Right-handed sneutrinos as asymmetric DM and neutrino masses from neutrinophilic Higgs bosons

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Abstract
We consider an extension of the Next-to-Minimal Supersymmetric Standard Model by three right-handed neutrinos and a pair of neutrinophilic Higgs superfields. The small neutrino masses arise naturally from a small vacuum expectation value of the additional Higgs fields (hence without lepton number violation), while the lightest right-handed sneutrinos can constitute asymmetric Dark Matter. The right-handed sneutrino and baryon asymmetries are connected through equilibrium processes in the early universe, explaining the coincidence of the DM and baryon abundances. We show that particle physics and astrophysical constraints are satisfied.
1 Introduction

A large part of the total matter of the universe manifests itself only through its gravitational effects. Although the evidence for its existence is strong, the nature of this Dark Matter (DM) is not yet clear. Observations of the Bullet Cluster [1] show a clear separation between the location of the luminous matter and the location of gravitating matter providing a confirmation of particle DM. Assuming the standard cosmological model and using observations of the cosmic microwave background, it is possible to determine with great accuracy the present abundance of DM [2,3].

In most scenarios, the DM abundance is a relic of annihilation processes during the early hot universe. These processes stopped when the expansion rate of the universe became larger than their rate and, since then, the DM density per comoving volume has remained constant. In order that the relic density fits the observed value, the annihilation cross section should be roughly of the order of weak interactions. Due to this fact Weakly Interacting Massive Particles (WIMPs) are preferred candidates for DM.

The Standard Model (SM) of particle physics does not include a particle with the required characteristics. Supersymmetric (R-parity conserving) extensions of the Standard Model provide two stable candidates, the neutralino and the sneutrino. However, even though they are both weakly interacting particles, their relic density does not automatically have the correct value. In the case of neutralinos it can vary over many orders of magnitude as a function of the unknown parameters of the theory. (Left-handed) sneutrinos are problematic due to their strong coupling to the Z bosons. This strong coupling results in a large annihilation cross section and hence in a too small relic density compared to the observed value, and in a scattering cross section off nuclei far above the current bounds from XENON100 [4].

The present baryonic matter abundance, on the other hand, is not sensitive to the baryon pair annihilation cross section. It originates from the baryon asymmetry, since all the anti-baryons have already been completely depleted. Astonishingly, both DM and baryon abundances are of the same order of magnitude ($\Omega_{DM}/\Omega_b \approx 5$). This coincidence of the values of these two quantities motivated the search for alternative mechanisms for the generation of the DM relic density, connected to the one of the baryons (see [5–9] for early discussions and [10,11] for reviews). If the DM particles carry a quantum number related to baryon number, the same mechanism might be responsible for the generation of their relic density via their asymmetry.

At first sight, sneutrinos are promising candidates for asymmetric DM (ADM) [12–17]. They carry a conserved quantum number, lepton number, such that they can share the asymmetry of charged leptons through processes which were in equilibrium in the early hot universe. Then, their asymmetry can have become related to the baryon asymmetry through sphaleron processes [19,20]. However, although the large annihilation cross section between sneutrinos and anti-sneutrinos is a good feature allowing for ADM, their self-annihilation (sneutrino–sneutrino or anti-sneutrino–anti-sneutrino annihilation) cross sections are also typically large, destroying the asymmetry [21].

1Higgsinos could also be ADM, but only if a number of strong constraints is satisfied, see [18].
However, non-zero neutrino masses suggest the existence of right-handed neutrinos (and sneutrinos). In the past years, a variety of models aiming at the explanation of neutrino masses have been proposed (see, for example, [22,23] and references therein). They can be categorized in two main classes, those which employ Majorana mass terms for right-handed neutrinos and those employing Dirac mass terms only. The former allow for various versions of seesaw mechanisms, amongst others the inverse seesaw which allows for electroweak scale right-handed neutrinos. The common characteristic of these models is the violation of lepton number by a Majorana mass term. Dirac neutrino masses are less well studied, though not less motivated. The simplest way to obtain Dirac masses for neutrinos is the introduction of Yukawa couplings to Higgs bosons, but with unnatural small Yukawa coupling constants. A more elegant way is the introduction of an additional Higgs field which couples only to right-handed neutrinos. Then, the smallness of neutrino masses is no longer due to small values of coupling constants, but can be due to a small vacuum expectation value (vev) of the new Higgs field.

The presence of right-handed sneutrinos opens new possibilities for sneutrino DM: right-handed sneutrinos with a small left-handed component may have at the same time a large pair annihilation cross section, but a negligible self-annihilation cross section. However, in scenarios as seesaw models which do not conserve lepton number, an asymmetry of sneutrinos is difficult to maintain due to oscillations between sneutrinos and anti-sneutrinos [24–26] (see also [27]). In the Dirac case with small Yukawa couplings, asymmetric DM faces the following difficulties: First, the annihilation cross section, proportional to the small couplings, is not adequate to eliminate the symmetric part of the DM, resulting in a large relic density unrelated to the asymmetry. Second, such small Yukawa couplings keep the right-handed neutrinos and sneutrinos out of equilibrium in the early universe and, as a result, the asymmetry of the sneutrinos was never related to the baryonic asymmetry. However, if the small neutrino masses originate from the small vev of an additional Higgs field (but with large Yukawa couplings), these difficulties are solved.

Such scenarios, often denoted as neutrinophilic Higgs doublet models, appeared first in [28,29] (for an earlier approach, but with Majorana neutrinos, see [30]). In these models, a $Z_2$ symmetry is spontaneously broken generating a small vev of the new Higgs scalar. This mechanism results also in a very light scalar with mass of the order of eV. Such light scalars have been ruled out [31,32] by astrophysical arguments. However, the $Z_2$ symmetry can be replaced by a global $U(1)$ symmetry in order to forbid Majorana masses, but which is broken explicitly so that a very light scalar is avoided [33]. The $U(1)$ symmetry makes very small explicit breaking terms (see below) more natural. The LHC phenomenology of this model is studied in [34], while in [35] a supersymmetric variant based on the MSSM and its phenomenology are examined. The additional Higgs doublets still allow the SUSY version to be embedded into a grand unified symmetry as verified in [36]. Furthermore, the additional Higgs doublets do not spoil proton stability since they couple only to leptons. In [37], the Higgs potential is studied for both SUSY and non-SUSY models. Scenarios for leptogenesis with neutrinophilic Higgs are discussed in [38–40]. Finally, the sneutrino of the SUSY
model of [35] has been used as DM candidate in [41, 42]. In particular, the possibility of ADM is also considered in [42], but with a trilinear soft coupling of the order of several TeV and a (related) very large annihilation cross section into monochromatic photons.

In this paper we consider the Next-to-Minimal Supersymmetric Standard Model (NMSSM) extended by a pair of neutrinophilic Higgs doublets and three generations of right-handed neutrino superfields. The NMSSM provides a natural solution of the $\mu$-problem of the MSSM by introducing a gauge singlet superfield $S$ (for reviews, see [43, 44]). In the NMSSM, the Higgs mass term $\mu$ in the superpotential of the MSSM is replaced by the term $\lambda S H_u \cdot H_d$. Otherwise the singlet plays no particular role in our model (its vev could be replaced by a constant dimensionful parameter), but due to the additional coupling $\lambda$ the lightest CP-even Higgs is naturally heavier than in the MSSM [45–48]. Therefore, a SM-like Higgs mass of $\sim 125$ GeV is much easier to explain [49, 50].

We are going to explore whether and under which circumstances this model can accommodate right-handed sneutrinos as ADM. We find that this is indeed possible under certain conditions. First, we note that the ordinary Higgs sector of the NMSSM is not affected by the introduction of the neutrophilic Higgses (henceforth $\nu$-Higgses). The scalar $\nu$-Higgses, however, have to be relatively heavy of $O(1)$ TeV such that the additional degrees of freedom of the Dirac neutrinos do not lead to $^4$He overabundance through their contribution to the expansion rate of the universe during big bang nucleosynthesis (BBN). On the other hand, light $\nu$-higgsinos are required for a large sneutrino–anti-sneutrino pair annihilation cross section which is necessary for the sneutrino relic density to be determined by its asymmetry.

In the next section we present the model and discuss constraints from lepton number violation and BBN. In Sec. 3 we explore the possibility for right-handed sneutrinos as asymmetric DM. In particular, in Sec. 3.1 we examine the connection between the sneutrino and baryon asymmetry via sphaleron processes, in Sec. 3.2 we study conditions from oscillations, self and pair annihilations, and in Sec. 3.3 we discuss possible DM signals and constraints from DM detection. Finally, we summarize our results in Sec. 4.

2 The model

We extend the NMSSM by three right-handed neutrino superfields $\tilde{\nu}_R$ and a pair of new Higgs doublets $\tilde{H}_u^\nu$ and $\tilde{H}_d^\nu$. These fields are charged under a new global $U(1)$ symmetry with charges $-1$, $+1$ and $-1$, respectively, while the usual NMSSM superfields remain uncharged. The superpotential is written as

$$W = W^{NMSSM} + y_\nu \tilde{L} \cdot \tilde{H}_u^\nu \tilde{\nu}_R + \lambda_\nu \tilde{S} \tilde{H}_u^\nu \cdot \tilde{H}_d^\nu, \quad (1)$$

where the Yukawa coupling $y_\nu$ and the superfields $\tilde{L}$ and $\tilde{\nu}_R$ should be understood as matrix and vectors, respectively, in flavor space. The corresponding soft SUSY
breaking masses and couplings are

\[-\mathcal{L}_{\text{soft}} = -\mathcal{L}_{\text{soft}}^{\text{NMSSM}} + m_{H_u}^2 |H_u^\nu|^2 + m_{H_d}^2 |H_d^\nu|^2 + m_{\nu_R}^2 |\nu_R|^2 + y_{\nu} A_\nu L \cdot H_u^\nu \nu_R^c + \lambda_{\nu} A_{\nu} S H_u^\nu \cdot H_d^\nu.\]

\(W^{\text{NMSSM}}\) and \(\mathcal{L}_{\text{soft}}^{\text{NMSSM}}\) are the superpotential and the soft terms of the \(Z_3\)-invariant NMSSM (see, e.g., [34]), respectively.

The new \(U(1)\) symmetry needs to be broken by the vev of the \(\nu\)-Higgs in order to give masses to the neutrinos. To this end we add to the Lagrangian (2) the additional soft terms

\[A_{\nu} S H_u \cdot H_d^\nu + A_{\nu} S H_u^\nu \cdot H_d.\]

These two terms do not correspond to terms in the superpotential. Since they break the \(U(1)\) explicitly, it is natural in the \(t\) Hooft sense for the trilinear couplings \(A_{\nu}\) to assume small values. Such small values can be obtained through higher dimensional operators involving SUSY and \(U(1)\) symmetry breaking spurion fields [35]. For instance, introducing a superfield \(\hat{X}\) with charge \(-1/2\) under \(U(1)\) and with \(\langle X \rangle = \theta^2 F + \sqrt{F}\) (see, e.g., [31] for similar mechanisms), a trilinear soft term can originate from the operator \(\frac{1}{\sqrt{M_{\text{Pl}}}} \hat{X} \hat{S} \hat{H}_u^\nu \cdot \hat{H}_d \|_F \sim \frac{F^{3/2}}{M_{\text{Pl}}} \hat{S} \hat{H}_u^\nu \cdot \hat{H}_d\). If \(F = m_1^2\) with \(m_1 \simeq \sqrt{\nu M_{\text{Pl}}}\) an intermediate scale where supersymmetry is broken, then \(\frac{F^{3/2}}{M_{\text{Pl}}} \sim 10^{-7}\) GeV, while the corresponding term in the superpotential is suppressed by several orders of magnitude.

The resulting vevs for the \(H_u^\nu, H_d^\nu\) fields have the form [30]

\[v_u^\nu \simeq \frac{A_{\nu}s}{m_{H_u}^\nu} v_d \quad \text{and} \quad v_d^\nu \simeq \frac{A_{\nu}s}{m_{H_d}^\nu} v_u,\]

respectively. Taking \(A_{\nu}s \simeq A_{\nu}s \sim 10^{-5}\) GeV\(^2\) and assuming soft masses \(m_{H_u}^\nu \sim m_{H_d}^\nu \sim \mathcal{O}(1)\) TeV (see below), then \(v_u^\nu \simeq v_d^\nu \sim \text{eV}\). Hence, the first extra term in the superpotential (1) will generate Dirac neutrino masses of the correct order for \(y_{\nu} \sim \mathcal{O}(1)\) [28,33].

The mass squared matrix of the sneutrinos, neglecting flavor indices, reads in the basis \((\tilde{\nu}_L, \nu_R)\)

\[M_\nu^2 = \begin{pmatrix} y_{\nu}^2 v_u^\nu + \frac{1}{2} g^2 (v_d^2 - v_u^2) + m_{\nu_L}^2 & y_{\nu} v_u^\nu (\lambda_{\nu} s + A_{\nu}) \\ y_{\nu}^2 v_u^\nu (\lambda_{\nu} s + A_{\nu}) & y_{\nu}^2 v_u^\nu (\lambda_{\nu} s + A_{\nu}) + m_{\nu_R}^2 \end{pmatrix}.\]

Taking into account the small value of \(v_u^\nu\), this matrix can be approximated by the diagonal form

\[M_\nu^2 \simeq \text{diag} \left[ \frac{1}{2} g^2 (v_d^2 - v_u^2) + m_{\nu_L}^2, m_{\nu_R}^2 \right].\]

We note that the mixing between the different sneutrino flavors in the right-handed sector is small, since it is proportional to the vev of the \(\nu\)-Higgs provided that \(m_{\nu_R} \gg v_u^\nu\) and flavour diagonal.

The \(\nu\)-Higgses form two nearly degenerate SU(2) doublets. (Since \(U(1)\) is not spontaneously broken, there are no Goldstone bosons.) These additional fields mix
very weakly with the standard Higgs fields due to their small vevs; in the following we will consider the new Higgs fields completely unmixed.

The mass matrices in the neutral sector are in the basis \((H_u^\nu, H_d^\nu)\) \[52\]

\[
M_{H^\nu}^2 = \begin{pmatrix}
\lambda_{\nu}^2 s^2 - \frac{1}{2} g^2 (v_d^2 - v_u^2) + m_{H_u^\nu}^2 & \pm \lambda_{\nu} (\lambda v_u v_d - \kappa s^2 + A_{\lambda\nu} s) \\
\lambda_{\nu}^2 s^2 + \frac{1}{2} g^2 (v_d^2 - v_u^2) + m_{H_d^\nu}^2 & \lambda_{\nu}^2 s^2 - \frac{1}{2} g^2 (v_d^2 - v_u^2) + m_{H_d^\nu}^2
\end{pmatrix}
\]

(6)

with plus (minus) signs in the off-diagonal element for the scalar (pseudoscalar), and

\[
M_{H^\nu^+}^2 = \begin{pmatrix}
\lambda_{\nu}^2 s^2 + \frac{1}{2} g^2 (v_d^2 - v_u^2) \cos 2\theta_W + m_{H_u^\nu}^2 & \lambda_{\nu} (\lambda v_u v_d - \kappa s^2 + A_{\lambda\nu} s) \\
\lambda_{\nu}^2 s^2 - \frac{1}{2} g^2 (v_d^2 - v_u^2) \cos 2\theta_W + m_{H_d^\nu}^2 & \lambda_{\nu}^2 s^2 + \frac{1}{2} g^2 (v_d^2 - v_u^2)
\end{pmatrix}
\]

(7)
in the charged \(\nu\)-Higgs sector. The neutral and charged \(\nu\)-higgsinos, forming Dirac fermions with masses \(\mu' = \lambda_{\nu} s\), are also practically unmixed with the neutralinos and the charginos of the NMSSM.

### 2.1 Constraints from lepton flavour violation and BBN

The charged Higgs \(H^\nu^+\) mediates the decay of the muon at one loop with a branching ratio \[53\]

\[
\text{BR} \left( \mu \to e\gamma \right) = \frac{\alpha_{\text{EM}}}{24\pi} \left( \frac{v}{v_{\mu}} m_{H^\nu^+} \right)^4 \left| \sum_j m_j^2 U_{ej} U_{\mu j} \right|^2,
\]

(8)

where \(U\) is the Pontecorvo–Maki–Nakagawa–Sakata (PMNS) matrix defined by \(|\nu_l\rangle = \sum_j U_{lj} |\nu_j\rangle\), with \(l = e, \mu, \tau\) and \(j = 1, 2, 3\) corresponding to the three mass eigenstates. The unitarity of the PMNS matrix allows to replace the sum in eq. \(8\) by

\[
\sum_j m_j^2 U_{ej}^* U_{\mu j} = -\Delta m_{21}^2 U_{e1}^* U_{\mu 1} + \Delta m_{32}^2 U_{e3}^* U_{\mu 3},
\]

where the mass squared differences are defined by \(\Delta m_{ij}^2 \equiv m_i^2 - m_j^2\) and \(m_i\) are the neutrino mass eigenvalues. Using the upper 90\% C.L. limit

\[
\text{BR} \left( \mu \to e\gamma \right) < 5.7 \times 10^{-13}
\]

(9)

from the MEG experiment \[54\] gives a lower bound on the charged \(\nu\)-Higgs mass,

\[
m_{H^\nu^+} \gtrsim \left( \frac{1\text{eV}}{v_{\nu}} \right) 300\text{ GeV},
\]

(10)

where we have used the standard values for \(\Delta m_{21}^2, \Delta m_{32}^2\) given in \[55\] and the elements of the PMNS matrix.

The additional degrees of freedom of the Dirac neutrinos contribute to the energy density and therefore to the expansion rate of the universe during Big Bang Nucleosynthesis (BBN). The abundance of \(^4\text{He}\) emerging from BBN depends on the Hubble expansion rate when processes like \(e^- + p \leftrightarrow n + \nu_e\) and \(e^+ + n \leftrightarrow p + \bar{\nu}_e\) were in equilibrium, since practically all the remaining neutrons (after these processes went out of equilibrium) were incorporated in helium nuclei. The larger the Hubble rate,
the faster (at a higher temperature) they are going out of equilibrium, resulting in a larger abundance of neutrons\(^2\).

In the epoch just before nucleosynthesis photons, electrons and left-handed neutrinos were in equilibrium\(^3\) at a common temperature \(T_{\gamma,n}\) (henceforth, the subscript \(n\) of temperatures will denote the epoch just before nucleosynthesis). Right-handed neutrinos remained in equilibrium as long as the processes \(\nu_R + \nu_R \leftrightarrow l + l\) (with \(l\) a charged lepton), mediated by the charged \(\nu\)-Higgses, were fast enough. However, even when the right-handed neutrinos go out of equilibrium at a temperature \(T_{R,d}\) (where the subscript \(d\) stands for decoupling), they continue to contribute to the total energy density of the universe with their own temperature \(T_R\) that is redshifting.

Usually, the helium abundance is parametrized by the effective number of degrees of freedom \(N_{\text{eff}}\) during BBN. Recently, Planck constrained this quantity to \(N_{\text{eff}} = 3.30 \pm 0.27\) at 68\% C.L.\(^3\). In the following we determine the lowest temperature \(T_{R,d}\) at which the right-handed neutrinos can decouple without \(N_{\text{eff}}\) exceeding the above limit, and subsequently we derive the necessary condition on the \(\nu_R + \nu_R \leftrightarrow l + l\) rate for this to occur.

Writing the energy density in the form
\[
\rho_n = \frac{\pi^2}{45} \left[ g_\gamma + \frac{7}{8} (g_e + N_{\text{eff}} g_\nu) \right] T_{\gamma,n}^4, \tag{11}
\]
\(N_{\text{eff}}\) is defined as \(N_{\text{eff}} = n_L + n_R \left( \frac{T_{R,n}}{T_{\gamma,n}} \right)^4\) with \(n_L\) (\(n_R\)) the number of left- (right-) handed neutrino generations. Taking \(n_L = n_R = 3\) we write \(N_{\text{eff}} = 3 + \Delta N_\nu\) with
\[
\Delta N_\nu = 3 \left( \frac{T_{R,n}}{T_{\gamma,n}} \right)^4. \tag{12}
\]

Applying entropy conservation separately for the decoupled species and the thermal bath\(^5\) one finds
\[
\frac{T_{R,n}}{T_{\gamma,n}} = \left( \frac{43}{9g(T_{R,d})} \right)^{1/3},
\]
where \(g(T_{R,d})\) is the number of degrees of freedom when the right-handed neutrinos decouple. Substituting the last relation into \((12)\), we obtain for the relation between the maximally allowed value of \(\Delta N_\nu^{\text{max}}\) and the temperature at which the right-handed neutrinos went out of equilibrium
\[
g(T_{d,R}) \geq \frac{43}{4} \left( \frac{3}{\Delta N_\nu^{\text{max}}} \right)^{3/4}. \tag{13}
\]
For \(\Delta N_\nu^{\text{max}} \lesssim 0.57\) at 1\(\sigma\), \(g(T_{R,d}) \geq 37.35\). This means that decoupling should have occurred before the quark-hadron phase transition when \(g = 51.25\) (just after the transition the number of degrees of freedom was \(g = 17.25\)). Assuming that the

\(^2\)The equilibrium density of neutrons falls as \(n_{eq}^n \sim \exp(-\Delta m T_{\gamma,n}) n_{eq}^p\) with \(\Delta m \equiv m_n - m_p\) the mass difference between neutron and proton.

\(^3\)The neutrinos decouple from the thermal plasma at a temperature \(T_d \gtrsim 1\text{ MeV}\) when \(H\) becomes larger than the rate of the processes \(\nu + \bar{\nu} \rightarrow e^+ + e^-\). Nevertheless, even after their decoupling but before the decoupling of the electrons, their temperature is the same as the one of photons since both are decreasing at the same rate.
QCD confinement temperature is roughly $T_c \simeq 200 \text{ MeV}$ [57] leads to the inequality $T_{R,d} \gtrsim 200 \text{ MeV}$.

Taking into account the approximate decoupling condition $n(T_d)\langle \sigma v \rangle(T_d) = H(T_d)$, one finds that the ratio of the decoupling temperatures of right- and left-handed neutrinos is

$$\left( \frac{T_{R,d}}{T_{L,d}} \right)^3 = \sqrt{\frac{g(T_{L,d})}{g(T_{R,d})}} \frac{\sigma_L}{\sigma_R},$$

with $\sigma_L$ and $\sigma_R$ the cross sections of the processes that were keeping left- and right-handed neutrinos, respectively, in equilibrium. Using $\frac{\sigma_L}{\sigma_R} = \left( \frac{2\sqrt{2}m_{H^{\nu+}}}{y_{\nu}^L y_{\nu}^R} \right)^4$ [33], where $U_{li}$ are again the elements of the PMNS mixing matrix, leads to the following bound on the charged Higgs mass and the couplings $y_{\nu}^L$

$$\frac{m_{H^{\nu+}}}{g_{\nu}^L} \gtrsim 3 \text{ TeV}.$$  

As we will explain later, the couplings $y_{\nu}^L$ cannot be very small, and as a consequence $m_{H^{\nu+}}$ has to be relatively large.

### 3 Right-handed sneutrinos as ADM

In the following we study more closely the rôle of right-handed neutrinos as ADM. We will use the notation $N \equiv \tilde{\nu}_{R1}$ with the index 1 denoting the lightest among the three right-handed neutrinos. Its mass $m_N$ is essentially its soft Susy breaking mass, and we safely assume that it is a pure state since its left-handed component is negligibly small.

#### 3.1 Asymmetry from sphaleron processes and the ADM mass

Since the sneutrinos carry a conserved charge (lepton number), it is possible that their relic density is not determined by the thermal mechanism but by their asymmetry. The asymmetry was related to the baryon asymmetry through equilibrium processes in the early universe. These allow to estimate the relation between the two asymmetries and, ultimately, to determine the mass range of the (right-handed sneutrino) DM that will provide the correct abundance.

If the $N, N^*$ annihilation is strong enough such that the less frequent species has been completely eliminated, the remaining abundance is the product of the charge density $\eta_N = |n_N - n_{N^*}|$ times its mass $m_N$ ($n$ denotes the number density). The relation between the DM relic density $\Omega_N$ and the baryonic relic density $\Omega_b$ is

$$\Omega_N = \frac{\eta_N m_N}{B m_p} \Omega_b,$$

where $m_p$ is the proton mass, which gives the desired result if $\eta_N$ is of same order of magnitude as the baryon charge density $B$. 

7
The charge density of a particle $X$ in kinetic equilibrium as a function of the temperature can be written as

$$\eta_X(T) = \frac{T^3}{6} g_X k(x) \mu_X T,$$

(17)

where we have assumed that $\mu_X / T \ll 1$, and $\mu_X$ is the chemical potential of the species $X$. $g_X$ is the number of internal degrees of freedom of the particle $X$, we defined $x \equiv \frac{m_X}{T}$, and

$$k(x) = \frac{6}{\pi^2} \int_x^\infty dy \frac{y^2 - x^2}{\sqrt{y^2 - x^2}} e^y \left(e^y \pm 1\right)^2.$$

(18)

In the above integral, the plus (minus) sign holds for fermions (bosons). In the ultra-relativistic limit $x \ll 1$, $k$ takes the values 1 for fermions and 2 for bosons, while in the opposite limit $x \gg 1$ it vanishes in both cases.

A sneutrino asymmetry can originate from primordial asymmetries in the baryonic or leptonic sectors. Although we will be agnostic about the exact mechanism that created these primordial asymmetries, the fact that certain processes were in equilibrium in the early universe can be used to relate $\eta_N$ to the baryon asymmetry $B$. We note that common mechanisms for thermal leptogenesis would not work in the present framework since the violation of lepton number is far too small (see the next section). However, other known mechanisms are possible, such as the Affleck-Dine mechanism [58].

In the absence of lepton number violating processes other than electroweak sphalerons, $\sum_{i=1}^3 (B/3 - L_i)$ is conserved. The relatively large Yukawa coupling constants of the neutrinos assure not only the equilibrium of the right-handed neutrinos with the thermal bath in the early universe, but also rapid flavor changing processes. As a result, lepton flavor equilibrium had been established and $B - L = \sum_{i=1}^3 (B/3 - L_i)$ is conserved. However, the sphaleron processes were still violating $B + L$. We are going to consider two cases [19]. In the first case we will assume that sphaleron processes were rapid only above the electroweak phase transition (EWPT), e.g., if the EWPT was strongly first order. In the second case, we will allow the sphaleron processes to violate $B + L$ also below the transition, until they went out of equilibrium because of the expansion of the universe.

In the first case we proceed along the lines of [20], where one can find a complete list of the equilibrium reactions in the MSSM and the relations between the chemical potentials. The reactions specific to the present model lead to the following equilibrium relations which have to be added to this list:

$$\mu_{L_i} + \mu_H = \mu_{\nu_i}, \quad \mu_H = \mu_{H_u} = \mu_{H_d},$$
$$\mu_{\tilde{L}_i} + \mu_{\tilde{H}} = \mu_{\tilde{\nu}_i}, \quad \mu_{\tilde{L}_i} + \mu_{\tilde{H}} = \mu_{\nu_i}, \quad \mu_{L_i} + \mu_{\tilde{H}} = \mu_{\tilde{\nu}_i},$$

(19)

where we have used the notation of [20], i.e. $\mu_{L_i}$ is the chemical potential of the left-handed leptons, $i$ is the flavor index, $\mu_{\nu_i}$ is the chemical potential of the right-handed neutrinos and tilde stands for the supersymmetric particles. The sneutrinos share the chemical potential with the neutrinos through the equilibrium of processes such as those of [Fig. 1]
Figure 1: Annihilation diagrams of right-handed sneutrinos $N$, $N^*$ into neutrinos and charged leptons.

Eliminating the chemical potentials using the sphaleron equilibrium relation and the fact that the total hypercharge of the universe vanishes, we can calculate the baryon charge $B$ and the DM leptonic charge $\eta_N$ as functions of the conserved difference $B - L$. We assume all the supersymmetric particles except the sleptons much heavier than the EWPT critical temperature $T_c$ and take massless SM particles. First, assuming light right-handed sneutrinos and masses $\sim 2T_c$ for the other sleptons, $B$ and $\eta_N$ are related by

$$B \simeq 0.14 \,(B - L) \quad \text{and} \quad \eta_N \simeq 0.10 \,(B - L), \quad \text{(light sleptons)} \quad (20a)$$

while for large slepton masses

$$B \simeq 0.18 \,(B - L) \quad \text{and} \quad \eta_N \simeq 0.12 \,(B - L). \quad \text{(heavy sleptons)} \quad (20b)$$

The DM mass, using (16) and $\Omega_N/\Omega_b \simeq 5.44 \, [3]$, has to be $m_N \sim 7.1 - 7.6 \, \text{GeV}$ (the smaller value corresponding to light sleptons).

In case the process induced by electroweak sphalerons were rapid also below the EWPT, the relations between the chemical potential are altered. First, due to the vacuum condensate of the neutral Higgs bosons, their chemical potentials have to vanish. However, the total hypercharge has no longer to be zero since $SU(2)_L$ has been broken, resulting in a non-zero chemical potential for the $W$ bosons. Generalizing the SM equilibrium processes of [19] we find, considering all the supersymmetric particles (except for the right-handed sneutrinos) as heavy,

$$B \simeq 0.18 \,(B - L) \quad \text{and} \quad \eta_N \simeq 0.10 \,(B - L). \quad \text{(heavy SUSY particles)} \quad (21a)$$

The resulting DM mass in this case has to be $m_N \sim 9.2 \, \text{GeV}$. However, allowing the left-handed sneutrinos to be light (with mass around the temperature at which the sphaleron processes went out of equilibrium), the value for the ratio $B/\eta_N$ becomes maximal:

$$B \simeq 0.31 \,(B - L) \quad \text{and} \quad \eta_N \simeq 0.07 \,(B - L). \quad \text{(light LH sneutrinos)} \quad (21b)$$

In this case the DM mass has to be larger, $m_N \simeq 23 \, \text{GeV}$.

Summarizing, depending on the sparticle spectrum and the nature of the EWPT, the DM mass can roughly be in the range

$$m_N \sim 7 \, \text{GeV} - 23 \, \text{GeV}.$$

(22)
The lowest value corresponds to light sleptons and a first order EWPT that terminated the sphaleron processes, while for the highest value the sphaleron processes have to continue to be in equilibrium for a short time after the EWPT and the left-handed sneutrinos have to be relatively light.

### 3.2 Constraints from oscillations, self and pair annihilation

In order that the current DM density to be determined by its asymmetry, a number of conditions have to be fulfilled. First, the annihilation of DM particles with antiparticles has to be strong enough so that one of them is completely depleted. However, in many cases, it is possible for a particle to oscillate into its antiparticle and vice versa. These oscillations, if rapid enough, might lead to a continuous repopulation of the depleted particles. As a result, however strong the pair annihilation cross section is, the antiparticles (or the particles) are never exhausted and, finally, the thermal mechanism is responsible for the relic density. Furthermore, if self-annihilation\footnote{the particle–particle or the antiparticle–antiparticle annihilation.} of DM particles occurs before the DM particles become non-relativistic, their asymmetry decreases rapidly due to this annihilation. If the self-annihilation does not freeze-out sufficiently fast, the thermal mechanism takes over again since there is no asymmetry left after the particle–antiparticle annihilation freeze-out. We will show that the sneutrino DM considered here can fulfill all these criteria for a successful asymmetric DM candidate.

Quantum mechanical oscillations occur between \( N \) and \( N^* \) if they are not the mass eigenstates. Then the rate of \( N - N^* \) conversion is approximately given by the mass difference \( \delta m \) of the two eigenstates. The conversion starts to be significant only at times larger than \( \delta m^{-1} \) or, expressed in terms of the temperature \( T \) of the universe, for \( T \lesssim T(\delta m) \) given by \[23\]

\[
T(\delta m) \sim \left( \frac{g_s^{1/2}}{h_{\text{eff}}} \sqrt{\frac{45}{4\pi^3}} M_{\text{Pl}} \delta m \right)^{1/2},
\]

where \( M_{\text{Pl}} \) is the Planck mass and \( g_s \) and \( h_{\text{eff}} \) are effective degrees of freedom (see, e.g., \[21\] for exact definitions).

A mass split appears if there exists a lepton number violating Majorana mass term \( m_M \); if \( m_M \ll m_D \) (\( m_D \) is the Dirac mass), the mass split can be written as \( \delta m \simeq m_M^2/m_D^2 \). The operator \( \frac{1}{M_{\text{Pl}}} |X^4SN^2|_F \), with \( X \) the superfield spurion whose vev brakes the \( U(1) \) symmetry, induces a tiny Majorana mass squared of the order \( m_M^2 \simeq 10^{-32} \text{GeV}^2 \), which corresponds to a mass difference \( \delta m \sim 10^{-33} \text{GeV} \) for \( m_D \sim 10 \text{GeV} \). With such a small value for the Majorana mass, the oscillations start very late in the history of the universe (see \[23\]), much later than the DM freeze-out (at \( T \sim 1 \text{GeV} \)), and do not affect the final DM density.

Upper bounds on the self-annihilation cross section have been derived in \[21\]. If the cross section does not obey these bounds, the asymmetry falls rapidly. In order that at least 90% of the asymmetry survives, the decoupling of self-annihilation should happen
before $x \equiv m_N/T \sim 5$ (we recall that the decoupling for WIMPs occurs at $x \sim 20 – 30$). However, in our case the possible annihilation of right-handed sneutrinos into two neutrinos through t- or u-channel exchange of neutral $\nu$-higgsinos is impossible due to the Dirac nature of the $\nu$-higgsinos. Furthermore, the left-handed components in $N$ and $N^*$ are sufficiently small, since they are induced only by the off-diagonal element of the mass matrix (1) and hence many orders of magnitude below the bound of [21]. Consequently, the self-annihilation cross sections of $N$ or $N^*$ are sufficiently small.

Now that we have shown that the asymmetry does not get destroyed by oscillations or self annihilations, the condition that remains to be satisfied is a sufficiently strong $N, N^*$ pair annihilation so that only the asymmetry survives as relic density. The dominant annihilation channels of right-handed sneutrinos are the annihilation into neutrinos and charged leptons (Fig. 1). The former proceeds through a t-channel neutral $\nu$-higgsino exchange, the latter by charged $\nu$-higgsino exchange. The thermal average of the cross section of these processes times velocity can be written as (to leading order in $x^{-1}$)

$$\langle \sigma v \rangle \simeq f \frac{y_\nu^4}{8\pi} \frac{m_N^2}{(m_N^2 + \mu'^2)^2} x^{-1},$$  

(24)

where the factor $f = 18$ counts the number of final states (9 neutrinos and 9 charged leptons) and we have assumed a common value $y_\nu$ for the coupling constants $y_{\nu i}^i$. The s-wave contribution is helicity suppressed and can be neglected (see also [59]).

In the usual symmetric DM case, the thermally averaged cross section during freeze out has to be of the order of the so-called thermal cross section, roughly given by [60]

$$\langle \sigma v \rangle_{\text{th}} \simeq \frac{3 \cdot 10^{-27}}{\Omega_{\text{DM}} h^2} \text{ cm}^3 \text{s}^{-1} \simeq 3 \cdot 10^{-26} \text{ cm}^3 \text{s}^{-1}$$  

(25)

from considerations using entropy conservation. In the asymmetric DM scenarios, the pair annihilation cross section must be equal to or larger than the thermal cross section; even if the cross section is much larger than its thermal value, the final density remains constant since annihilations become impossible due to the lack of $N$ or $N^*$. Examining eq. (24) at fixed $m_N$, the cross section decreases with increasing mass $\mu'$ of the $\nu$-higgsino.

[Fig. 2] shows the maximal value of the $\nu$-higgsino mass as function of the coupling constant $y_\nu$ for a sufficiently large annihilation cross section. The corresponding minimal allowed mass of the lightest scalar charged $\nu$-Higgs (from the condition that $N_{\text{eff}}$ is within the 1σ region determined by Planck) is shown along the upper axis. For a light right-handed sneutrino mass (7 GeV) the coupling constant has to be relatively large, $y_\nu \gtrsim 0.6$, which requires $m_{H^\pm} \gtrsim 1.8 \text{ TeV}$. Smaller values of $y_\nu$ (and lower bounds on the scalar $\nu$-Higgs mass) are allowed for heavier DM. For $m_N = 23 \text{ GeV}$, the smallest allowed value for the $\nu$-Higgs mass is $\sim 1 \text{ TeV}$.

We see that the right-handed sneutrino can have a relic density determined by its asymmetry if the $\nu$-higgsino is relatively light, while the scalar $\nu$-Higgs should be heavy. In terms of the coupling $\lambda_\nu$ in the NMSSM Lagrangian and the singlet vev $s$, the $\nu$-higgsino mass is given by $\mu' = \lambda_\nu s$ and hence small for small $\lambda_\nu$, whereas a heavy charged $\nu$-Higgs requires a large soft SUSY breaking mass.
Figure 2: The largest allowed $\nu$-higgsino mass as function of the coupling $y_{\nu}$ such that right-handed sneutrinos with mass 7 GeV (red –upper– line), 15 GeV (dashed line) or 23 GeV (blue –lower– line) have a large enough pair annihilation cross section such that their relic density is determined by their asymmetry. The corresponding lower limit on the lightest $\nu$-Higgs mass, derived from eq. (15), is indicated along the upper axis. The shaded area is excluded by chargino searches at LEP (e.g. [61]).

3.3 ADM Detection: prospects and constraints

Upper bounds on the ADM-nucleon scattering cross section originate from both the direct detection and observations of old neutron stars. Concerning the latter, if the cross section is too large, the accumulation of asymmetric DM inside the neutron stars can form a black hole which would potentially destroy the star. This is a specific feature of asymmetric DM, since in the common symmetric case the annihilation of DM with anti-DM prevents the accumulation. For bosonic asymmetric DM in the mass range $5 \text{ GeV} \lesssim m_{DM} \lesssim 16 \text{ GeV}$, nucleon–DM cross sections $\sigma \gtrsim 10^{-50}$ are excluded [62, 63]. However, the value of the cross section depends on the parameter space, particularly on the value of $A_{\nu}$ (see [41]), while for DM heavier than $\sim 16 \text{ GeV}$ there is no limit due to Hawking evaporation [62], letting a completely unconstrained mass range ($16 \text{ GeV} \lesssim m_{DM} \lesssim 23 \text{ GeV}$) for the ADM of this scenario.

Concerning direct detection, since right-handed sneutrinos couple only to neutrinophilic Higgses, there are no tree-level contributions to the scattering cross section.
of $N$ off nuclei. However, as it was pointed out in [41], the contribution of $Z$ exchange induced at one loop (with left-handed sneutrinos and $\nu$-Higgses running on the loop) may be significant. The value obtained in [41] ($\sim 10^{-45}$ cm$^3$s$^{-1}$), assuming a relatively large value for the trilinear soft coupling $A_\nu$, is at the lower bound of the current experimental direct detection sensitivity for a DM mass around 100 GeV. However, for the mass range considered here ($O(10)$ GeV), the upper limits on the scalar scattering cross section are much higher.

Concerning indirect detection, pure asymmetric DM does not give rise to detectable signals due to the absence of either DM particles or antiparticles. (The self-annihilation cross section is required to be too small to generate measurable signals.) However, the operators which break the $U(1)$ symmetry might also induce a very small mass difference $\delta m$ among the sneutrino and anti-sneutrino eigenstates. Even though the induced sneutrino – anti-sneutrino oscillations are slow enough in order not to affect the relic density, it may have led to the repopulation of the exhausted species if $\delta m^{-1}$ is smaller than the current age of the universe. This would be the case, e.g., for the value $\delta m \sim 10^{-33}$ GeV obtained in the scenario sketched below (23).

In case the exhausted species has been regenerated, the same $N, N^*$ annihilation processes (Fig. 1) that occurred in the early universe may happen today in galactic regions of large DM density, giving rise to leptonic charged cosmic rays and $\gamma$-rays. However, assuming that the excess of positrons observed amongst others by AMS-02 [64] originates from astrophysical sources, it constitutes an insurmountable background making the distinction of a potential DM signal from charged leptonic rays difficult. Concerning the diffuse photon radiation, we recall that the s-wave annihilation of $N, N^*$ is helicity suppressed. The low present-day velocity of DM particles leads to a low $\sigma v$, evading the bounds set by the Fermi collaboration [65]. Finally, as pointed out in [42], $N, N^*$ annihilation through a box loop with sleptons and charged $\nu$-Higgses can give rise to a monochromatic photon line with a large cross section proportional to $\left(\frac{M_{N^*}}{M_{N}}\right)^4$. The Fermi bound for a DM mass of $\sim 10$ GeV is quite severe, $\langle \sigma v \rangle_{\gamma\gamma} \lesssim 5 \cdot 10^{-29}$ cm$^3$s$^{-1}$ [66]. However, taking $A_\nu$ of the order of the EW scale ($\sim 100$ GeV), this bound is satisfied since $\langle \sigma v \rangle_{\gamma\gamma} \lesssim 10^{-29}$ cm$^3$s$^{-1}$.

4 Summary and outlook

In this paper we have presented an extension of the NMSSM introducing an additional pair of Higgs doublets with small vevs, explaining the smallness of neutrino masses and, at the same time, the present day coincidence of DM and baryon densities. The additional Higgses and the right-handed neutrinos are charged under a new $U(1)$ symmetry which is explicitly broken by soft SUSY breaking terms. This symmetry forces the new Higgses to couple in the superpotential only to right-handed neutrinos (so-called neutrinoophilic Higgses). The neutrinos have Dirac masses which are generated dynamically by the neutrinoophilic Higgs vev and hence naturally small.

We have shown that the right-handed sneutrinos can carry an asymmetry related to
the baryon asymmetry due to their conserved lepton number and equilibrium processes in the early universe. They can maintain their asymmetry at least until the freeze-out of sneutrino–anti-sneutrino annihilations. Therefore their relic density is determined by their asymmetry and of the correct value if their mass is $O(10)$ GeV, provided that the coupling constant $\lambda'$ is small compared to $\lambda$. However, the bound on the relativistic degrees of freedom during BBN set by the Planck collaboration requires large soft breaking mass for the neutrinophilic Higgs. At present this scenario satisfies constraints from DM detection experiments. Actually, the scattering cross section is too small in order to explain possible excesses observed in the CDMS, DAMA, CoGENT and CRESST-II experiments [67–70] in this mass range, which have been interpreted as possible evidence of DM. Still, direct detection is possible in the future once the sensitivity in the lower mass range is improved. On the other hand, neutrinoless double beta decay is impossible in this model and a future observation of this process would rule out the current scenario.

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