The maximal U(1)LU(1)LU(1)_L inverse seesaw from d=5d=5d=5 operator and oscillating asymmetric Sneutrino dark matter Postprint

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Abstract

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Preamble

The Maximal $\mathbf{U(1)}_L$ Inverse Seesaw from $d = 5$ Operator and Oscillating Asymmetric Sneutrino Dark Matter

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Abstract

The maximal $\mathrm{U(1)}_L$ supersymmetric inverse seesaw mechanism (MLSIS) provides a natural way to relate asymmetric dark matter (ADM) with neutrino physics. In this paper we point out that MLSIS is a natural outcome if one dynamically realizes the inverse seesaw mechanism in the next-to-minimal supersymmetric standard model (NMSSM) via the dimension-five operator $(N)^2S^2/M_*$, with S the NMSSM singlet developing a TeV-scale VEV; it slightly violates lepton number due to suppression by the fundamental scale M_*, thus preserving $\mathrm{U}(1)_L$ maximally. The resulting sneutrino is a distinguishable ADM candidate, oscillating and favored to have weak-scale mass. A fairly large annihilation cross section for such heavy ADM is available due to the presence of the singlet.

I. Introduction and Motivation

Origins of tiny but non-vanishing neutrino masses are of great interest. Among the various mechanisms, the inverse seesaw mechanism [1] gains special attention, mainly because it provides a natural, simple, and testable way to realize small neutrino masses at low energy without invoking suppressed couplings. Besides, this mechanism follows a symmetry principle: a tiny neutrino mass, which slightly breaks lepton number by two units, is closely related to the degree of lepton number symmetry $\mathrm{U}(1)_L$ violation. Such an observation yields deep implications for supersymmetric dark matter (DM) candidates in supersymmetric standard models (SSMs): if the inverse seesaw mechanism is realized while retaining maximal $\mathrm{U}(1)_L$, i.e., one attributes the lightness of neutrinos to $\mathrm{U}(1)_L$ violation to the maximum extent (we refer to Eq. (3) for a more detailed explanation), the lightest sneutrino can be an asymmetric dark matter (ADM) candidate. The resulting scenario is dubbed MLSIS—the maximal $\mathrm{U}(1)_L$ supersymmetric inverse seesaw.

Thus far, ADM [2-6] is the most attractive mechanism to understand the coincidence between the relic densities of dark and baryonic matter, $\Omega_{\mathrm{DM}} : \Omega_b \approx 5 : 1$. However, realizing the ADM scenario in SSMs usually requires substantial extension, for example, invoking higher-dimensional operators with new scales [3]. On the other hand, it was believed that low-scale supersymmetric type-I seesaw could provide a sneutrino as an economical ADM candidate [7], but this is rendered ordinary symmetric DM by the large $\mathrm{U}(1)_L$ violation effect [9]. In contrast, in MLSIS the degree of $\mathrm{U}(1)_L$ violation is by definition under control: it just regenerates the symmetric DM components via oscillation [10] but does not spoil the ADM picture. The oscillating sneutrino ADM is strongly favored to have mass around the weak scale instead of the conventional GeV scale [9].

Despite being a well-motivated scenario to embed ADM in SSMs and moreover providing a distinguishable ADM candidate, MLSIS can still be improved from two aspects in terms of model building. First, the origin of maximal lepton number, or minimal $\mathrm{U}(1)_L$ violation, is a concern. This is not a new problem but is inherited from the inverse seesaw mechanism; see some attempts to address this problem [11, 12]. As the central result of this paper, we find that the presence of a singlet S developing a TeV-scale vacuum expectation value (VEV) provides a quite simple solution via a dimension-five operator with a high cutoff scale. Such a singlet is furnished in the well-known next-to-minimal SSM (NMSSM) [13], which thus provides the basis for model building. Second, in the minimally realized MLSIS [9] the sneutrino ADM fails to annihilate away effectively, and again a singlet can help us cope with this problem.

The ordinary symmetric sneutrino dark matter has been studied extensively by many groups [8].

This work is organized as follows. In Section II, MLSIS is realized in the Z_3-NMSSM with a dimension-five operator. In Section III, we study the oscillating sneutrino asymmetric DM, focusing on annihilation. The conclusion is given in Section IV.

II. The Maximal $\mathrm{U}(1)_L$ Inverse Seesaw Based on NMSSM

Let us begin with a brief review of MLSIS, which was first proposed in Ref. [9]. In the minimal scenario, the superpotential is nothing but that of the supersymmetric inverse seesaw mechanism [14, 15]:

$$W_{\mathrm{IS}} = y_N H_u L N^c + m_N N N^c + \frac{M_N}{2} N^2$$

We follow the notation of Ref. [9]: the chiral superfields are denoted as $N^c = (\tilde{\nu}_R^c)$ and $N = (\tilde{\nu}_L')$, with ν_L and ν_R both carrying lepton number $+1$. The Majorana mass term is the source of $\mathrm{U}(1)_L$ violation, by two units. For simplicity, we consider the single-family case. In the flavor basis (ν_L, ν_L^c), the neutrino mass matrix is given by

$$M_{\text{inverse}} = \begin{pmatrix} 0 & m_D & 0 \\ m_D & 0 & m_N \\ 0 & m_N & M_N \end{pmatrix}$$

with Dirac neutrino mass $m_D = y_N \langle H_u^0 \rangle$. In order to avoid large non-unitarity, we impose the bound $K \equiv m_N/m_D \lesssim 10$ [16]. Then the lightest neutrino is dominated by the active ν_L with $\sin\theta_\nu \approx 1 - 1/2K^2 \approx 1$. The neutrino mass takes the form of double suppression

$$m_\nu^{\text{eff}} = -\frac{m_D^2 M_N}{m_N^2 + M_N^2} \approx -\frac{M_N}{K^2}$$

If K takes a value as small as possible, M_N should take the smallest value accordingly. So $\mathrm{U}(1)_L$ would be respected to the greatest extent, leading to maximal $\mathrm{U}(1)_L$.

The other two Weyl fermions $\nu_{2,3} \approx \frac{1}{\sqrt{2}}(\nu_R \pm \nu_L^c)$ are singlet-like and heavy. They have almost degenerate masses $|M_{2,3}| = \sqrt{m_N^2 + m_D^2} + \mathcal{O}(M_N) \approx m_N$ and form a pseudo-Dirac fermion.

A. Realizing MLSIS in NMSSM via a dimension-five operator

In MLSIS, M_N is required to be $\lesssim 10$ eV. Such a tiny mass scale implies that the $\mathrm{U}(1)_L$ breaking term may originate from a higher-dimensional operator, which resembles the understanding of active neutrino mass via the Weinberg operator $\mathcal{O}_{\mathrm{Win}} = (LH_u)^2/M_*$.

Owing to the fact that the weak scale $v_u \approx 246$ GeV is relatively low, to give realistic neutrino mass we need a somewhat peculiar scale $M_* \sim 10^{14}$ GeV, which is close but two orders of magnitude lower than the grand unification theory (GUT) scale $\sim 10^{16}$ GeV. It is even far less than another putative fundamental scale, the Planck scale $M_{\mathrm{Pl}} \sim 10^{18}$ GeV or the string scale that interpolates between them.

In the case under consideration, the situation becomes quite different and intriguing new possibilities open up. In order to construct a Weinberg operator-like term for the $\mathrm{U}(1)_L$ breaking mass term, a scalar singlet S is introduced; moreover, it develops a VEV $v_s \equiv \langle S \rangle$ so that we have the analogy $(LH_u)(LH_u) \to NNSS$. Now we have

$$m_\nu^{\mathrm{eff}} \sim \frac{\lambda_2 v_s^2}{K^2 M_*}$$

Given a multi-TeV v_s, M_* can be naturally identified as the GUT scale for operator coefficient $\lambda_2 \sim 1$. However, if v_s is merely at the sub-TeV scale, we need to allow a large coefficient $\lambda_2 \sim K^2$. In particular, if we have $v_s \sim \mathcal{O}(10)$ TeV, even $M_* = M_{\mathrm{Pl}}$ is possible.

In this article we prefer a lower v_s because then one can enjoy the benefits of NMSSM: enhancing the SM-like Higgs boson mass via the new quartic term $\lambda^2 |H_u H_d|^2$ without losing electroweak scale naturalness, i.e., keeping a smaller $\mu = \lambda v_s \sim \mathcal{O}(100)$ GeV [17-19].

In SSMs, such a singlet is very welcome. As is well known, the minimal SSM (MSSM) contains a unique mass parameter in the superpotential, i.e., the μ parameter of the mass term for Higgsinos $\mu H_u H_d$. It is expected to be around the weak scale, which is technically natural but the origin of such a low scale should be addressed. Among other solutions, the NMSSM provides a simple and attractive approach by upgrading μ to be a dynamical field, $\mu \to S$ [13]. As a bonus, S can also generate the supersymmetric Dirac mass term for the singlets N and N^c in MLSIS. So we propose the following scale-invariant (or Z_3-invariant) superpotential except for the dimension-five operator:

$$W = W_{\mathrm{NMSSM}} + (y_N L H_u N^c + \lambda_1 S N N^c) + \frac{\lambda_2}{M_*} S^2 N^2$$

The soft SUSY-breaking Lagrangian is

$$-\mathcal{L}_{\text{soft}} = m^2_{\tilde{L}}|\tilde{L}|^2+m^2_{\tilde{\nu}'_L}|\tilde{\nu}'_L|^2+m^2_{\tilde{\nu}_R}|\tilde{\nu}_R|^2+(y_N A_N H_u \tilde{L}\tilde{\nu}^*_R+B_m m_N \tilde{\nu}'_L\tilde{\nu}_R+B_M M_N(\tilde{\nu}'_L)^2+\text{h.c.})$$

The soft SUSY-breaking parameters A_N, A_1, etc., are assumed to be real and around the weak scale. The ordinary NMSSM sector with Z_3 symmetry takes the form

$$W_{\text{NMSSM}} = \lambda S H_u H_d + \frac{\kappa}{3}S^3$$

$$-\mathcal{L}^{\text{NMSSM}}_{\text{soft}} = m^2_{H_u}|H_u|^2 + m^2_{H_d}|H_d|^2 + m^2_S|S|^2 + (\lambda A_\lambda S H_u H_d + \frac{\kappa}{3}A_\kappa S^3 + \text{h.c.})$$

As usual, we insist on the perturbative bound on the dimensionless couplings, e.g. $\lambda \lesssim 0.7$.

After S develops a VEV, all the mass terms in the superpotential Eq. (1) are dynamically generated, just like μ:

$$m_N = \lambda_1 v_s, \quad M_N = \frac{\lambda_2 v_s^2}{M_*}$$

This simple model can provide all the elements we need and is the minimal model to dynamically realize the inverse seesaw mechanism because we do not introduce any new fields (only the NMSSM plus right-handed neutrinos).

We would like to stress that the inverse seesaw mechanism based on a dimension-five operator is bound to be realized with maximal preservation of $\text{U}(1)_L$. The reason is simple. From Eq. (3) and Eq. (9) one obtains $v_s \sim K\sqrt{m_\nu^{\text{eff}} M_*/\lambda_2}$. For $m_\nu^{\text{eff}} \sim 0.1$ eV, $M_* \sim M_{\text{GUT}}$ and a not very large λ_2, a large $K \gg 10$ would push v_s far above the TeV scale, hence losing the benefits of NMSSM stated before. Therefore, we want K to be as small as possible, giving rise to the MLSIS scenario.

B. Tentative UV completion

Since the small neutrino mass scale is simply a relic of fundamental scale physics, this inspires us to investigate possible models at the fundamental scale. We find that a new $\text{U}(1)'_R$ symmetry can guarantee the general form of our model. At the renormalizable level, the supersymmetric model described by Eq. (7) plus Eq. (5) possesses an accidental $\text{U}(1)_B\times$ $\text{U}(1)_L\times$ $\text{U}(1)_R$ symmetry with field charges assigned as

$$
\begin{array}{ll}
\mathrm{U}(1)_L: & H_u[0], H_d[0], S[0], L[1], E^c[-1], N^c[-1], N[1], ... \\
\mathrm{U}(1)_B: & H_u[0], H_d[0], S[0], Q[1], U^c[-1], D^c[-1], N[0], ... \\
\mathrm{U}(1)_R: & H_u[2/3], H_d[2/3], S[2/3], L[2/3], E^c[2/3], ... \\
\mathrm{U}(1)'_R: & H_u[2/3], H_d[2/3], S[2/3], L[1/3], E^c[4/3], N^c[4/3], N[1/3], \Phi[1], ...
\end{array}
$$

where the dots denote all other fields carrying the same charge 2/3. As a matter of fact, the $\mathrm{U}(1)_R$ charge assignment is not fixed according to this superpotential, and in the above we simply choose one example consistent with SU(5) GUT. Note that the Z_3 symmetry is simply an accidental result of $\mathrm{U}(1)_R$ symmetry, which forbids bare mass terms. At the dimension-five level, the operator S^2N^2 violates the global symmetry $\mathrm{U}(1)_L$ and $\mathrm{U}(1)_R$ simultaneously, but still leaves a discrete $Z_2^L \subset \mathrm{U}(1)_L$ invariance, $R' \equiv R - \frac{1}{3}L$. The R' charge assignment of various fields is presented in the last line of Eq. (10). In particular, if $\mathrm{U}(1)'_R$ is generated to all orders, it was found that as a consequence $\mathrm{U}(1)_B$ and the matter parity $Z_2^M \equiv (-1)^{3(B-L)}$ are conserved to all orders [20].

With such a $\mathrm{U}(1)'_R$ symmetry, we explore concrete UV completions of the low-energy model which contains dimension-five operators and thus hints at new physics. This is of concern since we will find that M_* tends to be far below the fundamental Planck scale.

We introduce a heavy singlet Φ carrying unit $\mathrm{U}(1)'_R$ charge, and thus it can (only) couple to S and N via the renormalizable term

$$W = \lambda_\Phi S N \Phi + \frac{M_\Phi}{2}\Phi^2$$

with $M_\Phi \sim M_*$. Now integrating out Φ via the F-flatness condition of Φ, namely $F_\Phi = M_\Phi \Phi + \lambda_\Phi S N = 0$, one obtains the operator

$$\frac{\lambda_\Phi^2}{2M_\Phi} S^2 N^2 \quad \text{with} \quad M_* = -M_\Phi / 2\lambda_\Phi^2$$

We would like to point out that, in the presence of three families of RHNs, one may arrange an accidental hierarchy among λ_Φ or (and) M_Φ such that one effective cutoff scale M_* is hierarchically larger than others, and consequently the corresponding M_N is much smaller than others. Later, we will see that this is helpful to realize the oscillating sneutrino ADM.

III. Oscillating Asymmetric Sneutrino Dark Matter

In this section we study the main phenomenology of MLSIS implemented in the NMSSM: oscillating asymmetric sneutrino dark matter. Although the main physics has been investigated in Ref. [9], there are still several differences between MLSIS with and without the singlet S; these will be our focus here. We briefly discuss the similarities like asymmetry transfer and symmetry regeneration in the first subsection; for more details, see Ref. [9]. For illustration, we show the thermal history and corresponding dynamics of sneutrino ADM in Fig. 1 [Figure 1: see original paper].

A. Profiles of the oscillating sneutrino ADM

A big bonus of maximal $U(1)_L$ is the presence of an ADM candidate: the sneutrino. Let us begin with an exact $U(1)_L$, thus strictly complex sneutrinos. In the basis $\Phi^T = (\tilde{\nu}_L, \tilde{\nu}_R, \tilde{\nu}'_L)^T$, the sneutrino mass-squared matrix is given by

$$\mathcal{M}^2_{\tilde{\nu}} \approx \begin{pmatrix} m^2_{\tilde{L}} & -m_D A_N + \mu m_D \cot\beta & -m_D m_N \\ -m_D A_N + \mu m_D \cot\beta & m^2_{\tilde{\nu}_R} + m^2_D & -B_m m_N \\ -m_D m_N & -B_m m_N & m^2_{\tilde{\nu}'_L} + m^2_N \end{pmatrix}$$

Among the three sneutrinos $\tilde{\nu}_{1,2,3}$ in the mass eigenbasis, the lightest sneutrino is denoted as $\tilde{\nu}_1$. Stringent DM direct detection requires that the left-handed sneutrino composition in $\tilde{\nu}_1$ should be very small, and hence we can make the approximation:

$$\tilde{\nu}'_L \approx -\sin\tilde{\theta}\,\tilde{\nu}_1 + \cos\tilde{\theta}\,\tilde{\nu}_2, \quad \tilde{\nu}_R \approx \cos\tilde{\theta}\,\tilde{\nu}_1 + \sin\tilde{\theta}\,\tilde{\nu}_2$$

with $\tilde{\theta}$ the mixing angle. $\tilde{\nu}_1$ gains asymmetry when it enters chemical equilibrium with the leptons; after equilibrium breaks at decoupling temperature T_d, the lepton asymmetry in ADM is $\eta_0 \sim f_{\rm ADM}(x_d)\eta_b$, with $\eta_b \approx 10^{-11}$ the baryon asymmetry. $f_{\rm ADM}(x_d)$ is a factor encoding thermal threshold effects; it tends to 1 in the relativistic limit $x_d = m_{\rm ADM}/T_d \ll 1$, while being exponentially suppressed in the opposite limit.

The above conventional picture of ADM may be spoiled by tiny $U(1)_L$ violation. It induces mixing between the CP-even and -odd components of $\tilde{\nu}_1 = \frac{1}{\sqrt{2}}({\rm Re}\,\tilde{\nu}_1 + i\,{\rm Im}\,\tilde{\nu}_1)$ and moreover splits their masses by an amount δm. Consequently, DM and anti-DM can oscillate into each other. If oscillation is significant during ADM freeze-out, ADM will turn out to be ordinary symmetric DM. The oscillation rate is very sensitive to δm, whose upper limit is strongly dependent on the ADM mass [22, 23]: ADM ~ 300 GeV can tolerate $\delta m \sim 10^{-5}$ eV; but for conventional GeV-scale ADM, δm is forced to be incredibly small, $\lesssim 10^{-10}$ eV. Therefore, we will consider a weak-scale sneutrino ADM.

However, even $\delta m \sim 10^{-5}$ eV is still hard to achieve in MLSIS. To see this, one can well approximate the mass splitting as

$$\delta m \approx \frac{m_N M_N \sin 2\tilde{\theta} - B_M M_N \sin^2\tilde{\theta}}{m_{\tilde{\nu}_1}}$$

As one can see, the natural scale of δm should not be far below M_N except for $\sin\tilde{\theta} \ll 1$. However, for $m_\nu^{\rm eff} \sim 0.1$ eV, one has $M_N \sim K^2 m_\nu^{\rm eff} \gg 10^{-5}$ eV. Therefore, it is likely that $m_\nu^{\rm eff}$ should be relaxed, say having a value $\lesssim 0.1$ eV. This is allowed with three families of RHNs and may be regarded as a prediction of MLSIS with sneutrino ADM.

B. Constraining λ_1 from charge washing-out

It has already been noted that viable sneutrino ADM in MLSIS needs the aid of a singlet to annihilate away the symmetric part through the term $\lambda_1 N N^c$ [9]. However, the magnitude of λ_1 is stringently constrained by the DM charge-violating scattering (CVS) process $\tilde{\nu}_1 \nu_1 \leftrightarrow \tilde{\nu}_1^* \bar{\nu}_1$, which is mediated by neutralinos and can maintain chemical equilibrium between ADM and light neutrinos down to quite low temperature T_d [21]. If T_d falls to the DM freeze-out temperature $T_f \sim m_{\rm DM}/20$, no asymmetry will remain.

To determine T_d, one must compare the Hubble expansion rate $H(T) \approx 5.5T^2/M_{\rm Pl}$ with the CVS reaction rate, which can be obtained from the effective Lagrangian

$$-\mathcal{L}_{\rm wash} = i\bar{\chi}_i \gamma^\mu \partial_\mu \chi_i + (y_{i1} \tilde{\nu}_1^* \bar{\chi}_i P_L \nu_1 + \text{h.c.})$$

with χ_i the five Majorana neutralinos in the NMSSM. They are related to the interaction eigenstates via $\chi_i = Z_{ij}\psi_j$ with $\psi = (\tilde{B}, \tilde{W}^3, \tilde{H}_d^0, \tilde{H}_u^0, \tilde{s})^T$. Approximately, the effective couplings y_{i1} are given by

$$y_{i1} \approx y_N \sin\theta_\nu \cos\tilde{\theta} Z_{4i} - \lambda_1 \cos\theta_\nu \cos\tilde{\theta} Z_{5i}$$

where the second term comes from the λ_1 term. The CVS rate is calculated to be

$$\Gamma_{\rm CVS} = 5 \times 10^3 \times \left(\frac{T}{m_{\tilde{\nu}_1}}\right)^4 \left(\frac{T}{100\,\text{GeV}}\right)^2 \sum_i |y_{i1}|^2$$

Now, the condition $\Gamma_{\rm CVS}(T_d) < H(T_d)$ gives the upper bound on couplings

$$\sum_i |y_{i1}|^2 \lesssim 0.41 x_d \left(\frac{M_i}{m_{\tilde{\nu}_1}}\right)^4 \frac{m_{\tilde{\nu}_1}}{100\,\text{GeV}} = 1.0\times 10^{-10} \left(\frac{M_i}{m_{\tilde{\nu}_1}}\right)^4 \left(\frac{x_d}{20}\right)^5 \left(\frac{m_{\tilde{\nu}_1}}{100\,\text{GeV}}\right)$$

In the above estimation, ADM is assumed to be relatively heavy, around the weak scale, but neutralinos are even heavier, with multi-TeV masses so as to suppress the CVS rate. Then, it is seen that $\sum_i |y_{i1}|^2 \lesssim 10^{-2}$ should be fulfilled.

But we typically need $\sum_i |y_{i1}|^2 \lesssim 10^{-3}$ if neutralinos merely have masses close to the ADM mass, which is probably true at least for Higgsinos, whose masses are mainly determined by the μ term and expected to lie around the weak scale for the sake of naturalness.

Now we investigate possible ways to obtain small $|\sum_i y_{i1}|^2$ and the difficulty therein. First, neutrinos in the decoupling limit, i.e., $\cos\theta_\nu \approx 1/K \ll 1$, help suppress the λ_1 contribution. However, by definition MLSIS needs K to be as small as possible, so we merely have $\cos\theta_\nu \sim 0.1$. Second, as long as $\tilde{\nu}_1$ is dominated by $\tilde{\nu}_L'$, all these couplings can be naturally small due to suppression from $\cos\tilde{\theta} \ll 1$. But such a situation will hamper attempts to decrease the mass splitting δm (see Eq. (14)). Of course, the smallness of δm can always be attributed to a small M_N, so the option $\tilde{\nu}_1 \approx \tilde{\nu}_L'$ serves as the last trick for avoiding large CVS.

C. Annihilating away the symmetric part

Now we are in a position to discuss sneutrino ADM symmetric annihilation. The interactions between sneutrinos and the NMSSM sector heavily rely on λ_1 as well as m_N; see Eq. (A1). We list the relevant terms for convenience:

$$\mathcal{L} \supset -\frac{i\lambda_1}{\sqrt{2}} a_s \tilde{\nu}_1^* \tilde{\nu}_2 + \frac{\lambda_1}{\sqrt{2}} \cos 2\tilde{\theta}\, h_s \tilde{\nu}_1^* \tilde{\nu}_2 + \left[\frac{1}{2} + \frac{\lambda_1 \kappa}{\sqrt{2}\lambda} \sin 2\tilde{\theta}\right] a_s |\tilde{\nu}_1|^2 + \left[\frac{\sqrt{2}\lambda_1 m_N}{\lambda v_s} - \frac{\lambda_1 \kappa}{\sqrt{2}\lambda} \sin 2\tilde{\theta}\right] h_s |\tilde{\nu}_1|^2 + \frac{\sin 2\tilde{\theta}}{2} (\lambda_1 A$$

Interactions involving $y_N = m_N/Kv_u \ll 1$ (to satisfy the CVS bound) are neglected. One may wonder if it is possible to obtain a large ADM annihilation cross section in the $\tilde{\theta} \to 0$ limit ($\tilde{\nu}_1 \approx \tilde{\nu}_R$), which is favored by small δm. Unfortunately, we cannot. In that limit, the CVS bound requires $\lambda_1 \lesssim \mathcal{O}(0.01)$, and thus all couplings in Eq. (19) are suppressed except for the massive coupling κm_N, which may be sizable due to large m_N. But this renders a large y_N, inconsistent with the CVS bound. In what follows we present a viable scenario characterized by large $\lambda_1 \sim \mathcal{O}(0.1)$ and small v_s at the sub-TeV scale.

Two ways are available to annihilate away the symmetric part with a cross section of at least a few pb [27]. One is annihilating into the lighter a_s/h_s pair via contact interactions, with cross sections $\sigma v \sim \lambda_1^4/(64\pi m_{\tilde{\nu}_1}^2)$. Thus it works for $\lambda_1 \gtrsim 0.3$ and lighter ADM with $m_{\tilde{\nu}_1} \lesssim 100$ GeV. The other is via s-channel h_s. Near the resonant enhancement region, the inclusive cross section is

$$\sigma v = \frac{4\pi\Gamma(h_s \to \tilde{\nu}_1 \tilde{\nu}_1^*)}{m_{\tilde{\nu}_1}^2 \Gamma_{h_s}^{\rm tot}} \left[1 - {\rm Br}(h_s \to \tilde{\nu}_1 \tilde{\nu}_1^*)\right]$$

Hence in principle it can easily reach $\mathcal{O}$(pb) as long as h_s does not dominantly decay into DM pairs. Actually, h_s, due to a sizable λ near 1, tends to dominantly decay into a pair of SM-like Higgs bosons or Higgsinos if kinematically accessible.

D. On the detection of sneutrino DM

Sneutrino DM can interact with quarks via the three Higgs bosons H_i, but the interaction strengths are expected to be fairly weak. One can see this from the last line in Eq. (19), where $\sin 2\tilde{\theta} \ll 1$ in order to satisfy the CVS bound, and thus the only sizable contribution is from the $\lambda_1 m_N$ term; moreover, this term is negligible unless h_s strongly mixes with the doublet component. We consider this case to examine the prospect for direct detection of ADM.

The H_i mediate DM-nucleon spin-independent (SI) scattering. Its cross section, normalized to DM-proton scattering, is conventionally written as $\sigma_{\rm SI} = 4a_p^2\mu_p^2/\pi$ with μ_p the proton-DM reduced mass. The effective proton-DM coupling a_p receives three contributions

$$a_{p,H_i} = \mu_{H_i}^{11}\left[\sum_{q=u,d,s} f_q^{(p)} g_{qqH_i} + \sum_{q=c,b,t} f_q^{(p)} g_{qqH_i}\right]$$

where $\mu_{H_i}^{11}$ are the massive couplings for $H_i|\tilde{\nu}_1|^2$; concretely, $\mu_{H_i}^{11} \approx \lambda_1 m_N O_{i3}/\lambda v_s$. The effective couplings are $g_{uuH_i} = O_{i2}/\sin\beta$ for up-type quarks and $g_{ddH_i} = O_{i1}/\cos\beta$ for down-type quarks, with O defined in Eq. (A2). The coefficients take values $f_u^{(p)} = 0.026$, $f_d^{(p)} = 0.023$, $f_s^{(p)} = 0.033$, and $\sum_{q=u,d,s} f_q^{(p)} = 0.684$ [24, 25]. With these one can parameterize a_{p,H_i} as

$$a_{p,H_i} = 4.0 \times 10^{-3} \times \mu_{H_i}^{11}\left(\frac{\sin\beta + 0.343}{\cos\beta}\right)$$

For DM around the weak scale like 300 GeV, currently the most stringent upper bound is $a_{p,H_i} \lesssim 1.6 \times 10^{-9}(\sigma_{\rm up}/10^{-9}\,{\rm pb})^{1/2}$ GeV^{-2}. $\sigma_{\rm up}$ is from LUX [26], about 10^{-9} pb, implying $a_{p,H_i} \lesssim 1.6 \times 10^{-9}$ GeV^{-2}. Typically, $\sigma_{\rm SI}$ here lies below the upper bound:

$$a_{p,H_1} \approx 0.8 \times 10^{-9}\left(\frac{200\,{\rm GeV}}{m_{\tilde{\nu}_1}}\right)\left(\frac{125\,{\rm GeV}}{m_{H_1}}\right)^2\left(\frac{0.03}{\rm mixing}\right)\left(\frac{\lambda_1 m_N}{10\,{\rm GeV}}\right)\,{\rm GeV}^{-2}$$

In this optimistic estimation, H_1 is the SM-like Higgs boson and "mixing"denotes the factor in the bracket of Eq. (22). The next round of detection may reach the sneutrino ADM. Of course, the most promising probe is from indirect detection, because our ADM possesses a large annihilation cross section today; it is totally different from most ADM scenarios except for decaying ones [28].

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IV. Conclusion

MLSIS provides an attractive way to relate ADM with neutrino physics. Such a scenario is a necessary outcome if one dynamically realizes the inverse seesaw mechanism in the NMSSM via the dimension-five operator $(N)^2 S^2/M_*$ to explain the origin of the smallness of lepton number violation. The sneutrino is a distinguishable ADM candidate, oscillating and favored to have weak-scale mass. A fairly large annihilation cross section for such heavy ADM is available due to the presence of the singlet.

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Appendix A: Relevant interactions of sneutrino DM with Higgs bosons

In studying sneutrino DM annihilation and its scattering with nucleons, the interactions with Higgs bosons are relevant. We collect the dominant terms from F-term and soft terms below:

$$\mathcal{L} \supset -|\kappa S^2 + \lambda H_u H_d + \lambda_1 \tilde{\nu}'_L \tilde{\nu}^*_R|^2 - \Bigg[\frac{\lambda_1}{\sqrt{2}} (v_d a_u + v_u a_d) \tilde{\nu}^*_1 \tilde{\nu}_2 + \frac{\lambda_1 \kappa}{\sqrt{2}\lambda} \sin 2\tilde{\theta} (h_s a_s - a_s h_s) \tilde{\nu}^*_1 \tilde{\nu}_2 + \Bigg(\frac{\sqrt{2}\lambda_1 m_N}{\lambda v_s} - \frac{\lambda_1 \kappa}{\sqrt{2}\lambda}$$

We have written the Higgs fields as $S = v_s + (h_s + i a_s)/\sqrt{2}$ and similarly for others. We have not transformed the Higgs fields into their mass eigenstates yet. Following the convention in Ref. [13], we use matrix O to rotate the CP-even Higgs bosons: $(H_1, H_2, H_3)^T = O(h_d, h_u, h_s)^T$, with H_i ordered by mass. As for the CP-odd Higgs bosons, we first work in the basis (A, a_s) with $A = \cos\beta\, a_u + \sin\beta\, a_d$; the Goldstone mode $G = -\cos\beta\, a_d + \sin\beta\, a_u$ is projected out. Then we diagonalize (A, a_s) using matrix P': $(A_1, A_2)^T = P'(A, a_s)^T$. Finally we have $a_d = P_{i1} A_i$, $a_u = P_{i2} A_i$, $a_s = P_{i3} A_i$, with $P_{i1} = \sin\beta P'_{i1}$, $P_{i2} = \cos\beta P'_{i1}$ and $P_{i3} = P'_{i2}$ $(i = 1, 2)$.

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Note: Figure translations are in progress. See original paper for figures.