Radiative Linear Seesaw Model, Dark Matter and $U(1)_{B-L}$

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In this paper we propose a radiated linear seesaw model where the naturally small term $\mu_L$ are generated at one-loop level and its soft-breaking of lepton number symmetry attributes to the spontaneous breaking (SSB) of B-L gauge symmetry. The value of $B-L$ charges for new particles are assigned to satisfy the anomalies cancelation. It is founded that some new particles may have exotic values of $B-L$ charge such that there exists residual $Z_2 \times Z'_2$ symmetry even after SSB of $B-L$ gauge symmetry. The $Z_2 \times Z'_2$ discrete symmetry stabilizes the these particles as dark matter candidates. In the model, two classes of inert fermions and scalars with different $B-L$ charges are introduced, leading to two-component dark matter candidates. The lepton flavor violation processes, the relic density of dark matter, the direct detection of dark matter and the phenomenology at LHC are investigated.

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I. INTRODUCTION

The origin of tiny but non-zero neutrino masses observed by neutrino oscillation experiments\cite{1} remains so far a mystery, and so provides us an opportunity to search the new physics beyond the standard model (SM). Perhaps the simplest scenario which may understand the neutrino puzzle is to introduce the Majorana mass that breaks the global $B - L$ symmetry though the dimension-5 Weinberg operator $\lambda LL\Phi\Phi/\Lambda$\cite{2}. This effective operator can be realized though various pathways which depends on the new physics scales the associated massages lie at. For instance in the case of widely known type-I seesaw mechanism\cite{3} with right-handed Majorana ingredients $N_R$ as massages, one needs the super-heavy masses for $N_R$ i.e $10^{14-16}$ GeV to fit the observed sub-eV neutrino mass. The right-handed neutrinos are too heavy to be detected at future experiments. In contrast, in the so-called low-scale scenarios, the small neutrino mass is not only due to the mass state with heavy mass but also to another naturally small mass parameter which breaks the lepton number symmetry. This is the basic ideas behind many schemes including type-II seesaw model\cite{4}, inverse seesaw model\cite{5} and linear seesaw model\cite{6,7}. In these models, the mass of massage particles can be lowered down to TeV or even hundreds GeV, a scale to be explored at collider experiments.

On the other hand, the Planck data has shown that 26% of the energy density of our universe is occupied by dark matter. In the view of particle physics, the weakly interacting massive particles (WIMPs) are the most promising dark matter candidates. In recent years, a class of models are proposed to incorporate the neutrino mass puzzle and the existence of dark matter in a unified framework. In these models, the neutrino masses are generated at loop level and the dark matter is naturally contained as an inert particle, where the $Z_2$ symmetry or a $U(1)_X$ symmetry is used to guarantee the stability of dark matter. The radiated generation of neutrino mass has been realized at one-loop level\cite{8-10}, two-loop level\cite{11-15} and three-loop level\cite{16}. The systematic analysis of one and two-loop realization for with possible topologies are performed in Ref\cite{17}.

In this paper, we proposed a radiated linear seesaw model where the lepton number violation is due to the spontaneous symmetry breaking (SSB) of $U(1)_{B-L}$ gauge symmetry, while the naturally small mass parameter is generated at one-loop level. The linear seesaw model was fist studied in the left-right theory with gauge group $SU(3) \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}\cite{6}$, and subsequently inspired on $SO(10)$ theory in the presence of gauge singlets\cite{7}. In linear seesaw scenario, $\Psi_R$ and $\Psi_L$ are added onto SM so that the Lagrangian is given by

$$\begin{align*}
L &= M_D \bar{\nu}_L \Psi_R + M_\Psi \bar{\Psi}_R \Psi_L + \mu_L \bar{\nu}_L \Psi_L + h.c
\end{align*}$$

(1)
The neutrino mass matrix in the basis of \((\nu_L, \Psi^c_R, \Psi_L)\) is

\[
M_\nu = \begin{pmatrix}
0 & M_D & \mu_L \\
M_D^T & 0 & M_{\Psi} \\
\mu_L^T & M_{\Psi}^T & 0
\end{pmatrix}
\]  

(2)

The light neutrino mass is given by \(m_\nu \sim \mu_L M_D M_\Psi^{-1}\). Note that the \(\mu_L\) violates the lepton number symmetry and plays the role of a naturally small parameter. Thus it seems natural that there exists a suppression mechanism where the \(\mu_L\) term is generated via loop diagram meanwhile the soft-breaking of lepton number symmetry may attribute to the SSB of \(B - L\) gauge symmetry. Moreover, when the WIMPs as the dark matter candidates are involved in the loop diagram, we can reasonably assume they are generated by a vacuum expectation value with respect to the SSB of \(B - L\) gauge symmetry at TeV scale. These are the main motivations of this work.

Following the spirit of Ref. [14], the value of \(B - L\) charges should be carefully assigned since the anomalies cancelation must be satisfied. It is founded that some new particles may have exotic values of \(B - L\) charge such that there exists residual \(Z_2 \times Z'_2\) symmetry even after SSB of \(B - L\) gauge symmetry. The \(Z_2 \times Z'_2\) discrete symmetry stabilizes the these particles from decaying to SM ingredients. Thus the lightest particle with the same exotic value of \(B - L\) charge can be a dark matter candidate. In practise, we introduce two classes of inert fermions and scalars to realize the model, leading to two component dark matter candidates.

The existence of new fermions and scalars provides rich phenomenon. Tiny neutrino masses are explained with one-loop induced linear-seesaw-like mechanism. The charged-scalar mediates lepton flavor violation (LFV) of charged leptons. The relic density and the direct detection of the two component dark matter are investigated. The properties of discovered SM Higgs will be changed by the new particles. And these new particles provide plenty of new signatures at LHC. Especially, multi-lepton signals with missing transverse energy \(E_T^\text{miss}\) can be used to test our model. We find that our model can satisfy current constraints from phenomenons mentioned above.

The rest of paper is organised as follows. In Sec. II we introduce the realization of radiative linear seesaw and multi-component dark matter from gauged \(U(1)_{B-L}\). In Sec. III we discuss the phenomenon of lepton flavor violation, dark matter and collider signatures. Conclusions are given in Sec. IV.
II. MODEL

A. Model Setup

In our model the neutrino masses are generated via the diagram depicted in Fig. 1. The new particles content and their charge assignment are listed in Table I. We add Weyl fermions \( \Psi_{Ri}, \Psi_{Li} \), right-handed Majorana neutrinos \( N_{R\alpha} \), and pairs of right-handed Majorana neutrinos \( N'_{R}, N''_{R} \) to the SM where \( i, \alpha \) and \( \beta \) are the generation indices. All the new fermions are singlets under SM gauge group. Five new scalars \( \eta_1, s_1, \eta_2, s_2 \) and \( \sigma \) are also added to SM. Because the new fermions are all SM singlets, the \( B-L \) gauge symmetry satisfies all anomaly cancellations except for \( \left[ U(1)_{B-L} \right] \times \left[ \text{Gravity} \right] \) and \( \left[ U(1)_{B-L} \right]^3 \) [18]. Considering the conditions for the absence of \( \left[ U(1)_{B-L} \right] \times \left[ \text{Gravity} \right]^2 \) and \( \left[ U(1)_{B-L} \right]^3 \) anomaly, one has

\[
3 + \left( -\frac{1}{2} \right) N_1 + x N_2 + (-1 - x) N_2 + (-1) N_{\psi} = 0 \tag{3}
\]

\[
3 + \left( -\frac{1}{2} \right)^3 N_1 + x^3 N_2 + (-1 - x)^3 N_2 + (-1)^3 N_{\psi} = 0 \tag{4}
\]

After solving the anomaly free condition, one obtains

\[
N_1 = 2, \quad N_2 = 1, \quad N_{\psi} = 1, \quad x = \sqrt{2} - \frac{1}{2} \tag{5}
\]

Thus we have the some inert particles classified into two parts. In the first class, there are two Majorana right-handed neutrinos \( (N_{R1}, N_{R2}) \) and the inert scalars \( (\eta_1, s_1) \) with their B-L charge being \( -\frac{1}{2} \). In the second class, we obtain a pair of Majorana right-handed neutrinos \( (N'_{R}, N''_{R}) \) along with the inert scalars \( (\eta_2, s_2) \) whose the B-L charges are irrational numbers \( (\sqrt{2} - \frac{1}{2}) \) or \( -\sqrt{2} - \frac{1}{2} \). One notices that the new particles with both \( -\frac{1}{2} \) and the irrational numbers can not decay into SM particles. Therefore the lightest particles belonging to the same class is stable and can be regarded as a dark matter candidate. The relevant Lagrangian for Yukawa sector is given by

\[
-L_Y = y_l \bar{L}_l \psi_{R} \tau_2 \Phi^* + y'_l \bar{\psi}_L \psi_{R} \sigma + h_{\alpha} \bar{N}_{R\alpha} \psi_L s_1 + f_{\alpha \beta} \bar{L}_l \bar{N}_{\alpha \beta} i \tau_2 \eta^*_1 + \frac{1}{2} Y_{\alpha} \bar{N}_{R\alpha} N_{R\alpha} \sigma \\
+ h \bar{N}_{R} \psi_L s_2 + f_{\beta} \bar{L}_l \bar{N}_{R} i \tau_2 \eta^*_2 + \frac{1}{2} Y_{\beta} \bar{N}_{R} N_{R} \sigma + h.c \tag{6}
\]
Without losing generality, we work in the basis where the mass term of $N_{R1,2}$ is diagonal. As for the mass term of $N'_{R}$ and $N''_{R}$, one can redefine the fields as

$$
\chi_1 = \frac{1}{\sqrt{2}}(N'_{R} + N''_{R}) \quad \chi_2 = \frac{i}{\sqrt{2}}(N'_{R} - N''_{R})
$$

so that

$$
\frac{1}{2}Y N''_{R} N'_{R} \sigma \rightarrow \frac{1}{2} Y (\chi_1 \chi_1 + \chi_2 \chi_2) \sigma
$$

Now we get two Majorana neutrino eigenstates having the same masses. Note that there is no interplay Yukawa terms between $N_{R1,2}$ and $(N'_{R}, N''_{R})$ because of the the B-L charges assignment they have.

The scalar potential in our model is given by

$$
V(\Phi, \sigma, \eta_1, s_1, \eta_2, s_2) = -\mu_\Phi^2 \Phi^\dagger \Phi + \lambda_\Phi(\Phi^\dagger \Phi)^2 - \mu_\sigma^2 |\sigma|^2 + \lambda_\sigma |\sigma|^4
$$

$$
+ \mu_{\eta_1}^2 |\eta_{1\dagger} \eta_1|^2 + \mu_{\eta_2}^2 |\eta_{2\dagger} \eta_2|^2 + \lambda_{\eta_1 \eta_2} \eta_{1\dagger} \eta_1 \eta_{2\dagger} \eta_2
$$

$$
+ \mu_{s_1}^2 |s_1|^2 + \mu_{s_2}^2 |s_2|^2 + \lambda_{s_1 s_2} |s_1| s_2|^2 + \lambda_{s_1 s_2} |s_1|^2 |s_2|^2
$$

$$
+ \lambda_{\eta_1 s_1}(\Phi^\dagger \Phi)(\eta_{1\dagger} \eta_1) + \lambda_{\eta_2 s_1}(\Phi^\dagger \Phi)(\eta_{2\dagger} \eta_2)
$$

$$
+ \lambda_{\eta_1 \eta_2}(\eta_{1\dagger} \eta_1)(\eta_{2\dagger} \eta_2) + \lambda_{\eta_1 \eta_2}(\eta_{1\dagger} \eta_1)(\eta_{2\dagger} \eta_2)
$$

$$
+ \lambda_{s_1 \eta_1} |s_1|^2 (\Phi^\dagger \Phi) + \lambda_{s_2 \eta_1} |s_2|^2 (\Phi^\dagger \Phi) + \lambda_{s_1 \eta_2} |s_1|^2 (\eta_{1\dagger} \eta_1)
$$

$$
+ \lambda_{s_2 \eta_2} |s_2|^2 (\eta_{2\dagger} \eta_2) + \lambda_{s_1 \sigma} |s_1|^2 |\sigma|^2 + \lambda_{s_2 \sigma} |s_2|^2 |\sigma|^2 + (\mu_1 s_1 \eta_{1\dagger} \eta_1 + \mu_2 s_2 \eta_{2\dagger} \eta_2 + h.c)
$$

where $\mu_\Phi^2, \mu_{\eta_1}^2, \mu_{\eta_2}^2, \mu_{s_1}^2$ and $\mu_{s_2}^2$ are taken as positive values and the value coupling constants $\mu_1$ and $\mu_2$ in trilinear terms can be set as positive by re-phasing $s_1$ and $s_2$. Notice that there is no terms like $s_1 \sigma^2$ or $s_2 \sigma^2$ appearing in the scalar potential. This has two fold meanings: First, the inert scalars $\eta_{1,2}$ and $s_{1,2}$ do not acquire the VEV after the SSB of $\Phi$ and $\sigma$; Second, there exists a residual $Z_2 \times Z'_2$ symmetry under which all the inert particles are odd even after the breakdown of B-L symmetry. Therefore the residual $Z_2 \times Z'_2$ symmetry stabilizes the inert particles, makes them to be two component dark matter candidates.

**B. Matrices of Scalar Particles**

After the SSB, the scalar $\Phi$ and $\sigma$ is parameterized as

$$
\Phi = \begin{pmatrix} G^+ \\ v_\phi + \phi_0 + iG_\phi \end{pmatrix} \quad \sigma = \frac{v_\sigma + \sigma_0 + iG_\sigma}{\sqrt{2}}
$$
where $v_\phi \simeq 246$GeV is the VEV of the SM higgs doublet scalar and $v_\sigma$ is responsible for the SSB of B-L symmetry [19]. The Nambu-Godstone bosons $G^+, G_\phi$ and $G_\sigma$ are absorbed by the longitudinal components of $W$, $Z$ and $Z'$ gauge bosons. For simplicity, we ignore the kinetic mixing between $U(1)_Y$ and $U(1)_{B-L}$ gauge boson [20]. Therefore the VEV $v_\sigma$ provides a mass of $U(1)_{B-L}$ gauge boson $Z'$ as $M_{Z'} = g_{B-L}v_\sigma$, where $g_{B-L}$ is the $U(1)_{B-L}$ gauge coupling constant. For the extra gauge boson $Z'$, LEP-II provides a combined bound $M_{Z'}/g_{B-L} > 7$ TeV [21], which is just the lower bound on $v_\sigma$. Then we obtain the mass matrix for CP-even scalars $\phi_0$ and $\sigma_0$

$$M^2(\phi_0, \sigma_0) = \begin{pmatrix} \cos \theta & \sin \theta \\ -\sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} M^2_h & 0 \\ 0 & M^2_H \end{pmatrix} \begin{pmatrix} \cos \theta & -\sin \theta \\ \sin \theta & \cos \theta \end{pmatrix}$$

(11)

where $h$ stands for the SM-like Higgs [22, 23] and $H$ is an extra CP-even Higgs boson [24–26] with the masses respectively as

$$M^2_h = \lambda_\phi v_\phi^2 + \lambda_\sigma v_\sigma^2 - \sqrt{(\lambda v_\phi^2 - \lambda_\sigma v_\sigma^2)^2 + \lambda_{0\Phi}^2 v_\phi^2 v_\sigma^2}$$

(12)

$$M^2_H = \lambda_\phi v_\phi^2 + \lambda_\sigma v_\sigma^2 + \sqrt{(\lambda v_\phi^2 - \lambda_\sigma v_\sigma^2)^2 + \lambda_{0\Phi}^2 v_\phi^2 v_\sigma^2}$$

(13)

and mixing angle $\theta$ determined as

$$\sin 2\theta = \frac{2\lambda_{1\Phi} v_\phi v_\sigma}{M^2_H - M^2_h}$$

(14)

On the other hand, the inert scalar $(\eta_1, s_1)$ and $(\eta_2, s_2)$ do not mix with $\Phi$ and $\sigma$ due to the residual $Z_2$ symmetry. The mass matrix for inert scalar fields are

$$M(\eta_1, s_1, \eta_2, s_2) = (\eta_1^\dagger, s_1^\dagger, \eta_2^\dagger, s_2^\dagger) \begin{pmatrix} M_{11} & M_{12} & 0 & 0 \\ M_{12} & M_{22} & 0 & 0 \\ 0 & 0 & M_{33} & M_{34} \\ 0 & 0 & M_{43} & M_{44} \end{pmatrix} (\eta_1, s_1, \eta_2, s_2)$$

(15)

where

$$M_{11} = \mu^2_{\eta_1} + \frac{1}{2} \lambda_{\phi\eta_1} v_\phi^2 + \frac{1}{2} \lambda_{\psi\eta_1} v_\psi^2 + \frac{1}{2} \lambda_{\sigma\eta_1} v_\sigma^2$$

$$M_{22} = \mu^2_{s_1} + \frac{1}{2} \lambda_{s_1\phi} v_\phi^2 + \frac{1}{2} \lambda_{s_1\sigma} v_\sigma^2$$

$$M_{33} = \mu^2_{\eta_2} + \frac{1}{2} \lambda_{\phi\eta_2} v_\phi^2 + \frac{1}{2} \lambda_{\psi\eta_2} v_\psi^2 + \frac{1}{2} \lambda_{\sigma\eta_2} v_\sigma^2$$

$$M_{44} = \mu^2_{s_2} + \frac{1}{2} \lambda_{s_2\phi} v_\phi^2 + \frac{1}{2} \lambda_{s_2\sigma} v_\sigma^2$$

$$M_{12} = M_{21} = \frac{\mu_1}{\sqrt{2}} v_\phi$$

$$M_{34} = M_{43} = \frac{\mu_2}{\sqrt{2}} v_\phi$$

(16)
There is also no mixing between \((\eta_1, s_1)\) and \((\eta_2, s_2)\), therefore a residual \(Z_2^\prime\) symmetry between the two classes can be realised. After diagonalizing the mass matrix, we obtain the mass eigenstates of inert scalars as

\[
\begin{pmatrix}
A_1^n \\
H_1^n
\end{pmatrix} = \begin{pmatrix}
\cos \theta_1 & -\sin \theta_1 \\
\sin \theta_1 & \cos \theta_1
\end{pmatrix} \begin{pmatrix}
\eta_{1,2}^0 \\
s_{1,2}^0
\end{pmatrix}, \quad \sin 2\theta_1 = \frac{\sqrt{2}\mu_{1,2}\nu_\phi}{M_{A_1,2}^2 - M_{H_1,2}^2}
\]

(17)

where

\[
M_{A_1,2}^2 = \frac{1}{2} \left( M_{\eta_1,2}^2 + M_{s_1,2}^2 + \sqrt{(M_{\eta_1,2}^2 - M_{s_1,2}^2)^2 + 2\mu_{1,2}^2\nu_\phi^2} \right)
\]

(18)

\[
M_{H_1,2}^2 = \frac{1}{2} \left( M_{\eta_1,2}^2 + M_{s_1,2}^2 - \sqrt{(M_{\eta_1,2}^2 - M_{s_1,2}^2)^2 + 2\mu_{1,2}^2\nu_\phi^2} \right)
\]

(19)

Here \(M_{\eta_1} \equiv M_{11}, M_{s_1} \equiv M_{22}, M_{\eta_2} \equiv M_{33}\) and \(M_{s_2} \equiv M_{44}\).

C. Neutrino Mass

\[
\langle \phi^0 \rangle \quad \nu_L \quad \psi_R \quad \psi_L \quad \langle \sigma \rangle \quad N_{R_{1,2}} \quad N_{L_{1,2}} \quad \langle \sigma \rangle
\]

\[
\langle \phi^0 \rangle \quad \nu_L \quad \psi_R \quad \psi_L \quad \langle \sigma \rangle \quad N_{R_{1,2}} \quad N_{L_{1,2}} \quad \langle \sigma \rangle
\]

FIG. 1. The one-loop diagrams for neutrino masses in our model.

As shown in Fig. 1, the tiny neutrino mass are generated by the linear seesaw mechanism except that the \(\mu_L\) terms are induced by a one-loop diagram. The effective mass matrix for active neutrinos depicted in Fig. 1 is expressed as

\[
M_{\nu ll'} = M_{\nu ll'}^I + M_{\nu ll'}^H
\]

(20)

where

\[
M_{\nu ll'}^I = \frac{\nu_\phi \sin \theta_1 \cos \theta_1}{16\pi^2\sqrt{2}M_\psi} \sum_{i=1}^2 h_{1,2} f_{ll'} M_i \left[ \frac{M_{A_2}^2}{M_i^2 - M_{A_1}^2} \ln \left( \frac{M_{A_2}^2}{M_i^2} \right) - \frac{M_{H_2}^2}{M_i^2 - M_{H_1}^2} \ln \left( \frac{M_{H_2}^2}{M_i^2} \right) \right] + (l \leftrightarrow l')
\]

\[
M_{\nu ll'}^H = \frac{\nu_\phi \sin \theta_2 \cos \theta_2}{16\pi^2\sqrt{2}M_\psi} y_l h f_{ll'} M_X \left[ \frac{M_{A_2}^2}{M_X^2 - M_{A_2}^2} \ln \left( \frac{M_{A_2}^2}{M_X^2} \right) - \frac{M_{H_2}^2}{M_X^2 - M_{H_2}^2} \ln \left( \frac{M_{H_2}^2}{M_X^2} \right) \right] + (l \leftrightarrow l')
\]

(21)
where $M_i (i=1,2)$ denotes the masses for $N_{R1}$ and $N_{R2}$; $M_\chi$ denotes the masses for the eigenstates of $N'_R$ and $N''_R$. Tiny neutrino masses can be obtained using the following benchmark points:

$$\mu_1 = \mu_2 = 0.1 \text{ GeV}, y = h = 0.0028, f = 0.01, M_\psi = 300 \text{ GeV}$$

$$M_{N_{R1}} = 149.5 \text{ GeV}, M_{N_{R2}} = 200 \text{ GeV}, M_\chi = 150 \text{ GeV}$$

$$M_{A_1} = 300 \text{ GeV}, M_{\eta_1^\pm} = 270 \text{ GeV}, M_{H_1^0} = 1000 \text{ GeV}$$

$$M_{A_2} = 700 \text{ GeV}, M_{\eta_2^\pm} = 690 \text{ GeV}, M_{H_2^0} = 62 \text{ GeV}$$

with the index of Yukawa couplings suppressed for simplicity. Then we get $M_\nu = 0.0164 \text{ eV} (\sqrt{\Delta m_{12}^2})$.

The benchmark point given above seems rather unusual since the $M_\nu$ becomes a rank-1 matrix as

$$M_\nu \sim D \equiv \begin{pmatrix} 1 & 1 & 1 \\ 1 & 1 & 1 \\ 1 & 1 & 1 \end{pmatrix}$$

(23)

which is obviously not consistent with the results of neutrino oscillation experiments. However, one reminds the expression of matrix (23) is just the so-called flavor democratic model studied by many authors and related to some flavor symmetries [28]. The matrix $D$ can be diagonalized as

$$V_\nu^T D V_\nu = \frac{1}{\sqrt{6}} \begin{pmatrix} \sqrt{3} & 1 & \sqrt{2} \\ -\sqrt{3} & 1 & \sqrt{2} \\ 0 & -2 & \sqrt{2} \end{pmatrix}$$

(24)

The unitary matrix $V_\nu$ corresponds to the democratic mixing pattern. Since $\tilde{D}$ contains a dominant non-zero element at $(3,3)$ position, the flavor democratic structure in $M_\nu$ in eqn(23) can be viewed as a good approximation for our rank-2 neutrino mass matrix which exhibiting the strong and normal order of neutrino mass spectrum, i.e $m_1 = 0 \ll m_2 = \sqrt{\Delta m_{12}^2} \ll m_3 = \sqrt{\Delta m_{13}^2}$. Thus the benchmark point we take is reasonable. It is noted in the flavor basis the democratic mixing matrix $V_\nu$ has already not been consistent with the PMNS matrix $U_{PMNS}$ measured by experiments, however the $V_\nu$ can be corrected by the charged lepton sector $V_l$ to fit the neutrino oscillation data [29].

The small values of $\mu_{1,2}$ lead to $\sin 2\theta_{1,2} = (3.8, 7.2) \times 10^{-5}$, which plays a key role in the suppression of tiny neutrino masses. The choice of values of $\mu_{1,2}$ is mainly for phenomenological consideration. First, in case of scalar DM, we have $Z - S - S^*$ coupling proportional to $\sin^2 \theta_{1,2}$. The spin independent elastic cross section of DM requires $\sin \theta_{1,2} < 0.05$ [30], setting an upper limit on $\mu_{1,2} \sim (10 \text{ GeV})$ for electroweak(EW) scale inert scalars. Second, in our benchmark point with inert particles at the EW-scale and Yukawa couplings of the order $10^{-2}$ or $10^{-3}$, rich phenomenon of new physics are expected for
LHC and LFV processes. On the other hand, larger values of $\mu_{1,2}$ is possible if we decrease the Yukawa couplings to be more smaller values. But this will predict a too small branch ratios for LFV processes. Too small Yukawa coupling also seems unnatural from the viewpoint of model building. Another solution is to increase the mass of $\psi$ or inert particle to TeV-scale, which is beyond the reach of LHC. For the DM candidate $N_{R1}(H^0)$, its mass is set to be about half of the mass of the s-channel mediator $H(h)$. Then one obtains the large enough DM annihilation cross section to account for the relic density. For the heavy dirac fermion $\psi$ and inert doublet scalars $\eta_1, A^0_1$, we choose their masses around 300 GeV, so that they are testable at LHC. The other inert scalars are around TeV-scale aiming to suppress the value of neutrino mass.

III. PHENOMENOLOGY

A. Lepton Flavor Violation

The Yukawa interactions of charged-scalar $\eta^\pm$ will contribute to the LFV processes of charged leptons. Detail studies on LFV processes in scotogenic models [9] have been carried out in Ref. [31, 32]. Currently, the most severe constraint coming from MEG collaboration on muon radiative decay with an upper limit $\text{BR}(\mu \rightarrow e\gamma) < 5.7 \times 10^{-13}$ (90% C.L.) [33]. In our model, the analytical branching ratio of $\mu \rightarrow e\gamma$ is calculated as [31, 34]:

$$\text{BR}(\mu \rightarrow e\gamma) = \frac{3\alpha_{\text{em}}}{64\pi G_F^2} \left| \sum_{i=1}^{2} \frac{f_{\mu i} f_{ei}^*}{M_{\nu_i}^2} F \left( \frac{M_{N_i}^2}{M_{\eta_1^i}^2} \right) + \frac{f_{\mu i} f_{ei}^*}{M_{\eta_2^i}^2} F \left( \frac{M_{N_i}^2}{M_{\eta_2^i}^2} \right) \right|^2,$$  

(25)

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

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

(26)

The benchmark point in Eqn[22] predicts $\text{BR}(\mu \rightarrow e\gamma) = 9.8 \times 10^{-14}$, which satisfies the current limit and is in the reach of future sensitivity [35]. The limits on $\tau$ observables are less stringent [36, 37]. With the much natural Yukawa structure in our benchmark point, the predicted $\text{BR}(\tau \rightarrow \mu\gamma) = 5.8 \times 10^{-13}$ is far beyond the future sensitivity. But on the other hand, hierarchal Yukawa structure $|f_{\tau e}| \lesssim |f_{\mu e}| \lesssim |f_{\tau \tau}|$ with $f_{\tau \tau} \sim \mathcal{O}(1)$ is still allowed from phenomenological point of view. In this case, fermionic dark matter candidate $F$ annihilation in the mass region between 2 GeV and 3 TeV through $t$-channel exchange of $\eta$ can satisfy dark matter relic density bound [32]. And as a consequence of hierarchal Yukawa structure, dark matter $F$ annihilates mainly into third-family leptons: $\tau^+\tau^-$ and $\nu_\tau\bar{\nu}_\tau$. 
B. Dark Matter

In our model, a multi-component dark matter scenario is possible due to the residual $Z'_2$ symmetry between two sets of new scalars and fermions. For instance, we choose the lightest fermion $N_{R1}$ (refer as $F$) in set one and the lightest scalar $H^0_2$ (refer as $S$) in set two as dark matter candidates. These dark matter candidates must satisfy two experimental constrains: (1) the dark matter relic density observed by Plank \cite{38} $\Omega_{DM}h^2 = 0.1193 \pm 0.0014$; (2) cross section for direct detection of dark matter scattering off nucleon set by LUX \cite{30}.

Theoretical calculation of dark matter relic density is well described in \cite{39}, and it is calculated with the help of packages Feynrules \cite{40} and micrOMEGAs \cite{41} in our analysis. Because the $t$-channel Yukawa portal may suffer constrains from LFV or neutrino masses, and in our benchmark point, the $t$-channel contribution to relic density is less than 1%, thus we will focus on the $s$-channel $h/H/Z'$ portal for simplicity.

Firstly, the relic density of fermion/scalar dark matter for one dark matter candidate is presented in Fig. 2, where we neglect the conversion $F \bar{F} \leftrightarrow SS^*$ between the two dark matter candidates. For the fermion dark matter $F$, the light $h$ portal can not acquire a sufficient annihilation cross section, while the heavy $H$ portal is still promising when $M_F \sim M_H/2$, which is because the suppression of large $v_\sigma = 8$ TeV. Anyway, the $Z'$ portal can easily satisfy the relic density when $M_F \sim M_{Z'}/2$. For the scalar dark matter $S$, it is dominantly made from the singlet scalar $s^0_2$, due to small mixing $\theta_2$. The relic density can easily be attained when $M_S \sim M_h/2$ and $M_S \sim M_H/2$, while the $Z'$ portal is not promising, mainly because the small $B-L$ charge of $S$ and suppression of heavy $M_{Z'}$.

Secondly, we take into account the conversion of two component dark matter $F \bar{F} \leftrightarrow SS^*$, which can be mediated by $s$-channel $h/H/Z'$, where the $H$-portal is expected to be the dominant one. Therefore, $HSS^*$ and $HF\bar{F}$ are the two most relevant couplings to study conversion. For simplicity, we further assume $\lambda_{\sigma \eta_2} = \lambda_{s_2 \sigma} = \lambda$, which determines $HSS^*$ coupling, and fix other parameters as discussed in FIG. 2 if not mentioned.

The dependence of $F/S$ relic density on $\lambda$ is shown in FIG. 3. For the fermion dark matter, when $M_F > M_S$, the larger $\lambda$ the larger $F \bar{F} \leftrightarrow SS^*$ annihilation rate, and therefor the smaller the relic density. It is clear that the $\Omega_Fh^2$ can differ by about one order of magnitude between $\lambda = -0.001$ and $\lambda = -0.05$. But when $M_F < M_S$, the effect of conversion $F\bar{F} \rightarrow SS^*$ is quite small. Noticing for $\lambda = -0.05$, the $H \rightarrow SS^*$ can greatly enhance the total decay width of $H$, which causes the increase of $\Omega_Fh^2$ around $M_F \sim M_H/2$. For the scalar dark matter, the increase of $\lambda$ will decrease the relic density significantly due to the increase of $SS^*$ annihilation. But, the friction of conversion $SS^* \rightarrow F\bar{F}$ keeps the same, since the Yukawa coupling $HFF\bar{F}$ is fixed by $M_F$. The arguments are true if we exchange the roles of $F$ and $S$, and
FIG. 2. Relic density of fermion (left) and scalar (right) dark matter as a function of $M_{DM}$ for one dark matter candidate. Here, we set the relevant scalar interaction coupling $\lambda_{\eta_2}^{\prime} = \lambda_{\eta_2}^{\prime} = \lambda_{s_2} - 0.001$, $\lambda_{s_2} = -0.001$. We also fix $M_h = 125\text{ GeV}$, $M_H = 300/600\text{ GeV}$, $\sin \theta = 0.3$, $M_{Z'} = 4\text{ TeV}$, and $v_{s} = 8\text{ TeV}$.

Another aspect of conversion is the masses of the two dark matter. The left(right) of FIG. 4 shows the $F(S)$ relic density for $M_{S(F)} = 60,150,300\text{ GeV}$ with $\lambda = -0.02$. In case of fermion dark matter, it is clear that the smaller the $M_S$ the larger $F\bar{F} \rightarrow SS^{*}$ annihilation rate, and therefore the smaller the relic
density. For relative heavy $M_S = 150, 300 \text{ GeV}$, the conversion has tiny effect on the $F$ relic density when $M_F < M_H/2$. In case of scalar dark matter, the dependence of $\Omega_S h^2$ on $M_F$ is a little complicated, since the $H F \bar{F}$ coupling is directly related with the $M_F$. For $M_S < 80 \text{ GeV}$, the effect of conversion is relatively small, and one expect that the larger the $M_F$, the smaller the $S$ relic density, which is mainly caused by the increase of $H F \bar{F}$ coupling. In medium mass region $80 < M_S < 200 \text{ GeV}$, the conversion effect would be dominant, thus the smaller the $M_F$, the smaller the $S$ relic density. In the high mass region, the conversion effect is comparable to $H F \bar{F}$ coupling effect, which makes the dependence of $\Omega_S h^2$ on $M_F$ nonlinear. In a word, the $H S S^*$ and $H F \bar{F}$ couplings play vital importance in dark matter conversion. The conversion can take place in both direction $F \bar{F} \rightarrow SS^*$ and $SS^* \rightarrow F \bar{F}$ when $M_F \sim M_S$, which can be obtained when both $F$ and $S$ are mainly annihilation through $H$-portal. If not the case, only the conversion of heavier one into light one is relevant \[42\].

\[\text{FIG. 4. The effect of two component dark matter conversion on } F/S \text{ relic density for fixed } \lambda = -0.02.\]

Finally, we discuss the constrains from direct detection of dark matter. The Current experiment constraints assume the existence of only one dark matter specie. However two-component dark matter candidates are predicted in our model. Therefore the contribution of cross section on nucleon for each specie should be rescaled by the fraction factor of relic density. We define the fraction of the mass density of $i$-th dark matter in case of multi-component dark matter [43, 44]:

$$\epsilon_i = \frac{\Omega_i h^2}{\Omega_{CDM} h^2},$$  

(27)
where \( i = F, S \) in our consideration. Thereafter, the up limit of direct detection is:

\[
\frac{\epsilon_F}{M_F} \sigma_{F-N} + \frac{\epsilon_S}{M_S} \sigma_{S-N} < \frac{\sigma_{exp}}{M_{DM}}
\]  

(28)

Here, \( \sigma_{F-N} (\sigma_{S-N}) \) denotes the scattering cross section of \( F (S) \) with a nucleon \( N \). The benchmark point in Eqn. 22 gives the spin-independent scattering cross section \( \sigma_{SI}^{F-N} = 1.10 \times 10^{-46} \text{ cm}^2 \) (\( \sigma_{SI}^{S-N} = 1.62 \times 10^{-44} \text{ cm}^2 \)) with \( \Omega_F h^2 = 1.12 \times 10^{-1} \) (\( \Omega_S h^2 = 2.64 \times 10^{-4} \)). Although the bare \( \sigma_{SI}^{S-N} \) is larger than the LUX upper constraint \( 1.1 \times 10^{-47} \text{ cm}^2/\text{GeV} \) [30], the contribution of scalar \( S \) to the scattering on nucleon is suppressed because of its small faction \( \epsilon_S = 2.21 \times 10^{-3} \). The value of expression on left of Eqn. 28 is \( 1.3 \times 10^{-48} \text{ cm}^2/\text{GeV} \), which is smaller than current LUX bound. Thus, the fermion dark matter is dominant in this scenario, while the scalar dark matter must be less than 4% to escape current LUX bound.

C. Collider Signatures

As shown in our bench mark point (Eqn. 22), the new particles are all at the electroweak scale, which makes them testable at LHC. Interactions between these new particles and SM Higgs \( h \) will of course modify the properties of \( h \), thus give us some indirect hints. Nowadays the most precise measurement of \( M_h \) is the combined results of ATLAS and CMS [45]:

\[
M_h = 125.09 \pm 0.21 \text{(stat.)} \pm 0.11 \text{(sys.) GeV}
\]  

(29)

Apparently, the extra of new scalars and fermions would change the decay rates of SM Higgs \( h \). For instance, mixing between \( h \) and additional scalar singlet \( H \) will modify tree-level \( h \) decays. And the additional charged-scalars \( \eta_{1,2}^\pm \) will contribute to the loop-induced decays as \( h \to \gamma \gamma \) [46].

It’s well known that, for Higgs-portal dark matter, upper limits on Higgs invisible decay is also interpreted an upper limits on dark matter-nucleon scattering cross section [47]. Direct measurement of Higgs invisible decay in associated with \( Z \) by ATLAS set an upper limit of 75% at 95% C.L. [48]. Combined analysis with Higgs signal strength give a more tight upper limit of 37% at 95% C.L. [49–51]. In the future, the weak boson fusion channel might have the ability to probe invisible decay to 2-3% with 3000 fb\(^{-1} \) at LHC [52]. Light dark matter candidates in our model will contribute the the invisible decays of \( h \) in our model. In case of scalar dark matter, the branching ratio of Higgs invisible decay, i.e., \( \text{BR}(h \to SS^*) \) is 1.7% for \( M_S = 62 \text{ GeV} \) with \( \lambda_{s2} = \lambda_{s2} = -0.001 \), \( \sin \theta = 0.3 \). On the other hand for fermion dark matter, \( \text{BR}(h \to FF) \) is \( 1.7 \times 10^{-6} \) for \( M_F = 62 \text{ GeV} \) with \( v_F = 8 \text{ TeV} \), \( \sin \theta = 0.3 \), although it might not be favored by the constrains on relic density of dark matter.

Then we discuss the mixing between \( h \) and \( H \). The analysis of signal strength of \( h \) constrains \( \sin^2 \theta < 0.23 \) at 95% C.L. [53, 54]. Direct search of \( H \) in the \( ZZ \) and \( WW \) channel now has push this limit down to
FIG. 5. Branching ratios of $H$ as a function of $M_H$ for $\sin \theta = 0.3$ (left) and $\sin \theta = 0.01$ (right) in our benchmark point in Eqn. 22.

$\sin^2 \theta < 0.1$ with no new physics contribute to the decays of $H$ [55]. Future hadron collider, i.e., HL-LHC, has the ability to probe $\sin^2 \theta \sim 4 \times 10^{-2}$, and lepton collider, i.e., CEPC, could reach $\sin^2 \theta \sim 2 \times 10^{-3}$ [56]. In Fig. 5, we show the branching ratios of $H$ for two values of $\sin \theta(0.3, 0.01)$. For a relatively large mixing angle $\sin \theta = 0.3$, the heavy neutral Higgs $H$ decays dominantly into SM particles. The branching ratio of invisible decay $H \rightarrow SS^*$ can reach 10% for $M_H \sim 165$ GeV, and we expect it becomes dominant when $M_H < 160$ GeV. The branching ratios of $H$ decaying into other new physical particles are below $10^{-3}$ in this case. While for $M_H \gg M_W$, it is well known that decays of $H$ into vector bosons are determined by their Goldstone nature, which implies:

$$BR(H \rightarrow hh) \approx BR(H \rightarrow ZZ) \approx \frac{1}{2}BR(H \rightarrow WW)$$  (30)

The asymptotic behavior of this relation is clear shown on left picture of Fig. 5. On the other aspect, for a relatively small mixing angle $\sin \theta = 0.01$, decays of $H$ into SM particles will be suppressed and decays into new particles will be greatly enhanced. $H \rightarrow SS^*$ is dominant when $M_H < 400$ GeV. The branching ratio of $\bar{\psi} \psi$ will reach about 0.25 when $M_H \sim 600$ GeV, which is comparable with $H \rightarrow W^+W^-$. In this case, $H \rightarrow FF$ is below 10% and $H \rightarrow \eta^+\eta^-/A_0^0A_1^0$ is below 2%.

The heavy neutral Higgs $H$ is testable for large mixing angle $\theta$. For example, the promising channels to probe heavy neutral Higgs $H$ would be $ZZ \rightarrow 4l$, $ZZ \rightarrow 2l2\nu$, $ZZ \rightarrow 2l2j$, $ZZ \rightarrow 2l2\tau$, $WW \rightarrow 2l2\nu$, $WW \rightarrow l\nu2j$, $hh \rightarrow 4b$, and $hh \rightarrow 2b2\gamma$ [57]. At the same time, we would like to mention that the heavy Higgs $H$ could enhance the di-Higgs production $hh$ [58] by a factor of 18 comparing to standard model
case [59]. For small mixing angle $\theta$, production of $H$ will be suppressed by this small $\theta$, thus makes it challenging to probe directly at colliders.

![Graph](image)

**FIG. 6.** Total decay width of $Z'$ as a function of $M_{Z'}$ (for fixed values of $g_{B-L}$), and $g_{B-L}$ (for fixed values of $M_{Z'}$).

Next we review the properties of $U(1)_{B-L}$ gauge boson $Z'$. With about 20 fb$^{-1}$ data at 8 TeV LHC, bound on $Z'$ has been push up to 2.95 TeV by CMS through the ratio $R_\sigma = \sigma(pp \to Z' \to \ell^+\ell^-)/\sigma(pp \to Z \to \ell^+\ell^-)$, where $\ell = e, \mu$ [60]. In our benchmark point, we choose $M_{Z'} = 4$ TeV and $g_{B-L} = 0.5$ ($\nu_\sigma = 8$ TeV), which can safely satisfy current experimental limits and can be tested at 14 TeV LHC with 100 fb$^{-1}$ [62, 63]. Fig. 6 shows the total decay width of $Z'$ as a function of $M_{Z'}$ and $g_{B-L}$. Depending on $g_{B-L}$, $\Gamma_{Z'}$ varies from a few to hundreds of GeV. For such large $\Gamma_{Z'}$, it can be directly measured by the leptonic final states at LHC [62, 63].

In table II we give the decay branching ratios of $Z'$ in our benchmark point. The dominant decay channels of $Z'$ are $q\bar{q}, l\bar{l}$, and $\nu_L\bar{\nu}_L$, while all of the new particle final states only account for about 20%. A distinct feature of $U(1)_{B-L}$ gauge boson $Z'$ is the definite relation between quark and lepton final states:

$$BR(Z' \to q\bar{q}) : BR(Z' \to l\bar{l}) \simeq 2 : 3$$ (31)

after summing over all flavors. This relation can be used to distinguish $U(1)_{B-L}$ gauge boson $Z'$ from $Z'$ in other models [20]. More practical on experiment, the B-L nature of $Z'$ can be tested if $BR(Z' \to b\bar{b})/BR(Z' \to \mu^+\mu^-) = 1/3$ is confirmed [13]. In our model with only left-handed light neutrinos, the dominant invisible decay channel of $Z'$ is $BR(Z' \to \nu_L\bar{\nu}_L)$, which is half of $BR(Z' \to l\bar{l})$. Further with dark matter candidate in our model, $Z'$ invisible decays get additional contributions from $Z'$ into dark
matter pairs. For instance, $\text{BR}(Z' \to \text{inv.})$ could be 0.2457, 0.1990, 0.1964 for $FF$, $FS$ and $SS$ dark matter separately. So precise measurement of $\text{BR}(Z' \to \text{inv.})$ will shed light on the nature of dark matter.

| $q\bar{q}$ | $t\bar{t}$ | $\nu_L\bar{\nu}_L$ | $\psi\bar{\psi}$ | $N_{R1}\bar{N}_{R1}$ | $N_{R2}\bar{N}_{R2}$ | $\chi_1\bar{\chi}_3$ | $\chi_2\bar{\chi}_2$ |
|------------|------------|-------------------|-----------------|----------------------|----------------------|------------------|------------------|
| 0.25       | 0.38       | 0.19              | 0.063           | 0.0077               | 0.0077               | 0.016            | 0.032            |
| $HH$       | $hh$       | $A_i^0A_1^0$      | $H_i^0H_1^0$    | $A_2^0A_2^0$         | $H_2^0H_2^0$         | $\eta_1^+\eta_1^-$| $\eta_2^+\eta_2^-$|
| 0.030      | 0          | 0.0076            | 0.0051          | 0.0010               | 0.0013               | 0.0076           | 0.0010           |

TABLE II. Branching ratios of $Z'$ in our benchmark point. Here, we set $M_H = 300$ GeV and $\sin \theta = 0$ for simplicity.

Another interesting feature of our model is the existence of heavy Dirac fermion $\psi$, thus there are no lepton number violation (LNV) decays as $\psi \to W^-l^+$. For $M_\psi < M_h$, the Higgs decay into a pair of light and heavy neutrinos, $h \to \bar{\nu}_\psi + \bar{\psi}\nu$ will open, which could increase $\Gamma_h$ by up to almost 30% and significantly affects Higgs searches at the LHC [61]. In this paper, we consider $M_\psi > M_h$. Therefore, decay channels of $\psi$ could be $W^+l^-$, $Z\nu$, $h\nu$, and if kinematically allowed $H\nu$, $A_1^0N_j$, $A_1^0\chi_j$, $H_i^0N_j$, $H_i^0\chi_j (i, j = 1, 2)$ are also possible. Due to tiny mixing angle $\theta_{1,2}$, branching ratios of $\psi \to A_i^0N_j, A_i^0\chi_j$ are negligible. In Fig. 7 we show the branching ratios of $\psi$. It is clear that $\psi$ will decays dominantly into standard model final states for comparable Yukawa couplings of $y$ and $h$. Approximately for $M_\psi \gg M_W$, we have:

$$\frac{1}{\cos^2 \theta} \text{BR}(\psi \to h\nu) \approx \text{BR}(\psi \to Z\nu) \approx \frac{1}{2} \text{BR}(\psi \to W^+l^-)$$ (32)
Decays of $\psi$ into new physical particles is small in this case. $\text{BR}(\psi \rightarrow H_2^0 \chi)$ is about 10%, once is kinematically opened. $\text{BR}(\psi \rightarrow H \nu)$ is suppressed by $\sin^2 \theta$, thus it is always much smaller. As shown in table II, $\text{BR}(Z' \rightarrow \psi \bar{\psi}) \approx 0.063$ for one generation in our model, so $\psi \bar{\psi}$ can be produced through $Z'$ portal. A possible promising signature is the tri-lepton channel [62]:

$$pp \rightarrow Z' \rightarrow \psi \bar{\psi} \rightarrow W^+ l^- + W^- l^+ \rightarrow 2l^\pm l^\mp jj + E_T$$  \hspace{2cm} (33)

The cross section of this tri-lepton signal is about 0.017 fb in our benchmark point, so the tri-lepton is only promising at future high-luminosity LHC. The mass of $\psi$ can be reconstructed using the transverse mass of two opposite sign leptons with missing transverse momentum [62]. Another feature of $\psi$ is the possible large mixing with $\nu_L$ comparing to canonical type-I seesaw [3]. As discussed in Sec. II C, the mixing $V_{\nu \psi}$ between $\nu_L$ and $\psi_L$ is $M_D/\sqrt{M_D^2 + M_{\psi'}^2}$. For $M_D \sim O(1)\text{GeV}$, $M_{\psi'} \sim O(100)\text{GeV}$, $V_{\nu \psi} \sim O(10^{-2})$.

Thereafter, $\psi$ could be largely associated produced with charged leptons through $W$ [64]:

$$pp \rightarrow W^* \rightarrow l^\pm \psi \rightarrow l^\pm + W^\pm l^\mp \rightarrow l^\pm l^\mp jj$$  \hspace{2cm} (34)

$$pp \rightarrow W^* \rightarrow l^\pm \psi \rightarrow l^\pm + W^\pm l^\mp \rightarrow l^\pm l^\mp l^\mp E_T$$  \hspace{2cm} (35)

The production cross section $\sigma(l^\pm \psi) = 350 \times |V_{\nu \psi}|^2$ fb in our benchmark point. And it might be promising at 14TeV LHC with about 100 fb$^{-1}$. Testability of this heavy Dirac neutrino $\psi$ is less promising than the heavy Majorana neutrinos with same mixing scale, since the latter could rise LNV signatures [65-68].

Finally, we discuss the decays of inert scalars and fermions. $N_{R1}$ and $H_2^0$ are dark matter candidate as in our benchmark point in Eqn. [22] Decays of $N_{R2}$ are dominated by $N_{R2} \rightarrow l^\pm \eta_1^{\mp*} \rightarrow l^\pm l^\mp N_{R1}$ and $N_{R2} \rightarrow \nu \bar{A}_1^0 \rightarrow \nu \nu N_{R1}$ through the Yukawa coupling $f$. Decays of $\chi_i$ are $\chi_i \rightarrow H_2^0 \psi^*$ with the off-shell $\psi^*$ further decaying into $W^+l^-/Z\nu/h\nu/H_2^0 \nu$. $A_1^0$ and $\eta_1^\pm$ mainly decay through the Yukawa coupling $f$, which leads to $A_1^0 \rightarrow \nu N_{Ri}$ and $\eta_1^\pm \rightarrow l^\pm N_{Ri}$. The heavy $Z_2$ odd scalar $H_2^0$ decays into $\psi N_{Ri}$ through Yukawa coupling $h_\alpha$ and into $h A_1^0$ through trilinear coupling $\mu_1$. Similar, decays of $A_2^0$ are $A_2^0 \rightarrow \nu \chi_i$ and $A_2^0 \rightarrow hH_2^0$, while decays of $\eta_2^\pm$ are $\eta_2^\pm \rightarrow l^\pm \chi_i$ and $\eta_2^\pm \rightarrow W^\pm H_2^0$.

| Particles | $\eta_1^+ \eta_1^-$ | $\eta_1^\pm A_1^0$ | $A_1^0 A_1^{0*}$ | $\eta_2^+ \eta_2^-$ | $\eta_2^\pm A_2^0$ | $A_2^0 A_2^{0*}$ |
|-----------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|
| $\sigma$ (in fb) | 5.8 | 16 | 3.6 | 0.089 | 0.31 | 0.075 |

TABLE III. Production cross sections for inert scalar doublets.

The inert scalar doublets can be pair production through DY process. In Table III we list the production cross sections for inert scalar doublets. Many signatures can be risen from the inert particles. In Ma’s scotogenic model [9], promising signals of the doublet scalar on collider are multi-lepton final states with
missing transverse energy $E_T$ \cite{69,71}. Similar signals can also be produced in our model, for example:

\begin{align*}
2l + E_T: & \eta_1^+ \eta_1^- \rightarrow l^+ N_{R1} + l^- N_{R1} \\
& : \eta_1^+ \eta_1^- \rightarrow W^+ H_2^{0*} + W^- H_2^0 \rightarrow l^+ \nu_l H_2^{0*} + l^- \bar{\nu}_l H_2^0 \\
3l + E_T: & \eta_1^+ A_1^0 \rightarrow l^+ N_{R2} + \nu N_{R1} \rightarrow l^+ l^\pm l^\mp N_{R1} + \nu N_{R1} \\
& : \eta_2^+ A_2^0 \rightarrow W^\pm H_2^{0(*)} + h H_2^0 \rightarrow l^\pm \nu_l H_2^{0(*)} + l^\pm \nu_l \nu_l H_2^0 \\
4l + E_T: & \eta_1^+ \eta_1^- \rightarrow l^+ N_{R2} + l^- N_{R1} \rightarrow l^+ l^\pm l^\mp N_{R1} + l^- N_{R1} \\
& : A_2^0 H_2^{0*} \rightarrow h H_2^0 H_2^{0*} \rightarrow l^+ l^- l^\mp l^- N_{R1} + l^- N_{R1} \\
\end{align*}

With much different decay topologies between our and Ma’s model, it would be distinguishable even with same signals. Apart from these multi-lepton signals, there are also some other interesting signals in our model, i.e.:

\begin{align*}
2l^\pm jj + E_T: & \eta_2^+ A_2^0 \rightarrow W^\pm H_2^{0(*)} + h H_2^0 \rightarrow l^\pm \nu_l H_2^{0(*)} + l^\pm \nu_l j j H_2^0 \\
3l^\pm \bar{b}b + E_T: & \eta_2^+ A_2^0 \rightarrow W^\pm H_2^{0(*)} + h H_2^0 \rightarrow l^\pm \nu_l H_2^{0(*)} + \bar{b}b H_2^0 \\
\end{align*}

The lepton number violation signal $2l^\pm jj + E_T$ suffers much lower SM background, thus might make this signal very promising on LHC. The $l^\pm \bar{b}b + E_T$ has a relatively large production rate due to $h \rightarrow b \bar{b}$ dominant in $h$ decay, so it might also be promising.
IV. CONCLUSIONS

In usual canonical seesaw mechanisms, it requires the heavy states with the scale of masses being grand unification scale to generated the small neutrino mass. In linear seesaw scenario with $m_\nu \simeq \mu_L M_D/M_\Psi$, the neutrino masses suffers a two fold suppression by both lepton number symmetry violating term $\mu_L$ and heavy mass $M_\Psi$. The linear seesaw model can lower the seesaw scale such that new physics may arise at TeV scale. In this work, we construct a radiated linear seesaw model where the naturally small term $\mu_L$ are generated at one-loop level and its soft-breaking of lepton number symmetry attributes to the SSB of B-L symmetry at TeV scale. To satisfy the anomalies cancelation, the value of B-L charges for inert particles are found to be exotic such that there exists residual $Z_2 \times Z_2'$ symmetry even after SSB of $B - L$ gauge symmetry. It is shown that the residual symmetry stabilizes the inert particles as dark matter candidates. In our model, we introduce two no-interplay classes of inert particles to realize the model such that the lightest inert particles belonging in each class play is the dark matter matter candidate. Therefore we have propose a two-component dark matter model. The seesaw scale of radiated linear seesaw scale can be as low as a few hundred GeV, leading to interesting phenomenology.

Given a benchmark point at electro-weak scale, we illustrate the main prediction of our model. For the Yukawa coupling $f_{l_i}(f_l)$ at 0.01 order, our benchmark point predicts $\text{Br}(\mu \rightarrow e\gamma) \sim 10^{-13}$, an order slightly under the current constraints and in the reach of the forthcoming experiment. The two component dark matter candidates are realized in our model. To account for the observed relic density, the annihilation of dark matter are dominant by the $s$-channel scalars $h/H$ or gauge boson $Z'$. For the fermion DM, we find that the $h$-channel is excluded. But it is still allowed for the $H$-channel and $Z'$-channel. On the contrary, for the scalar DM, the $Z'$-channel is excluded while the $h/H$-channel is allowed. And the heavy Higgs $H$ also plays a vital import rule in the conversion between fermion and scalar dark matter. Collider signatures of our model are also very rich. The precise measurements of SM Higgs $h$ will put tight constrain light scalar DM and heavy scalar $H$. With a relatively large mixing angle $\sin \theta = 0.3$, the $H \rightarrow ZZ, W^+W^-$ channels are testable at LHC. For the extra $Z$ boson and heavy lepton $\psi$, the tri-lepton channel of $Z \rightarrow \tilde{\psi}\psi$ is promising at HL-LHC. With a larger cross section, the associated production of $l^\pm \psi$ may be more promising. The inert doublet scalar can also produce multi-lepton channels. And some distinct channels, as $2\ell^\pm j j + E_T$, $\ell^\pm b b + E_T$, can be used to distinguish our model.

Finally, we would like to mention that the radiated linear seesaw model we proposed is the minimal version where only one $\Psi$ fermion mediator is included. In this scenario, $M_\nu$ is a rank-2 mass matrix and the lightest neutrino must be massless. However, more complicated scenarios exist, corresponding to other solutions of the anomaly free condition. Then one may obtain the rank-3 neutrino mass matrix.
Such scenarios predict more new particles with different $B - L$ charges, the model construction and the phenomenology deserve us further study.

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[1] Q. R. Ahmad et al. (SNO Collaboration), Phys. Rev. Lett. 89, 011301 (2002); K. Eguchi et al. (KamLAND Collaboration), Phys. Rev. Lett 90, 021802 (2003); M. H. Ahn et al. (K2K Collaboration), Phys. Rev. Lett 90, 041801 (2003); F. P. An et al. (DAYA-BAY Collaboration), Phys. Rev. Lett. 108, 171803 (2012); J. K. Ahn et al. (RENO Collaboration), Phys. Rev. Lett108, 191802 (2012).

[2] S. Weinberg, Phys. Rev. Lett. 43, 1566 (1979).

[3] P. Minkowski, Phys. Lett. B 67 (1977) 421; T. Yanagida, in Proceedings of the Workshop on Unified Theories and Baryon Number in the Universe, eds. O. Sawada et al., (KEK Report 79-18, Tsukuba, 1979), p. 95; M. Gell-Mann, P. Ramond, R. Slansky, in Supergravity, eds. P. Van Nieuwenhuizen et al., (North-Holland, 1979), p. 315; S. Glashow, in Quarks and Leptons, Cargèse, eds. M. Lévy et al., (Plenum, 1980), p. 707; R. N. Mohapatra, G. Senjanović, Phys. Rev. Lett. 44 (1980) 912.

[4] R. N. Mohapatra and G. Senjanovic, Phys, Rev. Lett. 44, 912 (1980); J. Schechter and J. W. F. Valle, Phys. Rev. D22, 2227 (1980); J. Schechter and J. W. F. Valle, Phys. Rev. D25, 774 (1982).

[5] D. Weyler and L. Wolfenstein, Nucl. Phys. B 218, 205 (1983); R. N. Mohapatra and J. W. F. Valle, Phys. Rev. D 34, 1642 (1986); E. Ma, Phys. Lett. B 191, 287 (1987)

[6] E. Akhmedov, M. Lindner, E. Schnapka, J. W. F. Valle, Phys. Lett. B 368, 270 (1996); E. Akhmedov, M. Lindner, E. Schnapka, J. W. F. Valle, Phys. Rev. D53, 2752 (1996);

[7] M. Malinsky, J. C. Romao, and J. W. F. Valle, Phys. Rev. Lett. 95, 161801 (2005).

[8] A. Zee, Phys. Lett. B 93, 389 (1980); A. Pilaftsis, Z. Phys. C 55, 275 (1992).

[9] E. Ma, Phys. Rev. D 73, 077301 (2006) [hep-ph/0601225].

[10] J. Kubo, E. Ma and D. Suematsu, Phys. Lett. B 642, 18 (2006); P. H. Gu and U. Sarkar, Phys. Rev. D 77, 105031 (2008); E. Ma and D. Suematsu, Mod. Phys. Lett. A 24, 583 (2009); D. Aristizabal Sierra, J. Kubo, J. Suematsu, D. Restrepo and O. Zapata, Phys.Rev. D79, 013011 (2009); S. Kanemura, T. Nabeiisma and H. Sugiyama, Phys. Lett. B703, 66 (2011); S. Kanemura, O. Seto, T. Shimo, Phys.Rev.D 84, 016004 (2011); S. Kanemura and H. Sugiyama, Phys. Rev. D86, 073006 (2012); H. Okada and T. Toma, Phys. Rev. D86, 033011 (2012); P. S. Bhupal Dev and A. Pilaftsis, Phys. Rev. D 86, 113001 (2012); D. Schmidt, T. Schwetz and T. Toma, Phys. Rev. D 85, 073009 (2012); M. Hirsch, R. Lineros, S. Morisi, J. Palacio, N. Rojas, et al, JHEP 1310, 149 (2013);
D. Restrepo, O. Zapata, and C. E. Yaguna, JHEP 1311, 011 (2013); S. S. Law and K. L. McDonald, JHEP 1309, 092 (2013); S. Kanemura, T. Matsui and H. Sugiyama, Phys. Lett. B 727, 151 (2013); V. Brdar, I. Picek, and B. Radovcic, Phys.Lett. B 728, 198 (2014); H. Okada, K. Yagyu, Phys. Rev. D 90, 035019(2014);

[11] A. Zee, Nucl. Phys. B 264, 99 (1986); K. S. Babu, Phys. Lett. B 203, 132(1988).

[12] E. Ma, Phys. Lett. B 662, 49 (2008) [arXiv:0708.3371 [hep-ph]].

[13] S. Kanemura, T. Nabeshima and H. Sugiyama, Phys. Rev. D 85, 033004 (2012) [arXiv:1111.0599 [hep-ph]].

[14] S. Kanemura, T. Matsui and H. Sugiyama, Phys. Rev. D 90, 013001 (2014) [arXiv:1405.1935 [hep-ph]].

[15] M. Aoki and T. Toma, JCAP 1409, 016 (2014) [arXiv:1405.5870 [hep-ph]].

[16] L. M. Krauss, S. Nasri and M. Trodden, Phys. Rev. D 67, 085002 (2003); K. Cheung and O. Seto, Phys. Rev. D 69, 113009 (2004); M. Aoki, S. Kanemura and O. Seto, Phys. Rev. Lett. 102, 051805 (2009); M. Gustafsson, J. M. No, and M. A. Rivera, Phys. Rev. Lett 110, 211802 (2013); J. N. Ng and A. de la Puente, Phys. Lett. B 727, 204 (2013); A. Ahriche, C.-S. Chen, K. L. McDonald, and S. Nasri, Phys. Rev. D 90, 015024 (2014); C.-S. Chen, K. L. McDonald and S. Nasri,Phys. Lett. B 05, 082 (2014).

[17] F. Bonnet, M. Hirsch, T. Ota and W. Winter, JHEP 1207, 153 (2012) [arXiv:1204.5862 [hep-ph]].

[18] E. Ma, Mod. Phys. Lett. A 17, 535 (2002) [hep-ph/0112232].

[19] S. Khalil, J. Phys. G 35, 055001 (2008) [hep-ph/0611205].

[20] P. Langacker, Rev. Mod. Phys. 81, 1199 (2009) [arXiv:0801.1345 [hep-ph]].

[21] G. Cacciapaglia, C. Csaki, G. Marandella and A. Strumia, Phys. Rev. D 74, 033011 (2006) [hep-ph/0604111].

[22] G. Aad et al. [ATLAS Collaboration], Phys. Lett. B 716, 1 (2012) [arXiv:1207.7214 [hep-ex]].

[23] S. Chatrchyan et al. [CMS Collaboration], Phys. Lett. B 716, 30 (2012) [arXiv:1207.7235 [hep-ex]].

[24] V. Barger, P. Langacker and G. Shaughnessy, Phys. Rev. D 75, 055013 (2007) [hep-ph/0611239].

[25] V. Barger, P. Langacker, M. McCaskey, M. J. Ramsey-Musolf and G. Shaughnessy, Phys. Rev. D 77, 035005 (2008) [arXiv:0706.4311 [hep-ph]].

[26] T. Robens and T. Stefaniak, Eur. Phys. J. C 75, no. 3, 104 (2015) [arXiv:1501.02234 [hep-ph]].

[27] E. Ma, Phys. Rev. D 80, 013013 (2009) [arXiv:0904.4450 [hep-ph]].

[28] H. Fritzsch and Z. Z. Xing, Phys. Lett. B 372, 265 (1996); Phys. Lett. B 413, 396 (1997); Phys. Lett. B 440, 313 (1998); Y. Koide, Mod. Phys. Lett. A 11 2849 (1996); Y. Koide, S. Fusaoka, Prog. Theor. Phys 97 459 (1997); Z. Z. Xing, D. Yang, S. Zhou, Phys. Lett. B 690, 304 (2010); S. Dev, S. Gupta, R. R. Gautam, Phys. Lett. B 702 28 (2011);

[29] S. K. Garg, S. Gupta, JHEP 1310, 128 (2013);

[30] D. S. Akerib et al. [LUX Collaboration], Phys. Rev. Lett. 112, 091303 (2014) [arXiv:1310.8214 [astro-ph.CO]].

[31] T. Toma and A. Vicente, JHEP 1401, 160 (2014) [arXiv:1312.2840 [hep-ph]].

[32] A. Vicente and C. E. Yaguna, JHEP 1502, 144 (2015) [arXiv:1412.2545 [hep-ph]].

[33] J. Adam et al. [MEG Collaboration], Phys. Rev. Lett. 110, 201801 (2013) [arXiv:1303.0754 [hep-ex]].

[34] J. Hisano, T. Moroi, K. Tobe and M. Yamaguchi, Phys. Rev. D 53, 2442 (1996) [hep-ph/9510309].

[35] A. M. Baldini, F. Cei, C. Cerri, S. Dussoni, L. Galli, M. Grassi, D. Nicolo and F. Raffaelli et al., arXiv:1301.7225 [physics.ins-det].
[36] B. Aubert et al. [BaBar Collaboration], Phys. Rev. Lett. 104, 021802 (2010) [arXiv:0908.2381 [hep-ex]].
[37] K. Hayasaka, K. Inami, Y. Miyazaki, K. Arinstein, V. Aulchenko, T. Aushev, A. M. Bakich and A. Bay et al., Phys. Lett. B 687, 139 (2010) [arXiv:1001.3221 [hep-ex]].
[38] P. A. R. Ade et al. [Planck Collaboration], Astron. Astrophys. 571, A16 (2014) [arXiv:1303.5076 [astro-ph.CO]].
[39] G. Bertone, D. Hooper and J. Silk, Phys. Rept. 405, 279 (2005) [hep-ph/0404175].
[40] N. D. Christensen and C. Duhr, Comput. Phys. Commun. 180, 1614 (2009) [arXiv:0806.4194 [hep-ph]]; A. Alloul, N. D. Christensen, C. Degrande, C. Duhr and B. Fuks, Comput. Phys. Commun. 185, 2250 (2014) [arXiv:1310.1921 [hep-ph]].
[41] G. Blanger, F. Boudjema, A. Pukhov and A. Semenov, Comput. Phys. Commun. 192, 322 (2015) [arXiv:1407.6129 [hep-ph]]; G. Belanger, F. Boudjema, A. Pukhov and A. Semenov, Comput. Phys. Commun. 185, 960 (2014) [arXiv:1305.0237 [hep-ph]]; G. Belanger, F. Boudjema, A. Pukhov and A. Semenov, Comput. Phys. Commun. 176, 367 (2007) [hep-ph/0607059].
[42] S. Esch, M. Klasen and C. E. Yaguna, JHEP 1409, 108 (2014) [arXiv:1406.0617 [hep-ph]].
[43] Q. H. Cao, E. Ma, J. Wudka and C.-P. Yuan, arXiv:0711.3881 [hep-ph].
[44] M. Aoki, M. Duerr, J. Kubo and H. Takano, Phys. Rev. D 86, 076015 (2012) [arXiv:1207.3318 [hep-ph]].
[45] G. Aad et al. [ATLAS and CMS Collaborations], Phys. Rev. Lett. 114, 191803 (2015) [arXiv:1503.07589 [hep-ex]].
[46] A. Arhrib, R. Benbrik and N. Gaur, Phys. Rev. D 85, 095021 (2012) [arXiv:1201.2644 [hep-ph]].
[47] S. Baek, P. Ko and W. I. Park, Phys. Rev. D 90, no. 5, 055014 (2014) [arXiv:1405.3530 [hep-ph]].
[48] G. Aad et al. [ATLAS Collaboration], Phys. Rev. Lett. 112, 201802 (2014) [arXiv:1402.3244 [hep-ex]].
[49] P. Bechtle, S. Heinemeyer, O. Stål, T. Stefaniak and G. Weiglein, JHEP 1411, 039 (2014) [arXiv:1403.1582 [hep-ph]].
[50] T. Corbett, O. J. P. Eboli, D. Goncalves, J. Gonzalez-Fraile, T. Plehn and M. Rauch, arXiv:1505.05516 [hep-ph].
[51] G. Aad et al. [ATLAS Collaboration], arXiv:1507.04548 [hep-ex].
[52] C. Bernaciak, T. Plehn, P. Schichtel and J. Tattersall, Phys. Rev. D 91, no. 3, 035024 (2015) [arXiv:1411.7699 [hep-ph]].
[53] P. P. Giardino, K. Kannike, I. Masina, M. Raidal and A. Strumia, JHEP 1405, 046 (2014) [arXiv:1303.3570 [hep-ph]].
[54] A. Falkowski, F. Riva and A. Urbano, JHEP 1311, 111 (2013) [arXiv:1303.1812 [hep-ph]].
[55] M. Pelliccioni [CMS Collaboration], arXiv:1505.03831 [hep-ex].
[56] S. Dawson, A. Gritsan, H. Logan, J. Qian, C. Tully, R. Van Kooten, A. Ajaib and A. Anastassov et al., arXiv:1310.8361 [hep-ex].
[57] D. Buttazzo, F. Sala and A. Tesi, arXiv:1505.05488 [hep-ph].
[58] J. Baglio, A. Djouadi, R. Gröber, M. Mihleitner, J. Quevillon and M. Spira, JHEP 1304, 151 (2013) [arXiv:1212.5581 [hep-ph]].
[59] C. Y. Chen, S. Dawson and I. M. Lewis, Phys. Rev. D 91, no. 3, 035015 (2015) [arXiv:1410.5488 [hep-ph]].
[60] V. Khachatryan et al. [CMS Collaboration], JHEP 1504, 025 (2015) [arXiv:1412.6302 [hep-ex]].
[61] J. H. Chen, X. G. He, J. Tandean and L. H. Tsai, Phys. Rev. D 81, 113004 (2010) [arXiv:1001.5215 [hep-ph]].
[62] L. Basso, A. Belyaev, S. Moretti and C. H. Shepherd-Themistocleous, Phys. Rev. D 80, 055030 (2009) [arXiv:0812.4313 [hep-ph]].
[63] L. Basso, A. Belyaev, S. Moretti, G. M. Pruna and C. H. Shepherd-Themistocleous, Eur. Phys. J. C 71, 1613 (2011) [arXiv:1002.3586 [hep-ph]].
[64] G. Bambhaniya, S. Goswami, S. Khan, P. Konar and T. Mondal, Phys. Rev. D 91, 075007 (2015) [arXiv:1410.5687 [hep-ph]].
[65] T. Han and B. Zhang, Phys. Rev. Lett. 97, 171804 (2006) [hep-ph/0604064].
[66] F. del Aguila, J. A. Aguilar-Saavedra and R. Pittau, JHEP 0710, 047 (2007) [hep-ph/0703261].
[67] A. Atre, T. Han, S. Pascoli and B. Zhang, JHEP 0905, 030 (2009) [arXiv:0901.3589 [hep-ph]].
[68] F. F. Deppisch, P. S. B. Dev and A. Pilaftsis, New J. Phys. 17, 075019 (2015) [arXiv:1502.06541 [hep-ph]].
[69] E. Dolle, X. Miao, S. Su and B. Thomas, Phys. Rev. D 81, 035003 (2010) [arXiv:0909.3094 [hep-ph]].
[70] X. Miao, S. Su and B. Thomas, Phys. Rev. D 82, 035009 (2010) [arXiv:1005.0090 [hep-ph]].
[71] M. Gustafsson, S. Rydbeck, L. Lopez-Honorez and E. Lundstrom, Phys. Rev. D 86, 075019 (2012) [arXiv:1206.6316 [hep-ph]].