$U(1)_{B-3L_\alpha}$ extended scotogenic models and single-zero textures of neutrino mass matrices

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In this paper, we propose a scotogenic model of neutrino masses with flavor-dependent $U(1)_{B-3L_\alpha}$ ($\alpha = e, \mu, \tau$) gauge symmetry, in which neutrinos are generated at one-loop level and a fermion dark matter is naturally accommodated. In this model, three one-zero-texture structures of neutrino mass matrix, denoted as patterns A, B and C, are successfully realized by appropriate charge assignment for inert fermions and scalar fields. For each predicted texture, a detailed numerical analysis is carried out to find the allowed regions of neutrino mixing and masses. We find three scenarios (A for normal mass hierarchy and B and C for inverted mass hierarchy) favored by the latest data of neutrino oscillation experiments and Planck 2018 limit on the sum of neutrino masses. Other phenomenologies such as lepton flavor violation, dark matter, and collider signatures are also discussed.

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I. INTRODUCTION

Thanks to the overwhelming evidence from astrophysics, cosmology, and neutrino oscillation experiments, it is well established that the new physics beyond Standard Model (SM) is needed to accommodate two missing spices: tiny but nonzero neutrino masses and dark matter (DM). The scotogenic model, originally proposed by Krauss *et al.* [1] and Ma [2], is one of the appealing ways to explain the above two issues simultaneously, where the small neutrino masses arise only at loop level and DM plays a role of mediated field propagating inside the loop diagram. For a review on various interesting realizations, see Ref. [3].

In addition to the neutrino masses and DM mystery, the SM gauge symmetries cannot explain how different flavors of leptons distinguish each other and exhibit the observed mixing patterns. A promising ansatz which can potentially solves the flavor puzzle is to add an extra $U(1)$ gauge symmetry to SM and permit the three flavors of fermions transform nonuniversally under the new symmetry. On the other hand, three right-handed (RH) neutrinos N_R s are usually considered as necessary ingredients of the neutrino mass model. Then, assuming family nonuniversal charges of N_R s, the resulting lepton flavor structure admits a

number of anomaly-free solutions. Along this line of thinking, the $U(1)_X$ gauge symmetry, with X being the linear combinations of baryon number B and individual lepton numbers L_α , is frequently considered in various contexts. Well-known examples are $B - 3L_\tau$ [4], $L_\mu - L_\tau$ [5–8], $B + 3(L_e - L_\mu - L_\tau)$ [9], etc. In Ref. [10], the neutrino mass matrices containing two-zero-texture [11,12] or two-zero-cofactor structures [13] are realized by applying the $U(1)_{aB-\sum x_\alpha L_\alpha}$ gauge symmetry to type-I seesaw scenario. In Ref. [14], more solutions are found in the type-I and/or -III seesaw framework. It is then natural to ask if the flavor-dependent $U(1)$ gauge symmetry is compatible with the scotogenic models and leads to predictive lepton mass textures. The first attempts was made in Ref. [15], where the $U(1)_{L_\mu-L_\tau}$ gauge symmetry was adopted to realize the type-C two-zero-texture structure [12] in the one-loop induced neutrino mass matrix. This model can also interpret the $R_{K^{(*)}}$ anomaly and AMS-02 positron excess [16]. However, the latest analysis points out that such texture structure is incompatible with the Planck 2018 bound on the sum of neutrino masses, i.e., $\Sigma m_i < 0.12$ eV [17]. One way out of this problem is to consider different $U(1)$ gauge symmetries. In Refs. [18,19], other viable two-zero-texture structures are realized based on $U(1)_{B-2L_\alpha-L_\beta}$ gauge symmetry.

In this paper, we apply the $U(1)_{B-3L_\alpha}$ ($\alpha = e, \mu, \tau$) gauge symmetry to the scotogenic model. It is pointed out that by pairing one of N_R s with ν_{L_α} , the $U(1)_{B-3L_\alpha}$ gauged symmetry is anomaly free and may predict a realistic mixing structure among three families of leptons [4]. In the scotogenic model, we find that the form of Yukawa

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interactions is restricted by the $U(1)_{B-3L_\alpha}$ charge assignments, leading to a one-zero-texture structure of neutrino mass matrix. For each predicted texture, an updated phenomenological analysis is carried out in view of the current neutrino oscillation and cosmological data. With all new particles around TeV scale, the scotogenic model with a $B - 3L_\alpha$ gauge symmetry leads to testable phenomenologies. Thus, in addition to the neutrino issue, we discuss the phenomenologies such as lepton flavor violations (LFVs) and DM candidate and highlights of collider signatures. Without loss of generality, the $U(1)_{B-3L_\mu}$ gauge symmetry is considered as a case study.

The rest of this paper is organized as follows. In Sec. II, we first discuss the realization of one-zero textures in the scotogenic model with $U(1)_{B-3L_\alpha}$ gauge symmetry. In Sec. III, the numerical analysis of neutrino masses and mixing is presented. In Sec. IV, other predictions such as the lepton-flavor-violation rate, dark matter, and collider physics are discussed. The conclusion is given in Sec. V.

II. MODEL SETUP

We extend the scotogenic model proposed by Ma [2] to be compatible with the $U(1)_{B-3L_\alpha}$ gauge symmetry. The particle content related to lepton mass generation is listed in Table I. In the model, three RH fermion singlets N_{Ri} ($i = 1, 2, 3$) and an inert scalar doublet field η are added to the SM. Among three N_R fields, one of them carries the same nonzero charge as SM lepton doublet L_α under $U(1)_{B-3L_\alpha}$ gauge symmetry. In addition to the SM Higgs field Φ , a singlet scalar S is added to develop the vacuum expectation value (VEV) $v_S/\sqrt{2}$ after the spontaneous breaking of $U(1)_{B-3L_\alpha}$ symmetry. The unbroken discrete Z_2 symmetry forbids the tree-level neutrino Yukawa interaction and assures the stability of DM. In the original scotogenic model, the interactions relevant to lepton mass generation are given by

$$-\mathcal{L} \supset f_{\alpha\beta} \bar{L}_\alpha \Phi \ell_{\beta R} + h_{\alpha i} \bar{L}_\alpha \eta^c N_{Ri} + \frac{M_{Nij}}{2} N_{Ri} N_{Rj} + \frac{1}{2} \lambda (\Phi^\dagger \eta)^2 + \text{H.c.} \quad (1)$$

TABLE I. Particle content and charge assignment under $SU(2)_L \times U(1)_Y \times U(1)_{B-3L_\alpha} \times Z_2$ with (α, β, γ) a permutations of (e, μ, τ) .

Group	Lepton fields									Scalar fields		
	L_α	$\ell_{\alpha R}$	L_β	$\ell_{\beta R}$	L_γ	$\ell_{\gamma R}$	N_{R1}	N_{R2}	N_{R3}	Φ	η	S
$SU(2)_L$	2	1	2	1	2	1	1	1	1	2	2	1
$U(1)_Y$	$-\frac{1}{2}$	-1	$-\frac{1}{2}$	-1	$-\frac{1}{2}$	-1	0	0	0	$\frac{1}{2}$	$\frac{1}{2}$	0
Z_2	+	+	+	+	+	+	-	-	-	+	+	+
$U(1)_{B-3L_\alpha}$	-3	-3	0	0	0	0	0	0	-3	0	0	3

The mass matrix M_N can be diagonalized by unitary matrix V satisfying

$$V^T M_N V = \hat{M}_N \equiv \text{diag}(M_{N_1}, M_{N_2}, M_{N_3}). \quad (2)$$

The neutrino mass is generated radiatively, as show in Fig. 1, and can be computed exactly, i.e.,

$$(M_\nu)_{\alpha\beta} = \frac{1}{32\pi^2} \sum_k h_{\alpha i} V_{ik} h_{\beta j} V_{jk} M_{Nk} \left[\frac{m_R^2}{m_R^2 - M_{Nk}^2} \log \left(\frac{m_R^2}{M_{Nk}^2} \right) - \frac{m_I^2}{m_I^2 - M_{Nk}^2} \log \left(\frac{m_I^2}{M_{Nk}^2} \right) \right] \quad (3)$$

where $m_0^2 \equiv (m_R^2 + m_I^2)/2$. In the so-called radiative see-saw scenario, the usual assumption of $m_R^2, m_I^2 \ll M_{Nk}^2$ is adopted. Then, the neutrino mass matrix in Eq. (3) simplifies as

$$(M_\nu)_{\alpha\beta} \simeq -\frac{\lambda v^2}{32\pi^2} \sum_k \frac{h_{\alpha i} V_{ik} h_{\beta j} V_{jk}}{M_k} \left[\log \left(\frac{M_k^2}{m_0^2} \right) - 1 \right], \quad (4)$$

where m_R and m_I are the masses of $\Re\eta^0/\sqrt{2}$ and $\Im\eta^0/\sqrt{2}$. In this paper, we focus on another interesting scenario given in Ref. [2]. i.e., $m_R^2, m_I^2 \gg M_{Nk}^2$. In this case, the expression of neutrino mass matrix is given by

$$(M_\nu)_{\alpha\beta} \simeq -\frac{1}{32\pi^2} \frac{\lambda v^2}{m_0^2} \sum_k h_{\alpha i} V_{ik} h_{\beta j} V_{jk} M_{Nk} = -\frac{1}{32\pi^2} \frac{\lambda v^2}{m_0^2} (h M_N h^T)_{\alpha\beta}. \quad (5)$$

Now, we demonstrate how the one-zero-texture structures of M_ν arise from the charge assignments given by Table. I. Taking $U(1)_{B-3L_\tau}$ gauge symmetry as an example, the Yukawa matrix h , charged lepton mass matrix M_l , and the mass matrix of N_R are given by

$$h = \begin{pmatrix} h_{e1} & h_{e2} & 0 \\ h_{\mu 1} & h_{\mu 2} & 0 \\ 0 & 0 & h_{\tau 3} \end{pmatrix} \quad M_l = \frac{v}{\sqrt{2}} \begin{pmatrix} f_{ee} & f_{e\mu} & 0 \\ f_{\mu e} & f_{\mu\mu} & 0 \\ 0 & 0 & f_{\tau\tau} \end{pmatrix} \\ M_N = \begin{pmatrix} M_{11} & M_{12} & M_{13} \\ M_{12} & M_{22} & M_{23} \\ M_{13} & M_{23} & 0 \end{pmatrix}. \quad (6)$$

Note that the (1, 3) and (2, 3) entries in M_N originate from the interactions of $y_{13} N_{R1} N_{R3} S$ and $y_{23} N_{R2} N_{R3} S$. Given the results in Eq. (6) and using Eq. (5), M_ν can be written in the form

FIG. 1. Radiative neutrino mass at one loop.

$$M_\nu = \begin{pmatrix} (M_\nu)_{11} & (M_\nu)_{12} & (M_\nu)_{13} \\ (M_\nu)_{12} & (M_\nu)_{22} & (M_\nu)_{23} \\ (M_\nu)_{13} & (M_\nu)_{23} & 0 \end{pmatrix}, \quad (7)$$

where

$$\begin{aligned} (M_\nu)_{11} &= h_{e1}^2 M_{11} + 2h_{e1}h_{\mu1}M_{12} + h_{\mu1}^2 M_{22} \\ (M_\nu)_{12} &= h_{e1}h_{e2}M_{11} + (h_{e2}h_{\mu1} + h_{e1}h_{\mu2})M_{12} + h_{\mu1}h_{\mu2}M_{22} \\ (M_\nu)_{13} &= h_{\tau3}(h_{e1} + h_{\mu1})M_{13} \\ (M_\nu)_{22} &= h_{e2}^2 M_{11} + 2h_{e2}h_{\mu2}M_{12} + h_{\mu2}^2 M_{22} \\ (M_\nu)_{23} &= h_{\tau3}(h_{e2} + h_{\mu2})M_{23} \end{aligned} \quad (8)$$

in the unit of $-\lambda v^2/32\pi^2 m_0^2$. The result given in Eq. (7) directly reflects the one-zero-texture structure of the Majorana neutrino mass matrix. To have a more clear relation, the tedious expressions in Eq. (7) would be further simplified with diagonal Yukawa matrix $h = \text{diag}(h_{e1}, h_{\mu2}, h_{\tau3})$; then,

$$M_\nu = \begin{pmatrix} h_{e1}^2 M_{11} & h_{e1}h_{\mu2}M_{12} & h_{e1}h_{\tau3}M_{13} \\ h_{e1}h_{\mu2}M_{12} & h_{\mu2}^2 M_{22} & h_{\mu2}h_{\tau3}M_{23} \\ h_{e1}h_{\tau3}M_{13} & h_{\mu2}h_{\tau3}M_{23} & 0 \end{pmatrix}, \quad (9)$$

TABLE II. Possible mass textures of charged leptons (M_l) and neutrinos (M_ν), where $\times$ denotes a nonzero entry.

Pattern	Group	Texture of M_l	Texture of M_ν
A	$U(1)_{B-3L_e}$	$\begin{pmatrix} \times & 0 & 0 \\ 0 & \times & \times \\ 0 & \times & \times \end{pmatrix}$	$\begin{pmatrix} 0 & \times & \times \\ \times & \times & \times \\ \times & \times & \times \end{pmatrix}$
B	$U(1)_{B-3L_\mu}$	$\begin{pmatrix} \times & 0 & \times \\ 0 & \times & 0 \\ \times & 0 & \times \end{pmatrix}$	$\begin{pmatrix} \times & \times & \times \\ \times & 0 & \times \\ \times & \times & \times \end{pmatrix}$
C	$U(1)_{B-3L_\tau}$	$\begin{pmatrix} \times & \times & 0 \\ \times & \times & 0 \\ 0 & 0 & \times \end{pmatrix}$	$\begin{pmatrix} \times & \times & \times \\ \times & \times & \times \\ \times & \times & 0 \end{pmatrix}$

in the unit of $-\lambda v^2/32\pi^2 m_0^2$. The diagonal Yukawa matrix h can be obtained by imposing a discrete Z_3 symmetry. The corresponding Z_3 charges are $L_e, e_R, N_{R1}, S \sim \omega, L_\mu, \mu_R, N_{R2} \sim \omega^2$, and $L_\tau, \tau_R, N_{R3} \sim 0$ with $\omega = \exp(i2\pi/3)$. Therefore, the form of M_ν is proportional to M_N with an equal Yukawa structure $h_{e1} = h_{\mu2} = h_{\tau3}$.

The other two textures, predicted by $U(1)_{B-3L_e}$ and $U(1)_{B-3L_\mu}$ gauge symmetries, are produced in a similar approach. We list all of them in Table II. The M_l and M_ν are diagonalized by unitary matrix V_l and V_ν ,

$$M_l = V_l M_l^D V_l^\dagger \quad M_\nu = V_\nu M_\nu^D V_\nu^T, \quad (10)$$

where $M_l^D \equiv \text{diag}(m_e, m_\mu, m_\tau)$ and $M_\nu^D \equiv \text{diag}(m_1, m_2, m_3)$. Then, the Pontecorvo-Maki-Nakagawa-Sakata matrix U_{PMNS} [20] is given by

$$U_{\text{PMNS}} = V_l^\dagger V_\nu \quad (11)$$

and can be parametrized as

$$U_{\text{PMNS}} = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix} \times \text{diag}(e^{2i\rho}, e^{2i\sigma}, 1). \quad (12)$$

Then, we have

$$V_\nu = V_l U_{\text{PMNS}}. \quad (13)$$

In terms of the zero-texture condition $(M_\nu)_{\alpha\beta} = 0$, we obtain the constraint condition equation

$$\sum_{i=1,2,3} (V_\nu)_{\alpha i} (V_\nu)_{\beta i} m_i = 0. \quad (14)$$

From Table. II, one can see that the charged lepton matrix M_l is nondiagonal. However, the degrees of freedom in V_l do not change neutrino parameter space determined by

one-zero texture in M_ν . To see this, let us take the $B - 3L_\tau$ case as an example, which, from Eq. (14), satisfies the condition

$$(V_\nu)_{\tau 1}(V_\nu)_{\tau 1}m_1 + (V_\nu)_{\tau 2}(V_\nu)_{\tau 2}m_2 + (V_\nu)_{\tau 3}(V_\nu)_{\tau 3}m_3 = 0. \quad (15)$$

Because of the $B - 3L_\tau$ symmetry, the charged lepton of τ does not mix with the ones of e and μ , leading to

$$(V_\nu)_{\tau i} = (U_{\text{PMNS}})_{\tau i}. \quad (16)$$

By simply replacing the $(V_\nu)_{\tau i}$ in Eq. (15) with $(U_{\text{PMNS}})_{\tau i}$, the neutrino parameter space is constrained uniquely by the one-zero texture of M_ν .

III. NEUTRINO MASS AND MIXINGS

The systematic numerical analysis of one-zero-texture structures was first presented in Ref. [21]. In this section, we update the predictions with latest global-fit results [22] of neutrino oscillation parameters at 3σ confidence level (CL):

normal mass hierarchy (NH):

$$\begin{aligned} \sin^2 \theta_{12}/10^{-1} &\in [2.75, 3.50], & \sin^2 \theta_{23}/10^{-1} &\in [4.28, 6.24], & \sin^2 \theta_{13}/10^{-2} &\in [2.044, 2.437] \\ \delta m^2 \equiv m_2^2 - m_1^2 &\in [6.79, 8.01] \times 10^{-5} \text{ eV}^2, & \Delta m^2 \equiv |m_3^2 - m_1^2| &\in [2.431, 2.622] \times 10^{-3} \text{ eV}^2 \end{aligned} \quad (17)$$

inverted mass hierarchy (IH):

$$\begin{aligned} \sin^2 \theta_{12}/10^{-1} &\in [2.75, 3.50], & \sin^2 \theta_{23}/10^{-1} &\in [4.33, 6.23], & \sin^2 \theta_{13}/10^{-2} &\in [2.607, 2.461] \\ \delta m^2 \equiv m_2^2 - m_1^2 &\in [6.79, 8.01] \times 10^{-5} \text{ eV}^2, & \Delta m^2 \equiv m_3^2 - m_2^2 &\in [2.413, 2.606] \times 10^{-3} \text{ eV}^2. \end{aligned} \quad (18)$$

Solving Eq. (14) and using Eq. (16), we obtain two ratios of neutrino mass eigenvalues

$$\begin{aligned} \kappa_{12} &\equiv \frac{m_1}{m_2} = \frac{\text{Re}(U_{a3}U_{b3})\text{Im}(U_{a2}U_{b2}e^{2i\sigma}) - \text{Re}(U_{a2}U_{b2}e^{2i\sigma})\text{Im}(U_{a3}U_{b3})}{\text{Re}(U_{a1}U_{b1}e^{2i\rho})\text{Im}(U_{a3}U_{b3}) - \text{Re}(U_{a3}U_{b3})\text{Im}(U_{a1}U_{b1}e^{2i\rho})} \\ \kappa_{23} &\equiv \frac{m_2}{m_3} = \frac{\text{Re}(U_{a1}U_{b1}e^{2i\rho})\text{Im}(U_{a3}U_{b3}) - \text{Re}(U_{a3}U_{b3})\text{Im}(U_{a1}U_{b1}e^{2i\rho})}{\text{Re}(U_{a2}U_{b2}e^{2i\sigma})\text{Im}(U_{a1}U_{b1}e^{2i\rho}) - \text{Re}(U_{a1}U_{b1}e^{2i\rho})\text{Im}(U_{a2}U_{b2}e^{2i\sigma})}. \end{aligned} \quad (19)$$

The results of Eq. (19) imply that the two mass ratios (κ_{12} and κ_{13}) are fully determined in terms of six neutrino mix parameters ($\theta_{12}, \theta_{23}, \theta_{13}, \delta, \rho, \sigma$). We further define the ratio of squared mass difference:

$$R_\nu \equiv \frac{\delta m^2}{|\Delta m^2|}. \quad (20)$$

The three neutrino masses $m_{1,2,3}$ are obtained from Eq. (19):

$$m_2 = \sqrt{\frac{\delta m^2}{1 - \kappa_{12}^2}} \quad m_1 = \kappa_{12}m_2 \quad m_3 = \frac{m_2}{\kappa_{23}}. \quad (21)$$

Then, the R_ν defined in Eq. (20) can be rewritten in terms of κ_{12} and κ_{23} ,

$$R_\nu = \frac{\kappa_{23}^2(1 - \kappa_{12}^2)}{1 - \kappa_{12}^2\kappa_{23}^2} \quad (22)$$

for normal mass hierarchy and

$$R_\nu = \frac{\kappa_{23}^2(1 - \kappa_{12}^2)}{\kappa_{23}^2 - 1} \quad (23)$$

for inverted mass hierarchy.

We now perform a numerical analysis for each texture structure in Table. II. A set of random number inputs is generated for the three mixing angles ($\theta_{12}, \theta_{23}, \theta_{13}$) in their 3σ range. Instead, we generate a random input of (δ, ρ, σ) in the range of $[0, 2\pi)$. From Eqs. (22) and (23), R_ν is fully determined by the neutrino parameters of $(\theta_{12}, \theta_{23}, \theta_{13}, \delta, \rho, \sigma)$. We require the input scattering points acceptable only when R_ν falls inside the 3σ range $[\delta m_{\text{min}}^2/\Delta m_{\text{max}}^2, \delta m_{\text{max}}^2/\Delta m_{\text{min}}^2]$. From Eq. (21), the three absolute scales of neutrino masses $m_{1,2,3}$ are obtained. We should also consider the robust bound on the sum of three neutrino masses $\Sigma_\nu (\equiv \Sigma m_i) < 0.12 \text{ eV}$ set by the Planck Collaboration at 95% CL [17]. However, a latest analysis based on the physically motivated neutrino mass models yields a looser upper bound $\Sigma_\nu < 0.26 \text{ eV}$ at 95% CL [23]. In the following analysis, we check the consistency with the both of the bounds. In addition to Σ_ν , we also calculate the effective Majorana neutrino mass

$$m_{ee} = |m_1 c_{12}^2 c_{13}^2 + m_2 s_{12}^2 c_{13}^2 e^{2i\rho} + m_3 s_{13}^2 e^{2i\sigma}|, \quad (24)$$

which can be explored by neutrinoless double beta decay ($0\nu\beta\beta$) experiments. Among the current experiments in operation, the KamLAND-Zen Collaboration[24] has set

FIG. 2. Allowed samples of texture A in the case of normal mass hierarchy. The light green bands represent the 3σ uncertainty of δ . The light blue bands represent the 1σ uncertainty of δ or θ_{23} .

the most stringent limit of $m_{ee} < (0.061 - 0.165)$ eV depending on the nuclear matrix element value of the ^{136}Xe decays, while next-generation experiments aim for an improved sensitivity of m_{ee} being up to 0.01 eV.

We present the allowed range of neutrino oscillation parameters for each viable pattern in Figs. 2–6. The $U(1)_{B-3L_e}$ gauge symmetry leads to pattern A of one-zero-texture structure (see Table II). we find it phenomenologically allowed only for the case of normal mass hierarchy. Since $(M_\nu)_{ee} = 0$, the predicted effective Majorana neutrino mass m_{ee} is exactly zero for the $0\nu\beta\beta$ decay. Figure 2 shows no strong correlation between (δ, θ_{23}) and θ_{23} covers all the 3σ allowed region. The mass spectrum admits a strong normal hierarchy with $\Sigma_\nu \simeq 0.06\text{--}0.07$ eV; meanwhile, a lower bound on the lowest neutrino mass $m_1 \geq 0.0015$ eV is achieved. The $B - 3L_\mu$ gauge symmetry leads to pattern B. It is found that both

normal mass hierarchy and inverted mass hierarchy are allowed to satisfy the neutrino oscillation data, and corresponding scanning results are presented in Figs. 3 and 4, respectively. In the case of normal mass hierarchy, the predicted value of θ_{23} is restricted to be less than $\pi/4$, which, though allowed at 3σ level, is disfavored at 1σ level. Besides, the value of Σ_ν turns out to be relative large, which is excluded by the Planck 2018 limit. In the case of inverted mass hierarchy, one can observe from the left panel of Fig. 4 that the sample points of (δ, θ_{23}) fully cover the 3σ and 1σ allowed region. Furthermore, the interesting correlation of Σ_ν versus δ , shown in the middle panel of Fig. 4, indicates that the Planck bound $\Sigma_\nu < 0.12$ is satisfied only for $\delta \in [0^\circ, 110^\circ] \cup [250^\circ, 360^\circ]$, a range almost excluded at 1σ level. Both normal mass hierarchy and inverted mass hierarchy of pattern B predict the $m_{ee} = \mathcal{O}(0.01)$ eV, which is promising to be detected in the forthcoming

FIG. 3. Allowed samples of texture B in the case of normal mass hierarchy. The horizontal blue dashed line in the middle panel is the upper bound imposed by the Planck Collaboration: $\sum_\nu m_i < 0.12$ eV. The horizontal blue dashed line in the right panel is the upper limit of m_{ee} set by the KamLAND-Zen Collaboration.

FIG. 4. Allowed samples of texture B in the case of inverted mass hierarchy.

FIG. 5. Allowed samples of texture C in the case of normal mass hierarchy.

FIG. 6. Allowed samples of texture C in the case of inverted mass hierarchy.

experiments. The $B - 3L_\tau$ gauge symmetry leads to pattern C. In Figs. 5 and 6, we present the allowed region of pattern C for normal mass hierarchy and inverted mass hierarchy, respectively. Before proceeding, one notes that the $\mu - \tau$ permutation transformation can relate pattern B to pattern C by swapping the indices e and τ of the entries and the neutrino oscillation parameters between M_ν (pattern B) and $\tilde{M}_\nu$ (pattern C) are given by

$$\theta_{12} = \tilde{\theta}_{12} \quad \theta_{13} = \tilde{\theta}_{13} \quad \theta_{23} = \frac{\pi}{2} - \tilde{\theta}_{23} \quad \delta = \tilde{\delta} \pm \pi. \quad (25)$$

In the case of normal mass hierarchy, due to the $\mu - \tau$ symmetry, we arrive at the θ_{23} mixing angle greater than $\pi/4$, which is located in the 1σ range. Meanwhile, we obtain a rather large value of Σ_ν (Fig. 5), as similar to pattern B, being disfavored by Planck 2018 results. We now discuss the pattern C with inverted mass hierarchy, which is particularly interesting for us. The left panel of Fig. 6 shows no obvious correlation between δ and θ_{23} , while the correlation (δ, Σ_ν) shows a specific geometrical shape in the middle panel. In particular, part of the allowed region to satisfy the Planck 2018 bound $\Sigma_\nu < 0.12$ eV fully covers the 1σ range of the Dirac- CP phase: $\delta \in [1.40\pi, 1.68\pi]$. Therefore, for the inverted mass hierarchy of pattern C, there exist allowed regions of neutrino parameters in which all the experimental constraints are satisfied. Finally, both normal mass hierarchy and inverted mass hierarchy of pattern C predict the m_{ee} with a lower bound greater than 0.01 eV, which will be explored by the next generation experiments.

IV. OTHER PHENOMENOLOGY

In this section, we take the $U(1)_{B-3L_\mu}$ extension of the scotogenic model for some other specific phenomenology studies. The results for $U(1)_{B-3L_{e,\tau}}$ are similar except for the corresponding gauge boson Z' [25]. The $U(1)_{B-3L_\mu}$ is spontaneously broken by the S VEV $\langle S \rangle = v_S/\sqrt{2}$ at TeV scale and generates the mass for the new gauge boson Z' ,

$$M_{Z'} = 3g'v_S, \quad (26)$$

where g' is the gauge coupling of $U(1)_{B-3L_\mu}$. The Large Electron Positron Collider (LEP) bound requires $M_{Z'}/g' \gtrsim 7$ TeV [26]. In unitary gauge, the neutral component of Z_2 even scalars is parametrized as

$$\Phi^0 = \frac{v + \phi}{\sqrt{2}}, \quad S = \frac{v_S + S^0}{\sqrt{2}}. \quad (27)$$

The neutral scalars ϕ and S^0 are correlated with the mass eigenstates h and H via a rotation

$$\begin{pmatrix} h \\ H \end{pmatrix} = \begin{pmatrix} \cos \alpha & -\sin \alpha \\ \sin \alpha & \cos \alpha \end{pmatrix} \begin{pmatrix} \phi \\ S^0 \end{pmatrix}, \quad (28)$$

where the mixing angle α can be treated as a free parameter. To satisfy various theoretical and experimental constraints, $\sin \alpha \lesssim 0.2$ should be satisfied [27].

A. Lepton flavor violation

It is well known that the Yukawa interaction $\bar{L}\eta^c N_R$ could induce lepton-flavor-violation process [28]. Systematic studies on all LFV processes in scotogenic models have already been performed in Ref. [29]. Currently, the MEG experiment has set the most stringent constraint on $\mu \rightarrow e\gamma$, which requires $\text{BR}(\mu \rightarrow e\gamma) < 4.2 \times 10^{-13}$ at 90% CL [30]. The future upgrade of MEG could push the limit down to about 6×10^{-14} [31]. Meanwhile, the limits on LFV τ decays are much looser, i.e., $\text{BR}(\tau \rightarrow \mu\gamma) < 4.4 \times 10^{-8}$ and $\text{BR}(\tau \rightarrow e\gamma) < 3.3 \times 10^{-8}$ [32]. Hence, large Yukawa coupling in the τ sector is viable if one imposes hierarchical Yukawa structure as $|h_{ei}| \lesssim |h_{\mu i}| \lesssim |h_{\tau i}|$ [33]. On the other hand, if one considers universal Yukawa structure, then it is clearly expected that the tightest constraint comes from $\mu \rightarrow e\gamma$. The corresponding branching ratio is calculated as [29]

$$\text{BR}(\ell_\alpha \rightarrow \ell_\beta \gamma) = \frac{3(4\pi^3)\alpha_{\text{em}}}{4G_F^2} |A_D|^2 \text{BR}(\ell_\alpha \rightarrow \ell_\beta \nu_\alpha \bar{\nu}_\beta), \quad (29)$$

with the dipole form factor

$$A_D = \sum_{i=1}^3 \frac{h'_{\beta i} h'^*_{\alpha i}}{2(4\pi)^2 M_{\eta^+}^2} F_2\left(\frac{M_{N_i}^2}{M_{\eta^+}^2}\right), \quad (30)$$

where $h' = hV_N$ is the Yukawa coupling in the mass eigenstates, and the loop function $F_2(x)$ is given by

$$F_2(x) = \frac{1 - 6x + 3x^2 + 2x^3 - 6x^2 \log x}{6(1-x)^4}. \quad (31)$$

The three-body decay $\ell_\alpha \rightarrow \ell_\beta \bar{\ell}_\beta \ell_\beta$ is another kind of LFV process. The current limit on $\text{BR}(\mu \rightarrow 3e)$ is 1.0×10^{-12} [34] and is expected to reach 10^{-16} in the future [35]. The limits on $\text{BR}(\tau \rightarrow 3e)$ and $\text{BR}(\tau \rightarrow 3\mu)$ are 2.7×10^{-8} and 2.1×10^{-8} [36]. Considering the simplified scenario with small Yukawa coupling, e.g., $h' \lesssim 0.01$, then the three-body decay $\ell_\alpha \rightarrow \ell_\beta \bar{\ell}_\beta \ell_\beta$ is dominant by the γ -penguins. The branching ratio is given by

$$\begin{aligned} \text{BR}(\ell_\alpha \rightarrow \ell_\beta \bar{\ell}_\beta \ell_\beta) &= \frac{3(4\pi)^2 \alpha_{\text{em}}^2}{8G_F^2} \left[|A_{ND}|^2 + |A_D|^2 \left(\frac{16}{3} \log\left(\frac{m_\alpha}{m_\beta}\right) - \frac{22}{3} \right) \right. \\ &\quad \left. - (2A_{ND}A_D^* + \text{H.c.}) \right] \times \text{BR}(\ell_\alpha \rightarrow \ell_\beta \nu_\alpha \bar{\nu}_\beta), \end{aligned} \quad (32)$$

FIG. 7. Predicted results for various LFV processes as a function of M_{η^+} . Here, we assume $h'_{ai} = h'$ and $M_{N_1} \simeq M_{N_{2,3}} = 200$ GeV for illustration.

with the nondipole form factor

$$A_{ND} = \sum_{i=1}^3 \frac{h'_{\beta i} h'^*_{\alpha i}}{6(4\pi)^2 M_{\eta^+}^2} G_2\left(\frac{M_{N_i}^2}{M_{\eta^+}^2}\right) \quad (33)$$

and the loop function $G_2(x)$

$$G_2(x) = \frac{2 - 9x + 18x^2 - 11x^3 + 6x^3 \log x}{6(1-x)^4}. \quad (34)$$

Note when the photonic dipole contribution dominates a simple relation is derived [37]

$$\begin{aligned} & \text{BR}(\ell_\alpha \rightarrow \ell_\beta \bar{\ell}_\beta \ell_\beta) \\ & \simeq \frac{\alpha_{\text{em}}}{3\pi} \left(\log\left(\frac{m_\alpha^2}{m_\beta^2}\right) - \frac{11}{4} \right) \times \text{BR}(\ell_\alpha \rightarrow \ell_\beta \gamma). \end{aligned} \quad (35)$$

Therefore, comparing with $\ell_\alpha \rightarrow \ell_\beta \gamma$, $\ell_\alpha \rightarrow \ell_\beta \bar{\ell}_\beta \ell_\beta$ is suppressed. Because of quite loose limits on three-body τ decays, we will focus on $\mu \rightarrow 3e$ in following.

Meanwhile, the planned $\mu - e$ conversion in nuclei experiments will push current limits on the conversion rates, e.g., $\text{CR}(\mu - e, \text{Ti}) < 4.3 \times 10^{-12}$ [38] down to about 10^{-18} [39]. The converting rate is then expressed as [40]

$$\text{CR}(\mu - e, \text{Ti}) = 2G_F^2 |A_R D + \tilde{g}_{LV}^p V^p|^2, \quad (36)$$

with

$$A_R = -\frac{\sqrt{2}}{8G_F} e A_D, \quad \tilde{g}_{LV}^p = \frac{\sqrt{2}}{G_F} e^2 A_{ND}, \quad (37)$$

and $D = 0.0864 M_\mu^{5/2}$, $V^p = 0.0396 M_\mu^{5/2}$. Because of the relative opposite sign between above A_R and $\tilde{g}_{LV}^p$ terms, the cancellation effect is expected.

FIG. 8. Current and future prospect excluded region of various LFV processes in the $h' - M_{\eta^+}$ plane (left) and in the $M_{N_1} - M_{\eta^+}$ plane (right).

The predicted results for LFV are shown in Fig. 7. For electroweak scale N_i and η^+ , a universal Yukawa coupling $h' = 0.01$ can escape the current limit but is within the future limit. Specifically speaking, the $\mu \rightarrow e\gamma$ process could probe $M_{\eta^+} \lesssim 360$ GeV with $M_{N_1} = 200$ GeV. Although the $\mu \rightarrow 3e$ process is suppressed, the great improvement of future experiment will set a tighter bound than $\mu \rightarrow e\gamma$, i.e., excluding $M_{\eta^+} \lesssim 520$ GeV. The cancellation effect for $\mu - e$ conversion is obvious around $M_{\eta^+} \sim 260$ GeV. Together with the other LFV process, $\mu - e$ conversion could exclude $M_{\eta^+} \lesssim 700$ GeV. As for the $\tau \rightarrow e\gamma/\mu\gamma$, the Yukawa coupling as large of $h' = 0.5$ is still allowed and reachable in future experiments. It is also clear that $h' = 0.005$ is beyond all future experimental reach.

The derived current and future excluded region of various LFV processes are shown in Fig. 8, where we have assumed $h_{ai} = h'$ for illustration. The limits from $\mu \rightarrow e\gamma$ and $\mu \rightarrow 3e$ become less strict when M_{η^+} is larger. For instance, current $\mu \rightarrow e\gamma$ sets the limit $h' \lesssim 0.01$ when $M_{\eta^+} \sim 200$ GeV and $h' \lesssim 0.04$ when $M_{\eta^+} \sim 1000$ GeV, and $\mu \rightarrow 3e$ will improve the limit by a factor of about two. As for $\mu - e$ conversion, large cancellation happens when $M_{N_1} \lesssim M_{\eta^+}$, and its limit is more stringent than $\mu \rightarrow 3e$ when $M_{\eta^+} < 200$ GeV or $M_{\eta^+} > 400$ GeV. Then if we fix $h' = 0.01$, $\mu \rightarrow e\gamma$ could exclude the region $M_{N_1} \lesssim 300$ GeV and $M_{\eta^+} \lesssim 250$ GeV. For EW-scale N_1 , $\mu \rightarrow 3e$ ($\mu - e$ conversion) will probe $M_{\eta^+} \lesssim 500(700)$ GeV.

B. Dark matter

The dark matter candidate in this model is N_1 under our consideration. As for annihilation of N_1 DM, there are three kinds of interactions:

[η -portal] The same Yukawa interaction $h'_{ai}\bar{L}_\alpha\eta^c N_i$ inducing LFV is also involved in N_1 annihilation. Because of tight constraints on h' from LFV, $N_1 N_1$ pair annihilation is suppressed for small values of h' , which leads to an excess of the observed relic density. For instance, if we set $h'_{ai} = 0.01$, then the η -portal annihilation cross section

$$\langle\sigma v\rangle \simeq \sum_{\alpha\beta} \frac{|h'_{a1}h'_{\beta 1}^*|^2 r_1^2 (1 - 2r_1 + 2r_1^2)}{24\pi M_{N_1}} \langle v^2 \rangle \sim 6 \times 10^{-33} \text{ cm}^3 \text{ s}^{-1}, \quad (38)$$

where $r_1 = M_{N_1}^2 / (M_\eta^2 + M_{N_1}^2)$. Therefore, the contribution of η -portal annihilation is less than 10^{-7} in our consideration. One pathway to overcome this conflict is to consider a hierarchical Yukawa structure or coannihilation mechanism [33,41]. Besides, it is also possible to employ the freeze-in production of N_1 , which requires tiny Yukawa coupling $h'_{a1} \sim 10^{-10}$ [42].

[H -portal] The Higgs portal interactions $N_1 N_1 H(h)$ can mediate s -channel $N_1 N_1$ annihilation [43,44]. Correct relic density can be easily obtained when $M_{N_1} \sim M_H/2$ for the electroweak (EW) scale M_H [45]. And the LFV constraints are escaped by assuming small $h' \lesssim 0.01$. In the meantime, the Higgs portal interactions also induce spin-independent DM-nucleon scattering, which might be probed by DM direct detection experiments.

[Z' -portal] Although only N_{R3} is charged under the $U(1)_{B-3L_\mu}$, mixings of singlet fermions N_R also induce an effective coupling of $N_1 \gamma^\mu \gamma_5 N_1 Z'_\mu$. The Z' -portal

FIG. 9. Dark matter relic density as a function of M_{N_1} for $\sin \alpha = 0.1$ (left) and $\sin \alpha = 0.02$ (right). The green line denotes the Planck result $\Omega h^2 = 0.120 \pm 0.001$ [17].

interaction can also mediate $N_1 N_1$ annihilation [46,47]. Meanwhile, the direct LHC search for the dilepton signature $pp \rightarrow Z' \rightarrow \ell^+ \ell^-$ requires that Z' should be best above the TeV scale [48,49]. Because of the Majorana nature of N_1 , the DM-nucleon scattering cross section mediated by Z' is suppressed by $v_{\text{rel}}^2 \sim 10^{-6}$ [50].

In this paper, we consider the scalar singlet H at the EW scale and Z' at the TeV scale. Hence, the contributions to DM relic density from Z' -portal are suppressed by the heavy Z' mass. The micrOMEGAs package [51] is employed for the calculation of DM relic density and DM-nucleon scattering cross section. The numerical results for relic

density are depicted in Fig. 9. It is clear that the SM Higgs portal cannot lead to correct relic density for both $\sin \alpha = 0.1$ and $\sin \alpha = 0.02$, while resonance condition $M_{N_1} \sim M_H/2$ is always required to generate correct relic density for the H -portal.

The spin-independent DM-nucleon scattering cross section is calculated as

$$\sigma^{\text{SI}} = \frac{4}{\pi} \left(\frac{M_p M_{N_1}}{M_p + M_{N_1}} \right)^2 f_p^2, \quad (39)$$

where M_p is the proton mass and the hadronic matrix element f_p is

FIG. 10. Left: Spin-independent cross section as a function of M_{N_1} . The black solid and dashed lines correspond to current XENON1T [53] and future LZ [54] limits, respectively. Right: Exclusion limits of XENON1T and LZ in the $M_{N_1} - \sin \alpha$ plane.

FIG. 11. Left: Theoretical cross section of $pp \rightarrow \eta^+ \eta^- \rightarrow \ell^+ \ell^- + \cancel{E}_T$ signature at 13 TeV LHC. Right: Excluded region in the $M_{N_1} - M_\eta$ plane by LHC direct search. The current LFV exclusion region is obtained by assuming $h' = 0.01$.

$$\frac{f_p}{M_p} = \sum_{q=u,d,s} f_{Tq}^p \frac{\alpha_q}{M_q} + \frac{2}{27} \left(1 - \sum_{q=u,d,s} f_{Tq}^p \right) \sum_{q=c,b,t} \frac{\alpha_q}{M_q}, \quad (40)$$

with the effective vertex

$$\frac{\alpha_q}{M_q} = -\frac{y_{N_1}}{\sqrt{2}v} \sin 2\alpha \left(\frac{1}{M_h^2} - \frac{1}{M_{H_1}^2} \right). \quad (41)$$

The parameters f_{Tq}^p can be found in Ref. [52]. Figure 10 shows the results for DM-nucleon scattering. It is clear that the choice of $\sin \alpha = 0.1$ is enough to escape current XENON1T limit, but most of the range is in the reach of LUX-ZEPLIN (LZ). Since σ^{SI} is sensitive to $\sin \alpha$, we further derive the corresponding experimental limits. We find the current XENON1T limit on $\sin \alpha$ is actually less stringent than the collider limit in Ref. [27]. Meanwhile, if no signal is observed at LZ, then $\sin \alpha \lesssim 0.05$ should be satisfied.

C. Collider signature

With all new particles around the TeV scale, various promising signatures can be probed by the LHC [19]. For simplicity, we consider a degenerate mass spectrum of the inert Higgs doublet η , i.e., $M_{\eta^\pm} = M_{\eta_R} = M_{\eta_I}$. One promising signature originates from the pair production of $\eta^\pm$ at LHC, i.e., $pp \rightarrow \eta^+ \eta^- \rightarrow \ell^+ N_1 + \ell^- N_1$. Since N_1 is the DM candidate, this signature is just $\ell^+ \ell^- + \cancel{E}_T$.¹ In the left

panel of Fig. 11, theoretical production cross section at 13 TeV is shown, which is calculated by the help of MADGRAPH5_aMC@NLO [55]. Meanwhile, exclusion region by the LHC direct searches is shown in the right panel [56–58]. Two viable regions are now allowed by direct search. One is the heavy scalar region $M_\eta > 700$ GeV, and the other is the degenerate region $150 \text{ GeV} \lesssim M_{N_1} \lesssim M_\eta < 700 \text{ GeV}$. Together with the parameter space in Fig. 8, we find that the heavy scalar region might be probed by both the LHC and $\mu - e$ conversion when $M_{N_1} \lesssim 200 \text{ GeV}$. For $M_{N_1} \gtrsim 200 \text{ GeV}$ or $M_{\eta^\pm} \gtrsim 800 \text{ GeV}$, only the LHC has the ability to probe. Meanwhile, we hope that the future upgrade $\mu \rightarrow 3e$ and LHC searches with compressed mass spectra will probe the degenerate region. Note that the LFV exclusion limits depend on the choice of Yukawa coupling h' , i.e., $h' = 0.01$, but the LHC searches do not.

Another promising signature is the dilepton signature from $pp \rightarrow Z' \rightarrow \ell^+ \ell^-$. Neglecting final state mass, the corresponding partial decay widths are given by

$$\Gamma(Z' \rightarrow f \bar{f}) = \frac{M_{Z'}}{24\pi} g^2 N_C^f (Q_{fL}^2 + Q_{fR}^2), \quad (42)$$

$$\Gamma(Z' \rightarrow SS^*) = \frac{M_{Z'}}{48\pi} g^2 Q_S^2, \quad (43)$$

TABLE III. Decay branching ratio of $U(1)_{B-3L_\mu}$ gauge boson Z' , where we show the lepton flavor individually.

$q\bar{q}$	e^+e^-	$\mu^+\mu^-$	$\tau^+\tau^-$	$\nu\nu$	NN	HH
0.090	0	0.405	0	0.202	0.202	0.101

¹The LHC limits are obtained with $\ell = e, \mu$, while there would be τ leptons in our scenario. Here, we directly take such limits as a conservative estimation, because for nonuniversal Yukawa structure, the τ final states can be suppressed under certain circumstances.

where N_C^f is the number of colors of the fermion f , i.e., $N_C^{l,\nu} = 1$, $N_C^q = 3$, and Q_X is the $U(1)_{B-3L_\mu}$ charge of particle X . The branching ratios of Z' are presented in Table III, and it is clear that the dominant decay channel is $Z' \rightarrow \mu^+\mu^-$. The $B-3L_\mu$ nature of Z' can be confirmed by determining the branching ratios as

$$\text{BR}(Z' \rightarrow b\bar{b}) : \text{BR}(Z' \rightarrow e^+e^-) : \text{BR}(Z' \rightarrow \mu^+\mu^-) : \text{BR}(Z' \rightarrow \tau^+\tau^-) = \frac{1}{3} : 0 : 9 : 0. \quad (44)$$

Such an intrinsic property is useful to distinguish the $U(1)_{B-3L_\mu}$ gauge boson from other kinds of gauge bosons [59,60].

ATLAS [61,62] and CMS [63,64] have already performed searches for the Z' by the dilepton signature. Since $Z' \rightarrow \mu^+\mu^-$ is the only dilepton final state, the CMS limit on $\mu^+\mu^-$ is more suitable for this model. The limit is on the ratio

$$R_\sigma = \frac{\sigma(pp \rightarrow Z' + X \rightarrow \mu^+\mu^- + X)}{\sigma(pp \rightarrow Z + X \rightarrow \mu^+\mu^- + X)}. \quad (45)$$

$$\begin{aligned} h' &\in [10^{-3}, 10^{-1}], & \lambda &\in [10^{-11}, 10^{-1}], & M_\eta &\in [M_{N_1}, 1000] \text{ GeV}, \\ \sin \alpha &\in [0.01, 0.2], & M_H &\in [20, 1000] \text{ GeV}, & v_S &\in [5, 10] \text{ TeV}. \end{aligned} \quad (46)$$

During the scan, we require that the IH neutrino oscillation parameters interpret texture B, $m_{ee} < 0.061$ eV, and dark matter relic density satisfies $\Omega h^2 \in [0.117, 0.123]$. Then, we apply current limits from $\sum_\nu$, LFV, and LHC

Theoretical value and experimental limits are shown in left panel of Fig. 12, which indicates that $M_{Z'} \gtrsim 3.0(4.4)$ TeV should be satisfied for $g' = 0.1(0.5)$. The exclusion region in the $g' - M_{Z'}$ plane is shown in the right panel of Fig. 12. The LEP limit is now weaker than the LHC limit for $M_{Z'} \lesssim 5$ TeV.

D. Combined analysis

Based on the discussion of previous benchmark scenario, we now perform a scan to search for combined allowed parameter space by the neutrino data, lepton-flavor violation, dark matter, and collider signature. For simplicity, we consider the simplified scenario discussed in Eq. (9) further with the assumption $h_{ai} = h'$. In this way, M_ν and M_N have exactly the same structure and are related by $M_\nu = -\frac{\lambda v^2 h'^2}{32\pi M_\eta^2} M_N$. Using neutrino oscillation data, one can obtain M_ν , and then M_N directly. Note for scenario $B-3L_\mu$ only the IH is allowed. Besides IH neutrino oscillation data discussed in Sec. III, we scan the other relevant parameters in the following region:

FIG. 12. Left: Predicted cross section ratios in $U(1)_{B-3L_\mu}$ and corresponding limit from LHC. Right: Allowed parameter space in the $g' - M_{Z'}$ plane.

FIG. 13. Scanning results for allowed parameter sets. The gray points are excluded by $\sum \nu > 0.12$ eV. The red points are excluded by LFV. The blue points are excluded by the LHC. The orange points are within future reach of LFV and LZ. The green points are beyond future sensitivity.

excluded and surviving points are shown in Fig. 13. Although the $\sum \nu < 0.12$ eV limit is stringent, only the Dirac phase shown in Fig. 4 has a specific favored region, i.e., $\delta \lesssim 100^\circ$ or $\delta \gtrsim 250^\circ$. LFV processes are sensitive to Yukawa coupling h' , i.e., $h' \lesssim 0.05$ at present for $M_{\eta^+} < 1$ TeV. With such small Yukawa coupling, the DM N_1 dominantly annihilates via the scalar singlet H to SM final states. Thus, it requires $M_{N_1} \sim M_H/2$, which is

clearly shown in the third panel of Fig. 13. Note that for large enough mixing, e.g., $\sin \alpha \sim 0.2$, the SM Higgs h portal is also possible. But the h portal will be fully excluded by LZ. As for the LHC, its limits are sensitive to the mass spectrum of M_{N_1} and M_{η^+} with our assumption. And the direct detection experiment is sensitive to $\sin \alpha$. Finally, we show one benchmark point that satisfies all constraints:

$$\begin{aligned}
 \theta_{12} &= 35.59^\circ, & \theta_{23} &= 46.52^\circ, & \theta_{13} &= 8.325^\circ, & \delta &= 349.3^\circ, & \rho &= 216.7^\circ, & \sigma &= 292.1^\circ, \\
 m_1 &= 0.05067 \text{ eV}, & m_2 &= 0.05136 \text{ eV}, & m_3 &= 0.009627 \text{ eV}, & \sum m_i &= 0.1117 \text{ eV}, \\
 h' &= 0.00365, & M_{N_1} &= 239 \text{ GeV}, & M_{\eta^+} &= 927 \text{ GeV}, & M_H &= 483 \text{ GeV}, & \sin \alpha &= 0.041.
 \end{aligned} \tag{47}$$

TABLE IV. Some main results of the one texture-zeros in the $U(1)_{B-3L_\alpha}$ scotogenic model.

Patten	Group	Hierarchy	Oscillation@ 3σ	Oscillation@ 1σ	$\sum_\nu < 0.12$ eV	M_{ee} (eV)
A	$U(1)_{B-3L_e}$	NH	✓	✓	✓	0
		IH	✗	✗	✗	✗
B	$U(1)_{B-3L_\mu}$	NH	✓	✗	✗	$\gtrsim 0.05$
		IH	✓	✓	✓	$\gtrsim 0.015$
C	$U(1)_{B-3L_\tau}$	NH	✓	✓	✗	$\gtrsim 0.033$
		IH	✓	✓	✓	$\gtrsim 0.015$

V. CONCLUSION

In this paper, we consider the flavor-dependent $U(1)_{B-3L_\alpha}$ extension of the scotogenic neutrino mass model. Within this framework, three kinds of one-zero-texture structures are realized in the neutrino mass matrix. Predictions of the textures with latest neutrino oscillation parameters are performed. The main results are summarized in Table IV. Therefore, we obtain three scenarios (patterns A-NH, B-IH, and C-IH) favored by current experimental limits. These three scenarios are distinguishable by the forthcoming experiments. Since future $0\nu\beta\beta$ experiments are hopefully able to probe inverted hierarchy, pattern A-NH is favored once no positive signature is observed. For pattern B-IH, although oscillation parameters at 1σ range and $\sum_\nu < 0.12$ eV are individually satisfied, they can hardly have common parameter space. So, future precise measurements of neutrino oscillation data (especially the leptonic Dirac phase δ) and the sum of neutrino mass are able to exclude pattern B-IH. The corresponding gauge group can be confirmed by discovering a flavored gauge boson Z' .

Then, we have discussed phenomenologies such as lepton-flavor violation, dark matter, and collider signatures.

For EW scale inert particles, the Yukawa coupling should be less than 0.01 to satisfy current tight constraints from $\mu \rightarrow e\gamma$. The DM candidate N_1 dominantly annihilates via the heavy scalar singlet H with mass condition $M_H \sim 2M_{N_1}$. The observable direct detection signature is also mediated by H for $\sin \alpha \gtrsim 0.05$. The promising signature of inert particles at the LHC is $pp \rightarrow \eta^+\eta^- \rightarrow \ell^+\ell^- + \cancel{E}_T$ with viable parameter space $M_{N_1} \lesssim M_\eta < 700$ GeV or $M_\eta > 700$ GeV. The $B - 3L_\alpha$ nature of gauge boson Z' can be confirmed by $\text{BR}(Z' \rightarrow b\bar{b}) : \text{BR}(Z' \rightarrow \ell_\alpha \bar{\ell}_\alpha) = \frac{1}{3} : 9$.

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- [1] L. M. Krauss, S. Nasri, and M. Trodden, *Phys. Rev. D* **67**, 085002 (2003).
 - [2] E. Ma, *Phys. Rev. D* **73**, 077301 (2006).
 - [3] Y. Cai, J. Herrero-García, M. A. Schmidt, A. Vicente, and R. R. Volkas, *Front. Phys.* **5**, 63 (2017).
 - [4] E. Ma, *Phys. Lett. B* **433**, 74 (1998).
 - [5] E. Ma, D. P. Roy, and S. Roy, *Phys. Lett. B* **525**, 101 (2002).
 - [6] S. Choubey and W. Rodejohann, *Eur. Phys. J. C* **40**, 259 (2005).
 - [7] W. Altmannshofer, S. Gori, M. Pospelov, and I. Yavin, *Phys. Rev. D* **89**, 095033 (2014).
 - [8] J. Heeck, M. Holthausen, W. Rodejohann, and Y. Shimizu, *Nucl. Phys. B* **896**, 281 (2015).
 - [9] J. Heeck and W. Rodejohann, *Phys. Rev. D* **85**, 113017 (2012).
 - [10] T. Araki, J. Heeck, and J. Kubo, *J. High Energy Phys.* **07** (2012) 083.
 - [11] P. H. Frampton, S. L. Glashow, and D. Marfatia, *Phys. Lett. B* **536**, 79 (2002).
 - [12] H. Fritzsch, Z. z. Xing, and S. Zhou, *J. High Energy Phys.* **09** (2011) 083.
 - [13] L. Lavoura, *Phys. Lett. B* **609**, 317 (2005).
 - [14] L. M. Cebola, D. Emmanuel-Costa, and R. González Felipe, *Phys. Rev. D* **88**, 116008 (2013).
 - [15] S. Baek, H. Okada, and K. Yagyu, *J. High Energy Phys.* **04** (2015) 049.
 - [16] Z. L. Han, R. Ding, S. J. Lin, and B. Zhu, *Eur. Phys. J. C* **79**, 1007 (2019).
 - [17] N. Aghanim *et al.* (Planck Collaboration), *arXiv:1807.06209*.
 - [18] T. Nomura and H. Okada, *arXiv:1806.09957*.

- [19] Z. L. Han and W. Wang, *Eur. Phys. J. C* **79**, 522 (2019).
- [20] B. Pontecorvo, *Zh. Eksp. Teor. Fiz.* **33**, 549 (1957); Z. Maki, M. Nakagawa, and N. Sakata, *Prog. Theor. Phys.* **28**, 870 (1962).
- [21] E. I. Lashin and N. Chamoun, *Phys. Rev. D* **85**, 113011 (2012).
- [22] I. Esteban, M. C. Gonzalez-Garcia, A. Hernandez-Cabezudo, M. Maltoni, and T. Schwetz, *J. High Energy Phys.* **01** (2019) 106.
- [23] A. Loureiro *et al.*, *Phys. Rev. Lett.* **123**, 081301 (2019).
- [24] A. Gando *et al.* (KamLAND-Zen Collaboration), *Phys. Rev. Lett.* **117**, 082503 (2016); **117**, 109903(A) (2016).
- [25] T. Han, J. Liao, H. Liu, and D. Marfatia, *J. High Energy Phys.* **11** (2019) 028.
- [26] G. Cacciapaglia, C. Csaki, G. Marandella, and A. Strumia, *Phys. Rev. D* **74**, 033011 (2006).
- [27] T. Robens and T. Stefaniak, *Eur. Phys. J. C* **76**, 268 (2016).
- [28] J. Kubo, E. Ma, and D. Suematsu, *Phys. Lett. B* **642**, 18 (2006).
- [29] T. Toma and A. Vicente, *J. High Energy Phys.* **01** (2014) 160.
- [30] A. M. Baldini *et al.* (MEG Collaboration), *Eur. Phys. J. C* **76**, 434 (2016).
- [31] A. M. Baldini *et al.*, [arXiv:1301.7225](https://arxiv.org/abs/1301.7225).
- [32] B. Aubert *et al.* (BABAR Collaboration), *Phys. Rev. Lett.* **104**, 021802 (2010).
- [33] A. Vicente and C. E. Yaguna, *J. High Energy Phys.* **02** (2015) 144.
- [34] U. Bellgardt *et al.* (SINDRUM Collaboration), *Nucl. Phys.* **B299**, 1 (1988).
- [35] A. Blondel *et al.*, [arXiv:1301.6113](https://arxiv.org/abs/1301.6113).
- [36] K. Hayasaka *et al.*, *Phys. Lett. B* **687**, 139 (2010).
- [37] E. Arganda and M. J. Herrero, *Phys. Rev. D* **73**, 055003 (2006).
- [38] C. Dohmen *et al.* (SINDRUM II Collaboration), *Phys. Lett. B* **317**, 631 (1993).
- [39] PRIME Working Group, Search for the $\mu \rightarrow e$ Conversion Process at an Ultimate Sensitivity of the Order of 10^{-18} with PRISM.
- [40] R. Kitano, M. Koike, and Y. Okada, *Phys. Rev. D* **66**, 096002 (2002); **76**, 059902(E) (2007).
- [41] D. Suematsu, T. Toma, and T. Yoshida, *Phys. Rev. D* **79**, 093004 (2009).
- [42] E. Molinaro, C. E. Yaguna, and O. Zapata, *J. Cosmol. Astropart. Phys.* **07** (2014) 015.
- [43] N. Okada and O. Seto, *Phys. Rev. D* **82**, 023507 (2010).
- [44] S. Kanemura, O. Seto, and T. Shimomura, *Phys. Rev. D* **84**, 016004 (2011).
- [45] R. Ding, Z. L. Han, L. Huang, and Y. Liao, *Chin. Phys. C* **42**, 103101 (2018).
- [46] N. Okada and S. Okada, *Phys. Rev. D* **93**, 075003 (2016); S. Okada, *Adv. High Energy Phys.* **2018**, 1 (2018).
- [47] N. Okada and S. Okada, *Phys. Rev. D* **95**, 035025 (2017).
- [48] Z. L. Han, W. Wang, and R. Ding, *Eur. Phys. J. C* **78**, 216 (2018).
- [49] Z. L. Han and W. Wang, *Eur. Phys. J. C* **78**, 839 (2018).
- [50] S. Baek, *Phys. Lett. B* **756**, 1 (2016).
- [51] G. Bélanger, F. Boudjema, A. Goudelis, A. Pukhov, and B. Zaldivar, *Comput. Phys. Commun.* **231**, 173 (2018); G. Bélanger, F. Boudjema, A. Pukhov, and A. Semenov, *Comput. Phys. Commun.* **192**, 322 (2015).
- [52] J. R. Ellis, A. Ferstl, and K. A. Olive, *Phys. Lett. B* **481**, 304 (2000).
- [53] E. Aprile *et al.* (XENON Collaboration), *Phys. Rev. Lett.* **121**, 111302 (2018).
- [54] D. S. Akerib *et al.* (LUX-ZEPLIN Collaboration), *Phys. Rev. D* **101**, 052002 (2020).
- [55] J. Alwall, R. Frederix, S. Frixione, V. Hirschi, F. Maltoni, O. Mattelaer, H.-S. Shao, T. Stelzer, P. Torrielli, and M. Zaro, *J. High Energy Phys.* **07** (2014) 079.
- [56] M. Aaboud *et al.* (ATLAS Collaboration), *Eur. Phys. J. C* **78**, 995 (2018).
- [57] G. Aad *et al.* (ATLAS Collaboration), *Eur. Phys. J. C* **80**, 123 (2020).
- [58] G. Aad *et al.* (ATLAS Collaboration), *Phys. Rev. D* **101**, 052005 (2020).
- [59] L. Basso, A. Belyaev, S. Moretti, and C. H. Shepherd-Themistocleous, *Phys. Rev. D* **80**, 055030 (2009).
- [60] E. J. Chun, A. Das, J. Kim, and J. Kim, *J. High Energy Phys.* **02** (2019) 093.
- [61] M. Aaboud *et al.* (ATLAS Collaboration), *J. High Energy Phys.* **10** (2017) 182.
- [62] G. Aad *et al.* (ATLAS Collaboration), *Phys. Lett. B* **796**, 68 (2019).
- [63] A. M. Sirunyan *et al.* (CMS Collaboration), *J. High Energy Phys.* **06** (2018) 120.
- [64] CMS Collaboration, CERN Report No. CMS-PAS-EXO-19-019, 2019.