures between the heavy Higgs doublets and the slepton sector: lepton number may then be violated in neutralino decays by the double nature of the scalar mediators.

We first focus on scenarios with non-vanishing scalar bilinears B_i. Considering that it is necessary to break lepton number in diagrams mediating neutralino decays, hence play on the double nature of the slepton-Higgs admixtures, contributions to four-fermion operators then necessarily exploit Higgs couplings to SM fermions: this implies a Yukawa suppression for all the decay amplitudes involving light fermions, adding to the scale and mixing suppression factors. Consequently, the neutralino partial widths conveyed by the B_i terms are typically very narrow. On the other hand, contrarily to the trilinear terms, the individual bilinear couplings do not specify the nature of the ensuing (lepton-number violating) decays, so that semi-leptonic channels dominated by the strange Yukawa coexist with leptonic channels triggered by the muon Yukawa and electromagnetic decays conveyed by chargino / slepton-Higgs loops. The latter escape the Yukawa suppression while the loop-factor suppression might be alleviated by the logarithmic dependence in case of hierarchical spectra.

We illustrate this discussion about neutralino decays mediated by the bilinear scalar terms with Fig. 4. In the upper left-hand quadrant, through $B_2 \neq 0$, we generate a mixing of magnitude $\sim 10^{-2}$ between heavy Higgs and smuons taking mass in the TeV range. The higher mass-window is dominated by semi-leptonic neutralino decays conveyed by the strange Yukawa coupling, in particular the $\eta^{(\prime)}\nu_\mu$ and, to a lesser extent (after V_{us}^{CKM} suppression), the $K^+\mu^-$ channels. Below threshold for these, the $\pi^0\nu_\mu$ and $\pi^+\mu^-$ decays can still contribute, though at a reduced rate. Similarly to the previous subsection, channels involving meson pairs in the final state contribute at a sizable level in the upper mass range: the $\bar{S}S$ operator mediates $\pi\pi$

as well as KK decays, while $\bar{u}s$ conveys charged $K\pi$ decays. We also observe that the specific pattern imprinted by the B_i couplings favors semi-leptonic scalar contributions over vector and tensor, so that decays into vector mesons are virtually absent. Leptonic decay channels are of comparable magnitude, which can be related to the comparative size of the muon and strange Yukawa couplings, and are mostly represented by the $\nu_\mu\mu^+\mu^-$ channel: it is remarkable that this specific decay is not straightforwardly generated with the antisymmetric LLE couplings, unless slepton mixing is present. $\nu_e e^+\mu^-$ and $\nu_\mu e^+e^-$ contribute to the neutralino decays at a far lesser extent. Finally, the radiative decays $\psi \to \gamma\nu_\mu, \gamma\bar{\nu}_\mu$, generated by chargino and Higgs-smuon loops, control the low mass-range.

In the upper right-hand corner of Fig. 4, we turn to the case of a dominant B_3, resulting in Higgs-stau mixing. The size of leptonic and semi-leptonic decays is not fundamentally modified by this replacement, with all leading channels now involving ν_τ instead of ν_μ (while the

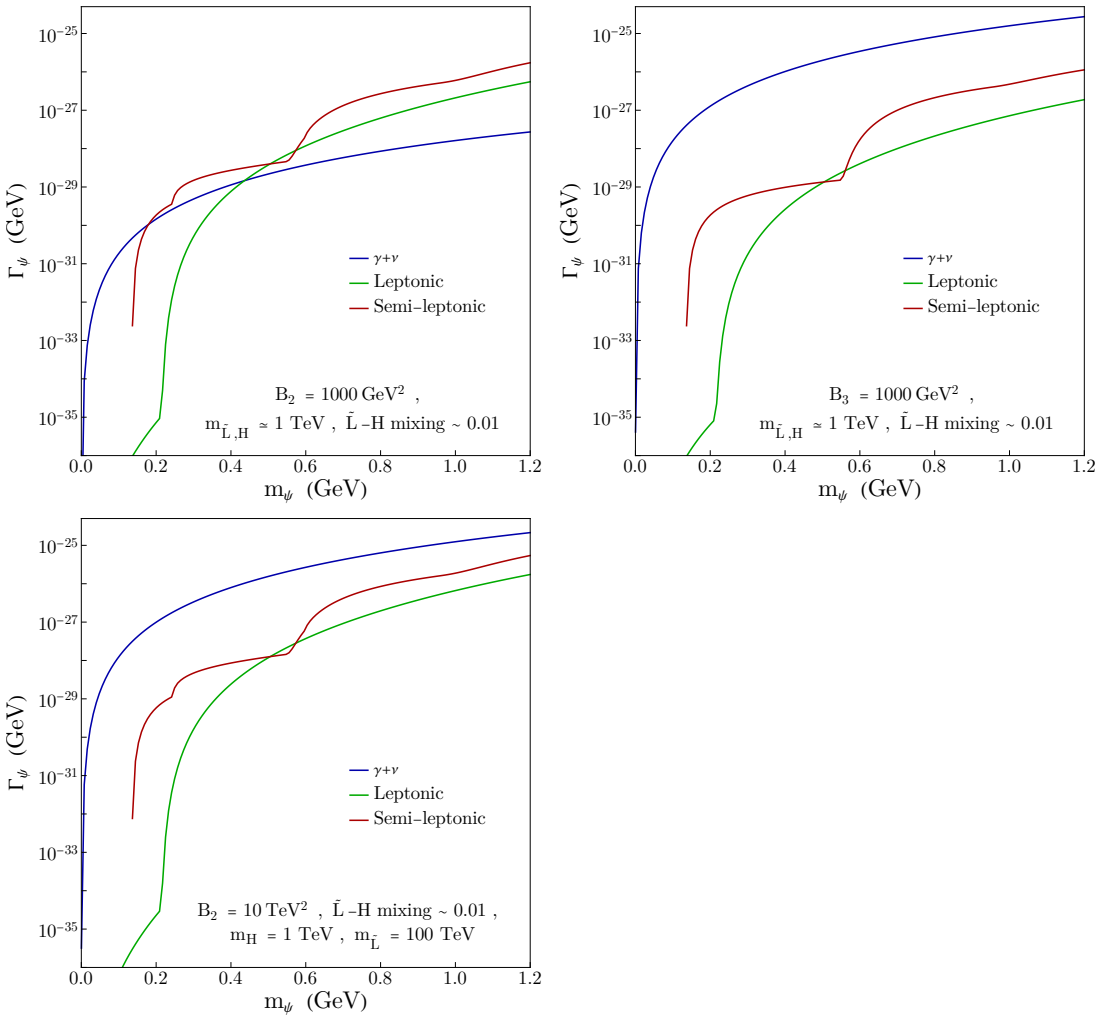

Figure 4: Neutralino decays mediated by bilinear scalar terms. Higgs-slepton mixing is maintained at the level of $\sim 10^{-2}$. *Left*: neutralino decays employ the non-vanishing B_2, resulting in Higgs-smuon mixing; all these scalars take mass in the TeV range in the upper plot while sleptons are made very heavy (100 TeV) in the lower plot. *Right*: B_3, generating Higgs-stau mixing, is chosen to dominate the decays. The blue, green and red curves respectively represent the inclusive photon+neutrino, leptonic and semi-leptonic widths.

charged channel are kinematically forbidden by the replacement $\mu \to \tau$). On the other hand, the electromagnetic dipole contributions are greatly enhanced by the additional mediation of τ + Higgs-stau loops, involving a sizable Yukawa coupling ($\tan\beta$ is set to the typical value of 10) as well as a large $\ln \frac{M_{\rm SUSY}}{m_\tau}$. As a consequence, the $\psi \to \gamma \nu_\tau, \gamma \bar{\nu}_\tau$ widths dominate the neutralino decays in all the considered mass-range.

In the lower plot of Fig. 4, we return to the case of a dominant B_2, but now with a decoupling SUSY spectrum, albeit the Higgs-slepton mixing is kept at the level of 10^{-2}. Leptonic and semi-leptonic decay channels receive contributions of comparable magnitude to their counterparts of the upper plot, since the Higgs-dominated mediator in the TeV range has about unchanged properties. On the other hand, the importance of the radiative channels $\psi \to \gamma \nu_\mu, \gamma \bar{\nu}_\mu$ is increased by the mass-splitting between scalar states, as cancellations between the smuon and Higgs loops no longer occur.

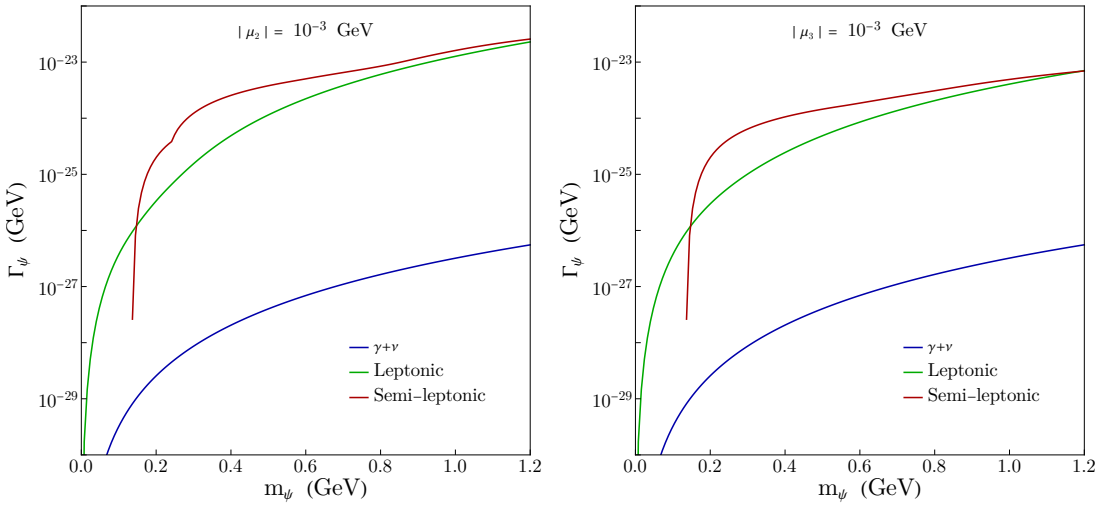

Figure 5: Neutralino decays mediated by bilinear 'fermionic' terms. *Left*: neutralino decays employ the non-vanishing $\mu_2 = 10^{-3}$, resulting in chargino-muon and neutralino-ν_μ mixing. *Right*: a non-vanishing μ_3 replaces μ_2. The blue, green and red curves respectively represent the inclusive photon+neutrino, leptonic and semi-leptonic widths.

In Fig. 5, we consider the impact of the supersymmetric bilinear mixing parameters μ_i on the decays of a light neutralino. The main effect rests with the appearance of chargino-lepton and neutralino-neutrino admixtures and the RpV / lepton-number breaking does not involve the mediation of heavy fields. Indeed, electroweak gauge bosons then become the main contributors to the dimension 6 operators of Eqs.(3-5) and the corresponding neutralino decays hardly depend on the mass of the sfermion fields. Nevertheless, large values of the μ_i parameters would generate unacceptable neutrino masses [20]. For small μ_i, the latter depend on secondary mixing of the neutrinos with higgsinos and gauginos, so that small mass values are easily explained. In practice, we will consider $|\mu_i| = 10^{-3}$ GeV, which typically produces neutrino masses under 10^{-10} GeV in our scenarios. We note that the bilinear 'fermionic' parameters also generate mixing of the heavy Higgs doublets with the right-handed sleptons (similarly to B_i with the left-handed sleptons). Yet, as the effect is proportional to the lepton masses (and the small μ_i), it hardly has a chance to reach phenomenological relevance for scalar fields in the TeV range or beyond.

In the plot on the left-hand side of Fig. 5, we set μ_2 to 10^{-3} GeV, thus generating mixings of the chargino states with the muon and of the neutralino states with ν_μ. The mediation of

the neutralino decays by electroweak gauge bosons means that, contrarily to the previously considered scenarios, the vector operators dominate in the current setup. Consequently, the production of vector mesons (mainly ρ's) controls the semi-leptonic decays in the higher mass window, generating large decays into pion pairs + muons or neutrinos. Decay channels involving single pions are still significant and control the lower mass range, where kinematically allowed. Leptonic decays are also sizable, which can be traced back to the gauge nature of the mediation. Channels violating muon number, such as $\nu_\mu \ell^+ \ell^-$ or $\nu_\mu \bar{\nu}_\ell \nu_\ell$, are determinant in our example, with decays involving a neutrino+antineutrino triplet controlling the low mass-range. Electromagnetic decays are subleading in this scenario and proceed essentially through a muon / W loop. The plot on the right-hand side, corresponding to the decays caused by $\mu_3 = 10^{-3}$ GeV, offers limited novelty, except for the disappearance of the channels with a charged lepton, since τ production is kinematically forbidden. The mediation through gauge fields ensures that the decay widths remain of comparable magnitude with respect to the previous case.

We thus observe that the choice of lepton-number violating bilinear parameter controlling the decays of the light bino-like state leaves a distinguishable imprint on the leading channels. This variety simply reflects the diversity of the accessible low-energy operators on which the lepton- and baryon-number violating parameters of the UV theory (in our case the RpV MSSM) can project.

6 Discussion and Conclusions

In this paper, we have systematically studied the baryon- or lepton-number violating decays of a light exotic fermion from the perspective of an EFT and connected this analysis to the specific case of a SUSY RpV UV completion. Here the light fermion is a light bino ψ, a well motivated scenario, with mass $m_\psi \leqslant m_\tau$. The interactions of the exotic particle with hadronic matter represent the main difficulty in controlling its phenomenology, but are of paramount importance in order to provide testable predictions for searches at collider and derive reliable limits.

In this respect, our study, though marking a step forward, still leaves ample room for improvement. In particular, baryon-number violating decays of the light fermion are at best qualitatively described by our chiral approach restricting to the fundamental baryon octet: going beyond this approximation certainly represents a challenge for hadronic physics, however. Turning to the semi-leptonic channels, our analysis greatly benefits from the inclusion of the two-meson + lepton final states through the form-factors of Ref. [61]: these decay modes indeed tend to overtake the channels with a single meson at energies of about 1 GeV. Corresponding effects can be approximated by the production of vector resonances only in scenarios where operators of vector type dominate, which, in the RpV context, only happens when the phenomenology is determined by the bilinear parameters μ_i. Extending our predictions for masses above 1.2 GeV however rests on uncertain theoretical ground, especially in the case of the determinant scalar operators. For all channels, we have neglected (non-hadronic) interactions between final state particles, which is certainly an illegitimate approximation in the immediate vicinity of decay thresholds.

Further improvements could also be considered from the perspective of the low-energy EFT itself, such as the inclusion of next-to-leading QCD logarithms in the group evolution. Higher-orders in the matching to the RpV MSSM would likewise help ascertain how much of the associated phenomenology can be absorbed within the definition of the tree-level couplings. The definition of multiple thresholds may also be justified in scenarios with hierarchical exotic sectors.

Another desirable feature would rest with a systematic assessment of the low-energy phenomenological constraints applying to the tested scenario. In the case of baryon-number violation, we have already studied the bounds resulting from nucleon decays in Ref. [48] and sketched how to apply limits from neutron-antineutron oscillations in the presence of a light exotic fermion in Appendix G. Further limits from di-nucleon decays [71] are also applicable. Nevertheless, associated constraints strongly depend on the modelization of the interactions of the new physics field with hadronic matter and require a good control over flavor-violating effects in corresponding observables, both in and beyond the SM. Concerning the lepton-number violating case, the various scenarios need to be put into the perspective of *e.g.* neutrino mass predictions or lepton-flavor transitions, often requiring the inclusion of loop effects.

Extension of our analysis for masses above the τ mass *a priori* appears as a difficult undertaking, since the interactions with charmed hadrons then need to be reliably described. A few individual channels, such as semi-leptonic decays involving a single $D^{(*)}$ meson, can be predicted (and have been considered in *e.g.* Ref. [33]), but would generally prove insufficient as soon as the phase space allows for further meson radiation.

Finally, quantitatively testing the existence of a light exotic field requires to control the production modes of this particle in experiment, which may be difficult when this production employs hadronic channels. A few reference formulae for calculable processes can be found in the appendix.

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A Dimension 6 operators for neutralino decay in two-component spinor notation

For clarity, we express all the operators listed in Eqs. (3-6), which potentially mediate a neutralino decay in two-component spinor notation.

- electromagnetic dipoles:

$$\mathcal{E}_i \equiv \frac{e}{16\pi^2}(\psi\sigma^{\mu\nu}\nu_i)F_{\mu\nu}\,,\ (i=1,2,3)\,, \tag{45}$$

- leptonic operators:

$$\begin{aligned}
&\widetilde{\mathcal{N}}_{ijk} \equiv (\psi\,\nu_i)(\nu_j\,\nu_k)\,,\\
&(i,j,k)\in\{(1,2,2),(1,3,3),(1,2,3),(2,1,1),(2,3,3),(2,1,3),(3,1,1),(3,2,2)\}\,,\\
&\mathcal{N}_{ijk} \equiv (\bar\psi\bar\sigma^{\mu}\nu_i)(\bar\nu_j\bar\sigma_{\mu}\nu_k)\,, \qquad (1\leqslant i\leqslant k\leqslant 3)\,,\\
&\mathcal{S}^{\nu e\,L}_{ijk} \equiv (\psi\,\nu_i)(e^c_j e_k)\,, \qquad \mathcal{S}^{\nu e\,R}_{ijk} \equiv (\psi\,\nu_i)(\bar e_j \bar e^c_k)\,,\\
&\mathcal{V}^{\nu e\,L}_{ijk} \equiv (\bar\psi\bar\sigma^{\mu}\nu_i)(\bar e_j\bar\sigma_{\mu}e_k)\,, \qquad \mathcal{V}^{\nu e\,R}_{ijk} \equiv (\bar\psi\bar\sigma^{\mu}\nu_i)(e^c_j\sigma_{\mu}\bar e^c_k)\,,\\
&\mathcal{T}^{\nu e}_{ijk} \equiv (\psi\sigma^{\mu\nu}\nu_i)(e^c_j\sigma_{\mu\nu}e_k)\,, \qquad (i=1,2,3\,,\ j,k=1,2)\,,
\end{aligned} \tag{46}$$

- semi-leptonic operators:

$$
\begin{aligned}
&\mathcal{S}^{eq\,LL}_{ijk} \equiv (\psi e_i)(d^c_j u_k)\,, && \mathcal{S}^{eq\,LR}_{ijk} \equiv (\psi e_i)(\bar{d}_j \bar{u}^c_k)\,,\\
&\mathcal{S}^{eq\,RL}_{ijk} \equiv (\bar{\psi}\bar{e}^c_i)(d^c_j u_k)\,, && \mathcal{S}^{eq\,RR}_{ijk} \equiv (\bar{\psi}\bar{e}^c_i)(\bar{d}_j \bar{u}^c_k)\,,\\
&\mathcal{V}^{eq\,LL}_{ijk} \equiv (\bar{\psi}\bar{\sigma}^\mu e_i)(\bar{d}_j \bar{\sigma}_\mu u_k)\,, && \mathcal{V}^{eq\,LR}_{ijk} \equiv (\bar{\psi}\bar{\sigma}^\mu e_i)(d^c_j \sigma_\mu \bar{u}^c_k)\,,\\
&\mathcal{V}^{eq\,RL}_{ijk} \equiv (\psi\sigma^\mu \bar{e}^c_i)(\bar{d}_j \bar{\sigma}_\mu u_k)\,, && \mathcal{V}^{eq\,RR}_{ijk} \equiv (\psi\sigma^\mu \bar{e}^c_i)(d^c_j \sigma_\mu \bar{u}^c_k)\,,\\
&\mathcal{T}^{eq\,L}_{ijk} \equiv (\psi\sigma^{\mu\nu} e_i)(d^c_j \sigma_{\mu\nu} u_k)\,, && \mathcal{T}^{eq\,R}_{ijk} \equiv (\bar{\psi}\bar{\sigma}^{\mu\nu} \bar{e}^c_i)(\bar{d}_j \bar{\sigma}_{\mu\nu} \bar{u}^c_k)\,,\\
& && (i,j=1,2\,,\ k=1)\,,\\
&\mathcal{S}^{\nu u\,L}_{ijk} \equiv (\psi\,\nu_i)(u^c_j u_k)\,, && \mathcal{S}^{\nu u R}_{ijk} \equiv (\psi\,\nu_i)(\bar{u}_j \bar{u}^c_k)\,,\\
&\mathcal{V}^{\nu u\,L}_{ijk} \equiv (\bar{\psi}\bar{\sigma}^\mu\,\nu_i)(\bar{u}_j \bar{\sigma}_\mu u_k)\,, && \mathcal{V}^{\nu u R}_{ijk} \equiv (\bar{\psi}\bar{\sigma}^\mu\,\nu_i)(u^c_j \sigma_\mu \bar{u}^k_k)\,,\\
&\mathcal{T}^{\nu u}_{ijk} \equiv (\psi\sigma^{\mu\nu}\,\nu_i)(u^c_j \sigma_{\mu\nu} u_k)\,, && (i=1,2,3\,,\ j,k=1)\,,\\
&\mathcal{S}^{\nu d\,L}_{ijk} \equiv (\psi\,\nu_i)(d^c_j d_k)\,, && \mathcal{S}^{\nu d R}_{ijk} \equiv (\psi\,\nu_i)(\bar{d}_j \bar{d}^c_k)\,,\\
&\mathcal{V}^{\nu d\,L}_{ijk} \equiv (\bar{\psi}\bar{\sigma}^\mu\,\nu_i)(\bar{d}_j \bar{\sigma}_\mu d_k)\,, && \mathcal{V}^{\nu d R}_{ijk} \equiv (\bar{\psi}\bar{\sigma}^\mu\,\nu_i)(d^c_j \sigma_\mu \bar{d}^k_k)\,,\\
&\mathcal{T}^{\nu d}_{ijk} \equiv (\psi\sigma^{\mu\nu}\,\nu_i)(d^c_j \sigma_{\mu\nu} d_k)\,, && (i=1,2,3\,,\ j,k=1,2)\,,
\end{aligned}
\tag{47}
$$

- hadronic operators:

$$
\begin{aligned}
&\mathcal{H}^{LL}_{ijk} \equiv \varepsilon_{\alpha\beta\gamma}(\psi u^{c\,\alpha}_i)(d^{c\,\beta}_j d^{c\,\gamma}_k)\,, && \mathcal{H}^{LR}_{ijk} \equiv \varepsilon_{\alpha\beta\gamma}(\psi u^{c\,\alpha}_i)(\bar{d}^\beta_j \bar{d}^\gamma_k)\,,\\
&\mathcal{H}^{RL}_{ijk} \equiv \varepsilon_{\alpha\beta\gamma}(\bar{\psi}\bar{u}^\alpha_i)(d^{c\,\beta}_j d^{c\,\gamma}_k)\,, && \mathcal{H}^{RR}_{ijk} \equiv \varepsilon_{\alpha\beta\gamma}(\bar{\psi}\bar{u}^\alpha_i)(\bar{d}^\beta_j \bar{d}^\gamma_k)\,,\\
& && (i=1\,,\ 1\leqslant j<k\leqslant 2)\,,\\
&\widetilde{\mathcal{H}}^{LL}_{ijk} \equiv \varepsilon_{\alpha\beta\gamma}(\psi d^{c\,\beta}_j)(u^{c\,\alpha}_i d^{c\,\gamma}_k)\,, && \widetilde{\mathcal{H}}^{LR}_{ijk} \equiv \varepsilon_{\alpha\beta\gamma}(\psi d^{c\,\beta}_j)(\bar{u}^\alpha_i \bar{d}^\gamma_k)\,,\\
&\widetilde{\mathcal{H}}^{RL}_{ijk} \equiv \varepsilon_{\alpha\beta\gamma}(\bar{\psi}\bar{d}^\beta_j)(u^{c\,\alpha}_i d^{c\,\gamma}_k)\,, && \widetilde{\mathcal{H}}^{RR}_{ijk} \equiv \varepsilon_{\alpha\beta\gamma}(\bar{\psi}\bar{d}^\beta_j)(\bar{u}^\alpha_i \bar{d}^\gamma_k)\,,\\
& && (i=1\,,\ j,k=1,2)\,.
\end{aligned}
$$

B Feynman Rules

Below, we list the relevant MSSM RpV couplings of Eq. (9). Mixing matrices are defined in appendix C.

B.1 Electroweak gauge couplings

$$
g_L^{Z\chi^0_i\chi^0_j} = -\frac{\sqrt{g_1^2+g_2^2}}{2}\left[N_{id}N^*_{jd} - N_{iu}N^*_{ju} + N_{i\nu_l}N^*_{j\nu_l}\right] = -\left(g_R^{Z\chi^0_j\chi^0_i}\right)^*\,, \tag{48a}
$$

$$
g_L^{Z\chi^+_i\chi^-_j} = \sqrt{g_1^2+g_2^2}\left[c_W^2\,U_{iw}U^*_{jw} + \tfrac{c_W^2-s_W^2}{2}(U_{id}U^*_{jd} + U_{ie_l}U^*_{je_l})\right] = -g_R^{Z\chi^-_j\chi^+_i}\,, \tag{48b}
$$

$$
g_R^{Z\chi^+_i\chi^-_j} = \sqrt{g_1^2+g_2^2}\left[c_W^2\,V^*_{iw}V_{jw} + \tfrac{c_W^2-s_W^2}{2}V^*_{iu}V_{ju} - s_W^2\,V^*_{ie_l}V_{je_l}\right] = -g_L^{Z\chi^-_j\chi^+_i}\,, \tag{48c}
$$

$$
g_L^{Zu^c_iu_j} = -\frac{\sqrt{g_1^2+g_2^2}}{2}\left(1-\tfrac{4}{3}s_W^2\right)\delta_{ij} = -g_R^{Zu_ju^c_i}\,, \tag{48d}
$$

$$g_R^{Zu_i^c u_j} = \frac{2}{3}\sqrt{g_1^2+g_2^2}\,s_W^2\,\delta_{ij} = -g_L^{Zu_j u_i^c}, \tag{48e}$$

$$g_L^{Zd_i^c d_j} = \frac{\sqrt{g_1^2+g_2^2}}{2}\left(1-\tfrac{2}{3}s_W^2\right)\delta_{ij} = -g_R^{Zd_j d_i^c}, \tag{48f}$$

$$g_R^{Zd_i^c d_j} = -\frac{\sqrt{g_1^2+g_2^2}}{3}s_W^2\,\delta_{ij} = -g_L^{Zd_j d_i^c}, \tag{48g}$$

$$\begin{aligned} g_L^{W^+\chi_i^0\chi_j^-} &= -g_2\left[N_{iw}U_{jw}^* + \tfrac{1}{\sqrt{2}}N_{id}U_{jd}^* + \tfrac{1}{\sqrt{2}}N_{i\nu_l}U_{je_l}^*\right] \\ &= \left(g_L^{W^-\chi_j^+\chi_i^0}\right)^* = -g_R^{W^+\chi_j^-\chi_i^0} = -\left(g_R^{W^-\chi_i^0\chi_j^+}\right)^*, \end{aligned} \tag{48h}$$

$$g_R^{W^+\chi_i^0\chi_j^-} = -g_2\left[N_{iw}^*V_{jw} - \tfrac{1}{\sqrt{2}}N_{iu}^*V_{ju}\right] = \left(g_R^{W^-\chi_j^+\chi_i^0}\right)^* = -g_L^{W^+\chi_j^-\chi_i^0} = -\left(g_L^{W^-\chi_i^0\chi_j^+}\right)^*, \tag{48i}$$

$$g_L^{W^-d_i^c u_j} = -\frac{g_2}{\sqrt{2}}V_{ji}^{\mathrm{CKM}\,*} = -g_R^{W^-u_j d_i^c} = \left(g_L^{W^+u_j^c d_i}\right)^* = -\left(g_R^{W^+d_i u_j^c}\right)^*, \tag{48j}$$

$$g_R^{W^-d_i^c u_j} = 0 = -g_L^{W^-u_j d_i^c} = \left(g_R^{W^+u_j^c d_i}\right)^* = -\left(g_L^{W^+d_i u_j^c}\right)^*. \tag{48k}$$

B.2 Scalar couplings

$$g_L^{S_m^0\chi_i^0\chi_j^0} = \frac{1}{2}(g_1N_{ib}^* - g_2N_{iw}^*)\left[X_{md}^*N_{jd}^* - X_{mu}^*N_{ju}^* + X_{m\tilde{N}_l}^*N_{j\nu_l}^*\right] + (i\leftrightarrow j) = \left(g_R^{S_m^0\chi_j^0\chi_i^0}\right)^*, \tag{49a}$$

$$\begin{aligned} g_L^{S_m^0\chi_i^+\chi_j^-} &= \frac{Y_e^l}{\sqrt{2}}V_{ie_l}^*\left[X_{m\tilde{N}_l}U_{jd}^* - X_{md}U_{je_l}^*\right] - \frac{g_2}{\sqrt{2}}\left[X_{md}^*V_{iw}^*U_{jd}^* + X_{mu}^*V_{iu}^*U_{jw}^* + X_{m\tilde{N}_l}^*V_{iw}^*U_{je_l}^*\right] \\ &\quad - \frac{\lambda_{lnp}}{\sqrt{2}}X_{m\tilde{N}_l}V_{ie_p}^*U_{je_n}^* = \left(g_R^{S_m^0\chi_j^+\chi_i^-}\right)^* = g_L^{S_m^0\chi_j^-\chi_i^+} = \left(g_R^{S_m^0\chi_i^-\chi_j^+}\right)^*, \end{aligned} \tag{49b}$$

$$g_L^{S_m^0u_i^cu_j} = -\frac{Y_u^i}{\sqrt{2}}X_{mu}\delta_{ij} = \left(g_R^{S_m^0u_j^cu_i}\right)^* = g_L^{S_m^0u_ju_i^c} = \left(g_R^{S_m^0u_iu_j^c}\right)^*, \tag{49c}$$

$$g_L^{S_m^0d_i^cd_j} = -\frac{Y_d^i}{\sqrt{2}}X_{md}\delta_{ij} - \frac{\lambda'_{lji}}{\sqrt{2}}X_{m\tilde{N}_l} = \left(g_R^{S_m^0d_j^cd_i}\right)^* = g_L^{S_m^0d_jd_i^c} = \left(g_R^{S_m^0d_id_j^c}\right)^*, \tag{49d}$$

$$\begin{aligned} g_L^{S_m^-\chi_i^+\chi_j^0} &= -\frac{g_1}{\sqrt{2}}(X_{mu}^{C\,*}V_{iu}^* + 2X_{mE_R^l}^{C\,*}V_{ie_l}^*)N_{jb}^* - \frac{g_2}{\sqrt{2}}X_{mu}^{C\,*}(V_{iu}^*N_{jw}^* + \sqrt{2}V_{iw}^*N_{ju}^*) \\ &\quad - Y_e^lV_{ie_l}^*(X_{mE_L^l}^{C\,*}N_{jd}^* - X_{md}^{C\,*}N_{j\nu_l}^*) - \lambda_{lnp}X_{mE_L^n}^{C\,*}V_{ie_p}^*N_{j\nu_l}^* \\ &= \left(g_R^{S_m^+\chi_i^-\chi_j^0}\right)^* = g_L^{S_m^-\chi_j^0\chi_i^+} = \left(g_R^{S_m^+\chi_j^0\chi_i^-}\right)^*, \end{aligned} \tag{49e}$$

$$\begin{aligned} g_R^{S_m^-\chi_i^+\chi_j^0} &= \frac{1}{\sqrt{2}}(X_{md}^{C\,*}U_{id} + X_{mE_L^l}^{C\,*}U_{ie_l})(g_1N_{jb} + g_2N_{jw}) - g_2U_{iw}(X_{md}^{C\,*}N_{jd} + X_{mE_L^l}^{C\,*}N_{j\nu_l}) \\ &\quad - Y_e^lX_{mE_R^l}^{C\,*}(U_{ie_l}N_{jd} - U_{id}N_{j\nu_l}) - \lambda_{lnp}^*X_{mE_R^p}^{C\,*}U_{ie_n}N_{j\nu_l} \\ &= \left(g_L^{S_m^+\chi_i^-\chi_j^0}\right)^* = g_R^{S_m^-\chi_j^0\chi_i^+} = \left(g_L^{S_m^+\chi_j^0\chi_i^-}\right)^*, \end{aligned} \tag{49f}$$

$$g_L^{S_m^-d_i^cu_j} = Y_d^iV_{ji}^{\mathrm{CKM}\,*}X_{md}^{C\,*} + \lambda'_{lni}V_{jn}^{\mathrm{CKM}\,*}X_{mE_L^l}^{C\,*} = \left(g_R^{S_m^+u_j^cd_i}\right)^* = g_L^{S_m^-u_jd_i^c} = \left(g_R^{S_m^+d_iu_j^c}\right)^*, \tag{49g}$$

$$g_R^{S_m^-d_i^cu_j} = Y_u^jV_{ji}^{\mathrm{CKM}\,*}X_{mu}^{C\,*} = \left(g_L^{S_m^+u_j^cd_i}\right)^* = g_R^{S_m^-u_jd_i^c} = \left(g_L^{S_m^+d_iu_j^c}\right)^*, \tag{49h}$$

$$g_L^{U_m u_i^c \chi_j^0} = -Y_u^i X^*_{mU_L^i} N^*_{ju} + \frac{2\sqrt{2}}{3} g_1 X^*_{mU_R^i} N^*_{jb} = \left(g_R^{U_m^* \chi_j^0 u_i}\right)^* = g_L^{U_m \chi_j^0 u_i^c} = \left(g_R^{U_m^* u_i \chi_j^0}\right)^*, \quad (49i)$$

$$g_R^{U_m u_i^c \chi_j^0} = -Y_u^i X^*_{mU_R^i} N_{ju} - \frac{1}{\sqrt{2}} X^*_{mU_L^i} \left(\tfrac{g_1}{3} N_{jb} + g_2 N_{jw}\right) = \left(g_L^{U_m^* \chi_j^0 u_i}\right)^* = g_R^{U_m \chi_j^0 u_i^c} = \left(g_L^{U_m^* u_i \chi_j^0}\right)^*, \quad (49j)$$

$$g_L^{D_m d_i^c \chi_j^0} = -Y_d^i X^*_{mD_L^i} N^*_{jd} - \frac{\sqrt{2}}{3} g_1 X^*_{mD_R^i} N^*_{jb} - \lambda'_{lni} X^*_{mD_L^n} N^*_{j\nu_l} = \left(g_R^{D_m^* \chi_j^0 d_i}\right)^* = g_L^{D_m \chi_j^0 d_i^c} = \left(g_R^{D_m^* d_i \chi_j^0}\right)^*, \quad (49k)$$

$$g_R^{D_m d_i^c \chi_j^0} = -Y_d^i X^*_{mD_R^i} N_{jd} - \frac{1}{\sqrt{2}} X^*_{mD_L^i} \left(\tfrac{g_1}{3} N_{jb} - g_2 N_{jw}\right) - \lambda'^*_{lin} X^*_{mD_R^n} N_{j\nu_l} = \left(g_L^{D_m^* \chi_j^0 d_i}\right)^* = g_R^{D_m \chi_j^0 d_i^c} = \left(g_L^{D_m^* d_i \chi_j^0}\right)^*, \quad (49l)$$

$$g_L^{U_m d_i^c \chi_j^-} = Y_d^i V_{ni}^{\mathrm{CKM}*} X^*_{mU_L^n} U^*_{jd} + \lambda'_{lpi} V_{np}^{\mathrm{CKM}*} X^*_{mU_L^n} U^*_{je_l} = \left(g_R^{U_m^* \chi_j^+ d_i}\right)^* = g_L^{U_m \chi_j^- d_i^c} = \left(g_R^{U_m^* d_i \chi_j^+}\right)^*, \quad (49m)$$

$$g_R^{U_m d_i^c \chi_j^-} = Y_u^n V_{ni}^{\mathrm{CKM}*} X^*_{mU_R^n} V_{ju} - g_2 V_{ni}^{\mathrm{CKM}*} X^*_{mU_L^n} V_{jw} = \left(g_L^{U_m^* \chi_j^+ d_i}\right)^* = g_R^{U_m \chi_j^- d_i^c} = \left(g_L^{U_m^* d_i \chi_j^+}\right)^*, \quad (49n)$$

$$g_L^{D_m^* u_i \chi_j^-} = Y_d^n V_{in}^{\mathrm{CKM}*} X_{mD_R^n} U^*_{jd} - g_2 V_{in}^{\mathrm{CKM}*} X_{mD_L^n} U^*_{jw} + \lambda'_{lnp} V_{in}^{\mathrm{CKM}*} X_{mD_R^p} U^*_{je_l} = \left(g_R^{D_m \chi_j^+ u_i^c}\right)^* = g_L^{D_m^* \chi_j^- u_i} = \left(g_R^{D_m u_i^c \chi_j^+}\right)^*, \quad (49o)$$

$$g_R^{D_m^* u_i \chi_j^-} = Y_u^i V_{in}^{\mathrm{CKM}*} X_{mD_L^n} V_{ju} = \left(g_L^{D_m \chi_j^+ u_i^c}\right)^* = g_R^{D_m^* \chi_j^- u_i} = \left(g_L^{D_m u_i^c \chi_j^+}\right)^*, \quad (49p)$$

$$g_L^{U_m^* d_i^c d_j^c} = -\lambda''_{nij} X_{mU_R^n} = \left(g_R^{U_m d_i d_j}\right)^* = -g_L^{U_m^* d_j^c d_i^c} = -\left(g_R^{U_m d_j d_i}\right)^*, \quad g_R^{U_m^* d_i^c d_j^c} = 0, \quad (49q)$$

$$g_L^{D_m^* u_i^c d_j^c} = -\lambda''_{inj} X_{mD_R^n} = \left(g_R^{D_m d_j u_i}\right)^* = g_L^{D_m^* d_j^c u_i^c} = \left(g_R^{D_m u_i d_j}\right)^*, \quad g_R^{D_m^* u_i^c d_j^c} = 0. \quad (49r)$$

C Mixing in the RpV MSSM

C.1 Mixing in the chargino / lepton sector

We work in the field basis where the sneutrino fields do not take a vacuum expectation value (vev) [67]. Then the Dirac mass-terms in the chargino/lepton sector read

$$-\mathcal{L} \ni (\tilde{w}^+, \tilde{h}_u^+, e_R^{cf}) \mathcal{M}_{\tilde{C}} (\tilde{w}^-, \tilde{h}_d^-, e_L^g)^T + h.c., \quad (50)$$

with the 5×5 matrix:

$$\mathcal{M}_{\tilde{C}} = \begin{bmatrix} M_2 & g_2 v_d & 0 \\ g_2 v_u & \mu & \mu_j \\ 0 & 0 & Y_e^i \delta_{ij} v_d \end{bmatrix}. \quad (51)$$

Here i, j correspond to the flavor indices. μ_j is the RpV bilinear coupling. For the mass matrix of Eq. (51) we perform a singular value decomposition [44] with a pair of unitary

matrices such that: $\mathcal{M}_{\tilde{C}} = V^T \mathrm{diag}(m_{\chi_i^\pm}) U$, from which we deduce the mass eigenstates: $\chi_i^+ = V_{iw}^* \tilde{w}^+ + V_{ih}^* \tilde{h}_u^+ + V_{ie_j}^* e_R^{cj}$, $\chi_i^- = U_{iw}^* \tilde{w}^- + U_{ih}^* \tilde{h}_d^- + U_{ie_j}^* e_L^j$.

In the hierarchical context $|M_2|, |\mu|, \big||M_2|-|\mu|\big| \gg g v, |\mu_j|, Y_e^g v_d$, we can expand the mixing matrices according to $W = \mathbf{I} + W^{(1)} + \ldots$, for $W = U, V$. The non-trivial elements of $W^{(1)}$ read:

$$
\begin{aligned}
U^{(1)}_{1d} &= \frac{g_2(v_d M_2^* + v_u \mu)}{|M_2|^2 - |\mu|^2}, & U^{(1)}_{2w} &= -\frac{g_2(v_d M_2 + v_u \mu^*)}{|M_2|^2 - |\mu|^2}, & U^{(1)}_{2e_l} &= \frac{\mu_l}{\mu}, \\
V^{(1)}_{1u} &= \frac{g_2(v_d \mu + v_u M_2^*)}{|M_2|^2 - |\mu|^2}, & V^{(1)}_{2w} &= -\frac{g_2(v_d \mu^* + v_u M_2)}{|M_2|^2 - |\mu|^2}, & U^{(1)}_{3_l d} &= -\frac{\mu_l^*}{\mu^*},
\end{aligned}
\tag{52}
$$

where the mass-indices 1, 2, 3_l respectively refer to mostly wino, higgsino and lepton states. Here $e_l + e\,\mu, \tau$ for $l = 1, 2, 3$.

C.2 Mixing in the neutralino / neutrino sector

Similarly to the chargino/lepton sector, electroweak-symmetry breaking and RpV mixing produce the 7×7 Majorana mass-matrix between neutral gauginos, higgsinos and neutrinos, in the basis $(\tilde{b}^0, \tilde{w}^0, \tilde{h}_d^0, \tilde{h}_u^0, \nu_L^f)$:

$$
\mathcal{M}_{\tilde{N}} = \begin{bmatrix}
M_1 & 0 & -\frac{g_1}{\sqrt{2}} v_d & \frac{g_1}{\sqrt{2}} v_u & 0 \\
0 & M_2 & \frac{g_2}{\sqrt{2}} v_d & -\frac{g_2}{\sqrt{2}} v_u & 0 \\
-\frac{g_1}{\sqrt{2}} v_d & \frac{g_2}{\sqrt{2}} v_d & 0 & -\mu & 0 \\
\frac{g_1}{\sqrt{2}} v_u & -\frac{g_2}{\sqrt{2}} v_u & -\mu & 0 & -\mu_g \\
0 & 0 & 0 & -\mu_g & 0
\end{bmatrix}.
\tag{53}
$$

This matrix is Takagi diagonalized [44] via a 7×7 unitary matrix N according to $\mathcal{M}_{\tilde{N}} = N^T \mathrm{diag}(m_{\chi^0}) N$, from which we deduce the mass eigenstates $\chi_i^0 = N_{ib}^* \tilde{b} + N_{iw}^* \tilde{w}^0 + N_{ih_d}^* \tilde{h}_d^0 + N_{ih_u}^* \tilde{h}_u^0 + N_{i\nu_f}^* \nu_L^f$, $i = 1, \ldots, 7$. Again, in a hierarchical context, the mixing matrix can be expanded as $N = N^{(0)} + N^{(1)} + N^{(2)} + \ldots$, with:

$$
N^{(0)} = \begin{bmatrix}
1 & 0 & 0 & 0 & 0 \\
0 & 1 & 0 & 0 & 0 \\
0 & 0 & \frac{1}{\sqrt{2}} & \frac{1}{\sqrt{2}} & 0 \\
0 & 0 & -\frac{1}{\sqrt{2}} & \frac{1}{\sqrt{2}} & 0 \\
0 & 0 & 0 & 0 & 1
\end{bmatrix},
\quad
\begin{aligned}
N^{(1)}_{1h_d} &= -\tfrac{g_1}{\sqrt{2}} \tfrac{M_1^* v_d + \mu v_u}{|M_1|^2 - |\mu|^2}, & N^{(1)}_{1h_u} &= \tfrac{g_1}{\sqrt{2}} \tfrac{M_1^* v_u + \mu v_d}{|M_1|^2 - |\mu|^2}, \\
N^{(1)}_{2h_d} &= \tfrac{g_2}{\sqrt{2}} \tfrac{M_2^* v_d + \mu v_u}{|M_2|^2 - |\mu|^2}, & N^{(1)}_{2h_u} &= -\tfrac{g_2}{\sqrt{2}} \tfrac{M_2^* v_u + \mu v_d}{|M_2|^2 - |\mu|^2}, \\
N^{(1)}_{3b} &= \tfrac{g_1}{2} \tfrac{(v_d - v_u)(M_1 - \mu^*)}{|M_1|^2 - |\mu|^2}, & N^{(1)}_{3w} &= \tfrac{g_2}{2} \tfrac{(v_u - v_d)(M_2 - \mu^*)}{|M_2|^2 - |\mu|^2}, \\
N^{(1)}_{4b} &= -\tfrac{g_1}{2} \tfrac{(v_d + v_u)(\mu^* + M_1)}{|M_1|^2 - |\mu|^2}, & N^{(1)}_{4w} &= \tfrac{g_2}{2} \tfrac{(v_d + v_u)(\mu^* + M_2)}{|M_2|^2 - |\mu|^2}, \\
N^{(1)}_{3\nu_l} &= \tfrac{\mu_l}{\sqrt{2}\mu} = -N^{(1)}_{4\nu_l}, & N^{(1)}_{5_l h_d} &= -\tfrac{\mu_l^*}{\mu^*},
\end{aligned}
\tag{54}
$$

and all other elements of $N^{(1)}$ are trivial. As we only consider bino- or neutrino-dominated fields in the external legs, we only provide the bino-neutrino mixing of second order:

$$
N^{(2)}_{1\nu_l} = -\frac{g_1}{\sqrt{2}} \frac{\mu_l}{M_1} \frac{M_1 v_u + \mu^* v_d}{|M_1|^2 - |\mu|^2}, \quad N^{(2)}_{5_l b} = -\frac{g_1}{\sqrt{2}} \frac{\mu_l^* v_d}{M_1^* \mu^*}.
\tag{55}
$$

In this description, the indices 1, 2, 3, 4 and 5_l refer to mostly bino, wino, a pair of higgsino and neutrino-like mass eigenstates.

C.3 Mixing in the Higgs / slepton sector

The neutral Higgs and sneutrino states mix at tree-level through the soft SUSY-breaking RpV parameters $B_i \equiv B_i^R + i\,B_i^I$ (after simplifying other mixing parameters via the minimization equations of the scalar potential). The 5×5 mass-blocks $\mathcal{M}_S^2$ and $\mathcal{M}_P^2$ for the CP-even and CP-odd states in the gauge eigenbases $(h_d^0, h_u^0, \frac{1}{\sqrt{2}}\mathrm{Re}[\tilde{N}_i])$ and $(a_d^0, a_u^0, \frac{1}{\sqrt{2}}\mathrm{Im}[\tilde{N}_i])$ (respectively) read:

$$\mathcal{M}_S^2 = \begin{bmatrix} M_A^2 s_\beta^2 + M_Z^2 c_\beta^2 & -(M_A^2+M_Z^2)s_\beta c_\beta & (B_i^R)^T t_\beta \\ -(M_A^2+M_Z^2)s_\beta c_\beta & M_A^2 c_\beta^2 + M_Z^2 s_\beta^2 & -(B_i^R)^T \\ (B_i^R)t_\beta & -(B_i^R) & \mathrm{Re}[M_{\tilde{N}}^2] \end{bmatrix}, \tag{56}$$

$$\mathcal{M}_P^2 = \begin{bmatrix} M_A^2 s_\beta^2 & M_A^2 s_\beta c_\beta & (B_i^R)^T t_\beta \\ M_A^2 s_\beta c_\beta & M_A^2 c_\beta^2 & (B_i^R)^T \\ (B_i^R)t_\beta & (B_i^R) & \mathrm{Re}[M_{\tilde{N}}^2] \end{bmatrix}, \tag{57}$$

where (B_i) should be understood as a column 3-vector and $M_{\tilde{N}}^2 = \left(m_{L\,mn}^2 + \mu_m^*\mu_n + M_Z^2\frac{c_{2\beta}}{2}\delta_{mn}\right)$ is the 3×3 mass subblock corresponding to the sneutrinos. $\mathrm{Im}[M_{\tilde{L}}^2]$ and B_i^I mix CP-even and CP-odd degrees of freedom and the mass eigenstates can be decomposed as $S_p^0 = X_{pd}^R h_d^0 + X_{pu}^R h_u^0 + \frac{X_{p\tilde{N}_i}^R}{\sqrt{2}}\mathrm{Re}[\tilde{N}_i] + X_{pd}^I a_d^0 + X_{pu}^I a_u^0 + \frac{X_{p\tilde{N}_i}^I}{\sqrt{2}}\mathrm{Im}[\tilde{N}_i]$, where $(X_{p\cdot}^R, X_{p\cdot}^I)$ forms an orthonormal eigenvector of the mass-matrix. We also define the notation $X \equiv X^R + i\,X^I$.

Assuming $M_A^2, M_{\tilde{N}\,ii}^2 \gg M_Z^2, |B_i|, |M_{\tilde{N}\,ij}^2|$ $(j\neq i)$, the mass eigenstates can be approximated by:

$$\begin{aligned} &h^0 \equiv c_\beta h_d^0 + s_\beta h_u^0, \;\; m_{h^0}^2 \approx M_Z^2 c_{2\beta}^2, && G^0 \equiv -c_\beta a_d^0 + s_\beta a_u^0, \;\; m_{G^0}^2 = 0, \\ &H^0 \equiv s_\beta h_d^0 - c_\beta h_u^0, \;\; m_{H^0}^2 \approx M_A^2 + M_Z^2 s_{2\beta}^2, && A^0 \equiv s_\beta a_d^0 + c_\beta a_u^0, \;\; m_{A^0}^2 = M_A^2, \\ &\tilde{N}_p^R \equiv \tfrac{1}{\sqrt{2}}\mathrm{Re}[\tilde{N}_p], \;\; \tilde{N}_p^I \equiv \tfrac{1}{\sqrt{2}}\mathrm{Im}[\tilde{N}_p], && m_{\tilde{N}_p}^2 \approx m_{L\,pp}^2 + |\mu_p|^2 + M_Z^2\tfrac{c_{2\beta}}{2}. \end{aligned} \tag{58}$$

First order corrections may be encoded as $S_p^0 = (\delta_{pm} + X_{pm}^{(1)})(S_m^0)^{(0)}$ with $(S_m^0)^{(0)}$ corresponding to the states of Eq. (58), where the non-trivial elements of $X^{(1)}$ are:

$$\begin{aligned} X_{h^0H^0}^{(1)} &= -\frac{M_Z^2 s_{4\beta}}{2M_A^2} = -X_{H^0h^0}^{(1)}, & X_{\tilde{N}_p^R\tilde{N}_q^I}^{(1)} &= -\frac{\mathrm{Im}[M_{\tilde{N}}^2]_{pq}}{m_{\tilde{N}_p}^2 - m_{\tilde{N}_q}^2} = -X_{\tilde{N}_q^I\tilde{N}_p^R}^{(1)}, \\ X_{\tilde{N}_p^R\tilde{N}_q^R}^{(1)} &= \frac{\mathrm{Re}[M_{\tilde{N}}^2]_{pq}}{m_{\tilde{N}_p}^2 - m_{\tilde{N}_q}^2} = -X_{\tilde{N}_q^R\tilde{N}_p^R}^{(1)}, & X_{\tilde{N}_p^I\tilde{N}_q^I}^{(1)} &= \frac{\mathrm{Re}[M_{\tilde{N}}^2]_{pq}}{m_{\tilde{N}_p}^2 - m_{\tilde{N}_q}^2} = -X_{\tilde{N}_q^I\tilde{N}_p^I}^{(1)}, \\ X_{H^0\tilde{N}_p^R}^{(1)} &= \frac{B_p^R/c_\beta}{M_A^2 - m_{\tilde{N}_p}^2} = -X_{\tilde{N}_p^RH^0}^{(1)}, & X_{H^0\tilde{N}_p^I}^{(1)} &= \frac{-B_p^I/c_\beta}{M_A^2 - m_{\tilde{N}_p}^2} = -X_{\tilde{N}_p^IH^0}^{(1)}, \\ X_{A^0\tilde{N}_p^I}^{(1)} &= \frac{B_p^R/c_\beta}{M_A^2 - m_{\tilde{N}_p}^2} = -X_{\tilde{N}_p^IA^0}^{(1)}, & X_{A^0\tilde{N}_p^R}^{(1)} &= \frac{B_p^I/c_\beta}{M_A^2 - m_{\tilde{N}_p}^2} = -X_{\tilde{N}_p^RA^0}^{(1)}. \end{aligned} \tag{59}$$

As a result, only the heavy Higgs doublets H^0, A^0 undergo RpV mixing with the sneutrinos at leading order.

Similarly, the mass matrix in the charged sector mixes charged Higgs and sleptons. In the

gauge basis $(H_d^-, H_u^-, \tilde{E}_L^i, (\tilde{E}_R^{c\,i})^*)$,

$$\mathcal{M}_C^2 = \begin{bmatrix} (M_A^2+M_W^2)s_\beta^2 & (M_A^2+M_W^2)s_\beta c_\beta & (B_i)^T t_\beta & (m_{e_i}\mu_i)t_\beta \\ (M_A^2+M_W^2)s_\beta c_\beta & (M_A^2+M_W^2)c_\beta^2 & (B_i)^T & (m_{e_i}\mu_i) \\ (B_i)^* t_\beta & (B_i)^* & M^2_{\tilde{E}_L} & M^2_{\tilde{E}_{LR}} \\ (m_{e_i}\mu_i^*)t_\beta & (m_{e_i}\mu_i^*) & M^2_{\tilde{E}_{RL}} & M^2_{\tilde{E}_R} \end{bmatrix}, \tag{60}$$

$$\begin{aligned} M^2_{\tilde{E}_L} &= \left(m^2_{L\,mn} + \mu^*_m\mu_n + (m^2_{e_m} + \tfrac{g_1^2-g_2^2}{4}v^2 c_{2\beta})\delta_{mn}\right), \\ M^2_{\tilde{E}_R} &= \left(m^2_{E\,mn} + \mu^*_m\mu_n + (m^2_{e_m} - \tfrac{g_1^2}{4}v^2 c_{2\beta})\delta_{mn}\right), \\ M^2_{\tilde{E}_{LR}} &= (A^{e\,*}_{mn}v_d - m_{e_m}\,t_\beta\,\mu\delta_{mn} - \mu_k\lambda^*_{kmn}v_u), \;\; M^2_{\tilde{E}_{RL}} = \left(M^2_{\tilde{E}_{LR}}\right)^\dagger. \end{aligned}$$

The mass eigenstates may then be written as $S_p^- = X^C_{pd}H_d^- + X^C_{pu}H_u^- + X^C_{p\tilde{E}_L^i}\tilde{E}_L^i + X^C_{p\tilde{E}_R^i}(\tilde{E}_R^{c\,i})^*$. In a hierarchical context, they may be approximated by:

$$\begin{aligned} &G^- = -c_\beta H_d^- + s_\beta H_u^-, \;\; m^2_{G^\pm} = 0, && H^- = s_\beta H_d^- + c_\beta H_u^-, \;\; m^2_{H^\pm} = M_A^2 + M_W^2, \\ &\tilde{E}_L^i, \;\; m^2_{\tilde{E}_L^i} = (M^2_{\tilde{E}_L})_{ii}, && (\tilde{E}_R^{c\,i})^*, \;\; m^2_{\tilde{E}_R^i} = (M^2_{\tilde{E}_R})_{ii}. \end{aligned} \tag{61}$$

Then, the non-trivial mixings at first order are given by:

$$\begin{aligned} X^{C(1)}_{H^-\tilde{E}_L^p} &= \frac{B_p/c_\beta}{M_A^2 - m^2_{\tilde{E}_L^p}} = -X^{C(1)}_{\tilde{E}_L^p H^-}, & X^{C(1)}_{H^-\tilde{E}_R^p} &= \frac{m_{e_p}\mu_p/c_\beta}{M_A^2 - m^2_{\tilde{E}_R^p}} = -\left(X^{C(1)}_{\tilde{E}_R^p H^-}\right)^*, \\ X^{C(1)}_{\tilde{E}_L^p\tilde{E}_L^q} &= \frac{(M^2_{\tilde{E}_L})^*_{pq}}{m^2_{\tilde{E}_L^p} - m^2_{\tilde{E}_L^q}} = -\left(X^{C(1)}_{\tilde{E}_L^q\tilde{E}_L^p}\right)^*, & X^{C(1)}_{\tilde{E}_R^p\tilde{E}_R^q} &= \frac{(M^2_{\tilde{E}_R})^*_{pq}}{m^2_{\tilde{E}_R^p} - m^2_{\tilde{E}_R^q}} = -\left(X^{C(1)}_{\tilde{E}_R^q\tilde{E}_R^p}\right)^*, \\ X^{C(1)}_{\tilde{E}_L^p\tilde{E}_R^q} &= \frac{(M^2_{\tilde{E}_{LR}})^*_{pq}}{m^2_{\tilde{E}_L^p} - m^2_{\tilde{E}_R^q}} = -\left(X^{C(1)}_{\tilde{E}_R^q\tilde{E}_L^p}\right)^*. \end{aligned} \tag{62}$$

C.4 Mixing in the squark sector

The sfermion mixing matrices can be written in the $(\tilde{F}_L, \tilde{F}_R^{c\,*})$ basis $(\tilde{F} = \tilde{U}, \tilde{D})$ as:

$$\mathcal{M}^2_{\tilde{F}} = \begin{bmatrix} M^2_{\tilde{F}_L} & M^2_{\tilde{F}_{LR}} \\ M^2_{\tilde{F}_{RL}} & M^2_{\tilde{F}_R} \end{bmatrix}, \quad \begin{cases} M^2_{\tilde{F}_L} \equiv m^2_{F_L} + m_f^2 + \frac{1}{2}\left(\frac{\mathcal{Y}_L^f}{2}g_1^2 - I_3^f g_2^2\right)(v_u^2 - v_d^2), \\ M^2_{\tilde{F}_R} \equiv m^2_{F_R} + m_f^2 + \frac{\mathcal{Y}_R^f}{4}g_1^2(v_u^2 - v_d^2), \\ M^2_{\tilde{F}_{LR}} \equiv A_f^* v_f - m_f\mu\frac{v_{f'}}{v_f} - \mu_i\lambda'^*_{ijk}v_u\delta_{\tilde{F}\tilde{D}}, \quad M^2_{\tilde{F}_{RL}} = M^{2\,\dagger}_{\tilde{F}_{LR}}, \end{cases} \tag{63}$$

where f is the fermion corresponding to the sfermion $\tilde{F}$, while f' is its $SU(2)_L$ partner. Then, $\mathcal{Y}^f_{L,R}$ are the associated hypercharges, I_3^f, the isospin. Finally, v_f denotes the vev. of the Higgs doublet to which the fermion f couples at tree-level. Each entry should again be understood as a 3×3 block. We write the mass eigenstates as $\tilde{F}_p = X_{p\tilde{F}_L^i}\tilde{F}_L^i + X_{p\tilde{F}_R^i}\tilde{F}_R^i$. It is also possible to work under the simplifying assumption that inter-generation and left-right mixings are suppressed compared to the diagonal hierarchy.

D Loop Integrals for the $\psi \to \bar{\nu}_i \gamma$ Transition

The loop integrals $\mathcal{I}^{V,S}_{L,R}$ of Eq. (20) can be written in terms of the standard Passarino-Veltman functions [68, 69]:

$$\begin{aligned}
\mathcal{I}^V_L(\tfrac{m^2_\psi}{M^2_W}, \tfrac{m^2_{\chi_j}}{M^2_W}) &= 2M^2_W \Big([C_{21} - C_{23} - C_{12}](p_\psi, -p_\gamma, m^2_{\chi_j}, M^2_W, M^2_W) \\
&\qquad + [C_{21} - C_{23} + 2C_{11} - C_{12} + C_0](-p_\psi, p_\gamma, M^2_W, m^2_{\chi_j}, m^2_{\chi_j}) \Big) , \qquad (64) \\
\mathcal{I}^V_R(\tfrac{m^2_\psi}{M^2_W}, \tfrac{m^2_{\chi_j}}{M^2_W}) &= 4M^2_W \Big(C_{11}(p_\psi, -p_\gamma, m^2_{\chi_j}, M^2_W, M^2_W) - [C_{11} + C_0](-p_\psi, p_\gamma, M^2_W, m^2_{\chi_j}, m^2_{\chi_j}) \Big) , \\
\mathcal{I}^S_L(\tfrac{m^2_\psi}{m^2_S}, \tfrac{m^2_f}{m^2_S}) &= m^2_S \Big([C_{11} + C_0](p_\psi, -p_\gamma, m^2_f, m^2_S, m^2_S) - C_{11}(-p_\psi, p_\gamma, m^2_S, m^2_f, m^2_f) \Big) , \\
\mathcal{I}^S_R(\tfrac{m^2_\psi}{m^2_S}, \tfrac{m^2_f}{m^2_S}) &= -m^2_S \Big([C_{21} - C_{23} + C_{11} - C_{12}](p_\psi, -p_\gamma, m^2_f, m^2_S, m^2_S) \\
&\qquad + [C_{21} - C_{23} + C_{11} - C_{12}](-p_\psi, p_\gamma, m^2_S, m^2_f, m^2_f) \Big) .
\end{aligned}$$

In the regime $m^2_\psi \ll M^2_W, m^2_S$ that we consider here, it is useful (in terms of numerical stability) to replace these functions by their limit with vanishing first argument:

$$\begin{aligned}
&\mathcal{I}^V_L(0,z) = -\tfrac{1}{1-z}\big[\tfrac{3}{2} - 2z + (2z-1) f_\gamma(z)\big] , \qquad \mathcal{I}^V_R(0,z) = 4\big(1 - f_\gamma(z)\big) , \qquad (65) \\
&\mathcal{I}^S_L(0,z) = -f_\gamma(z)/z , \qquad \mathcal{I}^S_R(0,z) = \tfrac{1}{2(1-z)}\big(f_\gamma(z) - \tfrac{1}{2}\big) , \qquad (66) \\
&f_\gamma(z) \equiv -\tfrac{z}{1-z}\big(1 + \tfrac{\ln z}{1-z}\big) .
\end{aligned}$$

E Wilson Coefficients at Leading order in RpV / Mixing

In this section, we express the Wilson coefficients at leading order in terms of the RpV parameters, in an expansion with respect to the mixing between gauge / flavor eigenstates and neglecting the Yukawa couplings. This operation provides a more intuitive view of the contributions of the RpV MSSM to the Wilson coefficients. We naturally assume that the light neutralino is bino-like. Nevertheless, the corresponding results do not accurately describe all arbitrary parameter choices of the model. In particular, bilinear RpV contributions (mixing leptons / sleptons and Higgs / Higgsinos, hence involving higher orders in mixing or Yukawa couplings) are then systematically discarded. In regimes where $|\lambda^{(\prime),(\prime\prime)}_{ijk}| \ll \left|\frac{\mu_{i'}}{\mu}\right|$ or if final states with charged leptons or hadrons are kinematically inaccessible, these bilinear parameters would however determine the decays of the light neutralino. For instance, decays into neutrinos are controlled by the following Wilson coefficients, involving two degrees of mixing:

$$C[\mathcal{N}_{ijk}] = -\frac{g_1 v_d (\mu_i \delta_{jk} + \mu_k \delta_{ij})}{4\sqrt{2}(1+\delta_{ik})\,\mu} \Big[\frac{1}{m^2_{\tilde{N}_j}} \Big(\frac{g_1^2}{M_1} + \frac{g_2^2}{M_2} \Big) + \frac{g_1^2 + g_2^2}{M_1 M_Z^2} \Big] . \qquad (67)$$

In addition, Yukawa contributions for the muon or the strange quarks may be competitive with mixing terms of higher order.

E.1 Leptonic operators

$$
\begin{aligned}
&C[\widetilde{\mathcal{N}}_{ijk}](\mu_0)\approx 0\,, \qquad C[\mathcal{N}_{ijk}](\mu_0)\approx 0\,, \qquad C[\mathcal{S}^{veR}_{ijk}](\mu_0)\approx 0\,,\\
&C[\mathcal{S}^{veL}_{ijk}](\mu_0)\approx -\frac{g_1}{\sqrt{2}}\Bigg[\lambda_{ikj}\Big(\frac{1}{m^2_{\tilde{N}_i}}-\frac{1}{2m^2_{\tilde{E}^k_L}}+\frac{1}{m^2_{\tilde{E}^j_R}}\Big)-\sum_{l\neq i}\frac{\lambda_{lkj}M^2_{\tilde{N}\,il}}{m^2_{\tilde{N}_i}m^2_{\tilde{N}_l}}+\sum_{l\neq k}\frac{\lambda_{ilj}M^2_{\tilde{E}_L\,lk}}{2m^2_{\tilde{E}^k_L}m^2_{\tilde{E}^l_L}}\\
&\qquad\qquad -\sum_{l\neq j}\frac{\lambda_{ikl}M^2_{\tilde{E}_R\,jl}}{m^2_{\tilde{E}^j_R}m^2_{\tilde{E}^l_R}}\Bigg]\,,\\
&C[\mathcal{V}^{veL}_{ijk}](\mu_0)\approx \frac{g_1\lambda_{ikl}}{2\sqrt{2}}\frac{M^2_{\tilde{E}_{LR}\,jl}}{m^2_{\tilde{E}^j_L}m^2_{\tilde{E}^l_R}}\,, \qquad C[\mathcal{V}^{veR}_{ijk}](\mu_0)\approx \frac{g_1\lambda_{ilj}}{\sqrt{2}}\frac{M^2_{\tilde{E}_{LR}\,lk}}{m^2_{\tilde{E}^l_L}m^2_{\tilde{E}^k_R}}\,,\\
&C[\mathcal{T}^{ve}_{ijk}](\mu_0)\approx \frac{g_1\lambda_{ikl}}{\sqrt{2}m^2_{\tilde{E}^j_R}}\Big[\delta_{jl}-(1-\delta_{jl})\frac{M^2_{\tilde{E}_R\,jl}}{m^2_{\tilde{E}^l_R}}\Big]+\frac{g_1\lambda_{ilj}}{2\sqrt{2}m^2_{\tilde{E}^k_L}}\Big[\delta_{kl}-(1-\delta_{kl})\frac{M^2_{\tilde{E}_L\,lk}}{m^2_{\tilde{E}^l_L}}\Big]\,.
\end{aligned}
\tag{68}
$$

Below, we omit the $\sum_{l\neq\cdot}$ or $(1-\delta_{l\cdot})$ indicating sfermion mixing.

E.2 Semi-leptonic operators

$$
\begin{aligned}
&C[\mathcal{S}^{vu\,L,R}_{ijk}](\mu_0)\approx 0\approx C[\mathcal{V}^{vu\,L,R}_{ijk}](\mu_0)\approx C[\mathcal{T}^{vu}_{ijk}](\mu_0)\,, \qquad C[\mathcal{S}^{vd\,R}_{ijk}](\mu_0)\approx 0\,,\\
&C[\mathcal{S}^{vd\,L}_{ijk}](\mu_0)\approx -\frac{g_1}{\sqrt{2}}\Big[\frac{\lambda'_{lkj}}{m^2_{\tilde{N}_i}}\big(\delta_{il}-\frac{M^2_{\tilde{N}\,il}}{m^2_{\tilde{N}_l}}\big)+\frac{\lambda'_{ilj}}{6m^2_{\tilde{D}^k_L}}\big(\delta_{kl}-\frac{M^2_{\tilde{D}_L\,lk}}{m^2_{\tilde{D}^l_L}}\big)+\frac{\lambda'_{ikl}}{3m^2_{\tilde{D}^j_R}}\big(\delta_{jl}-\frac{M^2_{\tilde{D}_R\,jl}}{m^2_{\tilde{D}^l_R}}\big)\Big]\,,\\
&C[\mathcal{V}^{vd\,L}_{ijk}](\mu_0)\approx -\frac{g_1\lambda'_{ikl}M^2_{\tilde{D}_{LR}\,jl}}{6\sqrt{2}m^2_{\tilde{D}^j_L}m^2_{\tilde{D}^l_R}}\,, \qquad C[\mathcal{V}^{vd\,R}_{ijk}](\mu_0)\approx \frac{g_1\lambda'_{ilj}M^2_{\tilde{D}_{LR}\,lk}}{3\sqrt{2}m^2_{\tilde{D}^l_L}m^2_{\tilde{D}^j_R}}\,,\\
&C[\mathcal{T}^{vd}_{ijk}](\mu_0)\approx \frac{g_1\lambda'_{ikl}}{3\sqrt{2}m^2_{\tilde{D}^j_R}}\Big[\delta_{lj}-\frac{M^2_{\tilde{D}_R\,jl}}{m^2_{\tilde{D}^l_R}}\Big]-\frac{g_1\lambda'_{ilj}}{6\sqrt{2}m^2_{\tilde{D}^k_L}}\Big[\delta_{lk}-\frac{M^2_{\tilde{D}_L\,lk}}{m^2_{\tilde{D}^l_L}}\Big]\,,
\end{aligned}
\tag{69}
$$

$$
\begin{aligned}
&C[\mathcal{S}^{eq\,LL}_{ijk}](\mu_0)\approx \frac{g_1}{\sqrt{2}}\Big[\frac{\lambda'_{lnj}V^{\mathrm{CKM}\,*}_{kn}}{m^2_{\tilde{E}^i_L}}\Big(\delta_{il}-\frac{M^2_{\tilde{E}_L\,li}}{m^2_{\tilde{E}^l_L}}\Big)+\frac{\lambda'_{ilj}V^{\mathrm{CKM}\,*}_{ml}}{6m^2_{\tilde{U}^k_L}}\Big(\delta_{km}-\frac{M^2_{\tilde{U}_L\,mk}}{m^2_{\tilde{U}^m_L}}\Big)\\
&\qquad\qquad +\frac{\lambda'_{ilm}V^{\mathrm{CKM}\,*}_{kl}}{3m^2_{\tilde{D}^j_R}}\Big(\delta_{jm}-\frac{M^2_{\tilde{D}_R\,jm}}{m^2_{\tilde{D}^m_R}}\Big)\Big]\,,\\
&C[\mathcal{S}^{eq\,LR}_{ijk}](\mu_0)\approx 0\,, \quad C[\mathcal{S}^{eq\,RL}_{ijk}](\mu_0)\approx \sqrt{2}g_1\lambda'_{lnj}V^{\mathrm{CKM}\,*}_{kn}\frac{M^2_{\tilde{E}_{LR}\,li}}{m^2_{\tilde{E}^l_L}m^2_{\tilde{E}^i_R}}\,, \quad C[\mathcal{S}^{eq\,RR}_{ijk}](\mu_0)\approx 0\,,\\
&C[\mathcal{V}^{eq\,LL}_{ijk}](\mu_0)\approx \frac{g_1\lambda'_{inm}V^{\mathrm{CKM}\,*}_{kn}}{6\sqrt{2}}\frac{M^2_{\tilde{D}_{LR}\,jm}}{m^2_{\tilde{D}^j_L}m^2_{\tilde{D}^m_R}}\,, \quad C[\mathcal{V}^{eq\,LR}_{ijk}](\mu_0)\approx \frac{\sqrt{2}g_1\lambda'_{inj}V^{\mathrm{CKM}\,*}_{mn}}{3}\frac{M^2_{\tilde{U}_{LR}\,mk}}{m^2_{\tilde{U}^m_L}m^2_{\tilde{U}^k_R}}\,,\\
&C[\mathcal{V}^{eq\,RL}_{ijk}](\mu_0)\approx 0\,, \qquad C[\mathcal{V}^{eq\,RR}_{ijk}](\mu_0)\approx 0\,, \qquad C[\mathcal{T}^{eq\,R}_{ijk}](\mu_0)\approx 0\,,\\
&C[\mathcal{T}^{eq\,L}_{ijk}](\mu_0)\approx \frac{g_1}{\sqrt{2}}\Big[\frac{\lambda'_{ilj}V^{\mathrm{CKM}\,*}_{mn}}{6m^2_{\tilde{U}^k_L}}\Big(\delta_{km}-\frac{M^2_{\tilde{U}_L\,mk}}{m^2_{\tilde{U}^m_L}}\Big)-\frac{\lambda'_{ilm}V^{\mathrm{CKM}\,*}_{kl}}{3m^2_{\tilde{D}^j_R}}\Big(\delta_{jm}-\frac{M^2_{\tilde{D}_R\,jm}}{m^2_{\tilde{D}^m_R}}\Big)\Big]\,.
\end{aligned}
\tag{70}
$$

E.3 Hadronic operators

$$C[\mathcal{H}^{LL}_{ijk}](\mu_0) \approx -\frac{2\sqrt{2}g_1\lambda''_{njk}}{3m^2_{\tilde{U}^i_R}}\Big[\delta_{in} - \frac{M^2_{\tilde{U}_R in}}{m^2_{\tilde{U}^n_R}}\Big], \qquad C[\mathcal{H}^{RL}_{ijk}](\mu_0) \approx -\frac{g_1\lambda''_{njk}}{3\sqrt{2}}\frac{M^2_{\tilde{U}_{LR} in}}{m^2_{\tilde{U}^i_L} m^2_{\tilde{U}^n_R}}, \quad (71)$$

$$C[\widetilde{\mathcal{H}}^{LL}_{ijk}](\mu_0) \approx \frac{\sqrt{2}g_1\lambda''_{ink}}{3m^2_{\tilde{D}^j_R}}\Big[\delta_{jn} - \frac{M^2_{\tilde{D}_R jn}}{m^2_{\tilde{D}^n_R}}\Big], \qquad C[\widetilde{\mathcal{H}}^{RL}_{ijk}](\mu_0) \approx -\frac{g_1\lambda''_{ink}}{3\sqrt{2}}\frac{M^2_{\tilde{D}_{LR} jn}}{m^2_{\tilde{D}^j_L} m^2_{\tilde{D}^n_R}}.$$

E.4 Electromagnetic dipole operators

Neglecting any degree of mixing:

$$\begin{aligned} C^{\text{MSSM}}[\mathcal{E}_i] \approx{}& \frac{g_1}{\sqrt{2}}\lambda_{ijj} m_{e_j}\Big[\frac{1}{m^2_{\tilde{E}^j_L}}\mathcal{I}^S_L\big(\tfrac{m^2_\psi}{m^2_{\tilde{E}^j_L}}, \tfrac{m^2_{e_j}}{m^2_{\tilde{E}^j_L}}\big) + \frac{2}{m^2_{\tilde{E}^j_R}}\mathcal{I}^S_L\big(\tfrac{m^2_\psi}{m^2_{\tilde{E}^j_R}}, \tfrac{m^2_{e_j}}{m^2_{\tilde{E}^j_R}}\big)\Big] \\ &- \frac{N_c g_1}{9\sqrt{2}}\lambda'_{ijj} m_{d_j}\Big[\frac{1}{m^2_{\tilde{D}^j_L}}\mathcal{I}^S_L\big(\tfrac{m^2_\psi}{m^2_{\tilde{D}^j_L}}, \tfrac{m^2_{d_j}}{m^2_{\tilde{D}^j_L}}\big) - \frac{2}{m^2_{\tilde{D}^j_R}}\mathcal{I}^S_L\big(\tfrac{m^2_\psi}{m^2_{\tilde{D}^j_R}}, \tfrac{m^2_{d_j}}{m^2_{\tilde{D}^j_R}}\big)\Big]. \end{aligned} \quad (72)$$

F Kinematic Integrals for the Three-body Leptonic Decays

The kinematic integrals of Eq. (29) are defined by ($x_j + x_k < 1$):

$$I_{\Omega,\Omega'}[x_j, x_k] = \int_{(x_j+x_k)^2}^1 ds\, f_{\Omega,\Omega'}(s)(1-s)^2\Big(1 - 2\tfrac{x_j^2+x_k^2}{s} + \tfrac{(x_j^2-x_k^2)^2}{s^2}\Big)^{1/2}, \quad (73)$$

where the functions $f_{\Omega,\Omega'}$ read (with $f_{\Omega',\Omega} = f_{\Omega,\Omega'}$):

$$\begin{aligned}
&f_{\mathcal{S}^L,\mathcal{S}^L}(s) = \tfrac{1}{2}(s - x_j^2 - x_k^2) = f_{\mathcal{S}^R,\mathcal{S}^R}(s), && f_{\mathcal{S}^L,\mathcal{S}^R}(s) = -x_j x_k = f_{\mathcal{S}^R,\mathcal{S}^L}(s), \\
&f_{\mathcal{V}^L,\mathcal{S}^L}(s) = \tfrac{x_j}{2}[s^{-1}(x_j^2 - x_k^2) - 1] = f_{\mathcal{V}^R,\mathcal{S}^R}(s), && f_{\mathcal{V}^L,\mathcal{S}^R}(s) = \tfrac{x_k}{2}(s^{-1}[x_j^2 - x_k^2] + 1) = f_{\mathcal{V}^R,\mathcal{S}^L}(s), \\
&f_{\mathcal{T}^L,\mathcal{S}^{L,R}}(s) = 0, && f_{\mathcal{V}^L,\mathcal{V}^R}(s) = 2x_j x_k, \\
&f_{\mathcal{V}^L,\mathcal{V}^L}(s) = \tfrac{1}{3}\big\{1 + 2s + \tfrac{x_j^2+x_k^2}{s}(1-s) - \tfrac{(x_j^2-x_k^2)^2}{s^2}(2+s)\big\} = f_{\mathcal{V}^R,\mathcal{V}^R}(s), \\
&f_{\mathcal{T}^L,\mathcal{V}^L}(s) = \tfrac{3x_j}{2}(1 + s^{-1}[x_k^2 - x_j^2]), && f_{\mathcal{T}^L,\mathcal{V}^R}(s) = \tfrac{3x_k}{2}(1 + s^{-1}[x_j^2 - x_k^2]), \\
&f_{\mathcal{T}^L,\mathcal{T}^L}(s) = \tfrac{2+s}{6}\big\{1 + \tfrac{x_j^2+x_k^2}{s} - 2\tfrac{(x_j^2-x_k^2)^2}{s^2}\big\}.
\end{aligned} \quad (74)$$

These can be written as linear combinations of the four base functions $f_m(s) = s^m$, $m \in \{-2, -1, 0, 1\}$; this decomposition can be used for the calculation of the integrals.

$$I_m[x_j,x_k] \equiv \int_{(x_j+x_k)^2}^1 ds\, f_m(s)(1-s)^2\left(1-2\frac{x_j^2+x_k^2}{s}+\frac{(x_j^2-x_k^2)^2}{s^2}\right)^{1/2}, \tag{75}$$

$$I_{-2}[x_j,x_k] = \left[\frac{x_j^2-x_k^2}{2}+\frac{x_j^2+x_k^2}{x_j^2-x_k^2}-\frac{x_j^2x_k^2}{(x_j^2-x_k^2)^3}\right]L_1+\left[1+\frac{x_j^2+x_k^2}{2}\right]L_2+\left[\frac{5}{2}+\frac{x_j^2+x_k^2}{2(x_j^2-x_k^2)^2}\right]S,$$

$$I_{-1}[x_j,x_k] = -\left[x_j^2-x_k^2+\frac{x_j^2+x_k^2}{2(x_j^2-x_k^2)}\right]L_1-\left[\frac{1}{2}+x_j^2+x_k^2-x_j^2x_k^2\right]L_2-\frac{1}{2}[5+x_j^2+x_k^2]S,$$

$$I_0[x_j,x_k] = \frac{x_j^2-x_k^2}{2}L_1+\frac{1}{2}\left[x_j^2+x_k^2+2x_j^2x_k^2(x_j^2+x_k^2-2)\right]L_2 + \frac{1}{6}\left[2+5(x_j^2+x_k^2)-x_j^4-x_k^4-10x_j^2x_k^2\right]S,$$

$$I_1[x_j,x_k] = x_j^2x_k^2\left[1-3(x_j^2+x_k^2)+x_j^4+x_k^4+3x_j^2x_k^2\right]L_2 + \frac{1}{12}\left[1-3(x_j^2+x_k^2)+3x_j^4+3x_k^4+32x_j^2x_k^2-x_j^6-x_k^6-29x_j^2x_k^2(x_j^2+x_k^2)\right]S,$$

$$L_1 \equiv \ln\tfrac{1+X_1}{1-X_1}\,,\quad X_1 \equiv \tfrac{x_j^2-x_k^2}{x_j^2+x_k^2-(x_j^2-x_k^2)^2}S\,,\quad L_2 \equiv \ln\tfrac{1+X_2}{1-X_2}\,,\quad X_2 \equiv -\tfrac{S}{1-x_j^2-x_k^2}\,,$$

$$S \equiv \left[1-2(x_j^2+x_k^2)+(x_j^2-x_k^2)^2\right]^{1/2}.$$

The functions I_m are regular in the limits $x_j \to x_k$ and $x_j, x_k \to 0$.

G Neutron-antineutron Oscillations

The chiral Lagrangian of Eq. (31) implies neutralino-baryon mixing, which can mediate baryon-antibaryon oscillations. Here, we focus on the very constrained case of the neutron [70]. Eq. (31) in particular contains the terms:

$$\mathcal{L}_\psi^{\chi PT} \ni -m_{\psi n}(\bar\Psi P_L n^0) - m^*_{\psi n^c}(\bar\Psi P_R n^0) + h.c.\,, \quad \begin{cases} m_{\psi n} \equiv -\left(\tilde\alpha\, C[\widetilde{\mathcal{H}}^{RL}_{111}]^* + \tilde\beta\, C[\widetilde{\mathcal{H}}^{RR}_{111}]^*\right)(\mu)\,, \\ m_{\psi n^c} \equiv -\left(\tilde\alpha\, C[\widetilde{\mathcal{H}}^{LR}_{111}] + \tilde\beta\, C[\widetilde{\mathcal{H}}^{LL}_{111}]\right)(\mu)\,. \end{cases} \tag{76}$$

We can express this in terms of a mass matrix for the neutralino-neutron-antineutron system:

$$-\frac{1}{2}(\psi\; n_L\; n_R^c)\begin{bmatrix} m_\psi & m_{\psi n} & m_{\psi n^c} \\ m_{\psi n} & m_{nn} & m_n \\ m_{\psi n^c} & m_n & m_{n^cn^c} \end{bmatrix}\begin{pmatrix}\psi \\ n_L \\ n_R^c\end{pmatrix} + h.c.\,, \tag{77}$$

where the entries m_{nn} and $m_{n^cn^c}$ emerge from dimension 9 (six-fermion) operators and correspond to a direct high-energy contribution to the neutron-antineutron mixing. At tree-level, these are mediated by gluino + squark or heavy neutralino + squark diagrams — see *e.g.* Ref. [71] for a recent review. They entail a scale-suppression double to that of $m_{\psi n}$ and $m_{\psi n^c}$. Relevant lattice evaluations of the hadronic matrix elements for the dimension 9 operators can be found in Ref. [72].

Diagonalization of the mass matrix of Eq. (77) generates a mass splitting between the neutron and antineutron. At leading order in $|m_{nn}, m_{n^cn^c}| \ll |m_{\psi n}, m_{\psi n^c}| \ll |m_n \pm m_\psi|$, the two mass eigenvalues for the hadron-dominated states read:

$$\begin{aligned} m_{n+} &\approx m_n + \tfrac{1}{2}\mathrm{Re}[m_{nn}+m_{n^cn^c}] + \tfrac{2m_n-m_\psi}{2(m_n-m_\psi)^2}\mathrm{Re}[m_{\psi n}+m_{\psi n^c}]^2 \\ &\quad + \tfrac{2m_n+m_\psi}{2(m_n+m_\psi)^2}\mathrm{Im}[m_{\psi n}+m_{\psi n^c}]^2\,, \\ m_{n-} &\approx -m_n + \tfrac{1}{2}\mathrm{Re}[m_{nn}+m_{n^cn^c}] - \tfrac{2m_n+m_\psi}{2(m_n+m_\psi)^2}\mathrm{Re}[m_{\psi n}-m_{\psi n^c}]^2 \\ &\quad - \tfrac{2m_n-m_\psi}{2(m_n-m_\psi)^2}\mathrm{Im}[m_{\psi n}-m_{\psi n^c}]^2\,. \end{aligned} \tag{78}$$

This results in a neutron-antineutron oscillation frequency of

$$\omega_{nn^c} \approx \mathrm{Re}[m_{nn}+m_{n^cn^c}]+\frac{2m_n-m_\psi}{2(m_n-m_\psi)^2}\mathrm{Re}[(m_{\psi n}+m^*_{\psi n^c})^2]-\frac{2m_n+m_\psi}{2(m_n+m_\psi)^2}\mathrm{Re}[(m_{\psi n}-m^*_{\psi n^c})^2], \quad (79)$$

which can be compared to the experimental limit on the oscillation time, from Super Kamiokande [70], $\tau_{nn^c} > 2.7\cdot 10^8$ s. Nevertheless, we stress that this particular limit critically depends on the connection between τ_{nn^c} and the O^{16} lifetime, which involves complex nuclear modeling. A comparable bound has been extracted from the deuteron lifetime [73] using EFT methods, but finding significant discrepancy with nuclear models [74]. Naturally, the existence of light new physics (in the presence of a light neutralino) may further modify corresponding analyses. More direct limits from neutron beam experiments are available in Ref. [75], with further improvement expected at the European Spallation Source [76].

Let us focus on the contributions originating in neutron-neutralino mixing, *i.e.* $m_{\psi n}$ and $m_{\psi n^c}$ of Eq. (76). In the RpV MSSM at leading order, $C[\widetilde{\mathcal{H}}^{LR}_{111}]$ and $C[\widetilde{\mathcal{H}}^{RR}_{111}]$ vanish while non-zero $C[\widetilde{\mathcal{H}}^{LL}_{111}]$ and $C[\widetilde{\mathcal{H}}^{RL}_{111}]$ imply flavor mixing in the squark sector. This suggests a source of suppression from flavor-mixing in addition to that originating in the RpV violation. In the approximation of suppressed flavor mixing (see Appendix E), we have:

$$C[\widetilde{\mathcal{H}}^{LL}_{111}](\mu_0) \approx \frac{\sqrt{2}g_1}{3}\lambda''_{1n1}\frac{M^2_{\tilde{D}_R 1n}}{m^2_{\tilde{D}^1_R}m^2_{\tilde{D}^n_R}}, \quad C[\widetilde{\mathcal{H}}^{RL}_{111}](\mu_0) \approx \frac{g_1}{3\sqrt{2}}\lambda''_{1n1}\frac{M^2_{\tilde{D}_{LR}1n}}{m^2_{\tilde{D}^1_L}m^2_{\tilde{D}^n_R}}. \quad (80)$$

Similarly, flavor-mixing is a necessary ingredient to generate non-trivial m_{nn} and $m_{n^cn^c}$, so that the existence of neutron-antineutron oscillations in the RpV MSSM strongly depends on the chosen flavor template.

H Hadronic Decay Widths

The mixing Lagrangian of Eq. (31) generates mixing terms between the neutralino and the baryons, as well as trilinear couplings with mesons:

$$\mathcal{L}_\chi \ni \Omega^{B_k}_{L,R}\bar{\psi}P_{L,R}B_k + i\,C^{\psi M_iB_j}_{L,R}\bar{\psi}P_{L,R}B_jM_i\,, \quad (81)$$

$$\Omega^{B_k}_L \equiv -(\tilde{\alpha}C^{LR*}_{mn}+\tilde{\beta}C^{RR*}_{mn})\mathrm{Tr}[E_{mn}\lambda_k], \quad \Omega^{B_k}_R \equiv -(\tilde{\alpha}C^{RL*}_{mn}+\tilde{\beta}C^{LL*}_{mn})\mathrm{Tr}[E_{mn}\lambda_k], \quad (82)$$

$$C^{\psi M_iB_j}_L \equiv -\frac{\tilde{\alpha}}{f_\chi}C^{LR*}_{mn}\mathrm{Tr}[E_{mn}(\lambda_i\lambda_j+\lambda_j\lambda_i)]-\frac{\tilde{\beta}}{f_\chi}C^{RR*}_{mn}\mathrm{Tr}[E_{mn}(\lambda_i\lambda_j-\lambda_j\lambda_i)], \quad (83)$$

$$C^{\psi M_iB_j}_R \equiv \frac{\tilde{\alpha}}{f_\chi}C^{RL*}_{mn}\mathrm{Tr}[E_{mn}(\lambda_i\lambda_j+\lambda_j\lambda_i)]+\frac{\tilde{\beta}}{f_\chi}C^{LL*}_{mn}\mathrm{Tr}[E_{mn}(\lambda_i\lambda_j-\lambda_j\lambda_i)]. \quad (84)$$

Here, $C^{LL,RR,LR,RL}_{mn}$ stand for the Wilson coefficients of the operators of $\mathcal{H}$- and $\widetilde{\mathcal{H}}$-types and can be straightforwardly identified by comparing with Eq. (31). We have written the baryon and meson octets as $B \equiv \lambda_j B_j$, $M \equiv \lambda_i M_i$, where λ_j are linear combinations of the Gell-Mann matrices, $B_j \in \{\Sigma^0,\Sigma^+,\Sigma^-,n^0,\Xi^0,p^+,\Xi^-,\Lambda^0\}$ and $M_i \in \{\pi^0,\pi^+,\pi^-,K^0,\bar{K}^0,K^+,K^-,\eta^0_8\}$.

Through the mixing terms of Eq. (81), the neutralino may decay into one baryon and one meson using the baryon-meson-baryon interaction (in four-component spinor notation):

$$\mathcal{L}_\chi \ni i\,C^{\bar{B}_kM_iB_j}\left(\bar{B}_k(-i\partial\!\!\!/M_i)\gamma_5B_j\right)+i\,\tilde{C}^{\bar{B}_kM_iB_j}M_i\left(\bar{B}_k\gamma_5B_j\right), \quad (85)$$

$$C^{\bar{B}_kM_iB_j} \equiv \frac{D-F}{f_\chi}\mathrm{Tr}[\lambda^\dagger_k\lambda_j\lambda_i]+\frac{D+F}{f_\chi}\mathrm{Tr}[\lambda^\dagger_k\lambda_i\lambda_j], \quad (86)$$

$$\tilde{C}^{\bar{B}_kM_iB_j} \equiv -\frac{2b_1}{f_\chi}\mathrm{Tr}[\lambda^\dagger_k(M_q\lambda_i+\lambda_iM_q)\lambda_j]-\frac{2b_2}{f_\chi}\mathrm{Tr}[\lambda^\dagger_k\lambda_j(M_q\lambda_i+\lambda_iM_q)], \quad (87)$$

where $M_q = \mathrm{diag}[m_u, m_d, m_s]$ is the quark mass matrix, $D \approx 0.8$, $F \approx 0.47$ while $b_{1,2}$ are ill-known (we will ignore them in practice). Further terms accounting for the weak interaction can be included: they are essential in the description of the hyperon decays — see *e.g.* Ref. [59]. We dispense with a detailed description here: they essentially matter in the immediate mass-vicinity of an hyperon or in situations where other sources of $s \to d$ transition are suppressed.

Then, evaluating the two topologies of Feynman diagrams contributing to the $\psi \to \bar{M}_i \bar{B}_j$ decay amplitude at tree-level, we obtain:

$$\mathcal{A}_\chi[\psi \to \bar{M}_i \bar{B}_j] = -\bar{v}_\psi(p_\psi) g_{L,R}^{\psi M_i B_j} P_{L,R} v_{B_j}(p_{B_j}), \tag{88}$$

$$\begin{aligned} g_L^{\psi M_i B_j} \equiv\ & C_L^{\psi M_i B_j} - C^{\bar{B}_k M_i B_j} \frac{\Omega_L^{B_k}(m_\psi^2 + m_{B_k} m_{B_j}) + \Omega_R^{B_k} m_\psi (m_{B_k} + m_{B_j})}{m_\psi^2 - m_{B_k}^2 + i m_{B_k}\Gamma_k} \\ & + \tilde{C}^{\bar{B}_k M_i B_j} \frac{\Omega_L^{B_k} m_{B_k} + \Omega_R^{B_k} m_\psi}{m_\psi^2 - m_{B_k}^2 + i m_{B_k}\Gamma_k}, \end{aligned} \tag{89}$$

$$\begin{aligned} g_R^{\psi M_i B_j} \equiv\ & C_R^{\psi M_i B_j} + C^{\bar{B}_k M_i B_j} \frac{\Omega_R^{B_k}(m_\psi^2 + m_{B_k} m_{B_j}) + \Omega_L^{B_k} m_\psi (m_{B_k} + m_{B_j})}{m_\psi^2 - m_{B_k}^2 + i m_{B_k}\Gamma_k} \\ & - \tilde{C}^{\bar{B}_k M_i B_j} \frac{\Omega_R^{B_k} m_{B_k} + \Omega_L^{B_k} m_\psi}{m_\psi^2 - m_{B_k}^2 + i m_{B_k}\Gamma_k}. \end{aligned} \tag{90}$$

I Semi-leptonic Three-body Decays

In the approximation where the lepton in the final state does not interact with the di-meson system, the decay width for an individual semi-leptonic process $\psi \to \bar{\ell}_i M_1 M_2$ can be derived from the amplitude of Eq. (41):

$$\begin{aligned} \Gamma[\psi \to \bar{\ell}_i M_1 M_2] = & \frac{m_\psi^3 B_0^2}{512(1+\delta_{M_1 M_2})\pi^3 f_\chi^2} \int_{(X_1+X_2)^2}^{(1-x_i)^2} d\tilde{s} \Big[1 - \tfrac{2}{\tilde{s}}(X_1^2 + X_2^2) + \tfrac{1}{\tilde{s}^2}(X_1^2 - X_2^2)^2\Big]^{1/2} \\ & \times \Big[1 - 2(\tilde{s} + x_i^2) + (\tilde{s} - x_i^2)^2\Big]^{1/2} \sum_{\Omega,\Omega'} C[\Omega] C[\Omega']^* \, g^{\bar{\ell}_i M_1 M_2}[\Omega, \Omega'](\tilde{s}), \end{aligned} \tag{91}$$

where $x_i \equiv m_{\ell_1}/m_\psi$, $X_{1,2} \equiv m_{M_{1,2}}/m_\psi$. This defines the kinematic integral $G^{M_1 M_2}[\Omega, \Omega']$ of Eq. (43). After application of Eq. (42), $g^{\bar{\ell}_i M_1 M_2}[\Omega, \Omega'](\tilde{s})$ can be related to the reduced form-factors $\mathcal{F}_\Omega^{M_1 M_2}(\tilde{s}) \equiv \sqrt{2} f_\chi \, F_\Omega^{M_1 M_2}(m_\psi^2 \tilde{s})/B_0$ for all relevant operator pairs (Ω, Ω') (we omit the permutations that can be obtained from $g^{\bar{\ell}_i M_1 M_2}[\Omega', \Omega] = g^{\bar{\ell}_i M_1 M_2}[\Omega, \Omega']^*$):

$$\begin{aligned} g^{\bar{\ell}_i M_1 M_2}[\mathcal{S}_{ijk}^{\ell q JK}, \mathcal{S}_{i\tilde{j}\tilde{k}}^{\ell q JK'}](\tilde{s}) &\equiv \tfrac{1}{8}(1 - \tilde{s} + x_i^2)\mathcal{F}_{S\,jk}^{M_1 M_2}(\tilde{s}) \big(\mathcal{F}_{S\,\tilde{j}\tilde{k}}^{M_1 M_2}(\tilde{s})\big)^*, \\ g^{\bar{\ell}_i M_1 M_2}[\mathcal{S}_{ijk}^{\ell q JK}, \mathcal{S}_{i\tilde{j}\tilde{k}}^{\ell q J+1K'}](\tilde{s}) &\equiv \tfrac{x_i}{4}\mathcal{F}_{S\,jk}^{M_1 M_2}(\tilde{s}) \big(\mathcal{F}_{S\,\tilde{j}\tilde{k}}^{M_1 M_2}(\tilde{s})\big)^*, \\ g^{\bar{\ell}_i M_1 M_2}[\mathcal{V}_{ijk}^{\ell q JK}, \mathcal{S}_{i\tilde{j}\tilde{k}}^{\ell q JK'}](\tilde{s}) &\equiv -\tfrac{y_j - y_k}{8}\big[\tfrac{1}{\tilde{s}}(1 - x_i^2) - 1\big]\mathcal{F}_{S\,jk}^{M_1 M_2}(\tilde{s}) \big(\mathcal{F}_{S\,\tilde{j}\tilde{k}}^{M_1 M_2}(\tilde{s})\big)^*, \\ g^{\bar{\ell}_i M_1 M_2}[\mathcal{V}_{ijk}^{\ell q JK}, \mathcal{S}_{i\tilde{j}\tilde{k}}^{\ell q J+1K'}](\tilde{s}) &\equiv -\tfrac{x_i}{8}(y_j - y_k)\big[\tfrac{1}{\tilde{s}}(1 - x_i^2) + 1\big]\mathcal{F}_{S\,jk}^{M_1 M_2}(\tilde{s}) \big(\mathcal{F}_{S\,\tilde{j}\tilde{k}}^{M_1 M_2}(\tilde{s})\big)^*, \\ g^{\bar{\ell}_i M_1 M_2}[\mathcal{T}_{ijk}^{\ell q J}, \mathcal{S}_{i\tilde{j}\tilde{k}}^{\ell q J'K'}](\tilde{s}) &\equiv 0, \end{aligned} \tag{92}$$

$$
\begin{aligned}
g^{\bar{\ell}_i M_1 M_2}[\mathcal{V}^{\ell q JK}_{ijk}, \mathcal{V}^{\ell q JK'}_{i\tilde{j}\tilde{k}}](\tilde{s}) \;\equiv\; & \tfrac{1}{8\tilde{s}^2}(y_j - y_k)(y_{\tilde{j}} - y_{\tilde{k}})\big[1 - \tilde{s} - x_i^2(2+\tilde{s}) + x_i^4\big]\mathcal{F}^{M_1M_2}_{S\,jk}(\tilde{s}) \\
& \times \big(\mathcal{F}^{M_1M_2}_{S\,\tilde{j}\tilde{k}}(\tilde{s})\big)^* + \tfrac{m_\psi^2}{24}\big[1 - \tfrac{2}{\tilde{s}}(X_1^2 + X_2^2) + \tfrac{1}{\tilde{s}^2}(X_1^2 - X_2^2)^2\big] \\
& \times \big[1 + \tilde{s} - 2\tilde{s}^2 + x_i^2(\tilde{s} - 2) + x_i^4\big]\mathcal{F}^{M_1M_2}_{V\,jk}(\tilde{s})\,\big(\mathcal{F}^{M_1M_2}_{V\,\tilde{j}\tilde{k}}(\tilde{s})\big)^*, \\
g^{\bar{\ell}_i M_1 M_2}[\mathcal{V}^{\ell q JK}_{ijk}, \mathcal{V}^{\ell q J+1K'}_{i\tilde{j}\tilde{k}}](\tilde{s}) \;\equiv\; & \tfrac{x_i}{4\tilde{s}}(y_j - y_k)(y_{\tilde{j}} - y_{\tilde{k}})\mathcal{F}^{M_1M_2}_{S\,jk}(\tilde{s})\,\big(\mathcal{F}^{M_1M_2}_{S\,\tilde{j}\tilde{k}}(\tilde{s})\big)^* \\
& - \tfrac{m_\psi^2}{4}x_i\tilde{s}\big[1 - \tfrac{2}{\tilde{s}}(X_1^2 + X_2^2) + \tfrac{1}{\tilde{s}^2}(X_1^2 - X_2^2)^2\big] \\
& \mathcal{F}^{M_1M_2}_{V\,jk}(\tilde{s})\,\big(\mathcal{F}^{M_1M_2}_{V\,\tilde{j}\tilde{k}}(\tilde{s})\big)^*, \\
g^{\bar{\ell}_i M_1 M_2}[\mathcal{T}^{\ell q J}_{ijk}, \mathcal{V}^{\ell q JK'}_{i\tilde{j}\tilde{k}}](\tilde{s}) \;\equiv\; & -\tfrac{m_\psi^3}{16}\tilde{s}\big(1 - \tilde{s} - x_i^2\big)\big[1 - \tfrac{2}{\tilde{s}}(X_1^2 + X_2^2) + \tfrac{1}{\tilde{s}^2}(X_1^2 - X_2^2)^2\big] \\
& \mathcal{F}^{M_1M_2}_{T\,jk}(\tilde{s})\,\big(\mathcal{F}^{M_1M_2}_{V\,\tilde{j}\tilde{k}}(\tilde{s})\big)^*, \\
g^{\bar{\ell}_i M_1 M_2}[\mathcal{T}^{\ell q J}_{ijk}, \mathcal{V}^{\ell q J+1K'}_{i\tilde{j}\tilde{k}}](\tilde{s}) \;\equiv\; & \tfrac{m_\psi^3}{16}x_i\,\tilde{s}\big(1 + \tilde{s} - x_i^2\big)\big[1 - \tfrac{2}{\tilde{s}}(X_1^2 + X_2^2) + \tfrac{1}{\tilde{s}^2}(X_1^2 - X_2^2)^2\big] \\
& \mathcal{F}^{M_1M_2}_{T\,jk}(\tilde{s})\,\big(\mathcal{F}^{M_1M_2}_{V\,\tilde{j}\tilde{k}}(\tilde{s})\big)^*, \\
g^{\bar{\ell}_i M_1 M_2}[\mathcal{T}^{\ell q J}_{ijk}, \mathcal{T}^{\ell q J}_{i\tilde{j}\tilde{k}}](\tilde{s}) \;\equiv\; & \tfrac{m_\psi^4}{96}\tilde{s}\big[2 - \tilde{s} - \tilde{s}^2 - x_i^2(4+\tilde{s}) + 2x_i^4\big] \\
& \times \big[1 - \tfrac{2}{\tilde{s}}(X_1^2 + X_2^2) + \tfrac{1}{\tilde{s}^2}(X_1^2 - X_2^2)^2\big]\mathcal{F}^{M_1M_2}_{T\,jk}(\tilde{s})\,\big(\mathcal{F}^{M_1M_2}_{T\,\tilde{j}\tilde{k}}(\tilde{s})\big)^*, \\
g^{\bar{\ell}_i M_1 M_2}[\mathcal{T}^{\ell q J}_{ijk}, \mathcal{T}^{\ell q J+1}_{i\tilde{j}\tilde{k}}](\tilde{s}) \;\equiv\; & -\tfrac{m_\psi^4}{16}x_i\,\tilde{s}^2\big[1 - \tfrac{2}{\tilde{s}}(X_1^2 + X_2^2) + \tfrac{1}{\tilde{s}^2}(X_1^2 - X_2^2)^2\big] \\
& \mathcal{F}^{M_1M_2}_{T\,jk}(\tilde{s})\,\big(\mathcal{F}^{M_1M_2}_{T\,\tilde{j}\tilde{k}}(\tilde{s})\big)^*,
\end{aligned}
$$

with $y_{j,k,\tilde{j},\tilde{k}} \equiv m_{q_{j,k,\tilde{j},\tilde{k}}}/m_\psi$; the chiralities J, J', K, K' take the values L, R, with the convention that $J+1 \neq J$. Setting baryon, meson, quark and lepton masses to their numerical values, the kinematic integrals $G^{\bar{\ell}_i M_m M_n}[\Omega, \Omega']$ can be seen as functions of the sole mass m_ψ.

J Partonic Decay Widths

For hadronic or semi-leptonic processes with a large phase space (*i.e.* far above thresholds), the partonic descriptions regains its legitimacy. We neglect quark and lepton masses below.

For the three hadronic decay channels $(i,j,k) \in \{(1,1,1),(1,1,2),(1,2,2)\}$ (and their CP-conjugates):

$$
\begin{aligned}
\Gamma[\psi \to \bar{u}_i \bar{d}_j \bar{d}_k] = & \frac{m_\psi^5}{1024(1+\delta_{jk})\pi^3}\Big\{\big|C[\mathcal{H}^{JK}_{ijk}]\big|^2 + \big|C[\widetilde{\mathcal{H}}^{JK}_{ijk}]\big|^2 + \big|C[\widetilde{\mathcal{H}}^{JK}_{ikj}]\big|^2 \\
& - \mathrm{Re}\big[C[\mathcal{H}^{JJ}_{ijk}]C[\widetilde{\mathcal{H}}^{JJ}_{ijk}]^*\big] - \mathrm{Re}\big[C[\mathcal{H}^{JJ}_{ijk}]C[\widetilde{\mathcal{H}}^{JJ}_{ikk}]^*\big] + \mathrm{Re}\big[C[\widetilde{\mathcal{H}}^{JJ}_{ijk}]C[\widetilde{\mathcal{H}}^{JJ}_{ikj}]^*\big]\Big\}.
\end{aligned}
\tag{93}
$$

For the semi-leptonic decay channels (and their CP-conjugates):

$$
\Gamma[\psi \to \bar{\ell}_i q_j \bar{q}_k] = \frac{m_\psi^5}{2048\pi^3}\Big\{\big|C[\mathcal{S}^{\ell q JK}_{ijk}]\big|^2 + 4\big|C[\mathcal{V}^{\ell q JK}_{ijk}]\big|^2 + 3\big|C[\mathcal{T}^{\ell q J}_{ijk}]\big|^2\Big\},
\tag{94}
$$

$$
(\ell_i, q_j, q_k) \in \{(\nu_i,u,u),(\nu_i,d,d),(\nu_i,s,s),(\nu_i,d,s),(\nu_i,s,d),(e_i,d,u),(e_i,s,u)\}.
$$

K Production from the Decays of SM Particles

Production in the decays of electroweak gauge bosons

Single ($\chi_i^0 \stackrel{!}{=} \psi$, $\chi_j^0 \stackrel{!}{=} \nu_f, \bar{\nu}_f$) or pair production ($\chi_i^0 \stackrel{!}{=} \psi \stackrel{!}{=} \chi_j^0$) can occur in decays of the Z-boson, relying on subleading neutralino components:

$$\Gamma[Z \to \chi_i^0 \chi_j^0] = \frac{M_Z}{24(1+\delta_{ij})\pi}\left[1 - 2\frac{m^2_{\chi_i^0}+m^2_{\chi_j^0}}{M_Z^2} + \frac{(m^2_{\chi_i^0}+m^2_{\chi_j^0})^2}{M_Z^4}\right]^{1/2} \tag{95}$$
$$\times \left\{\left(|g_L^{Z\chi_i^0\chi_j^0}|^2 + |g_R^{Z\chi_i^0\chi_j^0}|^2\right)\left[1 - \frac{m^2_{\chi_i^0}+m^2_{\chi_j^0}}{2M_Z^2} - \frac{(m^2_{\chi_i^0}+m^2_{\chi_j^0})^2}{2M_Z^4}\right] + 6\frac{m_{\chi_i^0} m_{\chi_j^0}}{M_Z^2}\mathrm{Re}\left[g_L^{Z\chi_i^0\chi_j^0}(g_R^{Z\chi_i^0\chi_j^0})^*\right]\right\}.$$

Single production in the decays of a W-boson:

$$\Gamma[W \to \psi e_i] = \frac{M_W}{24\pi}\left[1 - 2\frac{m^2_\psi+m^2_{e_i}}{M_W^2} + \frac{(m^2_\psi+m^2_{e_i})^2}{M_W^4}\right]^{1/2} \tag{96}$$
$$\times \left\{\left(|g_L^{W\psi e_i}|^2 + |g_R^{W\psi e_i}|^2\right)\left[1 - \frac{m^2_\psi+m^2_{e_i}}{2M_W^2} - \frac{(m^2_\psi+m^2_{e_i})^2}{2M_W^4}\right] + 6\frac{m_\psi m_{e_i}}{M_W^2}\mathrm{Re}\left[g_L^{W\psi e_i\chi_j^0}(g_R^{W\psi e_i})^*\right]\right\}.$$

Production in the decays of a neutral scalar (Higgs)

Single ($\chi_i^0 \stackrel{!}{=} \psi$, $\chi_j^0 \stackrel{!}{=} \nu_f, \bar{\nu}_f$) or pair production ($\chi_i^0 \stackrel{!}{=} \psi \stackrel{!}{=} \chi_j^0$) can occur in decays of the Higgs boson, relying on subleading neutralino components or sneutrino-Higgs mixing:

$$\Gamma[S_p^0 \to \chi_i^0 \chi_j^0] = \frac{m_{S_p^0}}{16(1+\delta_{ij})\pi}\left[1 - 2\frac{m^2_{\chi_i^0}+m^2_{\chi_j^0}}{m^2_{S_p^0}} + \frac{(m^2_{\chi_i^0}+m^2_{\chi_j^0})^2}{m^4_{S_p^0}}\right]^{1/2} \tag{97}$$
$$\times \left\{\left(|g_L^{S_p^0\chi_i^0\chi_j^0}|^2 + |g_R^{S_p^0\chi_i^0\chi_j^0}|^2\right)\left[1 - \frac{m^2_{\chi_i^0}+m^2_{\chi_j^0}}{m^2_{S_p^0}}\right] - 4\frac{m_{\chi_i^0} m_{\chi_j^0}}{m^2_{S_p^0}}\mathrm{Re}\left[g_L^{S_p^0\chi_i^0\chi_j^0}(g_R^{S_p^0\chi_i^0\chi_j^0})^*\right]\right\}.$$

Production in leptonic transitions

We collect the operators of Eq. (4), last line, and Eq. (44), $f \stackrel{!}{=} e$, in a single basis, allowing to describe single and pair production in leptonic transitions simultaneously ($i \geqslant j$, $n \geqslant m$):

$$\begin{aligned}
\mathcal{S}^{KJ}_{ijmn} &\equiv (\bar{X}_i^0 P_K X_j^0)(\bar{E}_m P_J E_n)\,, \qquad \mathcal{V}^{KJ}_{ijmn} \equiv (\bar{X}_i^0 \gamma^\mu P_K X_j^0)(\bar{E}_m \gamma_\mu P_J E_n)\,, \\
\mathcal{T}^{K}_{ijmn} &\equiv \frac{1}{4}(\bar{X}_i^0 \Sigma^{\mu\nu} P_K X_j^0)(\bar{E}_m \Sigma_{\mu\nu} P_K E_n)\,.
\end{aligned} \tag{98}$$

Then, the three-body decay of a charged lepton e_n into the lighter e_m and a light neutralino pair χ_i^0, χ_j^0 can be written as the following sum (with mass indices straightforwardly referring to the particle of corresponding index):

$$\Gamma[e_n \to e_m \chi_i^0 \chi_j^0] = \frac{1+\delta_{ij}}{512\pi m_n}\int_{(m_i+m_j)^2}^{(m_n-m_m)^2} ds \left[1 - 2\frac{m_i^2+m_j^2}{s} + \frac{(m_i^2-m_j^2)^2}{s^2}\right]\left[1 - 2\frac{s+m_m^2}{m_n^2} + \frac{(s-m_m^2)^2}{m_n^4}\right]$$
$$\times \sum_{\Omega,\Omega'} C[\Omega] C[\Omega']^*\, g^{e_m e_n \chi_i^0 \chi_j^0}[\Omega,\Omega'](s)\,, \tag{99}$$

where

$$
\begin{aligned}
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{S}_{ijmn}^{KJ}, \mathcal{S}_{ijmn}^{KJ}](s) &\equiv (s - m_i^2 - m_j^2)(m_n^2 + m_m^2 - s)\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{S}_{ijmn}^{KJ}, \mathcal{S}_{ijmn}^{K+1J}](s) &\equiv -2m_i m_j (m_n^2 + m_m^2 - s)\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{S}_{ijmn}^{KJ}, \mathcal{S}_{ijmn}^{KJ+1}](s) &\equiv 2m_n m_m (s - m_i^2 - m_j^2)\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{S}_{ijmn}^{KJ}, \mathcal{S}_{ijmn}^{K+1J+1}](s) &\equiv -4m_i m_j m_n m_m\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{V}_{ijmn}^{KJ}, \mathcal{S}_{ijmn}^{KJ}](s) &\equiv \tfrac{m_m m_i}{s}(s - m_i^2 - m_j^2)(m_n^2 - m_m^2 + s)\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{V}_{ijmn}^{KJ}, \mathcal{S}_{ijmn}^{K+1J}](s) &\equiv -\tfrac{m_m m_j}{s}(s + m_i^2 - m_j^2)(m_n^2 - m_m^2 + s)\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{V}_{ijmn}^{KJ}, \mathcal{S}_{ijmn}^{KJ+1}](s) &\equiv \tfrac{m_n m_i}{s}(s - m_i^2 + m_j^2)(m_n^2 - m_m^2 - s)\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{V}_{ijmn}^{KJ}, \mathcal{S}_{ijmn}^{K+1J+1}](s) &\equiv -\tfrac{m_n m_j}{s}(s + m_i^2 - m_j^2)(m_n^2 - m_m^2 - s)\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{T}_{ijmn}^{K}, \mathcal{S}_{ijmn}^{K'J}](s) &\equiv 0\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{V}_{ijmn}^{KJ}, \mathcal{V}_{ijmn}^{KJ}](s) &\equiv \big[(m_n^2 - m_m^2)^2 - s^2\big]\big[1 - \tfrac{(m_i^2 - m_j^2)^2}{s^2}\big] \\
&\quad -\tfrac{1}{3}\big[(m_n^2 - m_m^2)^2 - 2s(m_n^2 + m_m^2) + s^2\big] \\
&\quad \big[1 - 2\tfrac{m_i^2 + m_j^2}{s} + \tfrac{(m_i^2 - m_j^2)^2}{s^2}\big]\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{V}_{ijmn}^{KJ}, \mathcal{V}_{ijmn}^{K+1J}](s) &\equiv 4m_i m_j (m_n^2 + m_m^2 - s)\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{V}_{ijmn}^{KJ}, \mathcal{V}_{ijmn}^{KJ+1}](s) &\equiv -4m_n m_m (s - m_i^2 - m_j^2)\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{V}_{ijmn}^{KJ}, \mathcal{V}_{ijmn}^{K+1J+1}](s) &\equiv -16m_i m_j m_n m_m\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{T}_{ijmn}^{K}, \mathcal{V}_{ijmn}^{KK}](s) &\equiv -3\tfrac{m_i m_m}{s}(s - m_i^2 + m_j^2)(m_n^2 - m_m^2 + s)\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{T}_{ijmn}^{K}, \mathcal{V}_{ijmn}^{K+1K}](s) &\equiv -3\tfrac{m_j m_m}{s}(s + m_i^2 - m_j^2)(m_n^2 - m_m^2 + s)\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{T}_{ijmn}^{K}, \mathcal{V}_{ijmn}^{KK+1}](s) &\equiv 3\tfrac{m_i m_n}{s}(s - m_i^2 + m_j^2)(m_n^2 - m_m^2 - s)\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{T}_{ijmn}^{K}, \mathcal{V}_{ijmn}^{K+1K+1}](s) &\equiv 3\tfrac{m_j m_n}{s}(s + m_i^2 - m_j^2)(m_n^2 - m_m^2 - s)\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{T}_{ijmn}^{K}, \mathcal{T}_{ijmn}^{K}](s) &\equiv \tfrac{1}{3}\big[2(m_n^2 - m_m^2)^2 - s(m_n^2 + m_m^2) - s^2\big] \\
&\quad \big[1 + \tfrac{m_i^2 + m_j^2}{s} - 2\tfrac{(m_i^2 - m_j^2)^2}{s^2}\big]\,, \\
g^{e_m e_n \chi_i^0 \chi_j^0}[\mathcal{T}_{ijmn}^{K}, \mathcal{T}_{ijmn}^{K+1}](s) &\equiv -12m_i m_j m_n m_m\,.
\end{aligned}
$$

Production in hadronic tau decays

Two-body decays involving a pseudoscalar meson:

$$
\Gamma[\tau \to \psi M] = \frac{m_\tau}{32\pi}\Big\{\big(1 + \tfrac{m_\psi^2 - m_M^2}{m_\tau^2}\big)\big(|g_L^{\psi \bar{M} \tau}|^2 + |g_R^{\psi \bar{M} \tau}|^2\big) + 4\tfrac{m_\psi}{m_\tau}\mathrm{Re}\big[(g_L^{\psi \bar{M} \tau})^* g_R^{\psi \bar{M} \tau}\big]\Big\} \times \Big[1 - 2\tfrac{m_M^2 + m_\psi^2}{m_\tau^2} + \tfrac{(m_M^2 - m_\psi^2)^2}{m_\tau^4}\Big]^{1/2}, \quad (100)
$$

with the couplings defined in Eq. (36).

For a light ψ, we expect the three-body decay channel, involving a meson pair, to be com-

petitive:

$$\Gamma[\tau\to\psi M_1M_2]=\frac{m_\tau^3B_0^2}{512\pi^3f_\chi^2}\int_{(X_1+X_2)^2}^{(1-m_\psi/m_\tau)^2}d\tilde{s}\left[1-\tfrac{2}{\tilde{s}}(X_1^2+X_2^2)+\tfrac{1}{\tilde{s}^2}(X_1^2-X_2^2)^2\right]^{1/2}$$
$$\times\left[1-2(\tilde{s}+\tfrac{m_\psi^2}{m_\tau^2})+(\tilde{s}-\tfrac{m_\psi^2}{m_\tau^2})^2\right]^{1/2}\sum_{\Omega,\Omega'}C[\Omega]C[\Omega']^*\,g^{\psi M_1M_2}[\Omega,\Omega'](\tilde{s})\,, \quad (101)$$

where the relation $m_\tau^2 g^{\psi M_1M_2}[\Omega,\Omega'](\frac{s}{m_\tau^2})=m_\psi^2 g^{\bar{\ell}_3M_1M_2}[\Omega,\Omega'](\frac{s}{m_\psi^2})$ determines the functions in the integrand in terms of their counterparts of Eq. (92).

Production in baryon-number violating transitions

A baryon B_j may decay into a pseudoscalar meson $\bar{M}_i$ and a light neutralino according to:

$$\Gamma[B_j\to\psi\bar{M}_i]=\frac{m_{B_j}}{32\pi}\left\{\left(1+\tfrac{m_\psi^2-m_{M_i}^2}{m_{B_j}^2}\right)\left[\left|g_L^{\psi M_iB_j}\right|^2+\left|g_R^{\psi M_iB_j}\right|^2\right]+4\tfrac{m_\psi}{m_{B_j}}\mathrm{Re}\left[\left(g_L^{\psi M_iB_j}\right)^*g_R^{\psi M_iB_j}\right]\right\}$$
$$\times\left(1-2\tfrac{m_\psi^2+m_{M_i}^2}{m_{B_j}^2}+\tfrac{(m_\psi^2-m_{M_i}^2)^2}{m_{B_j}^4}\right)^{1/2}, \quad (102)$$

with the couplings of Eq. (32). In the case of a nucleon, $B_j\stackrel{!}{=}p^+,n^0$, with a light neutralino, $m_\psi\lesssim m_\mu$, the form factors computed on the lattice [56] may offer a quantitatively more accurate description of these transitions. Higher orders in the chiral perturbation theory (*e.g.* encoding the electroweak charged interactions) can be added in order to account for *e.g.* the $s\to d$ transition.

For a heavy quark, *i.e.* with a large available phase space for the decay products, the partonic description may offer a good approximation of the decays of a b-flavored hadron ($\bar{B}$, Λ_b^0, etc.) into a light ψ and light hadrons:

$$\begin{aligned}\Gamma[b\to\psi\bar{u}\bar{d}_j] &= \frac{m_b^5}{1024\pi^3}\Big\{\Big(\left|C[\mathcal{H}^{JK}_{1j3}]\right|^2+\left|C[\widetilde{\mathcal{H}}^{JK}_{1j3}]\right|^2+\left|C[\widetilde{\mathcal{H}}^{JK}_{13j}]\right|^2\\ &\quad+\mathrm{Re}\big[C[\widetilde{\mathcal{H}}^{JJ}_{1j3}]C[\widetilde{\mathcal{H}}^{JJ}_{13j}]^*-C[\mathcal{H}^{JJ}_{1j3}]\big(C[\widetilde{\mathcal{H}}^{JJ}_{1j3}]^*+C[\widetilde{\mathcal{H}}^{JJ}_{13j}]^*\big)\big]\Big)\mathcal{J}\big(\tfrac{m_\psi^2}{m_b^2}\big)\\ &\quad+4\frac{m_\psi}{m_b}\mathrm{Re}\Big[C[\widetilde{\mathcal{H}}^{JK}_{1j3}]C[\widetilde{\mathcal{H}}^{J+1K}_{13j}]^*-C[\mathcal{H}^{JJ+1}_{1j3}]C[\widetilde{\mathcal{H}}^{J+1J}_{1j3}]^*\\ &\quad-C[\mathcal{H}^{JJ}_{1j3}]C[\widetilde{\mathcal{H}}^{J+1J}_{13j}]^*+C[\widetilde{\mathcal{H}}^{JJ}_{1j3}]C[\widetilde{\mathcal{H}}^{J+1J}_{13j}]^*\Big]\mathcal{K}\big(\tfrac{m_\psi^2}{m_b^2}\big)\Big\}\,, \quad (103)\\ \mathcal{J}(x) &\equiv 1-8x+8x^3-x^4-12x^2\ln x\,,\\ \mathcal{K}(x) &\equiv 1+9x-9x^2-x^3+6x(1+x)\ln x\,.\end{aligned}$$

The same approach is evidently applicable to a top quark, provided higher-order operators do not compete with those of dimension 6 that we consider here (*i.e.* provided the scalar mediators are much heavier than the top quark in the RpV MSSM):

$$\begin{aligned}\Gamma[t\to\psi\bar{d}_j\bar{d}_k] &= \frac{m_t^5}{1024\pi^3}\Big\{\Big(\left|C[\mathcal{H}^{JK}_{3jk}]\right|^2+\left|C[\widetilde{\mathcal{H}}^{JK}_{3jk}]\right|^2+\left|C[\widetilde{\mathcal{H}}^{JK}_{3kj}]\right|^2\\ &\quad+\mathrm{Re}\big[C[\widetilde{\mathcal{H}}^{JJ}_{3jk}]C[\widetilde{\mathcal{H}}^{JJ}_{3kj}]^*-C[\mathcal{H}^{JJ}_{3jk}]\big(C[\widetilde{\mathcal{H}}^{JJ}_{3jk}]^*+C[\widetilde{\mathcal{H}}^{JJ}_{3kj}]^*\big)\big]\Big)\mathcal{J}\big(\tfrac{m_\psi^2}{m_b^2}\big)\\ &\quad+4\frac{m_\psi}{m_t}\mathrm{Re}\Big[C[\mathcal{H}^{JK}_{3jk}]C[\mathcal{H}^{J+1K}_{3jk}]^*-C[\mathcal{H}^{JJ}_{3jk}]C[\widetilde{\mathcal{H}}^{J+1J}_{3jk}]^*\\ &\quad-C[\mathcal{H}^{JJ}_{3jk}]C[\widetilde{\mathcal{H}}^{J+1J}_{3kj}]^*+C[\widetilde{\mathcal{H}}^{JJ+1}_{3jk}]C[\widetilde{\mathcal{H}}^{J+1J}_{3kj}]^*\Big]\mathcal{K}\big(\tfrac{m_\psi^2}{m_b^2}\big)\Big\}\,. \quad (104)\end{aligned}$$

A charm quark is likely too close to the threshold for baryon production for such an approximation to be meaningful in descriptions of the decays of D-mesons.

Production in leptonic and semi-leptonic meson decays

Two-body decays of a pseudoscalar meson $\bar{M}$ into ψ and a (anti)lepton $\bar{\ell}_i$:

$$\begin{aligned}\Gamma[\bar{M}\to\psi\bar{\ell}_i] &= \frac{m_M}{16\pi}\Big\{\big(1-\tfrac{m_{\ell_i}^2+m_\psi^2}{m_M^2}\big)\big(|g_L^{\psi\bar{M}\ell_i}|^2+|g_R^{\psi\bar{M}\ell_i}|^2\big)-4\tfrac{m_{\ell_i}m_\psi}{m_M^2}\mathrm{Re}\big[(g_L^{\psi\bar{M}\ell_i})^* g_R^{\psi\bar{M}\ell_i}\big]\Big\}\\ &\quad\times\sqrt{1-2\tfrac{m_{\ell_i}^2+m_\psi^2}{m_M^2}+\tfrac{(m_{\ell_i}^2-m_\psi^2)^2}{m_M^4}}\,,\end{aligned}\tag{105}$$

with the couplings defined in Eq. (36) (and possibly extending the operator basis to include third-generation fermions). For a neutral meson, pair production of ψ is obtained from a similar formula, up to a factor 2 (for identical particles) and the Wilson coefficients of the operators of Eq. (44) replacing those of Eq. (5) in the definition of the couplings: $C[\mathcal{S}_{ijk}^{\ell q\,LJ}]\to C[\dot{\mathcal{S}}_{jk}^{qJ}]$, $C[\mathcal{S}_{ijk}^{\ell q\,RJ}]\to C[\dot{\mathcal{S}}_{kj}^{qJ+1}]^*$, $C[\mathcal{V}_{ijk}^{\ell q\,LJ}]\to C[\dot{\mathcal{V}}_{jk}^{qJ}]$, $C[\mathcal{V}_{ijk}^{\ell q\,RJ}]\to -C[\dot{\mathcal{V}}_{kj}^{qJ}]^*$, $C[\mathcal{T}_{ijk}^{\ell q\,J}]\to 0$.

Similar decays are also accessible to a vector meson M^* with the couplings defined in Eq. (40):

$$\begin{aligned}\Gamma[M^*\to\psi\bar{\ell}_i] &= \frac{m_{M^*}}{48\pi}\Big[1-2\tfrac{m_\psi^2+m_{\ell_i}^2}{m_{M^*}^2}+\tfrac{(m_\psi^2-m_{\ell_i}^2)^2}{m_{M^*}^4}\Big]^{1/2}\\ &\times\Bigg\{\Big(|g_{V\,L}^{M^*\psi\ell_i}|^2+|g_{V\,R}^{M^*\psi\ell_i}|^2\Big)\Big[2-\tfrac{m_\psi^2+m_{\ell_i}^2}{m_{M^*}^2}-\tfrac{(m_\psi^2-m_{\ell_i}^2)^2}{m_{M^*}^4}\Big]+6\frac{m_{\ell_i}m_\psi}{m_{M^*}^2}\mathrm{Re}\big[g_{V\,L}^{M^*\psi\ell_i\,*}g_{V\,R}^{M^*\psi\ell_i}\big]\\ &+\Big[1-2\tfrac{m_\psi^2+m_{\ell_i}^2}{m_{M^*}^2}+\tfrac{(m_\psi^2-m_{\ell_i}^2)^2}{m_{M^*}^4}\Big]\Big[m_{M^*}^2\big(|g_{S\,L}^{M^*\psi\ell_i}|^2+|g_{S\,R}^{M^*\psi\ell_i}|^2\big)\big(1+\tfrac{m_\psi^2-m_{\ell_i}^2}{m_{M^*}^2}\big)\\ &-4\,m_{\ell_i}m_\psi\,\mathrm{Re}\big[g_{S\,L}^{M^*\psi\ell_i\,*}g_{S\,R}^{M^*\psi\ell_i}\big]+2m_{\ell_i}\mathrm{Re}\big[g_{V\,L}^{M^*\psi\ell_i\,*}g_{S\,L}^{M^*\psi\ell_i}+g_{V\,R}^{M^*\psi\ell_i\,*}g_{S\,R}^{M^*\psi\ell_i}\big]\\ &+2m_\psi\mathrm{Re}\big[g_{V\,L}^{M^*\psi\ell_i\,*}g_{S\,R}^{M^*\psi\ell_i}+g_{V\,R}^{M^*\psi\ell_i\,*}g_{S\,L}^{M^*\psi\ell_i}\big]\Big]\Bigg\}.\end{aligned}\tag{106}$$

This expression is again straightforwardly extended to pair production through the substitution of the Wilson coefficients and the application of a factor 2.

As soon as enough phase space is accessible, semi-leptonic decays become competitive. We restrict ourselves to a three-body decay involving pseudoscalar mesons in the initial and final states. Then, in analogy to Eq. (42), one can introduce the form-factors for a $M_1\to M_2$ transition ($s\equiv(p_1-p_2)^2$):

$$\begin{aligned}\big\langle M_2(p_2)|\,\bar{Q}_jQ_k\,|M_1(p_1)\big\rangle &\equiv \overline{F}_{S\,jk}^{M_1M_2}(s)\,,\\ \big\langle M_2(p_2)|\,\bar{Q}_j\gamma^\mu Q_k\,|M_1(p_1)\big\rangle &\equiv -(p_1-p_2)^\mu\tfrac{m_j-m_k}{s}\overline{F}_{S\,jk}^{M_1M_2}(s)\\ &\quad-\big[p_1^\mu+p_2^\mu-\tfrac{m_1^2-m_2^2}{s}(p_1-p_2)^\mu\big]\overline{F}_{V\,jk}^{M_1M_2}(s)\,,\\ \big\langle M_2(p_2)|\,\bar{Q}_j\Sigma^{\mu\nu}Q_k\,|M_1(p_1)\big\rangle &\equiv -i(p_1^\mu p_2^\nu-p_2^\mu p_1^\nu)\overline{F}_{T\,jk}^{M_1M_2}(s)\,.\end{aligned}\tag{107}$$

The decay width for $M_1\to M_2\psi\bar{\ell}_i$ then reads ($m_{1,2}\equiv m_{M_{1,2}}$):

$$\begin{aligned}\Gamma[M_1\to M_2\psi\bar{\ell}_i] &= \frac{1}{256\pi^3m_1}\int_{(m_\psi+m_{\ell_i})^2}^{(m_1-m_2)^2}ds\Big[1-2\tfrac{m_\psi^2+m_{\ell_i}^2}{s}+\tfrac{(m_1^2-m_2^2)^2}{s^2}\Big]^{1/2}\\ &\times\Big[1-2(\tfrac{s+m_2^2}{m_1^2})+\tfrac{(s-m_2^2)^2}{m_1^4}\Big]^{1/2}\sum_{\Omega,\Omega'}C[\Omega]C[\Omega']^*\,g^{M_2\psi\bar{\ell}_i}[\Omega,\Omega'](s)\,,\end{aligned}\tag{108}$$

where

$$
\begin{aligned}
g^{M_2\psi\bar{\ell}_i}[\mathcal{S}^{\ell qJK}_{ijk},\mathcal{S}^{\ell qJK'}_{i\tilde{j}\tilde{k}}](s) &\equiv \tfrac{1}{4}(s-m_\psi^2-m_{\ell_i}^2)\overline{F}^{M_1M_2}_{S\,jk}(s)\left(\overline{F}^{M_1M_2}_{S\,\tilde{j}\tilde{k}}(s)\right)^*,\\
g^{M_2\psi\bar{\ell}_i}[\mathcal{S}^{\ell qJK}_{ijk},\mathcal{S}^{\ell qJ+1K'}_{i\tilde{j}\tilde{k}}](s) &\equiv -\tfrac{m_\psi m_{\ell_i}}{2}\overline{F}^{M_1M_2}_{S\,jk}(s)\left(\overline{F}^{M_1M_2}_{S\,\tilde{j}\tilde{k}}(s)\right)^*,\\
g^{M_2\psi\bar{\ell}_i}[\mathcal{V}^{\ell qJK}_{ijk},\mathcal{S}^{\ell qJK'}_{i\tilde{j}\tilde{k}}](s) &\equiv -\tfrac{m_j-m_k}{4}\left[1-\tfrac{m_\psi^2-m_{\ell_i}^2}{s}\right]\overline{F}^{M_1M_2}_{S\,jk}(s)\left(\overline{F}^{M_1M_2}_{S\,\tilde{j}\tilde{k}}(s)\right)^*,\\
g^{M_2\psi\bar{\ell}_i}[\mathcal{V}^{\ell qJK}_{ijk},\mathcal{S}^{\ell qJ+1K'}_{i\tilde{j}\tilde{k}}](s) &\equiv \tfrac{m_{\ell_i}}{4}(m_j-m_k)\left[1+\tfrac{m_\psi^2-m_{\ell_i}^2}{s}\right]\overline{F}^{M_1M_2}_{S\,jk}(s)\left(\overline{F}^{M_1M_2}_{S\,\tilde{j}\tilde{k}}(s)\right)^*,\\
g^{M_2\psi\bar{\ell}_i}[\mathcal{T}^{\ell qJ}_{ijk},\mathcal{S}^{\ell qJ'K'}_{i\tilde{j}\tilde{k}}](s) &\equiv 0,\\
g^{M_2\psi\bar{\ell}_i}[\mathcal{V}^{\ell qJK}_{ijk},\mathcal{V}^{\ell qJK'}_{i\tilde{j}\tilde{k}}](s) &\equiv \tfrac{1}{4}(m_j-m_k)(m_{\tilde{j}}-m_{\tilde{k}})\left[\tfrac{m_\psi^2+m_{\ell_i}^2}{s}-\tfrac{(m_\psi^2-m_{\ell_i}^2)^2}{s^2}\right]\\
&\quad \left(\overline{F}^{M_1M_2}_{S\,\tilde{j}\tilde{k}}(s)\right)^*\overline{F}^{M_1M_2}_{S\,jk}(s)\\
&\quad +\tfrac{m_1^4}{12}\left[1-2\tfrac{s+m_2^2}{m_1^2}+\tfrac{(s-m_2^2)^2}{m_1^4}\right]\left[2-\tfrac{m_\psi^2+m_{\ell_i}^2}{s}-\tfrac{(m_\psi^2-m_{\ell_i}^2)^2}{s^2}\right]\\
&\quad \overline{F}^{M_1M_2}_{V\,jk}(s)\left(\overline{F}^{M_1M_2}_{V\,\tilde{j}\tilde{k}}(s)\right)^*,\\
g^{M_2\psi\bar{\ell}_i}[\mathcal{V}^{\ell qJK}_{ijk},\mathcal{V}^{\ell qJ+1K'}_{i\tilde{j}\tilde{k}}](s) &\equiv -\tfrac{m_{\ell_i}m_\psi}{2s}(m_j-m_k)(m_{\tilde{j}}-m_{\tilde{k}})\overline{F}^{M_1M_2}_{S\,jk}(s)\left(\overline{F}^{M_1M_2}_{S\,\tilde{j}\tilde{k}}(s)\right)^*\\
&\quad +\tfrac{m_{\ell_i}m_\psi}{2s}m_1^4\left[1-2\tfrac{s+m_2^2}{m_1^2}+\tfrac{(s-m_2^2)^2}{m_1^4}\right]\overline{F}^{M_1M_2}_{V\,jk}(s)\left(\overline{F}^{M_1M_2}_{V\,\tilde{j}\tilde{k}}(s)\right)^*,\\
g^{M_2\psi\bar{\ell}_i}[\mathcal{T}^{\ell qJ}_{ijk},\mathcal{V}^{\ell qJK'}_{i\tilde{j}\tilde{k}}](s) &\equiv -\tfrac{m_\psi}{8}m_1^4\left(1-\tfrac{m_\psi^2-m_{\ell_i}^2}{s}\right)\left[1-\tfrac{s+2m_2^2}{m_1^2}+\tfrac{m_2^2(3s-m_2^2)}{m_1^4}\right]\\
&\quad \overline{F}^{M_1M_2}_{T\,jk}(s)\left(\overline{F}^{M_1M_2}_{V\,\tilde{j}\tilde{k}}(s)\right)^*,\\
g^{M_2\psi\bar{\ell}_i}[\mathcal{T}^{\ell qJ}_{ijk},\mathcal{V}^{\ell qJ+1K'}_{i\tilde{j}\tilde{k}}](s) &\equiv -\tfrac{m_{\ell_i}}{8}m_1^4\left(1+\tfrac{m_\psi^2-m_{\ell_i}^2}{s}\right)\left[1-\tfrac{s+2m_2^2}{m_1^2}+\tfrac{m_2^2(3s-m_2^2)}{m_1^4}\right]\\
&\quad \overline{F}^{M_1M_2}_{T\,jk}(s)\left(\overline{F}^{M_1M_2}_{V\,\tilde{j}\tilde{k}}(s)\right)^*,\\
g^{M_2\psi\bar{\ell}_i}[\mathcal{T}^{\ell qJ}_{ijk},\mathcal{T}^{\ell qJ}_{i\tilde{j}\tilde{k}}](s) &\equiv \tfrac{m_1^4}{48}s\left[1-2\tfrac{s+m_2^2}{m_1^2}+\tfrac{(s-m_2^2)^2}{m_1^4}\right]\left[1+\tfrac{m_\psi^2+m_{\ell_i}^2}{s}-2\tfrac{(m_\psi^2-m_{\ell_i}^2)^2}{s^2}\right]\\
&\quad \overline{F}^{M_1M_2}_{T\,jk}(s)\left(\overline{F}^{M_1M_2}_{T\,\tilde{j}\tilde{k}}(s)\right)^*,\\
g^{M_2\psi\bar{\ell}_i}[\mathcal{T}^{\ell qJ}_{ijk},\mathcal{T}^{\ell qJ+1}_{i\tilde{j}\tilde{k}}](s) &\equiv \tfrac{m_{\ell_i}m_\psi}{8}m_1^4\left[1-2\tfrac{s+m_2^2}{m_1^2}+\tfrac{(s-m_2^2)^2}{m_1^4}\right]\overline{F}^{M_1M_2}_{T\,jk}(s)\left(\overline{F}^{M_1M_2}_{T\,\tilde{j}\tilde{k}}(s)\right)^*,
\end{aligned}
$$

where the sum runs over all semi-leptonic operators. As before, the case of ψ-pair production can be addressed by substituting ℓ_i by ψ in the expression above, hence employing the Wilson coefficients of the operators of Eq. (44), and introducing a factor 2 for identical particles in the final state.

When substantial phase space is accessible in the final state, *i.e.* for hadrons involving a heavy quark, the partonic description may again prove a useful approximation. Neglecting the masses of the light leptons and quarks, we can write the relevant decay widths for a bottom

quark:

$$\Gamma[b\to\psi\bar{\ell}_i q_j] = \frac{m_b^5}{2048\pi^3}\left\{\left(\left|C[\mathcal{S}_{ij3}^{\ell qJK}]\right|^2+4\left|C[\mathcal{V}_{ij3}^{\ell qJK}]\right|^2+3\left|C[\mathcal{T}_{ij3}^{\ell qJ}]\right|^2\right)\mathcal{J}\left(\tfrac{m_\psi^2}{m_b^2}\right)\right. \tag{109}$$
$$\left.+\,8\frac{m_\psi}{m_b}\mathrm{Re}\left[C[\mathcal{V}_{ij3}^{\ell qJK}]C[\mathcal{S}_{i3j}^{\ell qJK+1}]^*+3C[\mathcal{T}_{ij3}^{\ell qJ}]C[\mathcal{V}_{ij3}^{\ell qJJ+1}]^*\right]\mathcal{K}\left(\tfrac{m_\psi^2}{m_b^2}\right)\right\},$$
$$\Gamma[\bar{b}\to\psi\bar{\ell}_i\bar{q}_j] = \frac{m_b^5}{2048\pi^3}\left\{\left(\left|C[\mathcal{S}_{i3j}^{\ell qJK}]\right|^2+4\left|C[\mathcal{V}_{i3j}^{\ell qJK}]\right|^2+3\left|C[\mathcal{T}_{i3j}^{\ell qJ}]\right|^2\right)\mathcal{J}\left(\tfrac{m_\psi^2}{m_b^2}\right)\right.$$
$$\left.-\,8\frac{m_\psi}{m_b}\mathrm{Re}\left[C[\mathcal{V}_{i3j}^{\ell qJK}]C[\mathcal{S}_{i3j}^{\ell qJK}]^*+3C[\mathcal{T}_{i3j}^{\ell qJ}]C[\mathcal{V}_{i3j}^{\ell qJJ}]^*\right]\mathcal{K}\left(\tfrac{m_\psi^2}{m_b^2}\right)\right\},$$

with $\mathcal{J}$ and $\mathcal{K}$ the functions defined in Eq. (103) and the considered 'ℓq' operators extend the base of Eq. (5) to bottom quarks (the index '3' referring to b). The decays of a top quark would take exactly the same form (but the operators now include t instead of b). The decay widths of a b quark producing a pair of light neutralinos and a light quark $q_j = d,s$ (of negligible mass) read:

$$\Gamma[b\to\psi\psi q_j] = \frac{m_b^5}{1024\pi^3}\left\{\left(\left|C[\dot{\mathcal{S}}_{j3}^{dJ}]\right|^2+\left|C[\dot{\mathcal{S}}_{3j}^{dJ}]\right|^2+4\left|C[\dot{\mathcal{V}}_{j3}^{dJ}]\right|^2+4\left|C[\dot{\mathcal{V}}_{3j}^{dJ}]\right|^2\right)\widetilde{\mathcal{J}}\left(\tfrac{m_\psi^2}{m_b^2}\right)\right. \tag{110}$$
$$-\,8\frac{m_\psi^2}{m_b^2}\mathrm{Re}\left[C[\dot{\mathcal{S}}_{j3}^{dJ}]C[\dot{\mathcal{S}}_{3j}^{dJ+1}]+4C[\dot{\mathcal{V}}_{j3}^{dJ}]C[\dot{\mathcal{V}}_{3j}^{dJ}]\right]\widetilde{\mathcal{K}}\left(\tfrac{m_\psi^2}{m_b^2}\right)$$
$$+\,8\frac{m_\psi}{m_b}\mathrm{Re}\left[C[\dot{\mathcal{V}}_{j3}^{dJ}]C[\dot{\mathcal{S}}_{j3}^{dJ+1}]^*-C[\dot{\mathcal{V}}_{3j}^{dJ}]C[\dot{\mathcal{S}}_{3j}^{dJ}]^*\right.$$
$$\left.\left.-\,C[\dot{\mathcal{V}}_{j3}^{dJ}]C[\dot{\mathcal{S}}_{3j}^{dJ+1}]+C[\dot{\mathcal{V}}_{3j}^{dJ}]C[\dot{\mathcal{S}}_{j3}^{dJ}]\right]\widetilde{\mathcal{K}}\left(\tfrac{m_\psi^2}{m_b^2}\right)\right\},$$
$$\widetilde{\mathcal{J}}(x)\equiv\sqrt{1-4x}(1-14x-2x^2-12x^3)-48x^2(1-x^2)\ln\tfrac{1+\sqrt{1-4x}}{2\sqrt{x}},$$
$$\widetilde{\mathcal{K}}(x)\equiv\sqrt{1-4x}(1+5x-6x^2)-12x(1-2x+x^2)\ln\tfrac{1+\sqrt{1-4x}}{2\sqrt{x}}.$$

Again, this formula can be applied to the decays of a top quark (and $q_j \overset{!}{=} u,c$) up to the replacement of $\dot{\mathcal{S}},\dot{\mathcal{V}}_{jk}^{dJ}$ by $\dot{\mathcal{S}},\dot{\mathcal{V}}_{jk}^{uJ}$.

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