Global $U(1)_L$ Breaking in Neutrino 2HDM: From LHC Signatures to X-Ray Line

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Lepton number violation plays an essential role in many scenarios of neutrino mass generation and also provides new clues to search new physics beyond the standard model. We consider the neutrinoophilic two-Higgs-doublet model ($\nu$-2HDM) where additional right-handed neutral fermions $N_{Ri}$ and a complex singlet scalar $\sigma$ are also involved. In scalar sector, the global $U(1)_L$ symmetry is spontaneously broken, leading to Nambu-Goldstone boson, the Majoron $J$, accompanied by the Majorana neutrino mass generation. We find that the massless Majoron will induce large invisible Higgs decay, and current experiments have already set constraints on relevant parameters. For the first time, we point out that the $\nu$-2HDM with $N_{Ri}$ can be distinguished from other seesaw by the same sign tri-lepton signature $3\ell^\pm A j + E_T$. More interesting, for $O$(keV) scale Majoron, it is a good candidate of decaying dark matter to interpret the 3.5 keV and 511 keV line excesses by two different parameter spaces.

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I. INTRODUCTION

In standard model, the total lepton number is conserved at classical level, yet it is violated in many scenarios beyond the standard model. A widely discussed scenario of the lepton number violation (LNV) appears in the models for neutrino mass generation. To explain the no-zero but tiny neutrino mass, the dimension-5 effective operator $f(\Phi_L)(\Phi_L)/\Lambda$ is introduced so that the smallness of neutrino mass is attributed to the seesaw mechanism where lepton number is violated at a scale higher than electroweak scale.

The mechanism of LNV may play a key role in the dark side of our universe. The point is that the pseudo-Nambu-Goldstone boson (pNGB), the Majoron $J$, arises from the spontaneous breaking of global $U(1)_L$ symmetry and picks light mass from quantum gravitational effect. In Ref., Majoron as a keV dark matter (DM) candidate has been studied where high LNV scale (typically $10^3 - 10^6$ TeV) is required to guarantee the small coupling of Majoron with neutrinos and eventually produce a satisfactory DM relic density. Moreover, at one loop level there exists a sub-leading decay of the Majoron to two photons from its coupling to charged fermions, leading to further constraints from x- and $\gamma$-ray experiments. On the other hand, for a TeV LNV scale, the coupling of Majorons to standard model Higgs boson could be large. As a result, the new invisible decay modes of Higgs boson to Majorons is open and provide an interesting route to probe new physics at LHC. The possibility of Majoron as WIMP DM has also been studied where a soft $U(1)_L$ breaking term is added to generate the Majoron mass.

In this paper, we investigate the LNV effect in the context of neutrinoophilic two-Higgs-doublet model (ν-2HDM) where one scalar doublet $\Phi$ gives masses to standard model fermions, while the other scalar doublet $\Phi_{\nu}$ with small vacuum expectation value (VEV) generates the Dirac neutrino mass term. In fermion sector, the neutral right-handed fermion singlets $N_{Ri}$ are introduced to give a natural suppression for the light Majorana neutrino masses. Different from the conventional type-I seesaw model, lepton numbers of $N_{Ri}$ are set to be zero instead of one. In scalar sector, in addition to the SM doublet scalar $\Phi$, a doublet scalar $\Phi_{\nu}$ with lepton number $L = 1$ and a singlet scalar $\sigma$ with $L = 1/2$ are also required to produce the spontaneous LNV process. Hence the scheme we proposed can be called “122” seesaw model in comparison with the “123” seesaw model proposed in Ref. where the “3” denotes the triplet scalar $\Delta$ in type-II seesaw

The scale of LNV is still unknown, hence both low scale and high scale scenarios are considered in this work. In former case, the new massive particles are naturally with electroweak (EW) scale, and thus contribute rich phenomenon at LHC. For instance, a distinct same sign trilepton $3\ell^\pm 4j + E_T$ signature arising from the associated production of neutrinoophilic scalars is unique, and therefore making this model...
quite distinguishable. While for the massless Majoron, it will contribute to invisible decays of Higgs. By choosing certain parameters, we find that a large branching ratio of invisible Higgs decay is possible to escape current experimental constraints. In the scenario with high LNV scale, we postulate the existence of \( O(\text{keV}) \rightarrow O(\text{MeV}) \) Majoron particle, which serves as a late-decaying dark matter. We find that the Majoron can decay into two photons. Hence the current experimental results of X-ray background can set the emission line constraints on the relevant parameters. As already pointed out in Ref. [13], the 3.5 keV x-ray line observed by XMM-Newton observatory [18] can be naturally explained by \( J \rightarrow \gamma \gamma \). In addition, we further consider the 511 keV line from the galactic bulge observed by INTEGRAL experiment [19]. It is suggested that the 511 keV line can be originated from the annihilation of positronium [20–22] or radiative decaying of degenerate fermionic DM [23]. In our model, we suggest that the 511 keV emission line can be originated from the decay of Majoron into low energy electron-positron pairs \( J \rightarrow e^+e^- \). Then the positrons dissipate their kinetic energy by collisions with baryon galactic gas and eventually form the positronium with electrons in the cosmic dust. [24].

The paper is organized as follows. In Sec. II, we introduce the model and describe the details of the symmetry breaking. Possible constraints from astrophysics, lepton flavor violation, and direct collider searches are considered in Sec. III. In Sec. IV A, we discuss the contribution of massless Majoron to invisible decays of Higgs. Collider signatures, especially the LNV signatures, are carried out in Sec. IV B. In Sec. IV C, we consider the Majoron as decaying dark matter and X-ray sources, where the 3.5 keV and 511 keV line excesses are also interpreted. The conclusions are summarised in Sec. V.

II. THE 122 MAJORON MODEL

A. The Model

In addition to SM particles, we introduce a singlet scalar \( \sigma \), a neutrinoiphilic doublet scalar \( \Phi_\nu \), and neutral right-handed fermions \( N_{Ri} \). The representations of new particles are listed in Table. I, where the fields transform under not only SM gauge group but also global \( U(1)_L \) group. The lepton number assignment in Table. I forbids the interaction \( \bar{L}\Phi_\nu N_R \), so that only \( \Phi_\nu \) couples with \( N_R \). The quark and charged lepton sector, on the other hand, are the same as the ones in SM. Thus the FCNCs do not appear at tree level. The relevant interactions are

\[
\mathcal{L}_N = -y_L \bar{L} \Phi_\nu N_R + \frac{1}{2} \overline{N_R} m_N N_R + \text{h.c.}.
\]

(1)

Without loss of generality, we take the diagonal basis for charged leptons and \( N_R \).
TABLE I. New particles content under $G_{SM} \otimes U(1)_L$

| Field | Spin | $SU(3)_c$ | $SU(2)_L$ | $U(1)_Y$ | $U(1)_L$ |
|-------|------|-----------|-----------|----------|----------|
| $N_{Ri}$ | $1/2$ | 1         | 1         | 0        | 0        |
| $\Phi_\nu$ | 0    | 1         | 2         | $1/2$    | 1        |
| $\sigma$ | 0    | 1         | 1         | 0        | $1/2$    |

The complete scalar potential is given by

$$V = -\mu_2^2 \Phi^\dagger \Phi + \mu_3^2 \Phi_\nu^\dagger \Phi_\nu + \lambda_1 (\Phi^\dagger \Phi)^2 + \lambda_2 (\Phi_\nu^\dagger \Phi_\nu)^2 + \lambda_3 (\Phi^\dagger \Phi) (\Phi_\nu^\dagger \Phi_\nu) + \lambda_4 (\Phi^\dagger \Phi_\nu) (\Phi_\nu^\dagger \Phi)$$

$$-\mu_1^2 \sigma^\dagger \sigma + \beta_1 (\sigma^\dagger \sigma)^2 + \beta_2 (\Phi^\dagger \Phi) (\sigma^\dagger \sigma) + \beta_3 (\Phi_\nu^\dagger \Phi_\nu) (\sigma^\dagger \sigma) - k(\Phi^\dagger \Phi_\nu \sigma^2 + \text{h.c.}).$$

(2)

where after acquiring non-zero VEVs, the scalars are denoted as

$$\sigma = \frac{v_1 + R_1 + iI_1}{\sqrt{2}}, \quad \Phi = \left( \begin{array}{c} \phi^+ \\ v_2 + R_2 + iI_2 \\ \sqrt{2} \end{array} \right), \quad \Phi_\nu = \left( \begin{array}{c} \phi_\nu^+ \\ v_3 + R_3 + iI_3 \\ \sqrt{2} \end{array} \right).$$

(3)

The VEV of $\sigma$ breaks the global symmetry $U(1)_L$ spontaneously through the last term in Eq. 2 and also accounts for the generation of Majorana neutrino masses. The minimization conditions are given by

$$\mu_1^2 = -\frac{2 \beta_1 v_1^2 - \beta_2 v_2^2 v_1 - \beta_3 v_3^2 v_1 + 2 k v_1 v_2 v_3}{2 v_1}$$

$$\mu_2^2 = -\frac{2 \beta_1 v_2^2 - \beta_2 v_2^2 v_2 - \beta_3 v_2^2 v_3 + 2 k v_1 v_2 v_3}{2 v_2}$$

$$\mu_3^2 = -\frac{2 \beta_2 v_3^2 - \beta_3 v_2^2 v_3 - \beta_3 v_2^2 v_3 + k v_1^2 v_2}{2 v_3}$$

(4)

Taking the parameter set $\mu_{1,2,3}^2 > 0$ and $k \ll 1$, one can derive the VEV of $\Phi_\nu$ from Eq. 4 as following

$$v_3 \approx \frac{k v_1^2 v_2}{2 \mu_3^2 + \beta_3 v_1^2}.$$  

(5)

One notes that for $\mu_3 \ll v_1$, we have

$$v_3 \approx \frac{k}{\beta_3} v_2,$$

(6)

where $v_3$ is independent to the LNV scale $v_1$. For $\mu_3 \gg v_1$, we have

$$v_3 \approx \frac{k v_1^2}{\mu_3^2} v_2.$$  

(7)

Since $v_3$ is tightly related to tiny neutrino masses, then one expects the VEV hierarchy $v_3 \ll v_2$ in the condition of smallness of $k$ or $\mu_3 \gg v_1$. Notably, $k\Phi^\dagger \Phi_\nu \sigma^2$ term is the only source of $U(1)_L$ breaking, radiative corrections to $k$ are proportional to $k$ itself and are only logarithmically sensitive to the cutoff [16].

Thus, the VEV hierarchy $v_3 \ll v_2 \lesssim v_1$ is stable against radiative corrections [25].
B. Neutrino Masses and Mixing

Note that the term as $\lambda_5/2[(\Phi^\dagger \Phi)^2 + \text{h.c.}]$, which is allowed by discrete $\mathbb{Z}_2$ symmetry, is now forbidden by the global $U(1)_L$ symmetry in our model. As shown in Refs. [17, 26], such $\lambda_5$ term will contribute to one-loop induced neutrino masses, and the radiative induced neutrino masses would be dominant when $\lambda_5 v_5^2/(4\pi)^2 \gtrsim v_3^2$. Due to the forbiddance of such $\lambda_5$ term in our model, the neutrino masses are totally dominantly induced at tree-level as depicted in Fig. 1. Analogical to canonical Type-I seesaw [2], the mass matrix for light neutrinos can be written as

$$m_\nu = -\frac{v_3^2}{2} y_{mn}^{-1} y^T = \hat{m}_\nu U_{\text{PMNS}}^T,$$

where $\hat{m}_\nu = \text{diag}(m_1, m_2, m_3)$ is the diagonalized neutrino mass matrix, and $U_{\text{PMNS}}$ is the PMNS (Pontecorvo-Maki-Nakagawa-Sakata) matrix:

$$U_{\text{PMNS}} = \begin{pmatrix}
 c_{12} c_{13} & s_{12} c_{13} & s_{13} e^{i\delta} \\
 -s_{12} c_{23} - c_{12} s_{23} s_{13} e^{-i\delta} & c_{12} c_{23} - s_{12} s_{23} s_{13} e^{-i\delta} & s_{23} c_{13} \\
 s_{12} s_{23} - c_{12} c_{23} s_{13} e^{-i\delta} & -c_{12} s_{23} - s_{12} c_{23} s_{13} e^{-i\delta} & c_{23} c_{13}
\end{pmatrix} \times \begin{pmatrix}
 e^{i\varphi_1/2} & 0 & 0 \\
 0 & e^{i\varphi_2/2} & 0 \\
 0 & 0 & 1
\end{pmatrix}$$

Here, we use $c_{ij} = \cos \theta_{ij}$ and $s_{ij} = \sin \theta_{ij}$ for short, $\delta$ is the Dirac phase and $\varphi_1, \varphi_2$ are the two Majorana phases. In the following numerical discussion of the phenomenology, we take into account both normal (NH) and inverted hierarchy (IH), and use the latest best fit values of neutrino oscillation parameters in Ref. [27]. For simplicity, the Majorana phases $\varphi_{1,2}$ are neglected in the following numerical discussion. According to Eq. 8, the Yukawa matrix $y$ can be expressed in terms of quantities measured in neutrino masses.

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1 Early works on the global fit of neutrino oscillation can be found in Refs. [28].
oscillation experiments. Since the neutrino masses are induced by Type-I seesaw like mechanism in our model, we could adopt the Casas-Ibarra parametrization \cite{29} to express $y$ as:

$$ y = \frac{\sqrt{2}}{v_3} U_{\text{PMNS}} \sqrt{m_{\nu} R \sqrt{m_N}}, $$

(10)

where $R$ is a complex orthogonal matrix. In the minimal case for two massive neutrinos, $R$ can be expressed in terms of an angle $\omega$ \cite{30} as:

$$ R_{\text{NH}}^\omega = \begin{pmatrix} 0 & 0 & \sqrt{1 - \omega^2} \\ \sqrt{1 - \omega^2} & -\omega & \omega \\ \omega & \sqrt{1 - \omega^2} & 0 \end{pmatrix}, \quad R_{\text{IH}}^\omega = \begin{pmatrix} \sqrt{1 - \omega^2} & -\omega \\ \omega & \sqrt{1 - \omega^2} \\ 0 & 0 \end{pmatrix}, $$

(11)

for the normal (NH) and inverted (IH) hierarchy, respectively. Here in this work, we concentrate on the range $-1 < \omega < 1$. Typically, for $v_3 \sim 1 \text{MeV}$, $m_\nu \sim 0.1 \text{eV}$ and $m_N \sim 100 \text{GeV}$, we have $y \sim 0.01$.

C. Scalar Masses and Mixings

The squared mass matrix for neutral CP-even scalars in the weak basis $(R_1, R_2, R_3)$ is below

$$ M_R^2 = \begin{pmatrix} 2\beta_1 v_1^2 & \beta_2 v_1 v_2 - kv_1 v_3 & \beta_3 v_1 v_3 - kv_1 v_2 \\ \beta_2 v_1 v_2 - kv_1 v_3 & 2\lambda_1 v_2^2 + \frac{1}{2} kv_1^2 v_3^2 & (\lambda_3 + \lambda_4) v_2 v_3 - \frac{1}{2} k v_1^2 \\ \beta_3 v_1 v_3 - kv_1 v_2 & (\lambda_3 + \lambda_4) v_2 v_3 - \frac{1}{2} k v_1^2 & 2\lambda_2 v_3^2 + \frac{1}{2} k v_1^2 v_3^2 \end{pmatrix}. $$

(12)

The $M_R^2$ is diagonalized by orthogonal matrix $O^R$ as $O^R M_R^2 (O^R)^T = \text{diag}(m_{H_1}^2, m_{H_2}^2, m_{H_3}^2)$, where

$$ \begin{pmatrix} H_1 \\ H_2 \\ H_3 \end{pmatrix} = O^R \begin{pmatrix} R_1 \\ R_2 \\ R_3 \end{pmatrix}, $$

(13)

and $O^R$ is parameterized as

$$ O^R = \begin{pmatrix} c_{12} c_{13} & c_{13} s_{12} & s_{13} \\ -s_{12} c_{23} - c_{13} s_{23} & c_{12} c_{23} - s_{13} s_{23} & c_{13} s_{23} \\ s_{23} s_{12} - c_{12} c_{23} s_{13} & -c_{12} s_{23} - c_{23} s_{12} s_{13} & c_{13} c_{23} \end{pmatrix}, $$

(14)

with $c_{ij} = \cos \alpha_{ij}$ and $s_{ij} = \sin \alpha_{ij}$ for short. In our following discussion, we will always keep $H_2$ to be the discovered standard model (SM) like Higgs boson with $m_{H_2} = 125 \text{GeV}$ at LHC \cite{31,33}.

The squared mass matrix for CP-odd scalars in the basis of $(I_1, I_2, I_3)$ is given by

$$ M_I^2 = k \begin{pmatrix} 2 v_2 v_3 - v_1 v_3 & v_1 v_2 \\ -v_1 v_3 & \frac{1}{2} v_1^2 v_3^2 - \frac{1}{2} v_2^2 \\ v_1 v_2 & \frac{1}{2} v_1^2 \frac{1}{2} v_2^2 v_3^2 \end{pmatrix}, $$

(15)
The matrix $M_I^2$ is diagonalized as $O^I M_I^2 (O^I)^T = \text{diag}(0, 0, m_A^2)$, where

\[
\begin{pmatrix}
J \\
G^0 \\
A
\end{pmatrix} = O^I \begin{pmatrix}
I_1 \\
I_2 \\
I_3
\end{pmatrix}.
\]

(16)

As one could expect, two eigenstates with null masses are obtained, corresponding to the normal SM Goldstone boson $G^0$ and the Majoron $J$ generated from global LNV. The $m_A^2$ and matrix $O^I$ are given by

\[
m_A^2 = k \left( \frac{v_1^2 v_3^2 + 4 v_2^2 v_2^2 + v_1^2 v_2^2}{2 v_2 v_3} \right),
\]

(17)

\[
O^I = \begin{pmatrix}
cv_1 V^2 & 32cv_2 v_3^2 & -32cv_2^2 v_3 \\
0 & -\frac{4v_2}{V} & -\frac{4v_3}{V} \\
-\frac{2hv_2}{v_1} & b & -\frac{6v_2}{v_3}
\end{pmatrix},
\]

(18)

with

\[
V^2 = 16(v_2^2 + v_3^2)
\]

\[
c^{-2} = v_1^2 V^4 + 1024(v_2^2 v_3^2 + v_4^2 v_3^2)
\]

\[
b^2 = \frac{v_1^2 v_3^2}{v_2^2 v_1^2 + 4v_3^2 v_3^2 + v_3^2 v_1^2}
\]

(19)

Turning to the charged scalars, the associated squared mass matrix in the basis of $(\phi^\pm, \phi_{\nu}^\pm)$ is given by

\[
M_{H^\pm}^2 = \frac{1}{2} \begin{pmatrix}
kv_1^2 v_3^2 - \lambda_4 v_3^2 & \lambda_4 v_2 v_3 - k v_1^2 \\
\lambda_4 v_2 v_3 - k v_1^2 & k v_1^2 v_3 - \lambda_4 v_2^2
\end{pmatrix}
\]

(20)

Then we have $O^\pm M_{H^\pm}^2 (O^\pm)^T = \text{diag}(0, m_{H^\pm}^2)$, where

\[
\begin{pmatrix}
G^\pm \\
H^\pm
\end{pmatrix} = \begin{pmatrix}
c_\pm & s_\pm \\
-s_\pm & c_\pm
\end{pmatrix} \begin{pmatrix}
\phi^\pm \\
\phi_{\nu}^\pm
\end{pmatrix}
\]

(21)

with $c_\pm = v_2 / \sqrt{v_2^2 + v_3^2}$, $s_\pm = v_3 / \sqrt{v_2^2 + v_3^2}$ and the mass of $m_{H^\pm}$ given by

\[
m_{H^\pm}^2 = \frac{1}{2v_2 v_3} \left( v_2^2 + v_3^2 \right) (kv_1^2 - \lambda_4 v_2 v_3)
\]

(22)

Taking into account the smallness of $v_3$ and $k$ one notices from Eq. (12) that for the neutral scalars $H_3$ and $A$, the following mass relation holds approximately

\[
m_A^2 \simeq k v_1^2 v_2 / 2v_3 = [M_R^2]_{33} \simeq m_H^2
\]

(23)
In the same way, from Eq. (17) and (22) one derives the mass relation
\[ m_A^2 - m_{H^+}^2 \approx \frac{\lambda_4}{2} v_2^2, \]  
which implies that the differences between \( m_A \) and \( m_{H^+} \) can not be too large under perturbativity condition. For simplicity, we will assume that masses of neutrinophilic scalars are degenerate, i.e., \( m_{H^+} = m_{H_3} = m_A \equiv m_{\Phi_\nu} \).

### III. CONSTRAINTS

#### A. Theoretical Constraints

Using Eqs. (12), (17) and (22), we rewrite all the coupling constants \( \lambda_i \) and \( \beta_j \) in terms of mixing angles \( \alpha_{ij} \) and scalar masses

\[
\beta_1 = \frac{1}{2v_1^2} [M_R^2]_{11}
\]
\[
\beta_2 = \frac{1}{v_1 v_2} [M_R^2]_{12} + k \frac{v_3}{v_1}
\]
\[
\beta_3 = \frac{1}{v_1 v_3} [M_R^2]_{13} + k \frac{v_2}{v_3}
\]
\[
\lambda_1 = \frac{1}{2v_2^2} [M_R^2]_{22} - k \frac{v_1 v_3}{4v_2^2}
\]
\[
\lambda_2 = \frac{1}{2v_3^2} [M_R^2]_{33} - k \frac{v_1 v_2}{4v_3^2}
\]
\[
\lambda_3 = \frac{1}{v_2 v_3} [M_R^2]_{23} - \lambda_4 + k \frac{v_1^2}{2v_2 v_3}
\]
and

\[
\lambda_4 = \frac{1}{v_2 v_3} (k v_1^2 - \frac{2v_2}{v_2 + v_3} m_{H^+}^2)
\]
\[
k = \left( \frac{2v_2 v_3}{v_1^2 v_2^2 + v_3^2 v_2^2 + 4v_2^2 v_3^2} \right) m_A^2
\]

where \([M_R^2]_{ij}\) denotes the matrix elements of \( M_R^2 \).

The scalar potential is bounded from below if the quartic part of scalar potential is positive in the non-negative basis. In the following, we take the same procedure in Ref. [8] [9] [11]. Taking into account the fact of \( v_3 \ll v_1 \) and using Eq. (26), we derive the parameter \( k \) as

\[ k \approx m_A^2 \frac{2v_3}{v_1^2 v_2} \]  

Therefore, we have \( k \ll \lambda_i, \beta_i \) and the parameter \( k \) can be neglected with respect to other coupling constants. In this limit, the copositive criteria [34] can be applied to the quartic part of scalar potential to give
the boundedness condition as following

\[ \begin{align*}
\lambda_1 &> 0, \quad \lambda_2 > 0, \quad \beta_1 > 0 \\
x & = \lambda_3 + \theta(-\lambda_4)\lambda_4 + 2\sqrt{\lambda_1\lambda_2} > 0 \\
y & = \beta_2 + 2\sqrt{\lambda_1\beta_1}, \\
z & = \beta_3 + 2\sqrt{\lambda_2\beta_1} \\
\sqrt{\lambda_1\lambda_2\beta_1} + [\lambda_3 + \theta(-\lambda_4)\lambda_4]\sqrt{\beta_1} + \beta_2\sqrt{\lambda_2} + \beta_3\sqrt{\lambda_1} + \sqrt{xyz} &> 0
\end{align*} \]  

(28)

In addition, we set the values of coupling constants \( \lambda_i \) and \( \beta_j \) less than \( \sqrt{4\pi} \) to ensure the perturbative condition.

**B. Astrophysical Constraints**

![Graph](image)

**FIG. 2.** Allowed region of \( v_3 \) as a function of \( v_1 \) considering the constraint from Majoron-electron coupling.

Note that \( \rho = 1 \) at tree level in this \( \nu \)2HDM. The stringent constraint on \( v_3 \) comes from astrophysics, due to the contributions of Majoron-electron coupling \( g_{Jee} \) to supernova \[35\] and red giant cooling \[36\]. For a massless Majoron (or lighter than typical stellar temperatures), the Compton-like process \( \gamma + e \rightarrow J + e \) sets an upper bound for the \( g_{Jee} \) coupling as \[35, 36\]:

\[ |g_{Jee}| = |O_{12}^T \frac{m_e}{v_2}| \lesssim 1.4 \times 10^{-13}. \]  

(29)

Considering the profile of Majoron \[37\] in Eq. \[46\], we can translate this as a bound on the projection of the Majoron onto the doublet \( \Phi \) as \[38\]:

\[ |\langle J | \Phi \rangle| = \frac{2v_2v_3^2}{\sqrt{v_1^2(v_2^2 + v_3^2)^2 + 4v_2^2v_3^4 + 4v_2^4v_3^2}} \approx \frac{2v_3}{v_1v_2} \lesssim 6.7 \times 10^{-8}. \]  

(30)
where in above approximation, we have used the assumption that $v_3 \ll v_1, v_2$. So from Eq. 30 we expect that $v_3^{Max} \propto \sqrt{v_1}$. The allowed region of $v_3$ as a function of $v_1$ is presented in Fig. 2. For instance, $v_3 \lesssim 0.09\text{GeV}$ must be satisfied when $v_1 = 1000\text{GeV}$.

### C. Lepton Flavor Violation

![Graph showing BR($\mu \to e\gamma$) as a function of $\omega$ for four different values of $v_3$ and $m_{H^+}$ in NH (left panel) and IH (right panel).]

We find that the lepton flavor violating (LFV) processes would set a much more stringent lower bound on $v_3$ in models with heavy exotic leptons [39] than the canonical type-II seesaw [40, 41] as well as the Dirac neutrino scenario of $\nu$2HDM [42]. In this paper, we simply take the $\mu \to e\gamma$ process to illustrate such tight constraints, since the MEG experiment sets a severe upper limit as $\text{BR}(\mu \to e\gamma) < 5.7 \times 10^{-13}$ [43]. We also consider the future sensitivity of MEG experiment, which might be down to $6 \times 10^{-14}$ [44]. The branching ratio of $\mu \to e\gamma$ is calculated as [45]:

$$
\text{BR}(\mu \to e\gamma) = \frac{3\alpha}{64\pi G_F^2} \left| \sum_i y_{\mu i} y_{ei}^* \frac{m_{N_i}^2}{m_{H^+}^2} F \left( \frac{m_{N_i}^2}{m_{H^+}^2} \right) \right|^2,
$$

(31)

where the loop function $F(x)$ is:

$$
F(x) = \frac{1 - 6x + 3x^2 + 2x^3 - 6x^2 \ln x}{6(1 - x)^4}.
$$

(32)

In Fig. 3, we show the numerical results of $\text{BR}(\mu \to e\gamma)$ as a function of $\omega$ for $(v_3, m_{H^+}) = (1\text{MeV}, 300\text{GeV}), (1\text{MeV}, 600\text{GeV}), (3\text{MeV}, 300\text{GeV})$ and $(3\text{MeV}, 600\text{GeV})$ in both normal and inverted hierarchy. In case of
normal hierarchy, the present MEG bound \cite{43} requires $v_3 m_{H^+} \gtrsim 600\text{MeV} \cdot \text{GeV}$, meanwhile the future MEG sensitivity \cite{44} would push this bound up to $v_3 m_{H^+} \gtrsim 900\text{MeV} \cdot \text{GeV}$. On the other hand in case of inverted hierarchy, the bound on $v_3 m_{H^+}$ is slightly less stringent than the the bound of normal hierarchy. Briefly, we can conclude that to satisfy the LFV constraint, $v_3 \gtrsim \mathcal{O}(\text{MeV})$ is needed for $m_{H^+} \sim \mathcal{O}(\text{TeV})$.

In general, LFV processes depends on neutrino masses, mixing angles, Dirac phase, as well as Majorana phases. In our assumption with degenerate $N_R$ and real $R$ matrix, we obtain

$$\sum_i y_{\mu i} y_{ei}^* \propto U_{PMNS} \tilde{m}_\nu U_{PMNS}^\dagger = c_{12} c_{13} s_{12} c_{23} (m_2 - m_1) + c_{13} s_{13} s_{23} e^{-i\delta} [(m_3 - m_2) + c_{12}^2 (m_2 - m_1)]$$

(33)

Therefore, the $\mu \to e\gamma$ sets no constraint on the Majorana phases and the $R$ matrix. But for a large $s_{13}$, the branch ratio is sensitive to the Dirac phase $\delta$.

Comparing to the bound on type-II seesaw $v_\Delta m_{H^{++}} \gtrsim \frac{250}{130} \text{eV} \cdot \text{GeV} \\cite{40}$ and Dirac scenario of $\nu 2\text{HDM} v_3 m_{H^+} \gtrsim 600\text{MeV} \cdot \text{GeV}$ is about 6 orders of magnitude higher. Here we take Dirac and Majorana scenario of $\nu 2\text{HDM}$ to briefly estimate such great difference. From Eq. \ref{33}, it is clear that the constraint from LFV actually requires about the same order of the Yukawa coupling $y$, since the loop function $F(x)$ is of the same order in both Dirac and Majorana scenario if we also assume $m_N < m_{H^+}$. In Dirac scenario, $y^D \sim m_\nu/v_3^D$, while in Majorana scenario, $y^M \sim \sqrt{m_\nu m_N}/v_3^M$. For the same order of the Yukawa coupling, we could estimate that $v_3^M/v_3^D \sim \sqrt{m_N/m_\nu} \sim 10^6$ with $m_N \sim 10^2\text{GeV}$ and $m_\nu \sim 0.1\text{eV}$, which is just the result of the above discussion.

### D. Collider Constraints

The status of the Higgs singlet $H_1$ has been extensively studied in Refs. \cite{46,49}. We refer to Ref. \cite{49} for a more detail and updated study on the constraints of $H_1$. In the high mass region with $m_{H_1} > 130\text{GeV}$, the allowed value for $\sin \alpha_{12}$ as a function of $m_{H_1}$ is shown in FIG.1 of Ref. \cite{49}. For example, $\sin \alpha_{12} \lesssim 0.3$ is required with $m_{H_1} = 300\text{GeV}$. Although the invisible decay $H_1 \to JJ$ could affect the direct search bound to be less stringent, the indirect bound as from Higgs signal rate still requires $\sin \alpha_{12} < 0.36$ \cite{49}. So we will consider $\sin \alpha_{12} = 0.3, 0.2, 0.1$ with $m_{H_1} = 300\text{GeV}$ as our benchmark points for the high mass region in Sec. \text{IV} A.

However, in the low mass region with $m_{H_1} < 120\text{GeV}$, the stringent bound in Ref. \cite{49} is not applicable to our model. Mainly because the invisible decay $H_1 \to JJ$ is totally dominant in this region (see the detail}

\footnote{$v_\Delta$ is the vacuum expectation value of Higgs triplet $\Delta$, and $m_{H^{++}}$ is the mass of the doubly-charged scalar $H^{++}$ in $\Delta$.}
FIG. 4. Constraints on $\sin \alpha_{12}$ in the low mass region.

study in Sec. [IV A]. In FIG. 4, we show the constraints on $\sin \alpha_{12}$ in the low mass region coming from LHC Higgs signal rate [49], visible ($H_1 \rightarrow b\bar{b}$) [50] and invisible decaying ($H_1 \rightarrow JJ$) [51] Higgs at LEP through $ZH_1$ associated production. Note that for the constraint from $H_1 \rightarrow b\bar{b}$ we show the most severe case with $BR(H_1 \rightarrow b\bar{b}) = 1$. If we take into account a more realistic $BR(H_1 \rightarrow b\bar{b})$, the exclusion region will be $\sin \alpha_{12} > 0.4$ and less stringent than those from Higgs signal rate [10, 11]. For $m_{H_1} = 50$GeV, the most strict bound comes from invisible Higgs search at LEP with $\sin \alpha_{12} \lesssim 0.2$. Therefore, we will take $\sin \alpha_{12} = 0.2, 0.1, 0.05$ with $m_{H_1} = 50$GeV as our benchmark points for the low mass region in Sec. [IV A].

The collider signature of $\nu 2$HDM has been discussed in Refs. [16, 17, 52]. In the case of $m_N > m_{H^+}$, the dominant decay mode of $H^+$ could be $H^+ \rightarrow \ell^+ \nu$. The direct search for signature as $\ell^+ \ell^- + E_T$ at LHC has excluded the region of $m_{H^+} \lesssim 300$GeV [53, 54]. While in the case of $m_N < m_{H^+}$, the dominant decay mode of $H^+$ would be $H^+ \rightarrow \ell^+ N_{Ri}$ with the heavy Majorana neutrino $N_{Ri}$ further decaying into $\ell^\pm W^{\mp}, \nu Z$ and $\nu H_2$. A detail discussion and simulation at LHC of this case is still missing. Therefore, we consider the LEP bound on charged scalar, i.e., $m_{H^+} > 80$GeV [55]. And also to satisfy the constraints from electroweak precision tests (EWPT) [56], we further assume that the masses of neutrinoophilic doublet scalars are degenerate as $m_{H^+} = m_{H_3} = m_A = m_{\Phi_v}$.

Since the heavy Majorana neutrino $N_R$ also exists in canonical type-I seesaw [2], searches for $N_R$ are already well studied [57–66]. For more detail, see the recent review of neutrino and collider in Ref. [67] and references therein. Direct searches for $N_R$ at colliders have also been performed at LEP [68, 69] and LHC [70–72]. For $m_N < m_W$, LEP has excluded the mixing $V_{\ell N}$ between the heavy Majorana neutrino $N$ and
the neutrino of flavor $\nu_\ell$ with $|V_{\ell N}|^2 \gtrsim 2 \times 10^{-5}$ \cite{68}. For a more heavier $N_R$, LHC would give the most restrictive direct limits. For instance, at $m_N = 200\text{GeV}$ the limit is $|V_{\ell N}|^2 < 0.017$ and at $m_N = 500\text{GeV}$ the limit is $|V_{\ell N}|^2 < 0.017$. In $\nu$2HDM, the mixing $V_{\ell N}$ is predicted as \cite{72}:

$$V_{\ell N} = U_{\text{PMNS}}^\dagger m_\nu^{1/2} R_{\text{m}_{1/2}} \sim 10^{-6}$$

(34)

for EW-scale $m_N$, which is far below current limits.

IV. PHENOMENOLOGY

As pointed in introduction, the LNV scale is still unknown, hence both low scale and high scale scenarios are allowed. For low scale scenario, our model is a natural TeV-model, so it can be test at LHC as we will discuss in SEC. IV A and IV B. On the other hand, if the Majoron are assumed to be a DM candidate, a high LNV scale $v_1 \gtrsim 10^4 \text{TeV}$ are needed to satisfy constraints from WMAP. In high scale scenario, the natural way to get correct neutrino mass is keeping new scalar masses (i.e. $m_{H^+}, m_{H_2}$ and $m_A$) around $v_1$ scale. So new scalars in case of massive Majoron are out reach of LHC. The possible signatures of massive Majoron will be discussed in SEC. IV C.

A. Invisible Higgs Decay

Due to the existence of massless Majoron $J$, the Higgs scalar can decay into Majorons though $H_a \rightarrow JJ$ and $H_a \rightarrow H_b H_b \rightarrow 4J$ \cite{10,11,24,75}. The Higgs-Majoron couplings are derived as:

$$g_{H_a JJ} = \left( \frac{(O^I_{11})^2}{v_1} O^R_{1a} + \frac{(O^I_{12})^2}{v_2} O^R_{2a} + \frac{(O^I_{13})^2}{v_3} O^R_{3a} \right) m^2_{H_a},$$

(35)

where $O^R$ and $O^I$ are the mixing matrices for CP-even and CP-odd scalars in Eq. 14 and 18. The partial decay width of $H_a \rightarrow JJ$ is then given by:

$$\Gamma(H_a \rightarrow JJ) = \frac{1}{32\pi} \frac{g_{H_a JJ}^2}{m_{H_a}}.$$

(36)

Similar to the type-II seesaw case \cite{11}, the smallness of $v_3$ indicates that the neutrinophilic doublet $\Phi_\nu$ is basically decoupled. So we concentrate on the invisible decays of $H_1$ and $H_2$. Note that for light $m_{H_3/A} < m_{H_2}/2$, $H_3$ or $A$ could also mediate a sizable invisible decay of $H_2$\cite{76}. The trilinear coupling $H_2 H_1 H_1$ (see Eq. 55 in the appendix) contributes to invisible decay of SM Higgs $H_2$ with $H_1 \rightarrow JJ$. When $m_{H_1} < m_{H_2}/2$, the decay mode $H_2 \rightarrow H_1 H_1$ would be kinematically open. The partial decay width for $H_2 \rightarrow H_1 H_1$ is computed as:

$$\Gamma(H_2 \rightarrow H_1 H_1) = \frac{g_{H_2 H_1 H_1}^2}{32\pi m_{H_2}} \left( 1 - \frac{4m^2_{H_1}}{m^2_{H_2}} \right)^{1/2}.$$

(37)
For the decays of $H_1$ into SM particles, we refer Ref. [46] for a more detail description. In this paper, the benchmark points discussed in Sec. [III D] are used to illustrate the invisible decays of $H_1$ and $H_2$. Varying the parameters, we find that the most relevant parameters for the invisible Higgs decays are $\sin \alpha_{12}$, $v_1$, and $m_{H_1}$. For simplicity, we fix the value of the following parameters as:

\[
m_{H^+} = m_{H_3} = m_A = m_{\Phi_\nu} = 500 \text{ GeV},
\]
\[
v_3 = 2 \text{ MeV}, \sin \alpha_{13} = \sin \alpha_{23} = 2 \times 10^{-6},
\]

(38)

meanwhile we vary the values of parameters $\sin \alpha_{12}$, $v_1$, and $m_{H_1}$ as:

\[
\sin \alpha_{12} \in [0, 0.3], \; v_1 \in [500, 1500],
\]
\[
m_{H_1} \in [10, 100] \text{ GeV}, \; \text{for the low mass region},
\]
\[
\in [200, 1000] \text{ GeV}, \; \text{for the high mass region}.
\]

(39)

It is checked that the above region is allowed by the boundedness conditions in Eq. 28. A fully scanning the whole parameter space as done in Refs. [10, 11] is worthwhile but beyond the scope of this paper. In the following, we will give some qualitative discussion which is helpful to better understand the scanning results of Refs. [10, 11].

1. High Mass Region:

FIG. 5. Branching ratios of heavy singlet scalar $H_1$ as a function of $m_{H_1}$ for $\sin \alpha_{12} = 0.3, 0.2, 0.1$, respectively. Here, we also set $v_1 = 1000 \text{GeV}$.

First, we explore the high mass region $m_{H_1} > m_{H_2}$. In FIG. 5, we show the branching ratios of $H_1$ as a function of $m_{H_1}$ for $\sin \alpha_{12} = 0.3, 0.2, 0.1$ with $v_1 = 1000 \text{GeV}$. Clearly, the smaller $\sin \alpha_{12}$ is, the bigger $\text{BR}(H_1 \rightarrow JJ)$ is, thus the bigger invisible decay of $H_1$ is. For $\sin \alpha_{12} = 0.3$, $\text{BR}(H_1 \rightarrow JJ) \approx 0.14$ when $m_{H_1} > 350 \text{GeV}$, which is smaller than $\text{BR}(H_1 \rightarrow W^+W^-, H_2H_2, ZZ)$. While for $\sin \alpha_{12} = 0.1$,
BR($H_1 \to JJ$) $\approx 0.6$, which is the dominant decay channel and makes $H_1$ quite different from the real scalar singlet in Refs. [46–49]. If $m_{H_1} > 2m_{H_2}$, then $H_1 \to H_2H_2 \to 4J$ will also contribute the invisible decay of $H_1$, but BR($H_1 \to 4J$) is expected to be less than $0.3 \times 0.23^2 \approx 0.016$. On the other hand, BR($H_1 \to JJ$) is at least 0.14. Hence, the invisible decay of $H_1$ is dominant by $H_1 \to JJ$.

![Graphs](image)

FIG. 6. (a) Branching ratios of invisible $H_2$ decay as a function of $v_1$ for $\sin \alpha_{12} = 0.3, 0.2, 0.1$. (b) Relations between BR($H_1 \to \text{inv.}$) and BR($H_2 \to \text{inv.}$) for $\sin \alpha_{12} = 0.3, 0.2, 0.1$ by varying $v_1$ in the range of 500–1500GeV. (c) Relations between BR($H_1 \to \text{inv.}$) and BR($H_2 \to \text{inv.}$) for $v_1 = 500, 1000, 1500$GeV by varying $\sin \alpha$ in the range of 0–0.3. In all these figures, we have set $m_{H_1} = 300$GeV. The dashed and dotted line correspond to current (0.23) [77] and future limit ($\approx 0.1$) [78] on BR($H_2 \to \text{inv.}$).

In the high mass region with $m_{H_1} > m_{H_2}$, the invisible $H_2$ decay is dominant by $H_2 \to JJ$. In FIG. 6(a), we show BR($H_2 \to \text{inv.}$) vs. $v_1$ for $\sin \alpha_{12} = 0.3, 0.2, 0.1$. It is obvious that a smaller $\sin \alpha_{12}$ ($v_1$) will lead to smaller (larger) BR($H_2 \to \text{inv.}$) for fixed $v_1$ ($\sin \alpha_{12}$). Considering current bound on BR($H_2 \to \text{inv.}$)[77], $v_1 \geq 1200(800)$GeV is required for $\sin \alpha_{12} = 0.3(0.2)$. We then show the relations between BR($H_1 \to \text{inv.}$) and BR($H_2 \to \text{inv.}$) with $m_{H_1} = 300$GeV for $\sin \alpha_{12} = 0.3, 0.2, 0.1$ by varying $v_1$ in the range of 500–1500GeV in FIG. 6(b), and for $v_1 = 500, 1000, 1500$GeV by varying $\sin \alpha$ in the range of 0–0.3 in FIG. 6(c). From these figures, we expect a positive (negative) correlation between BR($H_1 \to \text{inv.}$) and BR($H_1 \to \text{inv.}$) for varying $v_1$ ($\sin \alpha_{12}$).

2. Low Mass Region:

Then, we study the low mass region $m_{H_1} < m_{H_2}$. In FIG. 7 we depict the branching ratios of $H_1$ as a function of $m_{H_1}$ for $\sin \alpha_{12} = 0.3, 0.2, 0.1$ with $v_1 = 1000$GeV. Different from the high mass region, $H_1 \to JJ$ is totally dominant in the low mass region for all allowed values of $\sin \alpha_{12}$. We expect BR($H_1 \to \text{inv.}$) $\gtrsim 0.9$. The dominant visible decay of $H_1$ is $H_1 \to b\bar{b}$, and BR($H_1 \to b\bar{b}$) is typically

\footnote{inv. is short for invisible.}
FIG. 7. Same as FIG. 5 but for $\sin \alpha_{12} = 0.2, 0.1, 0.05$ in the low mass region of $m_{H_1}$.

less than 0.1. So exotic $H_2$ decays as $H_2 \rightarrow H_1 H_1 \rightarrow b \bar{b} + \mathcal{E}_T$ as well as $H_2 \rightarrow H_1 H_1 \rightarrow 4b$ would be challenging at LHC [79] for these benchmark points.

FIG. 8. Same as FIG. 5 but for $\sin \alpha_{12} = 0.2, 0.1, 0.05$ with $m_{H_1} = 50$GeV.

In the low mass region with $m_{H_1} < m_{H_2}/2$, $H_2 \rightarrow H_1 H_1 \rightarrow 4J$ will contribute to $H_2 \rightarrow \text{inv.}$.

Fixed $m_{H_1} = 50$GeV, we present BR($H_2 \rightarrow \text{inv.}$) vs. $v_1$ in FIG. 8 (a), BR($H_2 \rightarrow \text{inv.}$) vs. BR($H_1 \rightarrow \text{inv.}$) for $\sin \alpha_{12} = 0.2, 0.1, 0.05$ in FIG. 8 (b), and BR($H_2 \rightarrow \text{inv.}$) vs. BR($H_1 \rightarrow \text{inv.}$) for $v_1 = 500, 1000, 1500$GeV in FIG. 8 (c). All the qualitative arguments in the high mass region are also applicable here. But due to contributions of $H_2 \rightarrow 4J$, bound on $v_1$ is slightly higher than it in the high mass case with same $\sin \alpha_{12}$. Since both $H_2 \rightarrow JJ$ and $H_2 \rightarrow 4J$ contribute to $H_2 \rightarrow \text{inv.}$, we quantize the contribution of $H_2 \rightarrow JJ$ to $H_2 \rightarrow \text{inv.}$ by defining:

$$R_{JJ} = \frac{\Gamma(H_2 \rightarrow JJ)}{\Gamma(H_2 \rightarrow \text{inv.})} = \frac{\Gamma(H_2 \rightarrow JJ)}{\Gamma(H_2 \rightarrow JJ) + \Gamma(H_2 \rightarrow 4J)}.$$ (40)

In FIG. 9 we plot the contour of $R_{JJ}$ in the $\sin \alpha_{12}$ vs. $v_1$ plane. For $\sin \alpha_{12} > 0.01$, $R_{JJ} > 0.5$, which
indicates that $H_2 \rightarrow JJ$ is the dominant contribution to invisible $H_2$ decays in quite a large parameter space. And from FIG. 9 we could conclude that the larger $v_1$ or the smaller $\sin \alpha$ is, the smaller the contribution of $H_2 \rightarrow JJ$ to $H_2 \rightarrow \text{inv.}$ is.

FIG. 9. Fraction of $H_2 \rightarrow JJ$ to total invisible $H_2$ decay with $m_{H_1} = 50\text{GeV}$ in the $\sin \alpha_{12}$ vs. $v_1$ plane.

B. Collider Signatures

Early papers on collider phenomenon of the $\nu$2HDM with heavy Majorana neutrino $N_{Ri}$ can be found in Refs. [14, 17, 80], and they mainly concentrate on the charged scalar $H^+$. Following Ref. [17], we give a brief discussion of the signatures at LHC by taking into account the contribution of neutral scalars $H_3$ and $A$ in neutrinophilic 2HDM. In FIG. 10 we show the cross section of pair and associate production of the neutrinoophilic doublet scalars at 14 TeV LHC. Typically for EW-scale $m_{\Phi\nu}$, the cross sections are at the order of $O(\text{fb})$. The cross section of associate production $H^+H_3/A$ is about twice larger than it of the pair production $H^+H^-$ or $H_3A$.

The decay properties of the neutrophilic doublet and the heavy Majorana neutrino are discussed in Ref. [17]. For $m_{\Phi\nu} < m_N$, the dominant decay mode of $H^+$ is $H^+ \rightarrow \ell^+\nu_i$ with $v_3 \lesssim O(\text{MeV})$ due to the mixing between light and heavy neutrino. In this case, the most promising signatures at LHC is $H^+H^- \rightarrow \ell^+\ell^- + \not{E}_T$ [52]. On the other hand for $m_{\Phi\nu} > m_N$, the dominant decay mode of $H^+$ is $H^+ \rightarrow \ell^+N_{Ri}$ with $v_3 \lesssim O(\text{GeV})$ [17]. The dominant decay mode of neutral scalars are reasonable to be $H_3 \rightarrow \nu_iN_{Ri}$ and $A \rightarrow \nu_iN_{Ri}$, since they have the same Yukawa coupling as $H^+$. The heavy
Majorana neutrino $N_{Ri}$ then decays as $N_{Ri} \rightarrow \ell^{\pm}W^{\mp}, N_{Ri} \rightarrow \nu_{\ell}Z, N_{Ri} \rightarrow \nu_{\ell}H_2$. For $m_N = 200\text{GeV}$, we have $\text{BR}(N_{Ri} \rightarrow \ell^{\pm}W^{\mp})=0.60$, $\text{BR}(N_{Ri} \rightarrow \nu_{\ell}Z)=0.28$, and $\text{BR}(N_{Ri} \rightarrow \nu_{\ell}H_2)=0.12$.

![GRAPH](image)

**FIG. 10.** Cross section of pair and associate production of the neutrophilic doublet scalars at 14 TeV LHC. We assume masses of the neutrophilic doublet scalars are degenerate as $m_{H^+} = m_{H_3} = m_A = m_{\Phi \nu}$.

| A : $H^+H^- \rightarrow \ell^+\ell^-N_{Ri}N_{Rj}$ | B : $H^\pm H_3/A \rightarrow \ell^\pm \nu N_{Ri}N_{Rj}$ | C : $H_3A \rightarrow \nu\nu N_{Ri}N_{Rj}$ |
|---|---|---|
| A.1: $\ell^+\ell^-\ell^\pm W^\mp \ell^\pm W^\mp (0.360)$ | B.1: $\ell^\pm \nu\ell^\pm W^\mp \ell^\pm W^\mp (0.360)$ | C.1: $\nu\nu\ell^\pm W^\mp \ell^\pm W^\mp (0.360)$ |
| A.2: $\ell^+\ell^-\ell^\pm W^\mp \nu Z (0.336)$ | B.2: $\ell^\pm \nu\ell^\pm W^\mp \nu Z (0.336)$ | C.2: $\nu\nu\ell^\pm W^\mp \nu Z (0.336)$ |
| A.3: $\ell^+\ell^-\ell^\pm W^\mp \nu H_2 (0.144)$ | B.3: $\ell^\pm \nu\ell^\pm W^\mp \nu H_2 (0.144)$ | C.3: $\nu\nu\ell^\pm W^\mp \nu H_2 (0.144)$ |
| A.4: $\ell^+\ell^-\nu Z \nu Z (0.079)$ | B.4: $\ell^\pm \nu\nu Z \nu Z (0.079)$ | C.4: $\nu\nu\nu Z \nu Z (0.079)$ |
| A.5: $\ell^+\ell^-\nu Z \nu H_2 (0.067)$ | B.5: $\ell^\pm \nu\nu Z \nu H_2 (0.067)$ | C.5: $\nu\nu\nu Z \nu H_2 (0.067)$ |
| A.6: $\ell^+\ell^-\nu H_2 \nu H_2 (0.014)$ | B.6: $\ell^\pm \nu\nu H_2 \nu H_2 (0.014)$ | C.6: $\nu\nu\nu H_2 \nu H_2 (0.014)$ |

**TABLE II.** Signals from pair and associate production of neutrophilic doublet $\Phi \nu$ with their branching ratios given in the parentheses. Here, we set $m_N = 200\text{GeV}$.

In this paper, we concentrate on the case of $m_{\Phi \nu} > m_N$. In TABLE II we summarize all the possible signatures (in $W^{\pm}, Z, H_2$ level) and classify them into three collum according to the production mechanism of $\Phi \nu$. With $W^{\pm}, Z, H_2$ further decaying, there are various possible signatures. Due to the existence of heavy Majorana neutrino $N_{Ri}$, we concentrate on LNV processes. The most interesting and distinct one is the same sign tri-lepton (SST) signature arising from B.1 of TABLE II:

$$B.1 \rightarrow \ell^\pm \nu \ell^\pm jj \ell^\pm jj \rightarrow 3\ell^\pm 4j + E_T$$

For simplicity, we set all the mixing angles among scalars to be zero here.
To our knowledge, such SST signature with $\Delta L = 3$ can only take place in this model, thus it could be used to distinguish this model from other seesaw models. There are also several same sign di-lepton (SSD) signatures with $\Delta L = 2$:

\begin{align*}
B.2 & \rightarrow \ell^+ \nu \ell^\pm jj \nu jj \rightarrow 2\ell^\pm 4j + E_T \\
B.3 & \rightarrow \ell^\pm \nu \ell^\pm jj \nu jj \rightarrow 2\ell^\pm 4j + E_T \\
C.1 & \rightarrow \nu \nu \ell^\pm jj \ell^\pm jj \rightarrow 2\ell^\pm 4j + E_T
\end{align*}

(41) (42) (43)

All these three processes contribute to the SSD signature $2\ell^\pm 4j + E_T$. And there is also a four lepton signature with $\Delta L = 2$:

\begin{align*}
A.1 & \rightarrow \ell^+ \ell^- \ell^\pm jj \ell^\pm jj \rightarrow 3\ell^\pm 4j \\
& \quad \quad \quad \quad \quad \quad \quad \text{(44)}
\end{align*}

In FIG. 11 we shows the theoretical cross section for the LNV signatures at 14 TeV LHC. The SSD signature $2\ell^\pm 4j + E_T$ has the largest cross section, but it also suffers a relative large background from $t\bar{t}W$. On the contrary, the four lepton signature $3\ell^\pm 4j$ has a relative clean background, but its cross section is the smallest. The SST signature $3\ell^\pm 4j + E_T$ seems very promising, since it is nearly background free. Thus it might be testable for $m_{\Phi\nu} \lesssim 700\text{GeV}$ with integrated luminosity of $300\text{fb}^{-1}$ at 14 TeV LHC. A fully discussion and simulation of these LNV signatures at LHC will be carried out in another paper [82].

![FIG. 11. Cross section of lepton number violation signatures at 14 TeV LHC. We also fix $m_N = 200\text{GeV}$.](image)
C. Majoron Dark Matter

Considering the non-perturbative gravitational effects, the Majoron \(J\) could get an \(O(\text{keV})\) mass \([6, 7]\), and play the role of decaying dark matter \([8, 9, 83]\). It is possible to realize EW-scale decaying \([12, 13]\) or stable dark matter \([84–86]\) in Majoron models. In this paper, we focus on \(O(\text{keV})\) Majoron and corresponding phenomenon.

For decaying Majoron dark matter, the present majoron density can be expressed as:

\[
\Omega h^2 = \beta \left( \frac{m_J}{1.25 \text{ keV}} \right) e^{-t_0/\tau_J},
\]

(45)

where \(h\) is the Hubble constant, \(t_0\) is the age of the universe, and \(\beta\) is in the range \(10^{-5} - 1\) corresponding to the majoron thermal history \([87]\). The decay mode of Majoron \(J\) is dominant by \(J \rightarrow \nu\nu\). Induced by the \(k\)-term in Eq. 2, the Majoron \(J\) has non-zero component along the SM and neutrinophilic doublet, and it is approximately given by:

\[
J \sim I_1 + \frac{2v_3^2}{v_1^2} I_2 - \frac{2v_3}{v_1} I_3
\]

(46)

According to this, we can derive the Majoron-neutrino coupling (to leading order) \([37]\):

\[
g_{J\nu_i\nu_j} = -\frac{2m_{\nu}^i}{v_1} \delta_{ij} + \ldots
\]

(47)

and the corresponding decay width:

\[
\Gamma_{J \rightarrow \nu\nu} = \frac{m_J}{2\pi} \sum_{i} (m_{\nu}^i)^2 \frac{v_1^2 v_3}{v_2^2} I_3
\]

(48)

The late decay \(J \rightarrow \nu\nu\) would produce too much power at large scales, thus spoiling the CMB anisotropy spectrum. WMAP third year data has set an upper limit \([8, 88]\):

\[
\Gamma_J < 6.4 \times 10^{-19} \text{ s}^{-1}, \text{ with } 0.12 \text{keV} < \beta m_J < 0.17 \text{keV}.
\]

(49)

From Eq. 48 it is clear that such limit can be easily satisfied as long as \(v_1\) is large enough. For instance, an \(O(\text{keV})\) Majoron requires \(v_1 \gtrsim O(10^4 \text{TeV})\) to satisfy the WMAP limit. In the following discussion, we take \(v_1\) to saturate the upper limit on \(J \rightarrow \nu\nu\). Since \(\beta \in [10^{-5}, 1]\), then \(m_J \sim 0.1 - 10^4 \text{keV}\. More interesting, the sub-leading decay mode of \(J\) is \(J \rightarrow \gamma\gamma\), which is mediated by charged fermions at one-loop level:

\[
\Gamma_{J \rightarrow \gamma\gamma} = \frac{\alpha^2 m_J^2}{64\pi^3} \left| \sum_f N_f Q_f^2 \frac{2v_3^2}{v_2^2 v_1} \left( -2T_3^f \right) \frac{m_J^2}{12m_f^2} \right|^2,
\]

(50)
where \( N_f, Q_f, T^f_3 \) and \( m_f \) are the color factor, electric charge, weak isospin and mass of SM fermion \( f \), respectively. Note from Eq. \( 50 \) that, \( \Gamma_{J \rightarrow \gamma \gamma} \) only depends on \( v_3 \) with fixed values of \( m_J \) and \( v_1 \). The predicted decay rate of \( J \rightarrow \gamma \gamma \) as a function of \( E_\gamma (= m_J/2) \) for different values of \( v_3 \) is shown in FIG. \( \ref{fig:12} \). It is clear that a larger \( v_3 \) leads to a larger \( \Gamma_{J \rightarrow \gamma \gamma} \). Since the performed line emission search has already excluded \( \Gamma_{J \rightarrow \gamma \gamma} \gtrsim \mathcal{O}(10^{-28}\text{s}^{-1}) \), a larger \( v_3 \) actually prefers a smaller \( E_\gamma \) for the survived \( \gamma \)-rays. Further considering the LFV bound on \( v_3 \gtrsim \mathcal{O}(\text{MeV}) \) with EW-scale \( m_\Phi \), the predicted \( E_\gamma \) is usually \( \lesssim 10\text{MeV} \), which covers the right region of \( m_J \) to satisfy the relic density of decaying dark matter \( \cite{87} \).

As discussed in Sec. \( \ref{sec:1} \) the Majoron DM is also a good candidate to explain several \( \text{keV} \)-line excesses. Here, we have chose two different benchmark points to interpret the observed \( 3.5\text{keV} \) and \( 511\text{keV} \) line excesses respectively. First, the direct decay mode \( J \rightarrow \gamma \gamma \) for \( \text{keV} \)-scale Majoron can be used to interpret the \( 3.5\text{keV} \) line excess with \( \cite{13,18} \):

\[
m_J \sim 7\text{keV}, \ \text{and} \ \Gamma_{J \rightarrow \gamma \gamma} \sim 10^{-28}\text{s}^{-1},
\]

which corresponds to benchmark point A in FIG. \( \ref{fig:12} \). Such requirement can be satisfied with \( v_3 \sim 10\text{GeV} \) and \( v_1 \sim 10^4\text{TeV} \). Note that \( v_3 \) in this range can satisfy the tight astrophysical constraints for \( v_1 \sim 10^4\text{TeV} \),
as well as the direct x-ray search bounds in FIG. [12] and WMAP limit in Eq. [49].

Second, for MeV-scale $m_J$, the decay mode $J \rightarrow e^+e^-$ is potential to explain the 511 keV line excess with the requirement [19–21]:

$$\Gamma_{J \rightarrow e^+e^-} \simeq 6.3 \frac{m_J}{1\text{MeV}} \times 10^{-27} \text{s}^{-1},$$

(52)

where we have assume that the Majoron DM $J$ accounts for all the observed DM relic density. In our model, the decay width of $J \rightarrow e^+e^-$ is given by:

$$\Gamma_J = \frac{m_J}{8\pi} \left| \frac{2v_3^2}{v_1v_2} \right| \frac{m_e^2}{m_J^2} \left( 1 - 4\frac{m_e^2}{m_J^2} \right)^{1/2}.$$ 

(53)

Combine Eq. [52] and [53] we have:

$$\frac{\Gamma_{J \rightarrow e^+e^-}}{\Gamma_{J \rightarrow e^+e^-}} = \left( \frac{v_3^2}{v_1v_2} \right)^2 \left( 1 - 4\frac{m_e^2}{m_J^2} \right)^{1/2} \times 1.7 \times 10^{35}.$$ 

(54)

Taking $m_J = 2\text{MeV}$, the required decay width can be obtained for $v_3 \sim 1\text{MeV}$ and $v_1 \sim 10^6\text{TeV}$. Meanwhile the WMAP limit on $\Gamma_J$ in Eq. [49] can be satisfied and the decay width of $\Gamma_{J \rightarrow \gamma\gamma}$ corresponding to benchmark point B in FIG. [12] is far below current direct x-ray limits. Note that to acquire $v_3 \sim 1\text{MeV}$, we also need $m_{\Phi_\nu} \gtrsim \text{TeV}$ to satisfy LFV constraints.

In principle, the discussions for invisible Higgs decay and LHC signatures in previous case for massless $J$ are still applicable for Majoron DM, since $J$ is still invisible at LHC and much lighter than electroweak scale. But with such large $v_1 \gtrsim 10^4 \text{TeV}$ to satisfy WMAP limit, the coupling of $H_aJ\bar{J}$ is so small, thus the branching ratio of invisible Higgs decay is tiny. On the other hand, the masses of $\Phi_\nu$ gets a large contribution from $\beta_3$-term in the scalar potential and would be much heavier than TeV-scale, thus beyond the reach of LHC.

V. CONCLUSION

In this paper, we propose a new model to realize the spontaneous violation of global $U(1)_L$ symmetry in the context of $\nu$-2HDM, where a neutrinophilic doublet scalar $\Phi_\nu$ with lepton number $L = 1$, a complex singlet scalar $\sigma$ with $L = 1/2$, and neutral right-handed fermion singlets $N_{Ri}$ with $L = 0$ are introduced in addition to SM particles. The global $U(1)_L$ symmetry is spontaneously broken by the VEV of $\sigma$, which leads to an (nearly) massless Majoron $J$ and also induces a small VEV of $\Phi_\nu$. Neutrino masses are generate at tree level type-I seesaw like diagram with the SM doublet $\Phi$ replaced by the neutrinophilic doublet $\Phi_\nu$. Due to the smallness of $\langle \Phi_\nu \rangle$, the model is naturally an $O(\text{TeV})$ scale seesaw, and thus detectable in the reach of LHC.
Constraints coming from astrophysics, lepton flavor violation, and direct collider searches are taking into account. The astrophysical constraints set an upper limit on VEV of $\Phi_2$, i.e., $v_3 \lesssim 0.09\text{GeV}$ for the LNV scale $v_1$ at 1 TeV. On the other hand, the LFV constraints set a lower limit on $v_3$, i.e., $v_3 \gtrsim 1\text{MeV}$ for $m_{\Phi_2} = 600\text{GeV}$. Due to the existence of heavy $N_R$ in Majorana case of $\nu$-2HDM, we explain the huge enhancement ($\sim 10^6$) of lower limit on $v_3$ comparing to Dirac case of $\nu$-2HDM. Based on various signals arising from new particles in our model, we investigate the direct search limits carried out by LEP and LHC as well. By choosing proper parameters, we find EW-scale new particles are allowed and some benchmark points are given to illustrate the phenomenological feature of the model.

For massless Majoron, two aspects of LHC signatures are studied: the invisible Higgs decays and LNV signatures. The invisible Higgs decays can be induced by $H_a \to JJ$ and $H_a \to H_b H_b \to 4J$. In the decoupling limit of $\Phi_2$, the two dominant variable that have impact on invisible Higgs decays are $\sin \alpha_{12}$ and $v_1$. Comparing several benchmark points, we conclude that the Majoron could induce large invisible Higgs decay and future experiments prefer smaller $\sin \alpha_{12}$ and larger $v_1$. The $\nu$-2HDM with $N_R$ has three kinds of LNV signatures. The most interesting and distinct one is the same sign tri-lepton signature $3\ell^\pm 4j + E_T$, which can be used to distinguish from other seesaw models. The other two LNV signatures $2\ell^\pm 4j + E_T$ and $3\ell^\pm \ell^\mp 4j$ are also promising to test this model at LHC.

Finally, the Majoron with $m_J \sim \mathcal{O}(\text{keV}) - \mathcal{O}(\text{MeV})$ mass is considered. In this case, the Majoron can serve as a good decaying dark matter candidate. To fulfill the CMB constraints on $\Gamma_{J\to\nu\nu}$, the LNV scale is required to be $\mathcal{O}(10^3 - 10^6\text{TeV})$. The sub-leading decay mode $J \to \gamma\gamma$ is also calculated and compared with current experiments. We find the current limits have already excluded some parameter space. Further, we point out that $J \to \gamma\gamma$ with $m_J \sim 7\text{keV}$ can explain the 3.5 keV line excess and $J \to e^+e^-$ with $m_J \sim \mathcal{O}(\text{MeV})$ can interpret the 511 keV line excess. Two different benchmark points are given to illustrate these two excesses respectively.

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The coupling of $H_2 H_1 H_1$:

\[
\frac{g H_2 H_1 H_1}{2} = 3\lambda_1 v_2 (O^{R}_{12}O^{R}_{22} + 3\lambda_1 v_3 (O^{R}_{13}O^{R}_{23} + 3\beta_1 v_1 (O^{R}_{11}O^{R}_{21} \\
+ \lambda_3 + \lambda_4 \frac{1}{2} \left[ (O^{R}_{12}O^{R}_{22}v_2 + (O^{R}_{12}O^{R}_{23}v_3 + 2O^{R}_{12}O^{R}_{13}(O^{R}_{23}v_2 + O^{R}_{22}v_3) \right) \\
+ \frac{\beta_2}{2} \left[ (O^{R}_{11}O^{R}_{21}v_1 + (O^{R}_{11}O^{R}_{23}v_3 + 2O^{R}_{11}O^{R}_{13}(O^{R}_{23}v_1 + O^{R}_{21}v_3) \right) \\
- k \frac{1}{2} \left[ (O^{R}_{11}O^{R}_{23}v_3 + (O^{R}_{11}O^{R}_{23}v_2 + 2O^{R}_{11}O^{R}_{12}(O^{R}_{21}v_3 + O^{R}_{23}v_1) + 2O^{R}_{11}O^{R}_{13}(O^{R}_{21}v_2 + O^{R}_{22}v_1) \right].
\]
\]

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