Right sneutrino with $\Delta L = 2$ masses as non-thermal dark matter

Avirup Ghosh, Tanmoy Mondal, Biswarup Mukhopadhyaya

*Regional Centre for Accelerator-based Particle Physics, Harish-Chandra Research Institute, HBNI, Chhatnag Road, Jhunsi, Allahabad - 211 019, India*

E-mail: avirupghosh@hri.res.in, tanmoymondal@hri.res.in, biswarup@hri.res.in

**Abstract:** We consider MSSM with right-chiral neutrino superfields with Majorana masses, where the lightest right-handed sneutrino dominated scalars constitutes non-thermal dark matter (DM). The $\Delta L = 2$ masses are subject to severe constraints coming from freeze-in relic density of such DM candidates as well as from sterile neutrino freeze-in. In addition, big-bang Nucleosynthesis and freeze-out of the next-to-lightest superparticle shrink the viable parameter space of such a scenario. We examine various $\Delta L = 2$ mass terms for families other than that $\Delta L = 2$ masses are difficult to reconcile with a right-sneutrino DM, unless there is either (a) a hierarchy of about 3 orders of magnitudes among various supersymmetry-breaking mass parameters, or, (b) strong cancellation between the higgsino mass and the trilinear supersymmetry breaking mass parameter for sneutrinos.

**Keywords:** Supersymmetry Phenomenology, Non-thermal Dark Matter
1 Introduction

While the search for the dark matter (DM) candidate(s) of our universe is on, one constantly feels the necessity of going beyond stereotypes in modelling physics beyond the standard model (SM) to accommodate the candidate particle(s). It is in this spirit that alternative candidates in small extensions of the minimal supersymmetric standard model (MSSM) have been explored. Scenarios with gravitino as warm DM are thus frequently discussed [1–10]. Another, quite minimalistic, extension is to extend the MSSM with a right-chiral neutrino superfield for each family, and postulate that one right-sneutrino-dominated scalar is the DM candidate. Various cosmological as well as phenomenological implications of this scenario have already been explored [10–18].

The interaction of such a particle with all other MSSM fields is proportional to the very small neutrino Yukawa coupling. Consequently, this scenario almost always leads to a non-thermal DM, with a long-lived next-to-lightest SUSY particle (NLSP). Such a long-lived NLSP mostly survives till it decouples from the thermal bath, and decays to the DM candidate thereafter. Usually, such scenarios are explored by assuming that neutrinos have just Dirac masses, in which cases their Yukawa coupling strengths are $O(10^{-13})$ [14, 19–26].

Some additional effects are operative if neutrinos have Majorana masses as well. This will require the right-handed neutrino fields to have $\Delta L = 2$ mass terms. The Type-I seesaw mechanism works here, thus requiring somewhat higher neutrino Yukawa couplings. This, however, entails the process of freeze-in of right sneutrinos while the NLSP (a stau, for example) is yet to decouple. Studies including those on similar non-SUSY theories have shown that such freeze-in contribution to the relic density impose rather strong constraints on the corresponding model parameters [27–30]. Since the freeze-in rates are proportional to the square of the Yukawa coupling(s), this effect is insignificant for Dirac neutrino scenarios [14, 24–26]. In a case with $\Delta L = 2$ mass terms present, however, the seesaw formula allows bigger Yukawa couplings, and thus the constraints tighten [31–35]. One has to also fit the neutrino mass and mixing patterns [36], and, quite seriously, the potential contribution of long-lived sterile neutrinos generated via the Dodelson-Widrow mechanism [37–44]. Thus the $\Delta L = 2$ Majorana masses, not only corresponding to the family where the right sneutrino DM candidate belongs but also for the other families, get seriously restricted, especially in view of the PLANCK data [45]. The way such restrictions arise is investigated in the present paper.

The right-handed Majorana mass matrix is assumed to be diagonal without loss of generality. Since the present neutrino oscillation experiments do not constrain mass of the lightest active neutrino we can constrain only two of the three heavy Majorana mass terms in the Majorana mass matrix.

In the pure Dirac case, the relic density of the non-thermal DM is obtained from post-freeze-out decay of the NLSP [14]. When freeze-in has a role to play, significant contribution to relic density also occurs in the decay of heavier superparticles still in thermal equilibrium [31]. The decay to the DM is governed essentially by the neutrino Yukawa coupling, even when it is driven by gauge interactions. Since such interactions will affect the (small but non-vanishing) left-sneutrino component in the DM, the deciding factor is the
sneutrino mixing angle, which essentially depends on the off-diagonal part of the sneutrino mass matrix. This block is proportional to the Yukawa coupling in usual parametrisations. Of course, the mixing is also reduced if one has the SUSY-breaking terms in the diagonal block are much larger than those in the off-diagonal block. This, however, requires a hierarchy of SUSY-breaking mass parameters, which is prima facie inexplicable, unless a cancellation between F-terms and soft terms in the off-diagonal block is engineered.

There can be several scales associated with the $\Delta L = 2$ masses, consistent with the Type-I seesaw mechanism. The commonly known GUT-scale seesaw won’t work for a case where non-thermal DM candidates are sought, as that would entail Yukawa couplings large enough for them to thermalise. Majorana masses in the electroweak scale [46–48], too, imply Yukawa couplings that could enhance the freeze-in rate unacceptably, unless a suppression in left-right mixing occurs via diagonal SUSY breaking terms that are several orders of magnitude larger than the off-diagonal ones. Though the situation is marginally better for $\mathcal{O}(1)$ GeV Majorana masses, a fine-tuned cancellation between the soft- and F-terms is required there, too. For hierarchical neutrino masses, the lightest neutrino mass being a free parameter, one can not easily constrain the corresponding $\Delta L = 2$ mass. We have considered this mass to be at the keV scale in our analysis, which does not hamper the generality of the constraints on the other Majorana masses derived here.

The scenario with all the three Majorana masses ranging from 500 MeV down to a few tens of keV is constrained by light element abundance. When all three masses are in keV scale, they all become warm dark matter, a scenario ruled out already due to overproduction of the warm keV scale dark matter via Dodelson-Widrow mechanism. We may have one $\Delta L = 2$ mass in the eV scale [49], as considered in [34], but such a situation brings in severe cosmological constraints [45, 50].

We establish freeze-in of right-sneutrino DM to be a major constraining factor with $\Delta L = 2$ masses. In such a case, not only the NLSP but also the rest of the MSSM spectrum contributes to freeze-in. This constraint is thus largely applicable to scenarios including various NLSP candidates. We focus on Stau and neutralino NLSP in particular.

Our paper has been organized as follows: In section 2 we discuss the model considered by us. Section 3 has been devoted to the discussions of constraints coming from freeze-in, NLSP freeze-out, low mass sterile neutrinos and Big-bang Nucleosynthesis (BBN). In section 4 we have shown how the constraints mentioned already can severely restrict the parameter space of such a scenario. In section 5 we have discussed the related issues of quasi-degenerate neutrino masses. We summarise and conclude in section 6. Finally the formula used by us have been tabulated in Appendices A and B.

### 2 Model

We consider the MSSM scenario augmented with three right chiral neutrino superfields ($\hat{N}_R$), which possess $\Delta L = 2$ mass terms. Hence the superpotential gets extended to the form [51–53]

$$
\mathcal{W} = \mathcal{W}_{MSSM} + Y^i_{\nu} \hat{H}_u \hat{L}^i \hat{N}_R^j + \frac{1}{2} M^{ij}_{N} \hat{N}_R^i \hat{N}_R^j
$$

(2.1)
where $\hat{H}_u$ is the higgs doublet that couples to up-type quarks, $\hat{L}$ is left-chiral SU(2) doublet lepton superfields and $\hat{N}_R$ is a right-handed neutrino superfield. While $Y_\nu$ is the neutrino Yukawa coupling matrix, $M_N$ is the $\Delta L = 2$ mass matrix for the heavy neutrinos, assumed, as already stated to be diagonal, since basis rotations in the right-handed neutrino sector is unlikely to affect our main conclusions.

Since we are not assuming any high scale mechanism of SUSY breaking we need to add all allowed SUSY-breaking terms phenomenologically. The relevant soft terms in this case are

$$L_{\text{soft}} = L_{\text{soft}}^{MSSM} - m^2_{N} \tilde{N}_R \tilde{N}_R + (m^2_{B} \tilde{N}_R \tilde{N}_R - T^i_{\nu} h_u \tilde{L}_i \tilde{N}_R^j + \text{h.c.}), \quad (2.2)$$

where $m_N (m_B)$ corresponds to $\Delta L = 0(2)$ susy-breaking sneutrino masses and $T_{\nu}$ is the coefficient of the well-known trilinear SUSY-breaking term. Left-right mixing in the sneutrino sector occurs, via the soft terms proportional to $T_{\nu}$, and the F-terms proportional to the higgsino mass parameter($\mu$), when $h_u (h_d)$ acquires the vacuum expectation values $v_u (v_d)$. In addition to $m_N$ and $m_B$, the F-term masses proportional to $M_N$ adds to the right-handed sneutrino masses. In our case, one requires $M_N$ not to exceed this SUSY-breaking scale, so that a right sneutrino may behave as a DM candidate.

2.1 Neutrino masses and mixings

An important constraint on the right sneutrino DM arises from the bearing it has on the mass and mixing patterns of light neutrinos. In this situation the simultaneous occurrence of a Dirac mass matrix ($m_D$) and a $\Delta L = 2$ mass matrix ($M_N$) leads to the standard Type-I seesaw expression for the light neutrino mass matrix:

$$m_\nu = -m_D M_N^{-1} m_D^T, \quad (2.3)$$

where $m_D = \frac{v_u}{\sqrt{2}} Y_\nu$ is the Dirac mass matrix. Working in the basis where the charged lepton Yukawa matrix is diagonal we need to incorporate a unitary matrix $U_{PMNS}$ to diagonalize $m_\nu$. Thus the diagonalized light neutrino mass matrix is

$$m_\nu^{\text{diag}} = U_{PMNS}^T m_\nu U_{PMNS}. \quad (2.4)$$

We parametrize $U_{PMNS}$ as,

$$U_{PMNS} = V \cdot \text{diag}(e^{i\alpha_1}, e^{i\alpha_2}, 1), \quad (2.5)$$

where $\alpha_1, \alpha_2$ are the Majorana phases and

$$V = \begin{bmatrix} c_{13} c_{12} & c_{13} s_{12} & s_{13} e^{-i\delta} \\ -c_{23} s_{12} - s_{23} s_{13} c_{12} e^{i\delta} & c_{23} c_{12} - s_{23} s_{13} s_{12} e^{i\delta} & s_{23} c_{13} \\ s_{23} s_{12} - c_{23} s_{13} c_{12} e^{i\delta} & -s_{23} c_{12} - c_{23} s_{13} s_{12} e^{i\delta} & c_{23} c_{13} \end{bmatrix},$$

where $c_{ij} = \cos \theta_{ij}$, $s_{ij} = \sin \theta_{ij}$ and $\delta$ is the Dirac CP phase.

To study the prospect of right-handed sneutrino being a non-thermal dark matter we need to appropriately fix $Y_\nu$. This in turn governs the decay of heavier sparticles to the
right-handed sneutrino DM, as will be discussed in section 3. Combining equations 2.3 into
equations 2.4,
\[ m_{\nu}^{\text{diag}} = -U_{PMNS}^T m_D \frac{1}{\sqrt{M_N}} \frac{1}{\sqrt{M_N}} m_D^T U_{PMNS}. \]  
(2.6)

Multiplying with \( \frac{1}{\sqrt{m_{\nu}^{\text{diag}}}} \) from left and right of equation 2.6 one obtains
\[ I = -\frac{1}{\sqrt{m_{\nu}^{\text{diag}}}} U_{PMNS}^T m_D \frac{1}{\sqrt{M_N}} \frac{1}{\sqrt{M_N}} m_D^T U_{PMNS} \frac{1}{\sqrt{m_{\nu}^{\text{diag}}}}. \]  
(2.7)

Using the Casas-Ibarra parametrization [54] one can introduce a complex orthogonal matrix
\( R \) as,
\[ R = i \frac{1}{\sqrt{M_N}} m_D U_{PMNS} \frac{1}{\sqrt{m_{\nu}^{\text{diag}}}}, \]  
(2.8)
so that equation 2.7 becomes \( RR^T = I \). Now by inverting equation 2.8, the neutrino
Yukawa coupling \( Y_\nu \) can be expressed as,
\[ Y_\nu = -i \frac{\sqrt{2}}{v_u} \sqrt{M_N} R \frac{1}{\sqrt{m_{\nu}^{\text{diag}}}} U_{PMNS}^\dagger. \]  
(2.9)

It is clear from equation 2.9 that \( R \) thus leads us to a Yukawa matrix that satisfies the
neutrino oscillation data by construction.\(^1\)

We have illustrated our main points for a normal hierarchy (NH) of neutrino masses.
To compute the Yukawa coupling matrix we need the \( U_{PMNS} \) and hence we have restricted
ourselves to the central value for the neutrino oscillation parameters \( \theta_{12} , \theta_{23}, \theta_{13} \) as given
by [36] and \( \delta \) is chosen to be \( \frac{3\pi}{2} \). Since the variation of Majorana phases does not change
the order of magnitude of the Yukawa couplings and our conclusion is not very sensitive
to such variation, we set \( \alpha_1, \alpha_2 \) to be \( \frac{\pi}{4} \). We have chosen the lightest neutrino mass
\( m_1 \) to be negligibly small (\( O(\sim 10^{-8} \text{ eV}) \)) while the masses of the other neutrinos are
determined by the solar(\( \Delta m^2_\odot \)) and atmospheric(\( \Delta m^2_\odot \)) mass-squared differences. The
lightest neutrino being effectively massless, masses of the relatively heavier neutrinos are
given by \( m_2 = \sqrt{\Delta m^2_\odot} = 0.0086 \text{ eV} \) and \( m_3 = \sqrt{\Delta m^2_\odot} = 0.05 \text{ eV} \) respectively. Since
one light neutrino is effectively massless, we have parametrized our \( R \) matrix as follows [55]:

\[ R = \begin{bmatrix} 1 & 0 & 0 \\ 0 & \cos \omega & \sin \omega \\ 0 & -\sin \omega & \cos \omega \end{bmatrix}, \]

where a non-vanishing \( R_{11} \) with appropriate \( m_{\nu}^{\text{diag}} \) reflects the postulated hierarchy. Our
analysis uses \( \omega = i \).

1Since for our present analysis only the order of magnitude of the Dirac Yukawa matrix elements mat-
ters, the Casas-Ibarra parametrization helps us fix the Yukawa matrix conveniently without any random
sampling.
2.2 Sneutrino masses and mixings

Turning now to the supersymmetric sector of the concerned model, the sneutrino mass matrix in the usual flavour basis \((\tilde{\nu}_L, \tilde{N}_R)\) is \([34, 35]\),

\[
M_{\tilde{\nu}} = \begin{pmatrix}
    m_{LL}^2 & m_{RL}^2 + \frac{v_u^2}{2} |Y_{\nu}|^2 & 0 & \frac{v_u}{\sqrt{2}} Y_{\nu}^1 M_N \\
    m_{RL}^2 & m_{RR}^2 & \frac{v_u}{\sqrt{2}} Y_{\nu}^0 M_N & -m_B^2 \\
    0 & \frac{v_u}{\sqrt{2}} Y_{\nu}^T M_N & m_{LL}^2 & m_{RL}^2 \\
    \frac{v_u}{\sqrt{2}} Y_{\nu} M_N & -m_B^2 & m_{RL}^2 & m_{RR}^2 \\
\end{pmatrix},
\]

where \(m_{LL}^2 = m_{iL}^2 + \frac{v_u^2}{2} |Y_{\nu}|^2 + \frac{m_Z^2}{2} \cos 2\beta\), \(m_{RL}^2 = -\mu^* \frac{v_u}{\sqrt{2}} Y_{\nu} + \frac{v_u}{\sqrt{2}} T_{\nu}\) and \(m_{RR}^2 = M_N^2 + \frac{v_u^2}{2} |Y_{\nu}|^2\) are the 3x3 mass matrices in the flavour basis. \(m_{iL}\) is the soft-mass term for charged lepton doublet \(\tilde{L}\). \(T_{\nu}\) will be parametrised as \(T_{\nu} = A_{\nu} Y_{\nu}, A_{\nu}\) being a mass-scale related to the scale of SUSY-breaking.

To study the sneutrino mass terms and phenomenology of sneutrino mass eigenstates originated from equation 2.10 it is convenient to introduce the real fields, \(\tilde{\nu}_L^i, \tilde{N}_R^i, \tilde{\nu}_R^i, \tilde{N}_L^i\) defined as follows \([34, 35, 53]\):

\[
\tilde{\nu}_L^i = \frac{1}{\sqrt{2}}(\tilde{\nu}_1^i + i \tilde{\nu}_2^i), \quad \tilde{N}_R^i = \frac{1}{\sqrt{2}}(\tilde{N}_1^i + i \tilde{N}_2^i),
\]

In the basis of these CP-even \((\tilde{\nu}_1^i, \tilde{N}_1^i)\) and CP-odd \((\tilde{\nu}_2^i, \tilde{N}_2^i)\) real sneutrino fields the sneutrino mass matrix is

\[
M_{\tilde{\nu}} = \begin{pmatrix}
    m_{LL}^2 & m_{RL}^2 + m_D M_N & 0 & 0 \\
    m_{RL}^2 + M_N^* m_D & m_{RR}^2 - m_B^2 & 0 & 0 \\
    0 & 0 & m_{LL}^2 & m_{RL}^2 + m_D M_N \\
    0 & 0 & m_{RL}^2 - M_N^* m_D & m_{RR}^2 + m_B^2 \\
\end{pmatrix}.
\]

The matrix in equation 2.12 can be diagonalized by an unitary rotation through an angle \(\theta_{\tilde{\nu}}\) given by,

\[
\tan 2\theta_{\tilde{\nu}} = \frac{2 (m_{RL}^2 \pm m_D M_N)}{m_{LL}^2 - (m_{RR}^2 \pm m_B^2)},
\]

where the top sign corresponds to mixing in the CP-even sector and bottom sign to CP-odd sector respectively. It is clear from the structure of the mixing matrix in (2.12) that the CP-even sneutrinos mix among themselves, as do the CP-odd ones. The decay of NLSP to LSP is proportional to \(\theta_{\tilde{\nu}}^2\) and hence to \(Y_{\nu}^2\). The BBN constraints suggests the NLSP lifetime \(\tau_{NLSP} \lesssim 100\) sec. there by posing a strong lower-limit on \(Y_{\nu}\). While the upper-limit on the \(Y_{\nu}\) has been set by correct relic density requirements.

Since we are considering the lightest sneutrino to be the LSP and also dominantly right-handed, its mass will be \(\sim \sqrt{m_{RR}^2 - m_B^2}\) and it will be CP-even in nature. Concomitantly, the lightest CP-odd sneutrino which is dominantly right-handed will have mass \(\sim \sqrt{m_{RR}^2 + m_B^2}\) and will also be a dark matter in addition to the CP-even LSP.\(^2\)

\(^2\)Since the only possible decay of the CP-odd lightest sneutrino is into LSP and the corresponding amplitude is suppressed by two factors of Yukawa coupling. Hence the lifetime of this CP-odd lightest sneutrino will always be larger than the age of the Universe.
have considered these lightest CP-even and CP-odd sneutrinos belong to the family which corresponds to the lightest active neutrino while the dominantly right-handed sneutrinos belonging to the other two families are heavy enough to decay before BBN. This assumption simplifies the calculation of relic density for the non-thermal sneutrino dark matter particle(s) which will be denoted as $\tilde{\nu}_{DM}$ collectively from now on.

3 Sources of constraints

While exploring the prospects of right-handed sneutrino being a non-thermal DM one has to take into account several constraints.

- The Yukawa interaction strength is bounded above by the consideration of freeze-in of sneutrinos.

- Post freeze-out decay of NLSP enhances the DM abundance and limits the parameter space following the observed density of DM.

- keV-MeV scale Majorana neutrino masses and the corresponding active-sterile mixings are constrained following BBN considerations, X-ray and Gamma-ray observations as well as the PLANCK data.

- Furthermore late decay of NLSP can mess up the observed light element abundances, thus implying a minimum value of the Yukawa couplings.

We discuss below these constraints in some detail.

3.1 Freeze-In of RH sneutrinos

A non-thermal DM candidate never reaches thermal equilibrium with the thermal plasma because of it’s low interaction strength. It is however, produced in the decays of the MSSM particles, mostly the NLSP, till the latter, being quasi-stable, freezes-out of the thermal bath. This is called the freeze-in process for the DM particles [31, 56–61]. The relic density of a non-thermal DM candidate generated via this process is proportional to the decay width of heavier superparticles into the $\tilde{\nu}_{DM}$. Since decay width is proportional to the square of the Yukawa coupling $Y_{\nu}$, the observed relic density as reported by the PLANCK collaboration [45] sets upper limits on $Y_{\nu}$ and $M_N$. The value $\Omega_{CDM} h^2 = 0.1199 \pm 0.0027$ has been used to obtain such limits.

In general, there exist a long-lived particle $\chi$, which is decoupled from the thermal bath. Despite the lack of equilibrium, $\chi$ can be produced from the decay of heavier bath particles, albeit via feeble interactions. Such production occurs mostly when the temperature $T$ drops below the mass of the decaying bath particles. For the sake of illustration, if a heavy particle $A$ in the thermal bath decays to lighter SM particle $B$ in the bath and $\chi$ then the Boltzmann equation governing the evolution of $\chi$ is given by,

$$ \frac{dn_{\chi}}{dt} + 3Hn_{\chi} = \int \frac{d^3p_A}{(2\pi)^3} \frac{m_A}{E_A} e^{-E_A/kT} \Gamma_A. $$

(3.1)
where $n_\chi$ is the comoving number density of the particle $\chi$, commonly known as *feebly interacting massive particle* (FIMP) and $H$ is the Hubble constant. Here $m_A$ and $E_A$ are the mass and energy of the decaying particle and $\Gamma_A$ is the decay width of the process $A \to B + \chi$.

By solving equation 3.1 for the yield of $\chi$ defined as, $Y_\chi = \frac{n_\chi}{s}$ where $s$ is the entropy density of the universe, one gets

$$Y_\chi \simeq \frac{45}{1.66 \times 4\pi^4} \frac{g_A M_{Pl} \Gamma_A}{m_A^2 g_* \sqrt{g_*}} \int_{x=0}^{x=\infty} K_1(x)x^3 dx.$$

If there are several heavy particle in the bath the decay of which may lead to the production of $\chi$ then one will have

$$Y_\chi \simeq \frac{45}{1.66 \times 4\pi^4} \frac{M_{Pl}}{g_* \sqrt{g_*}} \sum_{allA} \frac{g_A \Gamma_A}{m_A^2} \int_{x=0}^{x=\infty} K_1(x)x^3 dx,$$

where $x = \frac{m_A}{T}$ and $g_*^{s,\rho}$ are the number of degrees of freedom at the time of freeze-in i.e. $T \simeq m_A$ for entropy $s$ and energy density $\rho$ respectively.

In our specific context, the right-chiral sneutrino DM ($\tilde{\nu}_{DM}$) can be produced from the decays of higgsino dominated or gaugino dominated neutralinos ($\chi^0_i$), charginos ($\chi^\pm_i$) as well as left-chiral sleptons ($\tilde{l}_L$). The corresponding decay widths and amplitudes are tabulated in Appendices A and B respectively.

Let $g_*^{s}(g_*^\rho)$ be the number of degrees of freedom contributing to total entropy(energy) density. With $M_{Pl} = 10^{19}$ GeV, $g_*^s \simeq g_*^\rho \simeq 100$ [62] in equation 3.3, the freeze-in relic density of $\tilde{\nu}_{DM}$, denoted by $\Omega_{FI}$ is obtained

$$\Omega_{FI} h^2 = 2.755 \times 10^8 M_{\tilde{\nu}_{DM}} Y_{\tilde{\nu}_{DM}}$$

$$= 2.755 \times 10^8 \times 4 \times 10^{15} M_{\tilde{\nu}_{DM}} \frac{\Gamma_{\tilde{\nu}_{DM}}}{m_{susy}^2}$$

$$\simeq 10^{24} M_{\tilde{\nu}_{DM}} \frac{\Gamma_{\tilde{\nu}_{DM}}}{m_{susy}^2}$$

where $\Gamma_{\tilde{\nu}_{DM}} = \sum_{allA} g_A \Gamma_A \int_{x=0}^{x=\infty} K_1(x)x^3 dx$, using the appropriate degree of freedom $g_A$ for each of the decaying particles mentioned above, $M_{\tilde{\nu}_{DM}}$ is the mass of $\tilde{\nu}_{DM}$ and $m_{susy}$ is the generic value of the SUSY-breaking masses.

Freeze-in into $\tilde{\nu}_{DM}$ depends on the decay rates of charginos and neutralinos as well as sleptons into it. Considering neutralinos, one finds that any higgsino component in $\chi^0_i$ ($i=1,..4$) has a decay amplitude proportional to $Y_\nu$. If we further assume that the right sneutrino of the same family as that of the lightest active neutrino is the LSP and it has negligible mixing with other flavours, then the higgsino decay width to the LSP sneutrino is solely determined by the lightest neutrino Yukawa coupling $(Y_\nu)_{11}$. We have already seen that the lightest active neutrino mass is a free parameter in the normal hierarchical scenario. Therefore one can always have such a mass for it, when the contribution of any

\[ \text{For inverted hierarchy (IH), the conclusions are not very different if the sneutrino associated with the lightest active neutrino is the DM candidate. The assumption implicit here, which is not too artificial, is that the } M_N \text{ is only diagonal. The case of degenerate neutrino masses has been discussed in section 5.} \]
higgsino-dominated neutralino to sneutrino abundance is negligibly small compared to the observed DM abundance \([45]\). This choice considerably simplifies our analysis. The more general mixing scenario multiplies the number of free parameters but does not affect the constraint on the Majorana mass of the lightest neutrino. Throughout our analysis the Yukawa coupling corresponding to lightest active neutrino is \(\sim O(10^{-13})\), which ensures the contribution to sneutrino relic from higgsino decay is negligibly small (\(\sim 2\% - 3\%\) of total Relic).

For the freeze-in via slepton and gaugino-dominated neutralinos, on the other hand, the relevant factor is the sneutrino mixing angle \(\theta_{\tilde{\nu}}\). For a dominantly right-handed sneutrino dark matter, \(\theta_{\tilde{\nu}}\) is given by,

\[
\theta_{\tilde{\nu}} \sim \frac{Y_{\nu}(\mu v_d + \frac{v_u}{\sqrt{2}} A_{\nu}) \mp \frac{v_u}{\sqrt{2}} M_N Y_{\nu}}{m_{\tilde{l}_L}^2 - m_N^2}.
\] (3.5)

Freeze-in constrains \(\theta_{\tilde{\nu}}\) as a whole, and thus \(A_{\nu}\), too, is subjected to constraints \([24-26]\). We have varied \(A_{\nu}\) around Electroweak scale. A similar argument, based on equation 3.4, prompts us to keep \(M_{\tilde{\nu}_{DM}}\) not much higher than Electroweak scale.

For an order of magnitude estimate of parameters, let us consider \(M_{\tilde{\nu}_{DM}} \sim 100\) GeV. Then, from equation 3.4, it is clear that in order to restrain right sneutrino DM from being overabundant, one has to choose parameters in such a way that \(\frac{\Gamma_{\tilde{\nu}_{DM}}}{m_{\tilde{\nu}_{DM}}} \lesssim 10^{-27} \text{GeV}^{-1}\). Therefore, for different values of \(m_{\text{susy}}\), we have different limits on the mixing angle.

We illustrate a few cases below, keeping all non-strongly interacting superparticle masses (excepting higgsino mass parameter \(\mu\)) around the same scale, denoted by, \(m_{\text{susy}}\).

- For \(m_{\text{susy}} \sim 1\) TeV we must have \(\Gamma_{\tilde{\nu}_{DM}} \lesssim 10^{-21}\) GeV. As mentioned above we are looking for the least constrained scenario by choosing the lightest neutrino mass to be vanishingly small. Then the higgsino contribution to sneutrino relic is negligibly small and the main contribution comes from the decay of dominantly left-handed sfermions. The decay width of sfermions into \(\tilde{\nu}_{DM}\) is enhanced compared to those for gaugino and Higgsino decay by a phase-space factor \(\frac{m_{\text{susy}}^2}{m_{\text{EW}}^2}\). For \(m_{\text{susy}} \sim 1\) TeV this factor is \(\sim 100\). This in turn implies the mixing angle \(\theta_{\tilde{\nu}} \lesssim 3 \times 10^{-12}\). Since \(\theta_{\tilde{\nu}}\) is proportional to \(m_{\text{susy}}^{-2}\) the numerator on the right-hand side of equation 3.5 has to be \(\lesssim 3 \times 10^{-6}\) GeV\(^2\). There are two possible scenarios which give rise to such a small \(\theta_{\tilde{\nu}}\):

  - For a ‘natural’ scenario, each of the terms in the numerator of \(\theta_{\tilde{\nu}}\) has to be \(\sim 3 \times 10^{-6}\) GeV\(^2\). For Type I seesaw this implies that all the three Majorana neutrinos should be light, with \(M_N \lesssim O(100\) MeV\). This scenario is discussed in detail in section 4.1.

  - A relatively fine-tuned situation occurs for \(M_N \sim 500\) MeV or a few GeV. In this case, in order to obtain the correct relic density, one has to rely on the cancellation between \(\mu v_d\) and \(A_{\nu} v_u\). This situation is discussed in section 4.2.
Figure 1. Contribution to sneutrino relic coming from post freeze-out decay of Higgsino dominated $\chi_1^0$, Bino dominated $\chi_1^0$ and Stau NLSP for $m_{susy} = 1$ TeV. All strongly interacting sparticles are kept at 1.2–2.5 TeV in order to respect recent LHC bounds. For this plot CP-even DM mass have been taken to be 97 GeV while the mass of CP-odd sneutrino DM is 277 GeV.

- For a electroweak scale Majorana mass i.e. $M_N \simeq 100$ GeV, $v_u \sqrt{2} M_N Y_\nu \simeq 3 \times 10^{-2}$ GeV$^2$. In order to satisfy the constraint $\frac{\Gamma_{\tilde{\nu}_{DM}}}{m_{susy}} \lesssim 10^{-27}$ GeV$^{-1}$, one then requires $m_{susy} \simeq 100$ TeV or beyond. The corresponding results have been discussed in section 4.3.

3.2 Freeze-Out of NLSP

Till now we have considered only the contributions to sneutrino relic from heavier superparticles in thermal equilibrium. The NLSP is fairly long-lived, since its only possible decay into $\tilde{\nu}_{DM}$ is suppressed by the small mixing angle $\theta_\nu$. Thus, over and above the freeze-in process, a fair fraction of the NLSP population lives long enough to freeze-out of the thermal bath at an epoch to be decided by the Boltzmann equation [63, 64]. The post freeze-out decay of NLSP also raises $\tilde{\nu}_{DM}$ abundance to some extent. It should be noted here that since freeze-out is the only means of generating the DM abundance when the Yukawa coupling strength is very small, as happens for cases where neutrinos have just Dirac masses. The contribution to the $\tilde{\nu}_{DM}$ via NLSP freeze-out is given by [32, 33]

$$\Omega_{FO} h^2 = \frac{M_{\tilde{\nu}_{DM}}}{M_{NLSP}} \Omega_{NLSP} h^2,$$

where $\Omega_{NLSP} h^2$ is the density of NLSP at freeze-out relative to critical density $\rho_c = \frac{3H^2}{8\pi G}$. 

- 10 –
The contribution to relic density from the post freeze-out decay of \( \tilde{\tau}_R \) NLSP, higgsino dominated \( \chi_1^0 \) and bino dominated \( \chi_1^0 \) NLSP are shown in figure 1. For the purpose of illustration we have kept SUSY breaking masses of left-slepton, wino around 1 TeV while the strongly interacting superparticle masses are kept in the range 1.2 TeV – 2.5 TeV in order to be consistent with recent LHC bounds while the pseudoscalar mass \( m_{A^0} \) is kept at \( \simeq 2.5 \) TeV along with \( \tan \beta = 10.5 \). Such a choice of \( m_{A^0} \) prohibits higgsinos from reaching \( A \)-funnel region where its relic density is largely depleted. Thus throughout the region of \( M_{NLSP} \) depicted in figure 1 higgsino relic are determined by coannihilation with lightest chargino (\( \chi^+_1 \)) and second lightest neutralino (\( \chi^0_2 \)). While for stau (\( \tilde{\tau}_R \)) NLSP the relic density is determined by annihilation of staus into lepton pairs. The increase in the relic density with increase in mass of NLSP is attributed to decrease in s-channel annihilation cross-section. For a bino dominated \( \chi^0_1 \) NLSP the relic density is determined by t-channel anihilation into lepton-antilepton pairs. With increase in bino mass parameter the corresponding anihilation cross-section increases quadratically with \( M_{NLSP} \) and hence the relic density falls [65]. The figure 1 has been drawn with CP-even \( \tilde{\nu}_{DM} \) mass \( \simeq 97 \) GeV and CP-odd \( \tilde{\nu}_{DM} \) mass \( \simeq 277 \) GeV. The NLSP abundance at freeze-out has been calculated using micrOmegas-5.0 [66].

Clearly for bino dominated NLSP the post freeze-out decay of NLSP overproduces \( \tilde{\nu}_{DM} \) unless coanihilation comes into the picture. The situation is less binding for a Higgsino-dominated NLSP as well as dominantly right-handed stau NLSP. Thus soft masses corresponding to left-sleptons, bino and wino are taken to be in the same scale in our analysis.

### 3.3 Constraints on light sterile neutrino

The active-sterile neutrino mixing angle, is given by

\[
\sin^2 \omega_i = \sum_{\alpha} (Y_{\nu})_{\alpha i}^2 \frac{y_{\alpha i}^2}{M_{N,i}}.
\]

It is severely constrained for a MeV-keV scale sterile neutrino as shown in figure 2 [67–70]. The constraint emerges for a sterile neutrino with MeV scale, since otherwise it decays after BBN and disturbs the light element abundance via its hadronic decay modes. This region is shown in yellow in figure 2. The constraints apparently disappears if the lifetime of the sterile neutrino is larger than the age of the Universe(\( \simeq 10^{17} \) sec). However the mixing angle then suffers from a still stronger constraints from X-ray and gamma-ray observations attributed to the loop suppressed decay \( N_R \rightarrow \nu_L \gamma \) [41]. The violet region represents the parameter space thus ruled out. The orange region corresponds to the situation when the sterile neutrino abundance from the Dodelson-Widrow(DW) mechanism [37] is 1-100% of the currently measured DM relic density. The green region shows abundance of sterile neutrino exceeds the observed DM abundance. The red line in figure 2 denotes the mixing angle required to satisfy neutrino oscillation data via Type I seesaw mechanism. It is clear from the figure that for all three sterile neutrino masses \( \lesssim 10 \) keV the mixing angle satisfying oscillation data is large enough to overproduce(\( \Omega_N h^2 \gg \Omega_{CDM} h^2 \)) the sterile neutrinos via the DW mechanism. Additional constraints on sterile neutrino masses other than what shown in figure 2 come from the Lyman-\( \alpha \) forest observation [71–73] which eliminates \( M_N \lesssim 8 \) keV if the sterile neutrino is the only DM candidate. In our case sterile neutrino contribute only a fraction of the relic density, and hence the last mentioned...
constraint is somewhat relaxed.

3.4 Big Bang Nucleosynthesis

The presence of long-lived particles is a potential threat to the light element abundance, as predicted by standard Big-Bang Nucleosynthesis (BBN). Long-lived particles may give rise to non-thermal nuclear processes (non-thermal BBN) attributed to their late decay into energetic SM particles. It disturbs the observed light element abundance. As mentioned earlier, we have considered the possibility of a right-handed stau ($\tilde{\tau}_R$) and/or a neutralino ($\chi_1^0$) as the NLSP. A $\tilde{\tau}_R$ NLSP can have late decay into $\tilde{\nu}_{DM}$ and a $W^\pm$ boson via the small neutrino Yukawa coupling. The worst victim of subsequent hadronic decays of the $W^\pm$ is the deuterium abundance. A reasonable solution is to restrict the lifetime to $\mathcal{O}(100 \text{ sec})$ [14, 50, 74–77], a constraint we have imposed through out on the spectrum. The decay $\chi_1^0 \rightarrow \nu \tilde{\nu}_{DM}$ leads to highly energetic light neutrinos. Their scattering with background neutrinos generate energetic $e^\pm$ which subsequently produce energetic photons. These can change the light element abundances through photo-dissociation of light nuclei.

Figure 2. This plot depicts the available parameter space for a MeV-keV scale sterile neutrino [38]. Barring the constraints from Big-Bang Nucleosynthesis (Yellow patch implies sterile neutrino lifetime is between BBN and the current age of the Universe) the MeV-keV scale sterile neutrino behaves as a decaying WDM candidate which can decay into an active neutrino along with a energetic photon and hence X-ray, Gamma ray observations restrict the associated parameter space severely (Violet) [41]. For low enough masses ($\lesssim 10 \text{keV}$) the constraints on the present day energy density of the DM component of universe is proven to most stronger (Green). Orange patch signifies Dodelson-Widrow (DW) mechanism accounts for between 1% and all of dark matter (DM). red: mass-mixing angle combinations enforced by the seesaw mechanism.
The three-body decay $\chi_1^0 \rightarrow \tilde{\nu}_DM \nu h^{(*)}$, though phase space suppressed, causes additional tension [78]. All such problems are avoidable by imposing an upper limit of $O(100 \text{ sec})$ on the NLSP lifetime.

### 3.5 Other constraints

In addition to the aforementioned constraints, the following restrictions also have been adhered to in our numerical analysis:

- The mass of lightest higgs emerging from the spectrum calculation modulo various uncertainties, is in the range $123 \text{ GeV} < m_{h^0} < 128 \text{ GeV}$ [79, 80].
- The PMNS-driven contributions to FCNC processes like $\mu \rightarrow e\gamma$, $\tau \rightarrow \mu\gamma$ etc. are within the current limits [81].

### 4 Results

We use for illustration $m_N = 200 \text{ GeV}$, $m_B = 175 \text{ GeV}$, where $m_N$ and $m_B$ are defined in equation 2.2. This yields a CP-even state of mass $\simeq 97 \text{ GeV}$ and a CP-odd state of mass $\simeq 277 \text{ GeV}$. The two states split because of the $\Delta L = 2$ mass term proportional to $m_B$. As mentioned already, both of them serve as DM candidate at the present epoch. We have checked that different choices of $\tilde{\nu}_DM$ masses (around the electroweak scale) do not alter the constraints on the Majorana neutrino mass(es) significantly. We have considered four representative cases in terms of the diagonal entries of $M_N$. The corresponding results are presented in the subsequent subsections. The scenarios considered here are as follows:

(a) All three Majorana masses in the keV scale.
(b) Only the lightest Majorana mass in keV scale, the others being in the range 500MeV-GeV.
(c) The lightest Majorana mass in the keV scale and others in the electroweak scale.
(d) All three Majorana masses are in the electroweak scale. In this scenario, say for $m_1 \simeq 10^{-8} \text{ eV}$, the Yukawa coupling corresponding to the lightest neutrino will be $\approx 10^{-10}$ and the freeze-in contribution from higgsino decay will overclose the Universe. However, we can always tune the lightest active neutrino mass to reduce the Yukawa coupling to a desired value and then the scenario is exactly similar to the preceding case. Hence we will not discuss this scenario in detail.

For cases (b) and (c) one right-handed Majorana mass is kept at the keV scale. In such cases, the neutrino warm DM and the non-thermal DM in the form of $\tilde{\nu}_DM$ can coexist without being ruled out by freeze-in as well as freeze-out. However, it should be remembered that, in a hierarchical scenario, the lightest active neutrino mass being a free parameter the corresponding Majorana mass can in principle be made larger without changing the corresponding Yukawa coupling.
Figure 3. For $m_{\text{susy}} \approx 1\text{TeV}$ the dependence of freeze-in relic density with $\mu$ (left-plot) and $M_N$ (right-plot). All the Majorana masses have been taken in the keV scale, $A_\nu$ is varied in the EW scale and $\tan \beta = 10.5$ has been taken.

All the interaction of $\tilde{\nu}_{DM}$ depends on the order of its Yukawa coupling. In the first case, where all the Majorana masses are very small, the corresponding Yukawa couplings are also minuscule. Hence this scenario is very similar to the case of purely Dirac neutrinos. In the two remaining cases we have examined the effect of increase in the order of magnitude of Yukawa couplings on the freeze-in relic density of $\tilde{\nu}_{DM}$.  

4.1 keV scale Majorana mass

In this case we have kept all the three Majorana masses ($M_N$) of neutrinos in the keV scale as mentioned already. We have confined ourselves to the following illustrative values of SUSY parameters: $m_{\text{susy}} = 1\text{ TeV}$, $\tan \beta = 10.5$.

The largest Yukawa coupling for keV scale Majorana neutrinos is $\simeq O(10^{-10})$ and hence the term $\frac{v_u}{\sqrt{2}} M_N Y_{\nu}$ in the numerator of sneutrino mixing angle, $\theta_{\tilde{\nu}}$, is $\simeq 10^{-13}\text{GeV}^2$. This is much smaller than that required for $m_{\text{susy}} \approx 1\text{ TeV}$ as discussed in section 3.1. Hence the relic density is determined by the term $\frac{v_u}{\sqrt{2}} (A_\nu - \mu \cot \beta) Y_{\nu}$. Thus as the difference between $A_\nu$ and $\mu \cot \beta$ increases, the freeze-in relic density also goes up. The variation of freeze-in relic density with $\mu$ for three different values of $A_\nu$ with $M_N = 10$ keV is shown in figure 3 (left). As is clearly seen there for $A_\nu = 100$ GeV the relic density is minimum around $\mu \approx 1.05$ TeV. As one departs from this value of $\mu$, the relic density increases as a higher $\theta_{\tilde{\nu}}$ enhances the freeze-in process. The smooth increase in relic density is accounted for the low value of $Y_{\nu}$ multiplied to $(A_\nu - \mu \cot \beta)$. The term $\frac{v_u}{\sqrt{2}} M_N Y_{\nu}$ does not play a significant role in this case, when $A_\nu$ and $\mu$ are tuned to cancel up to five decimal places. A similar feature is also seen for $A_\nu = 50$ GeV. The relic density for $A_\nu = 20$ GeV is minimum for $\mu = 210$ GeV.

Similarly, for fixed $\mu$ and $A_\nu$, the relic density increases with the Majorana mass due to the term $\frac{v_u}{\sqrt{2}} M_N Y_{\nu}$ in the numerator of $\theta_{\tilde{\nu}}$, as per equation 3.5. The variation of relic density with $M_N$ is shown in figure 3 (right).
Figure 4. Show how the exact cancellation between supersymmetry conserving term $\mu \cot \beta$ and supersymmetry breaking term $A_\nu$ leads to correct freeze-in relic density for sneutrino with MeV-GeV scale Majorana masses. Left plot is for $\mu = 800$ GeV and right plot is for $\mu = 1200$ GeV. $\tan \beta = 10.5$ in both the cases.

Since our choices of Majorana masses ($M_N$) and Yukawa couplings ($Y_\nu$) are always consistent with neutrino oscillation data, the mixing angle of the heaviest active neutrino lies on the red line of figure 2. However, now there are three warm DM candidates which push the relic density to unacceptably large values via DW mechanism.

4.2 Majorana masses in the scale MeV-GeV

It is evident from figure 2 that, the simultaneous fulfillment of oscillation data, neutrino relic density constraints and BBN requirements imply that the two heavier sterile neutrinos must have masses $M_{H_2} \gtrsim 500$ MeV. For such a choice of parameters the largest Yukawa coupling is $\sim O(10^{-8})$ and hence the term $\frac{v_u}{\sqrt{2}} M_N Y_\nu \simeq 10^{-6}$ GeV$^2$, which is what necessary to obtain correct relic density following our discussion in section 3.1. However, one would also require the term $\frac{v_u}{\sqrt{2}} (A_\nu - \mu \cot \beta) Y_\nu$ to be on the same order of magnitude, namely, $O(10^{-6})$ GeV$^2$. For $A_\nu$ and $\mu$ in the electroweak scale, this is only possible if these two terms cancel between themselves. Thus one needs

$$A_\nu \simeq \mu \cot \beta,$$

in order to reproduce the correct relic density.

Figure 4 (left) shows the variation of relic density with $A_\nu$ for $\mu = 800$ GeV and $\tan \beta = 10.5$. One can see that for $\mu = 800$ GeV and $\tan \beta = 10.5$ relic density is minimized at $A_\nu = 76.2$ GeV which is in agreement with equation 4.1. The situation is quite similar for $\mu = 1200$ GeV as shown in figure 4 (right). Figure 5 depicts the variation of relic density with $\mu$ for $A_\nu = 100$ GeV, which minimizes around $\mu = 1.05$ TeV.

One must also notice that in situations when the two terms $\mu \cot \beta$ and $A_\nu$ cancel exactly, the mixing angle $\theta_{\tilde{\nu}}$ is dominated by the term $\frac{v_u}{\sqrt{2}} M_N Y_\nu$. This explains the increase in relic density with the heavier Majorana masses ($M_{H_2}$). Thus, for a fixed values $\mu$ and $\tan \beta$, the required $A_\nu$ to reproduce correct relic density is also constrained. The
rise in relic density as one moves away from the region $A_\nu \approx \mu \cot \beta$ is quite abrupt. This happens because, for $A_\nu$ and $\mu$ around the electroweak scale, $(A_\nu - \mu \cot \beta)$ is naturally around 100 GeV unless they are finely tuned.

### 4.3 Electroweak scale Majorana masses ($M_N \simeq 100$ GeV)

Here we have chosen the Majorana masses for heavier sterile neutrinos to be $\simeq 100$ GeV, while the lightest Majorana mass, $M_{N,1}$, kept at 10 keV. Constraints on the heavier Majorana masses do not really depend on $M_{N,1}$. For such large Majorana masses the largest Yukawa coupling is $O(10^{-7})$. Hence the term $\frac{v_u}{\sqrt{2}} M_N Y_\nu$ is $\simeq 10^{-3}$ GeV$^2$, which is about $3 \times 10^3$ times larger than the value required for allowing $m_{susy} \simeq 1$ TeV. Hence with electroweak scale Majorana neutrinos TeV scale SUSY breaking is disfavored for a right-handed sneutrino LSP. To reconcile electroweak scale Majorana masses with right sneutrino LSP, one has to increase the SUSY breaking scale to a much higher value ($m_{susy} \simeq 100$ TeV) since the sneutrino mixing angle decreases quadratically as $m_{susy}$ increases. In figure 6 the dependence of relic density on the variation of SUSY breaking scale is shown for $\mu$ and $A_\nu$, both around EW scale with $\tan \beta = 3.5$. It should be noted that, in such a case, the freeze-in constraint on $\theta_\tilde{\nu}$ requires $A_\nu$ and $M_{\tilde{\nu}_{DM}}$ to be still within a TeV. On the other hand, $\mu$ should also be within a TeV, not only for satisfying electroweak symmetry breaking conditions, but also to control freeze-in via $\tilde{\nu}_L \to \tilde{\nu}_R h$. However, we then face an inexplicable hierarchy of SUSY-breaking parameters, unless one engineer some highly contrived fine-tuning.

The monotonic fall in relic density with increasing $m_{susy}$ can be attributed to the dependence of $\theta_\tilde{\nu}$ on $1/m_{susy}^2$. Another interesting point to note is that the relic density for $\mu = 800$ GeV is larger for $A_\nu = 500$ GeV compared to, say, $A_\nu = 100$ GeV while the

![Figure 5](image)

**Figure 5.** Show the variation of freeze-in relic density with $\mu$ as $A_\nu$ is kept fixed. Only a restricted region allowed via the cancellation of supersymmetry conserving $\mu \cot \beta$-term and SUSY breaking term $A_\nu$. $\tan \beta = 10.5$ has been taken for illustration.
situation is opposite for $\mu = 1500$ GeV. In case of $\mu = 800$ GeV the term $(A_\nu - \mu \cot \beta)$ is larger for $A_\nu = 500$ GeV compared to $A_\nu = 100$ GeV, thereby enhancing $\theta_\tilde{\nu}$. The situation is exactly opposite for $\mu = 1500$ GeV.

5 Quasi-degenerate neutrino masses

It has been already mentioned that, for hierarchical neutrino masses, the free parameter in the form of the lightest active neutrino mass may be tuned to control the sneutrino abundance. For the quasi-degenerate scenario, the lightest neutrino mass is bound to be $\simeq 10^{-2}$ eV. In order to prevent higgsino decays from overproducing sneutrino dark matter via Yukawa interaction, one has to lower the Majorana mass corresponding to the lightest neutrino eigenstate. One thus finds oneself pushed down to eV-scale value for this Majorana mass. Such a situation has already been mentioned in [34]; however, recent limits from BBN and recombination [45, 50] strongly restricts this scenario. Although some wisely engineered ways of evading such constraints exist [82–88], we do not consider this possibility in greater detail.

6 Summary and conclusions

We have worked with an MSSM scenario augmented with three families of right-chiral neutrino superfields, and $\Delta L = 2$ terms in the superpotential as well as the scalar potential where masses for the light neutrinos have been generated following Type-I seesaw mechanism. The right-handed sneutrinos are non-thermal dark matter candidate as a result of small Yukawa couplings. We have restricted to hierarchical neutrino masses, latest constraints from Planck data on DM as well as data on the neutrino sector. We have gone beyond earlier studies where constraints on the neutrino Majorana masses were obtained, mostly within a degenerate neutrino scenario. While the eV-scale scenario proposed there is presently disfavoured, we have shown that the current picture admits of considerable
varieties as well as constraints in view of current observations. These include, on the theoretical side, constraints from both the NLSP freeze-out and freeze-in of right sneutrinos as well as sterile neutrinos. Our conclusions are:

- While the lightest neutrino mass and consequently the corresponding entry in the (diagonal) right-handed neutrino mass matrix is a free parameter, the remaining $\Delta L = 2$ mass terms are constrained rather tightly.

- $\Delta L = 2$ masses on the electroweak scale and above are possible only if there is a hierarchy of 2-3 orders of magnitude between DM mass $M_{\tilde{\nu}}^{\text{DM}}$, $\mu$ and $A_\nu$ on the one hand, and the remaining SUSY-breaking masses on the other.

- For $\Delta L = 2$ masses in the range 500 MeV-a few GeV, all constraints can be satisfied only if $A_\nu$ and the $\mu$-parameter are fine-tuned to 4-5 decimal places.

- Majorana masses in the few keV-500 MeV range are disfavoured by standard BBN, as the sterile neutrinos have lifetimes too long to maintain the observed light element abundances.

- If all three $\Delta L = 2$ entries in $M_N$ are in the keV range, then all three sterile neutrinos effectively constitute warm dark matter. Such a situation, however, is disfavoured by the Dodelson-Widrow mechanism where the freeze-in rate of the warm DM goes up unacceptably.

We thus conclude that the right-sneutrino DM scenario, if at all the picture of nature, is subject to rather severe constraints for practically all orders of magnitude of the $\Delta L = 2$ masses. The exception to this requires an inexplicable hierarchy among SUSY-breaking mass parameters which perhaps can be justified by going beyond MSSM.

7 Acknowledgements

We thank Raj Gandhi for helpful discussions and Saptarshi Roy for advice on numerical techniques. AG and BM acknowledges the hospitality of the Theoretical Physics Department, Indian Association for Cultivation of Science, Kolkata, where the last part of the project was carried out. This work was partially supported by funding available from the Department of Atomic Energy, Government of India, for the Regional Centre for Accelerator-based Particle Physics (RECAPP), Harish-Chandra Research Institute.
A Decay Widths

The decay widths of the channels that contribute to freeze-in of $\tilde{\nu}_D$ are as follows [89]:

$$\Gamma (\tilde{H}_u^0 \rightarrow \tilde{\nu}_R \nu_L) = \frac{1}{32\pi M_{\tilde{H}_u^0}} \beta_f (M_{\tilde{H}_u^0}, M_{\nu_L}) (M_{\tilde{H}_u^0}^2 + M_{\nu_L}^2 - M_{\tilde{\nu}_R}^2) |A \left( \tilde{H}_u^0 \rightarrow \tilde{\nu}_R \nu_L \right)|^2,$$

(A.1a)

$$\Gamma (\tilde{H}_u^+ \rightarrow \tilde{\nu}_R t^+) = \frac{1}{32\pi M_{\tilde{H}_u^0}} \beta_f (M_{\tilde{H}_u^+}, M_{\nu_L}) (M_{\tilde{H}_u^+}^2 + M_{\nu_L}^2 - M_{\tilde{\nu}_R}^2) |A \left( \tilde{H}_u^+ \rightarrow \tilde{\nu}_R t^+ \right)|^2,$$

(A.1b)

$$\Gamma (\tilde{\nu}_L \rightarrow \tilde{\nu}_R h) = \frac{1}{32\pi M_{\tilde{\nu}_L}} \beta_f (M_{\tilde{\nu}_L}, M_h) |A \left( \tilde{\nu}_L \rightarrow \tilde{\nu}_R h \right)|^2,$$

(A.1c)

$$\Gamma (\tilde{\nu}_L \rightarrow \tilde{\nu}_R Z) = \frac{1}{32\pi M_{\tilde{\nu}_L}} \frac{M_{\tilde{\nu}_R}^4}{M_Z^2} \beta_f (M_{\tilde{\nu}_L}, M_Z) |A \left( \tilde{\nu}_L \rightarrow \tilde{\nu}_R Z \right)|^2,$$

(A.1d)

$$\Gamma (\tilde{\nu}_L \rightarrow \tilde{\nu}_R W^-) = \frac{1}{32\pi M_{\tilde{\nu}_L}} \frac{M_{\tilde{\nu}_R}^4}{M_W^2} \beta_f (M_{\tilde{\nu}_L}, M_W) |A \left( \tilde{\nu}_L \rightarrow \tilde{\nu}_R W^- \right)|^2,$$

(A.1e)

$$\Gamma (\tilde{\nu}_L \rightarrow \tilde{\nu}_R W^0) = \frac{1}{32\pi M_{\tilde{\nu}_L}} \frac{M_{\tilde{\nu}_R}^4}{M_{W^0}^2} \beta_f (M_{\tilde{\nu}_L}, M_{W^0}) |A \left( \tilde{\nu}_L \rightarrow \tilde{\nu}_R W^0 \right)|^2,$$

(A.1f)

$$\Gamma (\tilde{\nu}_L \rightarrow \tilde{\nu}_R W^+ + \tilde{\nu}_R l^+) = \frac{1}{32\pi M_{\tilde{\nu}_L}} \frac{M_{\tilde{\nu}_R}^4}{M_{W^+}^2} \beta_f (M_{\tilde{\nu}_L}, M_{W^+}) |A \left( \tilde{\nu}_L \rightarrow \tilde{\nu}_R W^+ \right)|^2,$$

(A.1g)

where $\beta_f (M_A, M_B) = \left( 1 - 2 \frac{M^2_{\tilde{\nu}_R} + M^2_B}{M^2_A} + \frac{(M^2_{\tilde{\nu}_R} - M^2_B)^2}{M^4_A} \right)^{1/2}$ and $A_{A \rightarrow B \tilde{\nu}_R}$ is the amplitude for the decay $A \rightarrow B \tilde{\nu}_R$. We have tabulated all the relevant amplitudes in B.

B Amplitudes

In this section we give the expressions for the amplitudes for all the decays considered in section 3.1. In the expressions that follows we have denoted the $U_{PMNS}$ as $U_0$ and sneutrino mixing matrix as $Z_\nu$. $U_V$ will denote the full $6 \times 6$ matrix that diagonalizes the $6 \times 6$ neutrino mass matrix given by,

$$M_{\nu} = \begin{bmatrix} 0 & m_D^T \\ m_D & M_N \end{bmatrix}$$

(B.1)

The index $i, j$ denotes the superparticle generation appearing in the vertex while $I$ denotes the same for R-parity even particles. For $\tilde{\nu}_D M$ the index $i = 6$ need to be used.
B.1 Interactions with Gauginos

1. \(A(\tilde{B}^0 \to \tilde{\nu}_R^i \nu_L^j) = \frac{i}{2} g_1 P_L \sum_{k=1}^3 U_{01k}^* Z_{\tilde{\nu}^i k}\)

2. \(A(\tilde{W}^0 \to \tilde{\nu}_R^i \nu_L^j) = -\frac{i}{2} g_2 P_L \sum_{k=1}^3 U_{01k}^* Z_{\tilde{\nu}^i k}\)

3. \(A(\tilde{H}_u^0 \to \tilde{\nu}_R^i \nu_L^j) = -\frac{i}{2} g_2 P_L \left( \sum_{k,l=1}^3 Z_{\tilde{\nu}^i k}^* Y_{\nu^l} U_{V^l} U_{V^l} + \sum_{k,l=1}^3 Z_{\tilde{\nu}^i k}^* Y_{\nu^l} U_{V^l} \right)\)

4. \(A(\tilde{H}_u^- \to \tilde{\nu}_R^i \nu_L^j) = -\frac{i}{2} g_2 P_L \sum_{k=1}^3 Z_{\tilde{\nu}^i k}^* Y_{\nu^l} U_{V^l}\)

B.2 Interactions with sfermions

Sneutrino LSP can decay from sleptons and heavier sneutrinos via SM gauge bosons or SM higgs which we have discussed below:

B.2.1 Interactions with Gauge-Bosons

1. \(A(\tilde{e}_L^j \to \tilde{\nu}_R^i W^\mu) = \frac{i}{2} Z_{\tilde{\nu}^i j}^* (p_{\tilde{\nu}^i} - p_{\tilde{e}_L})^\mu\)

2. \(A(\tilde{\nu}_H^j \to \tilde{\nu}_R^i Z^\mu) = \frac{1}{2} (g_2 \cos \theta_w + g_1 \sin \theta_w) \sum_{p,k=1}^3 Z_{\tilde{\nu}^i j}^* Z_{\tilde{\nu}^i k}^* Y_{\nu^p k} + \mu^\ast \sum_{p,k=1}^3 Z_{\tilde{\nu}^i j}^* Z_{\tilde{\nu}^i k}^* Y_{\nu^p k}\)

where \(H\) denotes the heavier sneutrino eigenstates and hence \(j\) runs from 1 to 5. One has to keep in mind that this vertex is non-zero only when the CP nature of the two participating sneutrinos (\(\tilde{\nu}_H^j\) and \(\tilde{\nu}_R^i\)) are opposite.

B.2.2 Interactions with SM higgs

1. For a heavier CP-even (or CP-odd) sneutrino (\(\tilde{\nu}_H^j\)) decaying into a lighter CP-even (or CP-odd) sneutrino (\(\tilde{\nu}_R^i\)),

\[
A(\tilde{\nu}_H^j \to \tilde{\nu}_R^i h) = \frac{i}{2\sqrt{2}} \left( \mu \sum_{p,k=1}^3 \left[ Z_{\tilde{\nu}^i j}^* Z_{\tilde{\nu}^i (3+p)}^* + Z_{\tilde{\nu}^i j}^* Z_{\tilde{\nu}^i (3+p)}^* \right] Y_{\nu^p k}^* + \mu^* \sum_{p,k=1}^3 \left[ Z_{\tilde{\nu}^i j}^* Z_{\tilde{\nu}^i (3+p)}^* + Z_{\tilde{\nu}^i j}^* Z_{\tilde{\nu}^i (3+p)}^* \right] Y_{\nu^p k}^* \right) - 2(g_1^2 + g_2^2) v d \sum_{k=1}^3 Z_{\tilde{\nu}^i j}^* Z_{\tilde{\nu}^i k}^* \]
2. For a heavier CP-even (or CP-odd) sneutrino ($\tilde{\nu}_H^j$) decaying into a lighter CP-odd (or CP-even) sneutrino ($\tilde{\nu}_i^*$),

$$A\left(\tilde{\nu}_H^j \rightarrow \tilde{\nu}_i^* h\right) = \frac{1}{2\sqrt{2}} \left( \mu \sum_{p,k=1}^{3} \left[ -Z_{\tilde{\nu}^j_H}^* Z_{\tilde{\nu}^i}^* (3+p) + Z_{\tilde{\nu}^j_H}^* (3+p) Z_{\tilde{\nu}^i} Z_{\tilde{\nu}^i}^* \right] Y_{\nu^* pk} + \mu^* \sum_{p,k=1}^{3} \left[ Z_{\tilde{\nu}^j_H}^* Z_{\tilde{\nu}^i}^* (3+p) - Z_{\tilde{\nu}^j_H}^* (3+p) Z_{\tilde{\nu}^i} Z_{\tilde{\nu}^i}^* \right] Y_{\nu pk} \right)$$

For higgs-mediated sneutrino decays also $j$ runs from 1 to 5 as in the case of Z-mediated processes.

In diagonalizing sneutrino mass matrix and calculating all the Feynman Rules we have used SARAH [90].

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