

Low scale seesaw mechanism with local lepton number

Hridoy Debnath^{*} and Pavel Fileviez Pérez[†]

*Physics Department and Center for Education and Research in Cosmology and Astrophysics (CERCA),
Case Western Reserve University, Cleveland, Ohio 44106, USA*



(Received 17 July 2023; accepted 6 September 2023; published 11 October 2023)

We discuss a class of theories for Majorana neutrinos where the total lepton number is a local gauge symmetry. These theories predict a dark matter candidate from anomaly cancellation. We discuss the properties of the dark matter candidate and using the cosmological bounds, we obtain the upper bound on the lepton number symmetry breaking scale. The dark matter candidate has unique annihilation channels due to the fact that the theory predicts a light pseudo-Goldstone boson, the Majoron, and one can obtain the correct relic density in a large fraction of the parameter space. In this context, the seesaw scale is below the $\mathcal{O}(10^2)$ TeV scale and one can hope to test the origin of neutrino masses at current or future colliders. We discuss the lepton number violating Higgs decays and the possibility to observe lepton number violation at the Large Hadron Collider.

DOI: [10.1103/PhysRevD.108.075009](https://doi.org/10.1103/PhysRevD.108.075009)

I. INTRODUCTION

The origin of the charged fermion and gauge boson masses in the Standard Model (SM) of particle physics is well understood. After the discovery of the Brout-Englert-Higgs (BEH) boson at the Large Hadron Collider we know that spontaneous symmetry breaking plays a fundamental role in nature. It is well known that the SM does not provide a mechanism to generate neutrino masses and one needs to look for a more complete theory. Today, we know nothing about the nature of neutrinos; they could be Majorana or Dirac fermions. In the case of Majorana neutrinos the total lepton number ($\ell = \ell_e + \ell_\mu + \ell_\tau$ in our notation) symmetry has to be broken in two units, while in the Dirac case the total lepton number is conserved or broken in more than two units.

One of the leading ideas for the origin of neutrino masses is related to the canonical seesaw mechanism [1–4], where the SM is extended including at least two copies of right-handed neutrinos. In this context, the SM neutrino masses are suppressed by a new mass scale, $M_\nu \approx Y_\nu^T M_N^{-1} Y_\nu v_0^2/2$, where Y_ν is the Yukawa coupling between the SM neutrinos and the new right-handed neutrinos, M_N defines a new Majorana mass scale and $v_0 = 246$ GeV is the constant value of the BEH field. This simple mechanism

does not predict the scale where the new mass scale M_N is generated; it could be $M_N \approx 10^{14-15}$ GeV if $Y_\nu \sim 1$. If the origin of neutrino masses is related to the so-called canonical scale, $M_{\text{seesaw}} \sim 10^{14}$ GeV, there is no way to have direct access to the mechanism of Majorana neutrino masses at current or future collider experiments.

The scale for total lepton number violation could be around the TeV scale if $Y_\nu \sim 10^{-6}$ and one could hope to test the origin of Majorana neutrino masses at colliders. In this case one can look for lepton number violating (LNV) signatures with same sign leptons and multijets as pointed out in Ref. [5]. See Ref. [6] for a detailed discussion about lepton number violation at colliders and Ref. [7] for a discussion of different ideas in physics beyond the Standard Model that could predict lepton number violation.

The total lepton number is a key symmetry to understanding the origin of neutrino masses. In the SM the total lepton symmetry is a conserved global symmetry at the classical level but broken by the $SU(2)_L$ instantons in three units. One can promote the total lepton number to a local gauge symmetry to understand the link between spontaneous symmetry breaking and neutrino masses. In Refs. [8–10] the authors studied simple anomaly-free theories based on local $U(1)_\ell$. See also Ref. [11] for an earlier discussion. The simple theories in Refs. [9,10] predict the following:

- (1) a new sector needed for anomaly cancellation,
- (2) a dark matter candidate from anomaly cancellation,
- (3) the total lepton number must be broken below the multi-TeV scale, and
- (4) possible lepton number violating signatures at colliders.

In this article, we study for the first time a simple theory based on local total lepton number with Majorana neutrinos.

^{*}hxd253@case.edu

[†]pxf112@case.edu

Published by the American Physical Society under the terms of the Creative Commons Attribution 4.0 International license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

We discuss in detail the properties of the dark matter candidate discussing each annihilation channel. This theory predicts the existence of a pseudo-Nambu-Goldstone boson (Majoron) associated with the spontaneous breaking of lepton number [12]. We show that the dark matter candidate can annihilate through different channels containing the Majoron and one can satisfy the relic density constraints in different regions of the parameter space compared to the study in Ref. [13]. Using the cosmological bounds on the dark matter relic density we point out the upper bound on the mass of the new neutral gauge boson associated with the total lepton number gauge symmetry.

Since the new symmetry breaking scale is below the $\mathcal{O}(10^2)$ TeV scale, we discuss the possibility to observe lepton number violation from Higgs decays at the Large Hadron Collider. In the case of the SM-like Higgs one can have a large number of LNV events because the branching ratio, $\text{BR}(h \rightarrow NN)$, could be relatively large. We also show the predictions for the LNV signatures from the other Higgs decays. In this case, one cannot have many events at the LHC due to the fact that the production cross section is small if one uses the bounds on the mixing between the SM-like Higgs and the new Higgses coming from direct detection dark matter experiments. The theory studied in this article provides an interesting correlation between the origin of neutrino masses, the nature of dark matter, and possible exotic signatures at colliders.

This article is organized as follows: In Sec. II we discuss a simple theory for Majorana neutrino masses where the total lepton number is a local gauge symmetry spontaneously broken at the low scale. In Sec. III we discuss the main properties of the fermionic dark matter candidate predicted from anomaly cancellation in the theory. The lepton number violating Higgs decays at the Large Hadron Collider are discussed in Sec. IV, and we summarize our main findings in Sec. V.

II. LEPTON NUMBER AS A LOCAL GAUGE SYMMETRY

We can have a simple theory for massive neutrinos based on the gauge symmetry [8–10]

$$SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes U(1)_\ell,$$

where $U(1)_\ell$ is defined as the local total lepton number. It is well known that the total lepton number is a symmetry conserved in the Standard Model at the classical level but broken by three units at the quantum level by $SU(2)_L$ instantons. In order to define an anomaly-free theory based on the above gauge symmetry, we need to add new fermions together with right-handed neutrinos (ν_R). In Table. I we list the extra fermions needed for anomaly cancellation as proposed in Ref. [9]. See also the other possibilities in Refs. [8,10].

TABLE I. Fermions needed for anomaly cancellation with $\ell_1 - \ell_2 = -3$ [9].

| Fields | $SU(3)_C$ | $SU(2)_L$ | $U(1)_Y$ | $U(1)_\ell$ |
|---|-----------|-----------|----------------|-------------|
| $\Psi_L = \begin{pmatrix} \Psi_L^0 \\ \Psi_L^- \end{pmatrix}$ | 1 | 2 | $-\frac{1}{2}$ | ℓ_1 |
| $\Psi_R = \begin{pmatrix} \Psi_R^0 \\ \Psi_R^- \end{pmatrix}$ | 1 | 2 | $-\frac{1}{2}$ | ℓ_2 |
| η_R | 1 | 1 | -1 | ℓ_1 |
| η_L | 1 | 1 | -1 | ℓ_2 |
| χ_R | 1 | 1 | 0 | ℓ_1 |
| χ_L | 1 | 1 | 0 | ℓ_2 |

The Lagrangian of this theory is given by

$$\begin{aligned} \mathcal{L} = & \mathcal{L}_{\text{SM}} + i\bar{\nu}_R \not{\partial} \nu_R - g_\ell (\bar{\ell}_L \gamma^\mu \ell_L + \bar{e}_R \gamma^\mu e_R + \bar{\nu}_R \gamma^\mu \nu_R) Z_\mu^\ell \\ & - \frac{1}{4} Z_{\mu\nu}^2 Z^{\ell,\mu\nu} + \mathcal{L}_K^L + \mathcal{L}_Y^L - V(H, S). \end{aligned} \quad (1)$$

Here $\ell_L \sim (\mathbf{1}, \mathbf{2}, -1/2, 1)$ and $e_R \sim (\mathbf{1}, \mathbf{1}, -1, 1)$ are the SM leptonic fields, \mathcal{L}_{SM} is the SM Lagrangian, and $V(H, S)$ contains the new terms in the scalar potential. Here we neglect the kinetic mixing term between the two Abelian gauge symmetries for simplicity. The kinetic terms for the new fields can be written as

$$\begin{aligned} \mathcal{L}_K^L = & i\bar{\Psi}_L \not{D} \Psi_L + i\bar{\Psi}_R \not{D} \Psi_R + i\bar{\eta}_L \not{D} \eta_L + i\bar{\eta}_R \not{D} \eta_R \\ & + i\bar{\chi}_L \not{D} \chi_L + i\bar{\chi}_R \not{D} \chi_R + (D_\mu S)^\dagger (D^\mu S), \end{aligned} \quad (2)$$

and the new Yukawa interactions read as

$$\begin{aligned} -\mathcal{L}_Y^L = & y_1 \bar{\Psi}_L H \eta_R + y_2 \bar{\Psi}_R H \eta_L + y_3 \bar{\Psi}_L \tilde{H} \chi_R + y_4 \bar{\Psi}_R \tilde{H} \chi_L \\ & + y_\Psi \bar{\Psi}_L \Psi_R S^* + y_\eta \bar{\eta}_R \eta_L S^* + y_\chi \bar{\chi}_R \chi_L S^* + \text{H.c.} \end{aligned} \quad (3)$$

Here $H \sim (\mathbf{1}, \mathbf{2}, 1/2, 0)$ is the Standard Model Higgs, and $\tilde{H} = i\sigma_2 H^*$. Here we do not consider the case when $\ell_1 \neq -\ell_2$ and assume $\ell_i \neq \pm 1$. Notice that when we have fractional values for ℓ_i one makes sure that higher-dimensional operators do not affect the stability of our dark matter candidate. The case $\ell_1 = -\ell_2 = -3/2$ is interesting but less generic. In this case the Yukawa interactions $\chi_L \chi_L S^*$ and $\chi_R \chi_R S$ are allowed by the gauge symmetry, and the dark matter candidate is a Majorana fermion. In this article we focus on the most generic cases with $\ell_1 - \ell_2 = -3$, $\ell_1 \neq -\ell_2$ and $\ell_i \neq \pm 1$. Notice that the condition $\ell_i \neq \pm 1$ is needed to avoid the mixing between the χ_L , χ_R and the right-handed neutrinos.

The new Higgs quantum numbers are determined by anomaly cancellation condition, $\ell_1 - \ell_2 = -3$, and the above Yukawa interactions. Therefore, the new Higgs transforms as

$$S \sim (\mathbf{1}, \mathbf{1}, 0, 3). \quad (4)$$

For some studies in this context see Refs. [8,13–17].

In this theory we can generate Majorana neutrinos using the following interactions:

$$-\mathcal{L}_\nu \supset Y_\nu \bar{\ell}_L \tilde{H} \nu_R + \lambda_R \nu_R^T C \phi \nu_R + \text{H.c.} \quad (5)$$

Here $\phi \sim (\mathbf{1}, \mathbf{1}, 0, -2)$ is a new Higgs field needed to implement the seesaw mechanism. After the local lepton number symmetry, $U(1)_\ell$, is broken, the theory has an accidental global $U(1)$ symmetry:

$$\begin{aligned} \Psi_L &\rightarrow e^{i\theta} \Psi_L, & \Psi_R &\rightarrow e^{i\theta} \Psi_R, \\ \eta_L &\rightarrow e^{i\theta} \eta_L, & \eta_R &\rightarrow e^{i\theta} \eta_R, \\ \chi_L &\rightarrow e^{i\theta} \chi_L, & \chi_R &\rightarrow e^{i\theta} \chi_R. \end{aligned}$$

Therefore, the lightest field in the new sector is stable. In order to have a consistent scenario for cosmology we assume that the lightest new stable field is the neutral Dirac field: $\chi = \chi_L + \chi_R$. Therefore, one can say that this theory predicts a cold dark matter candidate from anomaly cancellation and its stability is a natural consequence from spontaneous symmetry breaking. Notice that field $\Psi = \Psi_L^0 + \Psi_R^0$ is ruled out as a dark matter candidate because it has a large coupling to the Z gauge boson and one cannot satisfy the dark matter direct detection constraints.

A. Higgs sector

The Higgs sector is composed of the SM Higgs, H , and the new Higgses, S and ϕ . The scalar potential in this theory is given by

$$\begin{aligned} V(H, S, \phi) &= -m_H^2 H^\dagger H + \lambda(H^\dagger H)^2 - m_S^2 S^\dagger S + \lambda_S (S^\dagger S)^2 \\ &\quad - m_\phi^2 \phi^\dagger \phi + \lambda_\phi (\phi^\dagger \phi)^2 + \lambda_1 (H^\dagger H) S^\dagger S \\ &\quad + \lambda_2 (H^\dagger H) \phi^\dagger \phi + \lambda_3 (S^\dagger S) \phi^\dagger \phi. \end{aligned} \quad (6)$$

Notice that this scalar potential has the global symmetry $O(4)_H \otimes U(1)_\phi \otimes U(1)_S$. The scalar fields in this theory can be written as

$$H = \begin{pmatrix} h^+ \\ \frac{1}{\sqrt{2}}(v_0 + h_0)e^{i\sigma_0/v_0} \end{pmatrix}, \quad (7)$$

$$S = \frac{1}{\sqrt{2}}(v_s + h_s)e^{i\sigma_s/v_s}, \quad (8)$$

and

$$\phi = \frac{1}{\sqrt{2}}(v_\phi + h_\phi)e^{i\sigma_\phi/v_\phi}. \quad (9)$$

After the spontaneous symmetry breaking, the constant part of the scalar potential can be written as

$$\begin{aligned} V(v_0, v_L, v_\phi) &= -\frac{1}{2}m_H^2 v_0^2 + \frac{\lambda}{4}v_0^4 - \frac{1}{2}m_S^2 v_s^2 + \frac{\lambda_S}{4}v_s^4 \\ &\quad - \frac{1}{2}m_\phi^2 v_\phi^2 + \frac{\lambda_\phi}{4}v_\phi^4 + \frac{\lambda_1}{4}v_0^2 v_s^2 \\ &\quad + \frac{\lambda_2}{4}v_0^2 v_\phi^2 + \frac{\lambda_3}{4}v_\phi^2 v_s^2, \end{aligned} \quad (10)$$

and the minimization conditions read as

$$-m_H^2 v_0 + \lambda v_0^3 + \frac{1}{2}\lambda_1 v_s^2 v_0 + \frac{1}{2}\lambda_2 v_\phi^2 v_0 = 0, \quad (11)$$

$$-m_\phi^2 v_\phi + \lambda_\phi v_\phi^3 + \frac{1}{2}\lambda_2 v_0^2 v_\phi + \frac{1}{2}\lambda_3 v_s^2 v_\phi = 0, \quad (12)$$

and

$$-m_S^2 v_s + \lambda_S v_s^3 + \frac{1}{2}\lambda_1 v_0^2 v_s + \frac{1}{2}\lambda_3 v_\phi^2 v_s = 0. \quad (13)$$

In this theory, the mass matrix for the CP -even Higgses, in the (h_0, h_s, h_ϕ) basis, can be written as

$$M_{\text{even}}^2 = \begin{pmatrix} 2\lambda v_0^2 & \lambda_1 v_0 v_s & \lambda_2 v_0 v_\phi \\ \lambda_1 v_0 v_s & 2\lambda_S v_s^2 & \lambda_3 v_s v_\phi \\ \lambda_2 v_0 v_\phi & \lambda_3 v_s v_\phi & 2v_\phi^2 \lambda_\phi \end{pmatrix}. \quad (14)$$

In our notation the physical Higgses, (h, H_1, H_2) , are defined as

$$\begin{pmatrix} h_0 \\ h_s \\ h_\phi \end{pmatrix} = U \begin{pmatrix} h \\ H_1 \\ H_2 \end{pmatrix}. \quad (15)$$

There are three CP -odd Higgses in this theory; two of them are Goldstone's bosons eaten by the neutral gauge bosons. Notice that the CP -odd Higgses, σ_s and σ_ϕ , masses are protected by the a shift symmetry, $\sigma_i \rightarrow \sigma_i + c$, where c is a constant. This shift symmetry is broken by the dimensional five term in the scalar potential:

$$V(H, S, \phi) \supset \lambda_M \frac{S^2 \phi^3}{\Lambda} + \text{H.c.} \quad (16)$$

This term is allowed by all the symmetries of the theory and tells us that there is only one Nambu-Golstone boson in the new sector. Notice that this term breaks the global symmetry, $U(1)_\phi \otimes U(1)_S$, to a new $U(1)$ symmetry. The CP -odd mass matrix in the (σ_s, σ_ϕ) basis can be written as

$$M_{\text{odd}}^2 = \frac{\lambda_M}{\sqrt{2}\Lambda} \begin{pmatrix} 2v_\phi^3 & 3v_s v_\phi^2 \\ 3v_s v_\phi^2 & \frac{9}{2}v_s^2 v_\phi \end{pmatrix}, \quad (17)$$

with eigenvalues

$$M_{G_\ell}^2 = 0 \quad \text{and} \quad M_J^2 = \frac{\lambda_M v_\phi}{2\sqrt{2}\Lambda} (4v_\phi^2 + 9v_s^2). \quad (18)$$

In this theory, the Z_ℓ mass can be written as

$$M_{Z_\ell}^2 = g_\ell^2 (9v_s^2 + 4v_\phi^2). \quad (19)$$

Notice that the mass of the leptophilic gauge boson can be used as the seesaw scale in this theory because it tells us about the scale where the lepton number is spontaneously broken.

Using Eqs. (18) and (19) one can estimate the Majoron mass as a function of the ratio between the new gauge boson mass and gauge coupling: $M_J^2 = \lambda_M M_{Z_\ell}^3 \cos \beta / (4\sqrt{2}g_\ell^3 \Lambda)$. One finds that $M_J \sim 10^{-4}$ GeV when $\lambda_M \sim 1$, $M_{Z_\ell}/g_\ell \sim 10$ TeV, $\cos \beta \sim 1$, and $\Lambda \sim M_{\text{Pl}}$, with M_{Pl} being the Planck scale. As we will discuss, in our case the seesaw scale will be below $\mathcal{O}(10^2)$ TeV. Therefore, the Majoron will decay fast into neutrinos and it cannot be a dark matter candidate. See for example Ref. [18] for a detailed discussion of the Majoron properties. For the cosmological bounds on the Majoron interactions with neutrinos, see for example Refs. [19,20].

The CP -odd Higgs eigenstates are defined by

$$\begin{pmatrix} \sigma_s \\ \sigma_\phi \end{pmatrix} = \begin{pmatrix} \cos \beta & \sin \beta \\ -\sin \beta & \cos \beta \end{pmatrix} \begin{pmatrix} G_\ell \\ J \end{pmatrix}, \quad (20)$$

where

$$\tan 2\beta = \frac{12v_s v_\phi}{4v_\phi^2 - 9v_s^2}. \quad (21)$$

B. Fermionic Fields

1. Neutral Dirac fermions

The mass matrix for the neutral fermions can be written, in the $(\chi_L^0 \Psi_L^0)$ and $(\chi_R^0 \Psi_R^0)$ basis, as

$$-\mathcal{L} \supset (\overline{\chi_R^0} \overline{\Psi_R^0}) \mathcal{M}_0 \begin{pmatrix} \chi_L^0 \\ \Psi_L^0 \end{pmatrix} + \text{H.c.}, \quad (22)$$

where

$$\mathcal{M}_0 = \frac{1}{\sqrt{2}} \begin{pmatrix} y_\chi v_s & y_3 v_0 \\ y_4 v_0 & y_\Psi v_s \end{pmatrix}. \quad (23)$$

One can diagonalize the mass matrix as follows:

$$\mathcal{M}_0^{\text{diag}} = N_R^\dagger \mathcal{M}_0 N_L. \quad (24)$$

where the neutral fields are related by the N_L and N_R mixing matrices as given by

$$\begin{pmatrix} \chi_L^0 \\ \Psi_L^0 \end{pmatrix} = N_L \begin{pmatrix} \chi_{1L}^0 \\ \chi_{2L}^0 \end{pmatrix}, \quad \text{and} \quad \begin{pmatrix} \chi_R^0 \\ \Psi_R^0 \end{pmatrix} = N_R \begin{pmatrix} \chi_{1R}^0 \\ \chi_{2R}^0 \end{pmatrix}. \quad (25)$$

In the limit when $v_s \gg v_0$, the mass of the dark matter candidate is given as $M_\chi = y_\chi v_s / \sqrt{2}$.

2. Charged fermions

The mass matrix for the charged fermions can be written in the $(\eta_R^- \Psi_R^-)$ and $(\eta_L^- \Psi_L^-)$ basis as

$$-\mathcal{L} \supset (\overline{\eta_R^-} \overline{\Psi_R^-}) \mathcal{M}_C \begin{pmatrix} \eta_L^- \\ \Psi_L^- \end{pmatrix} + \text{H.c.}, \quad (26)$$

where

$$\mathcal{M}_C = \frac{1}{\sqrt{2}} \begin{pmatrix} y_\eta v_s & y_1 v_0 \\ y_2 v_0 & y_\Psi v_s \end{pmatrix}. \quad (27)$$

In our notation, the mass matrix can be diagonalized as follows:

$$\mathcal{M}_C^{\text{diag}} = V_R^\dagger \mathcal{M}_C V_L, \quad (28)$$

while the physical charged fields are related by the V_L and V_R matrices

$$\begin{pmatrix} \eta_L^- \\ \Psi_L^- \end{pmatrix} = V_L \begin{pmatrix} \chi_{1L}^- \\ \chi_{2L}^- \end{pmatrix}, \quad \text{and} \quad \begin{pmatrix} \eta_R^- \\ \Psi_R^- \end{pmatrix} = V_R \begin{pmatrix} \chi_{1R}^- \\ \chi_{2R}^- \end{pmatrix}. \quad (29)$$

3. Neutrino masses

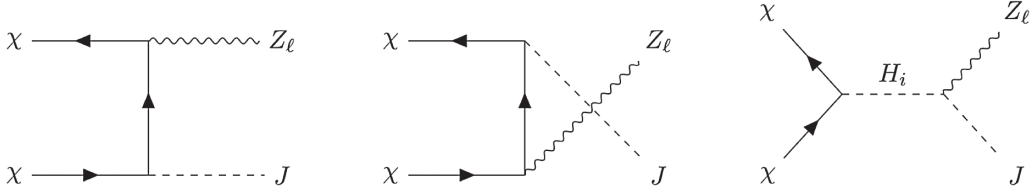
In this theory, the neutrino masses are generated through the type I seesaw mechanism and the SM neutrino mass matrix is given by

$$M_\nu = \frac{v_0^2}{2} Y_\nu M_N^{-1} Y_\nu^T, \quad (30)$$

where

$$M_N = \sqrt{2} \lambda_R v_\phi = \frac{\lambda_R}{\sqrt{2}} \frac{M_{Z_\ell}}{g_\ell} \cos \beta. \quad (31)$$

Here we used $v_\phi = v \cos \beta / 2$ and $v_s = v \sin \beta / 3$. Therefore, in this theory the upper bound on the seesaw scale is determined by the ratio M_{Z_ℓ}/g_ℓ and the perturbative bound on the Yukawa coupling λ_R .


 FIG. 1. Feynman graphs for the annihilation channel $\bar{\chi}\chi \rightarrow Z_\ell J$.

III. DARK MATTER CONSTRAINTS

We have discussed above that this theory predicts a dark matter from anomaly cancellation. In this context, the dark matter is a Dirac fermion, $\chi = \chi_L + \chi_R$. After symmetry breaking χ_L and χ_R have the same quantum numbers. The dominant dark matter annihilation channels are

$$\bar{\chi}\chi \rightarrow e_i e_i, \nu\nu, Z_\ell Z_\ell, H_i H_j, Z_\ell H_i, Z_\ell J, J H_i.$$

Notice that the existence of the annihilation channels containing the pseudo-Goldstone boson J , the Majoron, is quite unique. These channels will allow us to obtain the correct relic density in regions of the parameter space far from the Z_ℓ resonance.

One can compute the relic density using [21]

$$\Omega_{\text{DM}} h^2 = \frac{1.05 \times 10^9 \text{ GeV}^{-1}}{J(x_f) M_{\text{Pl}}}, \quad (32)$$

where the function $J(x_f)$ can be written as

$$J(x_f) = \int_{x_f}^{\infty} \frac{g_*^{1/2}(x) \langle \sigma v \rangle (x)}{x^2} dx, \quad (33)$$

where g_* is the total number of degrees of freedom at freeze-out and the thermal average cross section times velocity is given by

$$\langle \sigma v \rangle (x) = \frac{x}{8M_\chi^2 K_2^2(x)} \int_{4M_\chi^2}^{\infty} \sigma \times (s - 4M_\chi^2) \sqrt{s} K_1 \left(\frac{x\sqrt{s}}{M_\chi} \right) ds. \quad (34)$$

Here $x = M_\chi/T$, $K_1(x)$, and $K_2(x)$ are the modified Bessel functions. The freeze-out parameter x_f can also be computed by using

$$x_f = \ln \left(\frac{0.038 g M_{\text{Pl}} M_\chi \langle \sigma v \rangle (x_f)}{\sqrt{g_* x_f}} \right), \quad (35)$$

where g is the effective number of degrees of freedom of the dark matter particle and $M_{\text{Pl}} = 1.22 \times 10^{19}$ GeV. In order to understand the importance of each annihilation channel, we study the contribution of each channel independently, and later taking into account all relevant annihilation channels we find the upper bound on the lepton number symmetry breaking scale.

In this theory, we have unique annihilation channels since the dark matter can annihilate into a light Majoron, J . Then, we start our discussion taking into account only the annihilation channel $\bar{\chi}\chi \rightarrow Z_\ell J$. This channel has three contributions, see Fig. 1.

In Fig. 2 (left panel) we show the allowed region when only this channel is included in the calculation of the relic density. As one can appreciate, there is an upper bound on the Z_ℓ mass around 30 TeV. This annihilation channel can

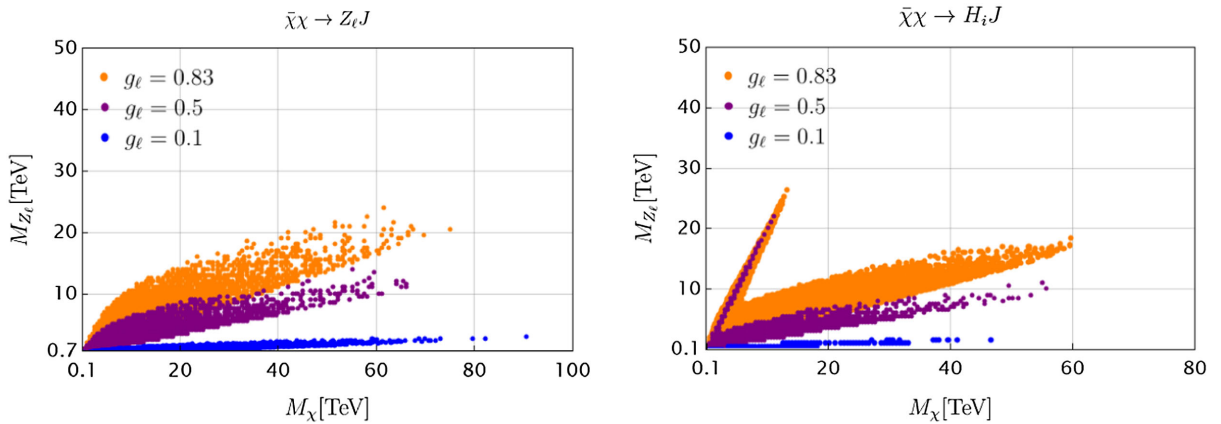


FIG. 2. Allowed regions by the cosmological bounds on dark matter relic density [22] and perturbative bound on the y_χ coupling. In the left panel, we have the annihilation into $Z_\ell J$, while in the right panel one has the $H_i J$ channel. We show the results for different values of the gauge coupling g_ℓ and assume that the mass of the new Higgses are equal to 1 TeV.

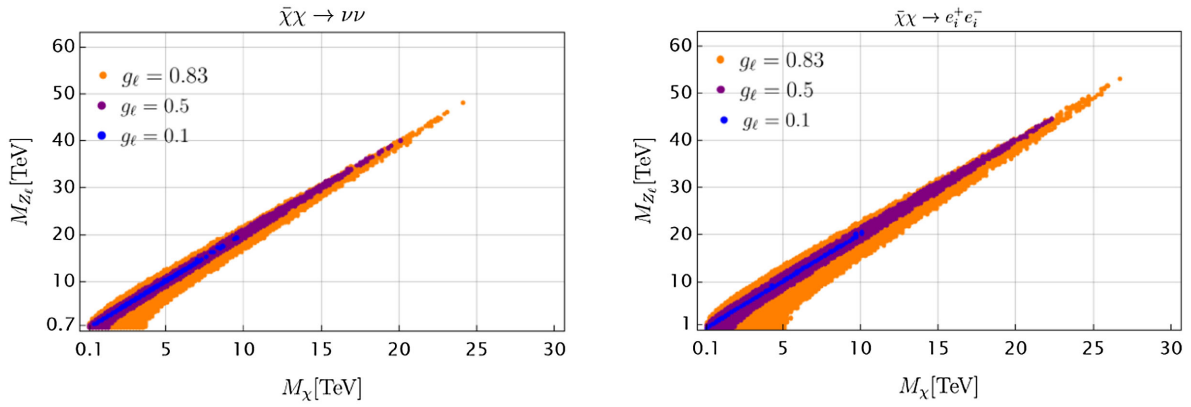


FIG. 3. Allowed regions by the cosmological bounds on dark matter relic density [22] and perturbative bound on the y_χ coupling, when only the annihilation into leptons is included. In the left panel we consider the annihilation into neutrinos, while in the right panel we have the annihilation into charged leptons.

give us the correct relic density in a large fraction of the parameter space. We show the results for three different values of the gauge coupling, $g_\ell = 0.83$ (in orange), $g_\ell = 0.5$ (in purple), and $g_\ell = 0.1$ (in blue). Notice that when we decrease the value of the gauge coupling g_ℓ from 0.83 to 0.1, the allowed region is smaller because the cross section for the process $\chi\bar{\chi} \rightarrow Z_\ell J$ is smaller for smaller g_ℓ and one obtains the correct relic density only in a reduced region of the parameter space. For all numerical examples, we use the values $\ell_1 = 1/2$ and $\ell_2 = 7/2$. The main results are very similar if we choose different values for ℓ_1 and ℓ_2 charges. This channel is important because one can have a large annihilation cross section when $M_\chi > M_{Z_\ell}/2$ because the Majoron mass is very small and one does not need to rely on a resonance to achieve the correct relic density. We show similar results for the annihilation channel $\chi\bar{\chi} \rightarrow H_i J$ in the right panel in Fig. 2. However, in this case one has also the allowed solutions when you are close to the Z_ℓ resonance because one of the processes is $\chi\bar{\chi} \rightarrow Z_\ell^* \rightarrow H_i J$, and then when $M_\chi \sim M_{Z_\ell}/2$ one can have the large contribution of

the Z_ℓ resonance. The different colored regions correspond to the different values of the gauge couplings as in the left panel. The maximal value of g_ℓ is determined by the perturbative bound on the $S^\dagger S Z_\ell Z_\ell$ coupling, which gives us that $g_\ell \leq 0.83$.

In Fig. 3 we show the allowed parameter when we include only the dark matter annihilation into leptons. In the left panel we have the annihilation into neutrinos, while in the right panel we have the annihilation into charged leptons. These channels can give the correct relic density only when one has the Z_ℓ resonance, i.e., $M_\chi \sim M_{Z_\ell}/2$. The different colored regions correspond to different values of the gauge coupling. As one expects, when the gauge coupling is smaller it is more difficult to find solutions allowed by the relic density constraints because the annihilation cross section is smaller and one typically obtains too much relic density for smaller couplings. Therefore, when the coupling is smaller one has allowed solutions only when the dark matter mass is very close to $M_{Z_\ell}/2$.

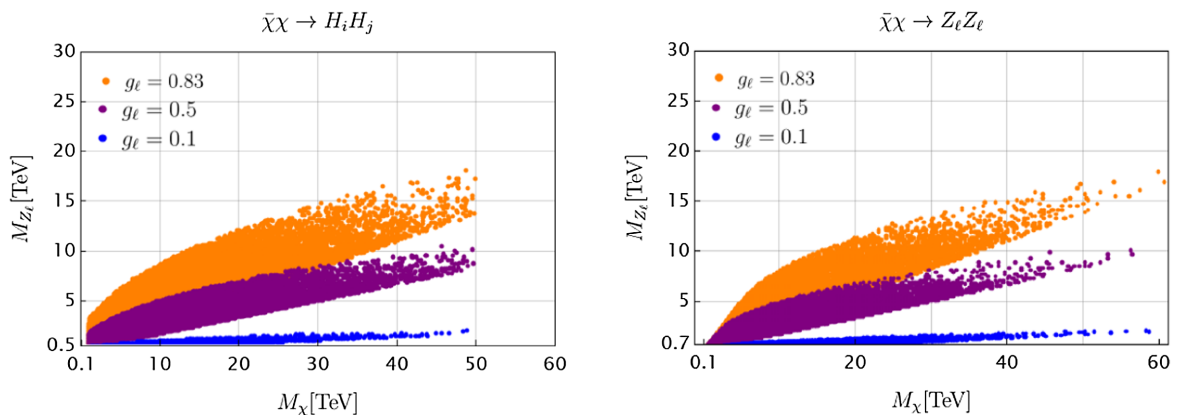


FIG. 4. Allowed regions by the cosmological bounds on dark matter relic density [22] and perturbative bound on the y_χ coupling, when only the annihilation into two Higgses or two gauge bosons are included. In the left panel we consider the annihilation into two Higgses, while in the right panel we have the annihilation into two gauge bosons. Here we use $M_{H_1} = M_{H_2} = 1$ TeV.

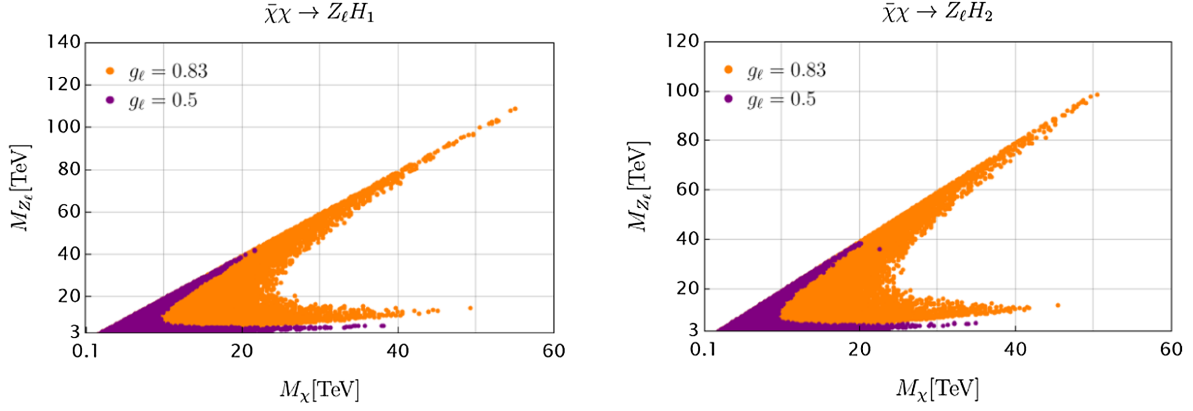


FIG. 5. Allowed regions by the cosmological bounds on dark matter relic density [22] and perturbative bound on the y_χ coupling when only the annihilations into a Higgs and a gauge boson are included. In the left panel we have $\bar{\chi}\chi \rightarrow Z_\ell H_1$, while in the right panel we show the results for $\bar{\chi}\chi \rightarrow Z_\ell H_2$. Here we use $M_{H_1} = M_{H_2} = 1$ TeV.

In Fig. 4 we show the same results for the $H_i H_j$ and $Z_\ell Z_\ell$ annihilation channels. When these channels are kinematically allowed, one can have a relic density in agreement with experiments in a large fraction of the $M_{Z_\ell} - M_\chi$ plane. The main contributions for the annihilation cross sections for these channels are from the u and t channels. Since the masses of Z_ℓ and the new Higgses must be well above the electroweak scale, these processes are mainly relevant when the dark matter mass is above the electroweak scale.

In Fig. 5, we show the allowed parameter space when one has only the $Z_\ell H_1$ (left panel) or the $Z_\ell H_2$ (right panel) channels. In this case, one has two main regions, the region around the Z_ℓ resonance and the region when $2M_\chi \gg M_{Z_\ell} + M_{H_i}$. Here we show the allowed solutions by the relic density constraints only for two values of the gauge couplings, $g_\ell = 0.5$ (in orange) and $g_\ell = 0.83$ (in purple), because there are no allowed solutions when $g_\ell = 0.1$. Notice that in this case one has two main regions;

below the Z_ℓ resonance and the u and t channels allow us to obtain solutions far from the resonance.

In Fig. 6 the allowed $M_{Z_\ell} - M_\chi$ plane is shown by considering only the $\chi\bar{\chi} \rightarrow JJ$ (left panel) and $\chi\bar{\chi} \rightarrow NN$ (right panel) annihilation channels. In this theory, the $\chi\bar{\chi} \rightarrow JJ$ annihilation channel is velocity suppressed, and the allowed parameter space by the relic density does not satisfy the collider bounds on the Z_ℓ mass. The $\chi\bar{\chi} \rightarrow NN$ channel can satisfy both the relic density and collider bounds only in the resonance region. These channels are clearly unique because they are present in the theory predicts the existence of the Majoron and the right-handed neutrinos. The colored regions correspond to the same values of the gauge couplings as in the previous figures. Notice that in the case when we study the annihilation into two Majorons, one has only the u and t contributions. For the annihilation into two right-handed neutrinos we have also $\chi\bar{\chi} \rightarrow Z_\ell^* \rightarrow NN$, and then we can have also the allowed region close to the Z_ℓ resonance.

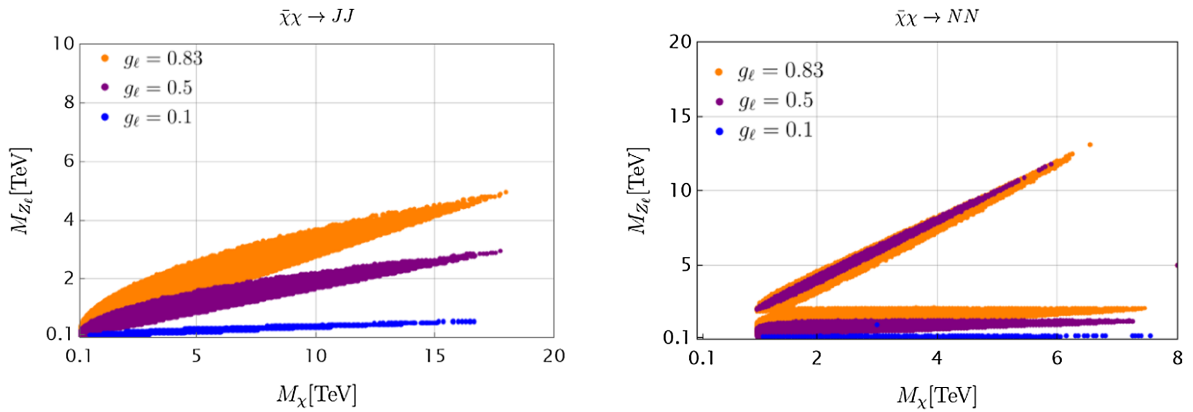


FIG. 6. Allowed region by the cosmological relic density [22] and perturbative bound on the Yukawa coupling y_χ . In the left-panel we show the annihilation to Majorons, while in the right panel we have the annihilation into right-handed neutrinos. For illustration, we used $M_N = 1$ TeV.

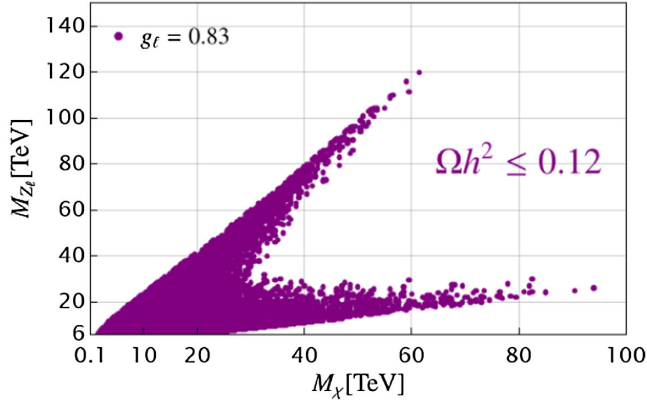


FIG. 7. Allowed region by the cosmological relic density [22] and perturbative bound on the Yukawa coupling y_χ . Here we include all relevant annihilation channels. For illustration, we use $M_{H_1} = M_{H_2} = 1$ TeV, and the mixing between the heavy Higgses is $\theta_H = \pi/6$.

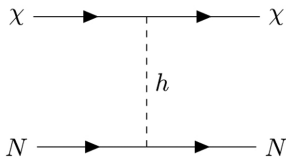


FIG. 8. Feynman diagram for the dark matter–nucleon cross section mediated by the SM-like Higgs boson.

Finally, in Fig. 7 we include all annihilation channels discussed above and show the allowed parameter space by the relic density constraints. As one expects, there are two main regions: (a) The Z_ℓ resonance (less generic) and (b) the region when the dark matter is far from the Z_ℓ resonance and the annihilation channels with the Majoron and the new Higgses are allowed. The second region is more generic because it does not rely in any particular relation of the masses of the fields involved. Clearly, if we

think about the most generic allowed solutions the upper bound on the leptophilic gauge boson mass is around 30 TeV. See Appendix B for the Feynman graphs for each different annihilation channel. The numerical results in Fig. 7 are in agreement with the perturbative unitarity bounds [23]. Notice that these results tell us that the seesaw scale in this theory is below the $\mathcal{O}(10^2)$ TeV scale and one can hope to test the origin of neutrino masses at colliders.

In this theory, the main contribution to the spin-independent dark matter–nucleon cross section is mediated by the SM-like Higgs as we show in Fig. 8. This cross section is given by

$$\sigma_{SI} = \frac{72}{\sqrt{2}} \frac{G_F}{\pi M_h^4} \frac{g_\ell^2 M_\chi^2 \sin^2 \theta}{M_{Z_\ell}^2 \sin^2 \beta} \frac{M_\chi^2 m_N^4}{(M_\chi + m_N)^2} f_N^2. \quad (36)$$

Here m_N is the nucleon mass, G_F is the Fermi constant, and $f_N = 0.3$ [24] is the effective Higgs–nucleon–nucleon coupling. Here $\sin \theta = U_{21} U_{11}$, where the U_{ij} are the elements of the mixing matrix in Eq. (15). Notice that if we include the kinetic mixing between the two Abelian gauge groups the new neutral gauge boson can also mediate this scattering process. See Ref. [25] for the calculation of this cross section at one-loop level. This contribution to the cross section for the dark matter–nucleon scattering is highly suppressed by the mass of the new gauge boson but it is independent of the mixing angle θ . In this theory, the new gauge boson Z_ℓ mediates the dark matter–electron scattering but it is very suppressed by the ratio g_ℓ/M_{Z_ℓ} .

In Fig. 9 we show the numerical results for the spin-independent dark matter–nucleon cross section for different values of g_ℓ and $\sin \theta$. In the case when $g_\ell = 0.83$, the orange and green lines represent the spin-independent cross section for $\sin \theta = 0.01$ and $\sin \theta = 0.1$, respectively. Notice that all numerical results presented in Fig. 9 are valid for the different values of the input parameters, g_ℓ , M_χ , and M_{Z_ℓ} , where one obtains the correct relic density,

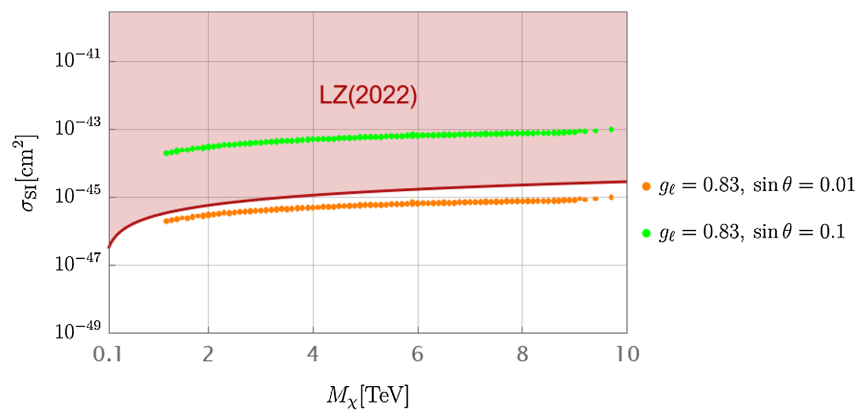


FIG. 9. Numerical values for the spin-independent cross section for dark matter–nucleon scattering. The red shaded region is excluded by the LZ experiment [26]. For $g_\ell = 0.83$, orange and green lines represent the spin-independent cross section for $\sin \theta = 0.01$ and $\sin \theta = 0.1$, respectively. We have used $\sin \beta = 0.54$ (when $v_s = v_\phi$) and $f_N = 0.3$.

i.e., $\Omega_\chi h^2 = 0.12$. As one can appreciate, in this simple theory, one can satisfy the LZ experimental bounds [26] when $\sin \theta$ is below 0.01, which is also in agreement with the experimental bounds of the SM Higgs mixing angle, $\sin \theta < 0.3$. Therefore, one can easily satisfy the current experimental bounds from direct detection experiments.

IV. LEPTON NUMBER VIOLATING SIGNATURES AT THE LHC

As we have discussed above, in this theory the seesaw scale can be below the $\mathcal{O}(10^2)$ TeV scale in agreement with the cosmological bounds on the dark matter relic density, and one can hope to test the mechanism responsible for Majorana neutrino masses at colliders. In this article, we will focus on the testability at the Large Hadron Collider (LHC). The collider bounds on the leptophilic gauge boson, Z_ℓ , are very strong and it will be very difficult to produce it at the LHC. However, the Standard Model-like Higgs and new Higgses, H_1 and H_2 , decays can give rise to lepton number violating signatures. For a previous study of lepton number violating signatures from Higgs decays see for example Refs. [27,28].

One can produce the Higgses at the LHC through gluon fusion and look for the decays into right-handed neutrinos that give rise to lepton number violating signatures with two same-sign charged leptons and four jets:

$$pp \rightarrow h, H_i \rightarrow N_k N_k \rightarrow e_j^\pm e_m^\pm 4j. \quad (37)$$

A. LNV signatures from h decays

The SM-like Higgs can decay into two right-handed neutrinos, if kinematically allowed, through the mixing between the Higgses. The decay width can be written as

$$\Gamma(h \rightarrow NN) = \frac{3}{4\pi} \frac{M_N^2 g_\ell^2 \sin^2 \theta}{M_{Z_\ell}^2 \cos^2 \beta} M_h \left(1 - \frac{4M_N^2}{M_h^2}\right)^{3/2}, \quad (38)$$

and the branching ratio is given by

$$\text{BR}(h \rightarrow NN) = \frac{\Gamma(h \rightarrow NN)}{\cos^2 \theta \Gamma_{\text{SM}} + \Gamma(h \rightarrow NN)}. \quad (39)$$

Here we are assuming for simplicity that all right-handed neutrinos have the same mass. In this case, we can estimate the number of events as given by

$$N_h(e_j^\pm e_m^\pm 4j) = \mathcal{L} \times \sigma(pp \rightarrow h) \times \text{BR}(h \rightarrow NN) \times 2 \times \text{BR}(N \rightarrow e_j^\pm jj) \times \text{BR}(N \rightarrow e_m^\pm jj). \quad (40)$$

In Fig. 10 we show the number of LNV events from the SM-like Higgs decays using the HL-LHC luminosity $\mathcal{L} = 3000 \text{ fb}^{-1}$, $\sigma(pp \rightarrow h) = 54.7 \text{ pb}$, $\cos \beta = 0.84$, and $\sin \theta = 0.01$. Notice that the mixing angle θ has to

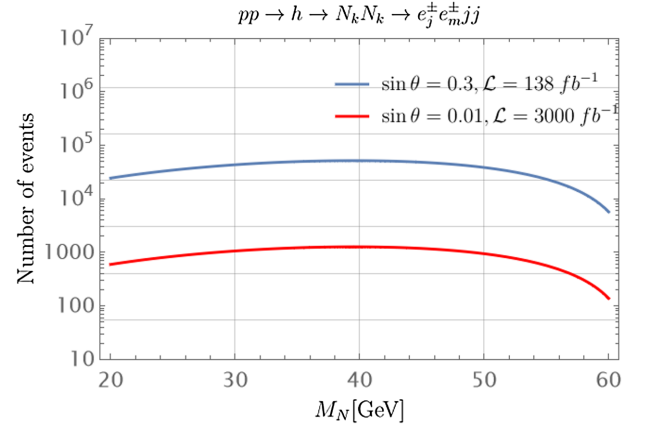


FIG. 10. Lepton number violating events from the SM-like Higgs decays as a function of right-handed neutrino masses (M_N). The red line shows the predictions when the luminosity is $\mathcal{L} = 3000 \text{ fb}^{-1}$, $\cos \beta = 0.84$, $\text{BR}(N_k \rightarrow e_i^\pm jj) = 0.5$, and $\sin \theta = 0.01$. The blue line corresponds to the case when $\mathcal{L} = 138 \text{ fb}^{-1}$ and $\sin \theta = 0.3$.

be small, $\sin \theta \leq 0.01$, to be in agreement with the dark matter direct detection constraints. The branching ratio $\text{BR}(h \rightarrow NN)$ is around 10^{-5} in the scenarios shown in Fig. 10. Since the dark matter direct detection bounds are only valid in a fraction of the parameter space, we show in Fig. 10 the number of events when the luminosity $\mathcal{L} = 138 \text{ fb}^{-1}$ and $\sin \theta = 0.3$. This simple estimation tells us that the HL-LHC could have access to a large number of LNV signatures coming from the SM-like Higgs decays. These signatures can be quite exotic because the light right-handed neutrinos can be long lived and then the LNV decays can give rise to displaced vertices. Clearly, these signatures are quite unique but the predictions depend on the mixing angle θ between the Higgses that could be smaller. The main reducible background for these signals is $pp \rightarrow t\bar{t}W^\pm \rightarrow W^\pm W^\pm jjb\bar{b}$ but by using a set of kinematic cuts one can reduce the background quite effectively. See for example the discussion in Ref. [29] for details.

B. LNV Signatures from H_i decays

The new Higgses can be produced at the LHC though the mixing with the SM-like Higgs and their decays can give rise to LNV signatures. In this case the number of events is given by

$$N_H(e_j^\pm e_m^\pm 4j) = \mathcal{L} \times \sigma(pp \rightarrow H_i) \times \text{BR}(H_i \rightarrow N_k N_k) \times 2 \times \text{BR}(N_k \rightarrow e_j^\pm W^\mp) \times \text{BR}(N_k \rightarrow e_m^\pm W^\mp) \times \text{BR}^2(W^\mp \rightarrow jj), \quad (41)$$

where the hadronic decay of the W boson is $\text{BR}(W^\mp \rightarrow jj) \simeq 2/3$. The right-handed neutrino decay width for $N_i \rightarrow e_j^- W^+$ is given by

$$\Gamma(N_i \rightarrow e_j^- W^+) = \frac{g_2^2}{64\pi M_W^2} |V_{ji}|^2 M_{N_i}^3 \left(1 + 2 \frac{M_W^2}{M_{N_i}^2}\right) \times \left(1 - \frac{M_W^2}{M_{N_i}^2}\right)^2. \quad (42)$$

The matrix V defining the mixing between the right-handed and left-handed neutrinos can be written as [30]

$$V = V_{\text{PMNS}} m^{1/2} R M^{-1/2}, \quad (43)$$

where V_{PMNS} is the Pontecorvo–Maki–Nakagawa–Sakata (PMNS) mixing matrix, $m = \text{diag}(m_1, m_2, m_3)$ is the matrix of the light neutrino masses, and

$M = \text{diag}(M_{N_1}, M_{N_2}, M_{N_3})$ is the matrix for the heavy neutrino masses, and the R matrix is complex and orthogonal. For a detailed study of the $\text{BR}(N_k \rightarrow e^\pm W^\mp)$ taking into account the neutrino constraints and using the freedom in the matrix R , see Ref. [31].

The cross section reads $\sigma(pp \rightarrow H_i) = \sin^2 \theta \sigma(pp \rightarrow H_i)_{\text{SM}}$, where $\sigma(pp \rightarrow H_i)_{\text{SM}}$ is the predicted cross section in the Standard Model changing the Higgs mass. Since this cross section is suppressed by the mixing angle and $\sin \theta \leq 0.01$, to be in agreement with dark matter direct detection constraints, one cannot expect a large number of events even at the HL-LHC.

In the above equation, the branching ratio for the Higgs decays into two right-handed neutrinos is given by

$$\text{BR}(H_i \rightarrow N_k N_k) = \frac{\Gamma(H_i \rightarrow N_k N_k)}{\sin^2 \theta \Gamma(H_i \rightarrow \text{SM SM}) + \Gamma(H_i \rightarrow \bar{\chi}\chi) + \Gamma(H_i \rightarrow N_k N_k) + \Gamma(H_i \rightarrow hh)}, \quad (44)$$

where $H_i \rightarrow \text{SM SM}$ are the decays into SM particles excluding the SM-like Higgs. Since in our case $\sin \theta \leq 0.01$ these decays are very suppressed. The decay width for $H_i \rightarrow N_k N_k$ is given by

$$\Gamma(H_i \rightarrow N_k N_k) = \frac{3}{4\pi} \frac{M_N^2 g_\ell^2 U_{3i}^2}{M_{Z_\ell}^2 \cos^2 \beta} M_{H_i} \left(1 - \frac{4M_N^2}{M_{H_i}^2}\right)^{3/2}, \quad (45)$$

the decay width for $H_i \rightarrow \bar{\chi}\chi$ is

$$\Gamma(H_i \rightarrow \bar{\chi}\chi) = \frac{9}{8\pi} \frac{M_\chi^2 g_\ell^2 U_{2i}^2}{M_{Z_\ell}^2 \sin^2 \beta} M_{H_i} \left(1 - \frac{4M_\chi^2}{M_{H_i}^2}\right)^{3/2}, \quad (46)$$

and the decay width for $H_i \rightarrow hh$ is

$$\Gamma(H_i \rightarrow hh) = \frac{\lambda_{H_i hh}^2}{32\pi} \left(\frac{M_{Z_\ell}}{6g_\ell}\right)^2 \frac{1}{M_{H_i}} \left(1 - \frac{4M_h^2}{M_{H_i}^2}\right)^{1/2} \quad (47)$$

where

$$\lambda_{H_i hh} = 2\lambda_1 U_{2i} \sin \beta + 3\lambda_2 U_{3i} \cos \beta.$$

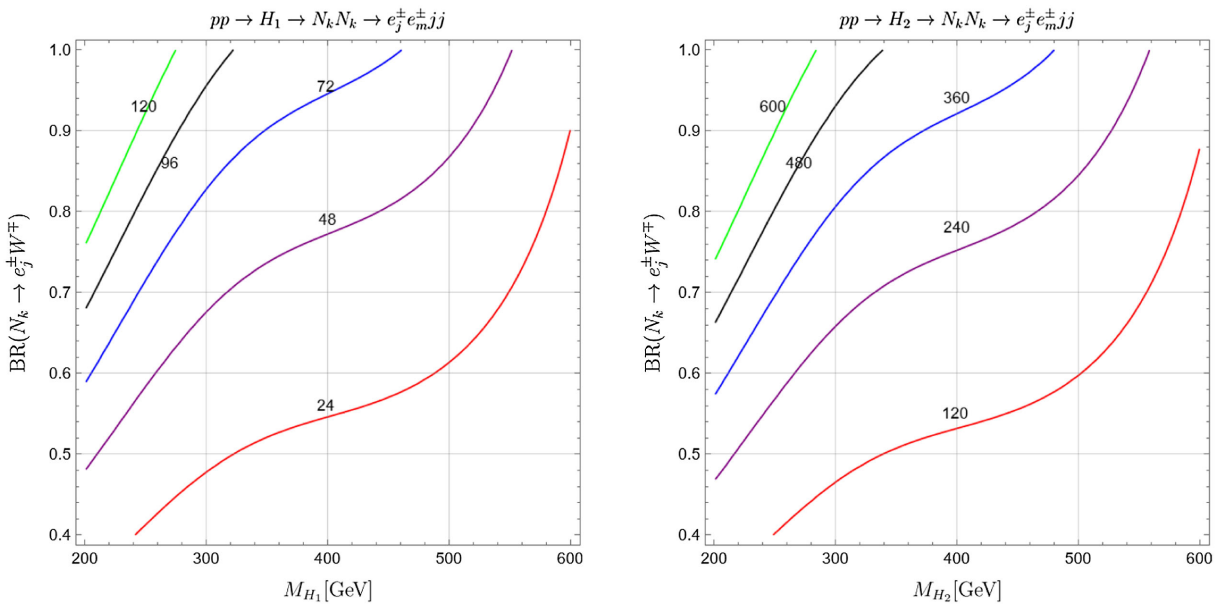


FIG. 11. Number of lepton number violating events as a function of heavy Higgs masses (M_{H_i}) and $\text{BR}(N_k \rightarrow e_j^\pm W^\mp)$. Here we used luminosity $\mathcal{L} = 3000 \text{ fb}^{-1}$ and $\cos \beta = 0.84$, $M_N = M_\chi = 100 \text{ GeV}$, $\sin \theta = 0.01$, $\sin \theta_H = 0.5$, and $M_{Z_\ell}/g_\ell = 7 \text{ TeV}$.

In Fig. 11 we show the predictions for the LNV events at the HL-LHC as a function of heavy Higgs mass, M_{H_i} , and $\text{Br}(N_i \rightarrow e^\pm W^\mp)$. For illustration, we use $\mathcal{L} = 3000 \text{ fb}^{-1}$, $\cos\beta = 0.84$, $M_N = M_\chi = 100 \text{ GeV}$, and $\sin\theta = 0.01$. Here we neglected the decays into two SM-like Higgses to show the most optimistic scenarios. Unfortunately, the number of events from the heavy Higgses are not too large at the LHC due to the fact that the production cross section is suppressed. Therefore, the best hope at the LHC is to look for the lepton number violating signatures from the SM-like Higgs decays into two right-handed neutrinos.

V. SUMMARY

The origin of neutrino masses is one of the most pressing issues in particle physics. In this article, we discussed a class of theories for Majorana neutrinos where the total lepton number is a local gauge symmetry. In order to define an anomaly-free theory based on the total lepton number one needs to add extra fermions including the right-handed neutrinos. These theories predict a fermionic dark matter candidate from anomaly cancellation. The properties of the dark matter candidate were discussed in great detail taking into account all annihilation channels. The dark matter candidate has unique annihilation channels due to the fact that the theory predicts a very light pseudo-Goldstone boson, the Majoron, and one can obtain the correct relic density in a large fraction of the parameter space. Using the cosmological bounds on the dark matter relic density we pointed out the upper bound on the total lepton number symmetry breaking scale.

These theories provide a unique scenario predicting a low scale seesaw mechanism for Majorana neutrino masses since the seesaw scale is below the $\mathcal{O}(10^2)$ TeV scale. We have shown that one could hope to test the origin of neutrino masses at current or future colliders. We discussed the lepton number violating Higgs decays and the possibility to observe lepton number violation at the Large Hadron Collider. We have shown that the SM-like Higgs decays can provide a unique window to the origin of neutrino masses if the lepton number violating decays can be observed in the near future.

ACKNOWLEDGMENTS

The work of H. D. and P. F. P. is supported by the U.S. Department of Energy, Office of Science, Office of High Energy Physics, under Award No. DE-SC0024160. This work made use of the High Performance Computing Resource in the Core Facility for Advanced Research Computing at Case Western Reserve University.

APPENDIX A: FEYNMAN RULES

In our study, the dark matter candidate is the Dirac spinor, $\chi = \chi_L + \chi_R$. Neglecting the mixing between χ and

the neutral components of the fermionic $SU(2)$ doublets, we list the simplified Feynman rules relevant to study the dark matter annihilation channels and direct detection cross section:

$$\bar{\chi}\chi Z_\ell: -ig_\ell\gamma^\mu(\ell_2 P_L + \ell_1 P_R), \quad (\text{A1})$$

$$\bar{\chi}\chi H_i: i\frac{3g_\ell M_\chi}{M_{Z_\ell}}\frac{U_{2i}}{\sin\beta}, \quad (\text{A2})$$

$$\bar{\chi}\chi J: i\frac{3g_\ell}{2M_{Z_\ell}}P_\mu^J\gamma^\mu\gamma^5, \quad (\text{A3})$$

$$Z_\mu^\ell H_i J: i(2P_\mu^J)(3g_\ell \sin\beta U_{2i} + 2g_\ell \cos\beta U_{3i}), \quad (\text{A4})$$

$$NNJ: i\frac{g_\ell}{2M_{Z_\ell}}P_\mu^J\gamma^\mu\gamma^5, \quad (\text{A5})$$

$$\nu\nu J: -i\frac{g_\ell}{2M_{Z_\ell}}P_\mu^J\gamma^\mu\gamma^5, \quad (\text{A6})$$

$$\nu\nu Z_\ell^\mu: i\frac{g_\ell}{2}\gamma^\mu\gamma^5, \quad (\text{A7})$$

$$NNZ_\ell^\mu: -i\frac{g_\ell}{2}\gamma^\mu\gamma^5, \quad (\text{A8})$$

$$\bar{e}e Z_\ell^\mu: -ig_\ell\gamma^\mu, \quad (\text{A9})$$

$$H_i NN: i\frac{g_\ell M_N}{M_{Z_\ell} \cos\beta}U_{3i}, \quad (\text{A10})$$

$$H_i Z_\ell Z_\ell: 2ig_\ell M_{Z_\ell}(2U_{3i} \cos\beta + 3U_{2i} \sin\beta)g_{\mu\nu}. \quad (\text{A11})$$

We have used $v_s = \frac{1}{3}v \sin\beta$ and $v_\phi = \frac{1}{2}v \cos\beta$ to obtain the above Feynman rules. Here we are neglecting the small mixing angle between the SM neutrinos and the right-handed neutrinos. Notice that we are working in the basis where the J interactions are invariant under the shift symmetry. We redefine the fields as follows:

$$\nu_R \rightarrow e^{-i\sigma_\phi/2v_\phi}\nu_R, \quad (\text{A12})$$

$$e_R \rightarrow e^{-i\sigma_\phi/2v_\phi}e_R, \quad (\text{A13})$$

$$\ell_L \rightarrow e^{-i\sigma_\phi/2v_\phi}\ell_L, \quad (\text{A14})$$

$$\chi_L \rightarrow e^{i\sigma_s/2v_s}\chi_L, \quad (\text{A15})$$

$$\chi_R \rightarrow e^{-i\sigma_s/2v_s}\chi_R. \quad (\text{A16})$$

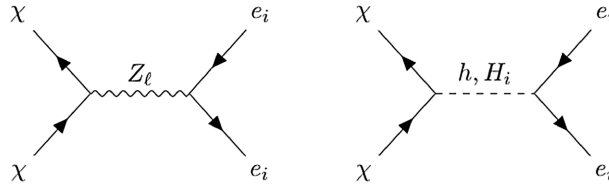
Notice that the Majoron, defined as $J = \cos\beta\sigma_\phi + \sin\beta\sigma_s$, couples to fermions as a pseudo-Goldstone boson.

APPENDIX B: DARK MATTER ANNIHILATION CHANNELS

The annihilation cross sections were calculated using FEYNCALC [32–34]. Our dark matter candidate has the following annihilation channels:

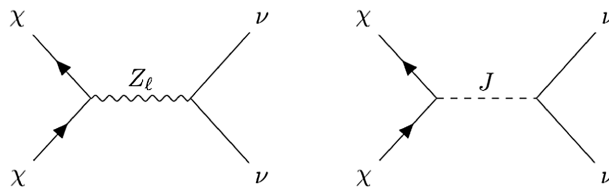
(i)

- $\bar{\chi}\chi \rightarrow \bar{e}_i e_i$



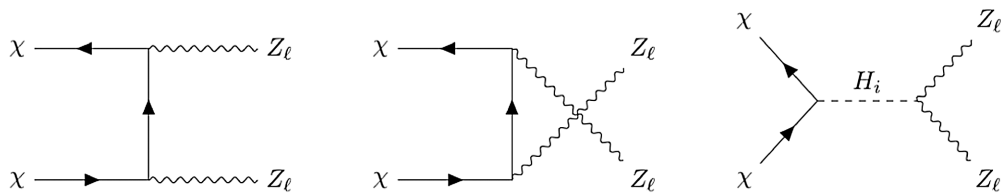
(ii)

- $\bar{\chi}\chi \rightarrow \nu\nu$



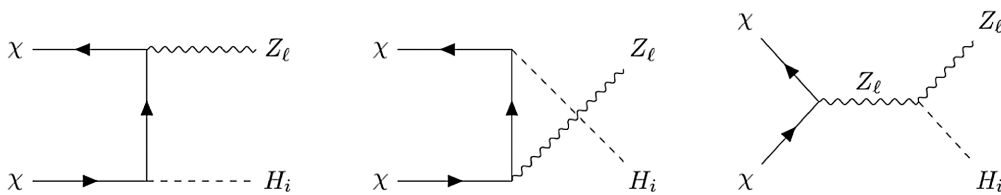
(iii)

- $\bar{\chi}\chi \rightarrow Z_\ell Z_\ell$



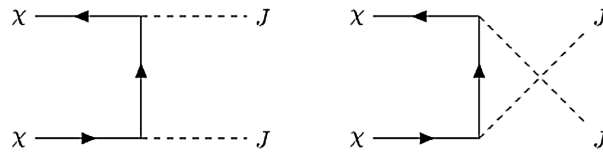
(iv)

- $\bar{\chi}\chi \rightarrow Z_\ell H_i$



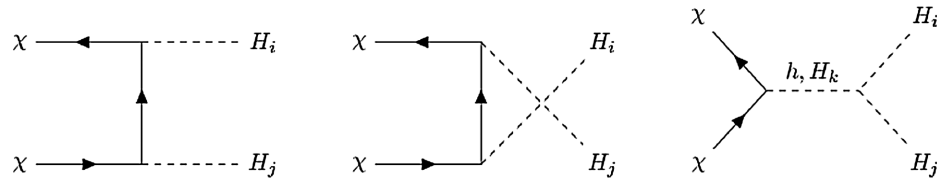
(v)

• $\bar{\chi}\chi \rightarrow JJ$



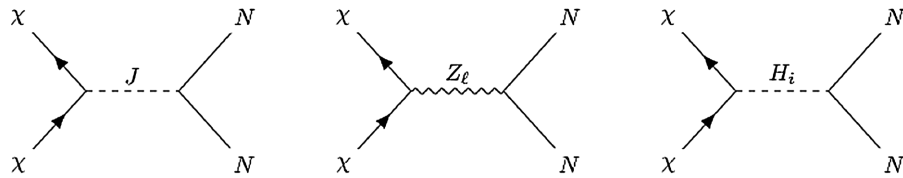
(vi)

• $\bar{\chi}\chi \rightarrow H_i H_j$



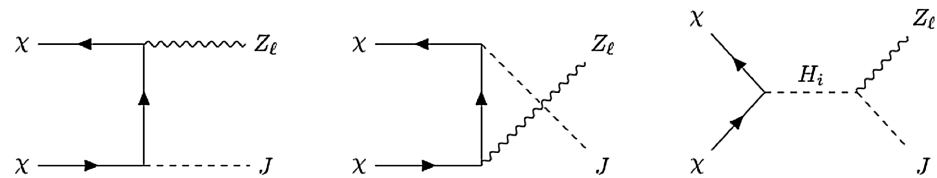
(vii)

• $\bar{\chi}\chi \rightarrow NN$



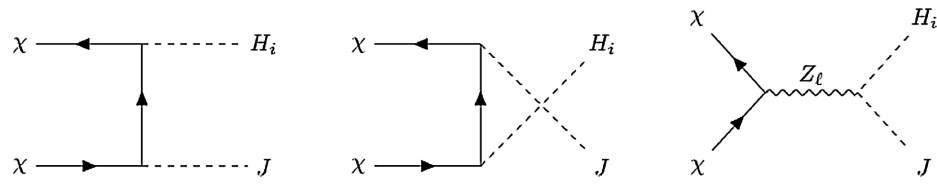
(viii)

• $\bar{\chi}\chi \rightarrow Z_\ell J$



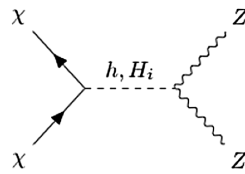
(ix)

• $\bar{\chi}\chi \rightarrow H_i J$



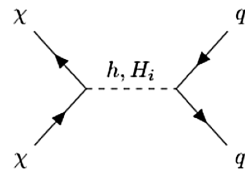
(x)

• $\bar{\chi}\chi \rightarrow ZZ$



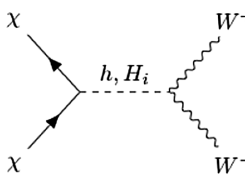
(xi)

• $\bar{\chi}\chi \rightarrow \bar{q}q$



(xii)

• $\bar{\chi}\chi \rightarrow W^+W^-$



- [1] Peter Minkowski, $\mu \rightarrow e\gamma$ at a rate of one out of 10^9 muon decays?, *Phys. Lett.* **67B**, 421 (1977).
- [2] Murray Gell-Mann, Pierre Ramond, and Richard Slansky, Complex spinors and unified theories, *Conf. Proc. C* **790927**, 315 (1979).
- [3] Rabindra N. Mohapatra and Goran Senjanovic, Neutrino Mass and Spontaneous Parity Nonconservation, *Phys. Rev. Lett.* **44**, 912 (1980).
- [4] Tsutomu Yanagida, Horizontal gauge symmetry and masses of neutrinos, *Conf. Proc. C* **7902131**, 95 (1979).
- [5] Wai-Yee Keung and Goran Senjanovic, Majorana Neutrinos and the Production of the Right-Handed Charged Gauge Boson, *Phys. Rev. Lett.* **50**, 1427 (1983).
- [6] Yi Cai, Tao Han, Tong Li, and Richard Ruiz, Lepton number violation: Seesaw models and their collider tests, *Front. Phys.* **6**, 40 (2018).
- [7] Pavel Fileviez Perez *et al.*, On baryon and lepton number violation, [arXiv:2208.00010](https://arxiv.org/abs/2208.00010).
- [8] Pavel Fileviez Perez and Mark B. Wise, Breaking local baryon and lepton number at the TeV scale, *J. High Energy Phys.* **08** (2011) 068.
- [9] Michael Duerr, Pavel Fileviez Perez, and Mark B. Wise, Gauge Theory for Baryon and Lepton Numbers with Leptoquarks, *Phys. Rev. Lett.* **110**, 231801 (2013).
- [10] Pavel Fileviez Perez, Sebastian Ohmer, and Hiren H. Patel, Minimal theory for lepto-baryons, *Phys. Lett. B* **735**, 283 (2014).
- [11] Robert Foot, Girish C. Joshi, and H. Lew, Gauged baryon and lepton numbers, *Phys. Rev. D* **40**, 2487 (1989).
- [12] Y. Chikashige, Rabindra N. Mohapatra, and R. D. Peccei, Are there real goldstone bosons associated with broken lepton number?, *Phys. Lett.* **98B**, 265 (1981).
- [13] Pavel Fileviez Pérez, Clara Murgui, and Alexis D. Plascencia, Neutrino-dark matter connections in gauge theories, *Phys. Rev. D* **100**, 035041 (2019).
- [14] Pedro Schwaller, Tim M.P. Tait, and Roberto Vega-Morales, Dark matter and vectorlike leptons from gauged lepton number, *Phys. Rev. D* **88**, 035001 (2013).
- [15] Marcela Carena, Ying-Ying Li, Tong Ou, and Yikun Wang, Anatomy of the electroweak phase transition for dark sector induced baryogenesis, *J. High Energy Phys.* **02** (2023) 139.
- [16] Eric Madge and Pedro Schwaller, Leptophilic dark matter from gauged lepton number: Phenomenology and gravitational wave signatures, *J. High Energy Phys.* **02** (2019) 048.
- [17] Pavel Fileviez Perez, New paradigm for baryon and lepton number violation, *Phys. Rep.* **597**, 1 (2015).
- [18] Julian Heeck and Hiren H. Patel, Majoron at two loops, *Phys. Rev. D* **100**, 095015 (2019).
- [19] Miguel Escudero and Samuel J. Witte, A CMB search for the neutrino mass mechanism and its relation to the Hubble tension, *Eur. Phys. J. C* **80**, 294 (2020).
- [20] Stefan Sandner, Miguel Escudero, and Samuel J. Witte, Precision CMB constraints on eV-scale bosons coupled to neutrinos, *Eur. Phys. J. C* **83**, 709 (2023).
- [21] Paolo Gondolo and Graciela Gelmini, Cosmic abundances of stable particles: Improved analysis, *Nucl. Phys.* **B360**, 145 (1991).
- [22] N. Aghanim *et al.* (Planck Collaboration), Planck 2018 results. VI. Cosmological parameters, *Astron. Astrophys.* **641**, A6 (2020); *Astron. Astrophys.* **652**, C4(E) (2021).
- [23] Kim Griest and Marc Kamionkowski, Unitarity Limits on the Mass and Radius of Dark Matter Particles, *Phys. Rev. Lett.* **64**, 615 (1990).
- [24] Martin Hoferichter, Philipp Klos, Javier Menéndez, and Achim Schwenk, Improved Limits for Higgs-Portal Dark Matter from LHC Searches, *Phys. Rev. Lett.* **119**, 181803 (2017).
- [25] Joachim Kopp, Viviana Niro, Thomas Schwetz, and Jure Zupan, DAMA/LIBRA and leptonically interacting dark matter, *Phys. Rev. D* **80**, 083502 (2009).
- [26] J. Aalbers *et al.* (LZ Collaboration), First Dark Matter Search Results from the LUX-ZEPLIN (LZ) Experiment, *Phys. Rev. Lett.* **131**, 041002 (2023).
- [27] Elena Accomando, Luigi Delle Rose, Stefano Moretti, Emmanuel Olaiya, and Claire H. Shepherd-Themistocleous, Novel SM-like Higgs decay into displaced heavy neutrino pairs in $U(1)'$ models, *J. High Energy Phys.* **04** (2017) 081.
- [28] Frank F. Deppisch, Wei Liu, and Manimala Mitra, Long-lived heavy neutrinos from Higgs decays, *J. High Energy Phys.* **08** (2018) 181.
- [29] Pavel Fileviez Perez, Tao Han, and Tong Li, Testability of type I seesaw at the CERN LHC: Revealing the existence of the B-L symmetry, *Phys. Rev. D* **80**, 073015 (2009).
- [30] J.A. Casas and A. Ibarra, Oscillating neutrinos and $\mu \rightarrow e, \gamma$, *Nucl. Phys.* **B618**, 171 (2001).
- [31] Pavel Fileviez Pérez and Alexis D. Plascencia, Probing the nature of neutrinos with a new force, *Phys. Rev. D* **102**, 015010 (2020).
- [32] Vladyslav Shtabovenko, Rolf Mertig, and Frederik Orellana, New developments in FeynCalc 9.0, *Comput. Phys. Commun.* **207**, 432 (2016).
- [33] Vladyslav Shtabovenko, Rolf Mertig, and Frederik Orellana, FeynCalc 9.3: New features and improvements, *Comput. Phys. Commun.* **256**, 107478 (2020).
- [34] R. Mertig, M. Böhm, and A. Denner, FeynCalc: Computer-algebraic calculation of Feynman amplitudes, *Comput. Phys. Commun.* **64**, 345 (1991).