

Hidden from view: Neutrino masses, dark matter, and TeV-scale leptogenesis in a neutrinophilic two-Higgs-doublet model

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We consider a simple extension of the Standard Model providing dark matter and a TeV-scale seesaw mechanism that also allows for viable leptogenesis. In addition to the Standard Model degrees of freedom, the model contains a neutrinophilic Higgs doublet, a scalar singlet, and six singlet fermions (including three right-handed Majorana neutrinos) that are charged under a local $U(1)'$ gauge symmetry. We show how the $U(1)'$ charge assignments and the choice of scalar potential can lead to a TeV-scale seesaw mechanism and $\mathcal{O}(1)$ neutrino Yukawa couplings in a straightforward way. While this scenario has all the ingredients one would expect for significant experimental signatures, including several new TeV-scale degrees of freedom, we find that most distinctive features associated with neutrino mass generation, leptogenesis and the dark sector are likely to remain inaccessible in the absence of additional lepton flavor symmetries.

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I. INTRODUCTION

Although the Standard Model (SM) is in spectacular agreement with the results of most terrestrial experiments, it is certainly incomplete. Apart from theoretical considerations, such as the hierarchy problem, the origin of electric charge quantization, and the “near miss” for gauge unification in the SM, several observations point to the need for physics beyond the Standard Model (BSM). The observation of neutrino oscillations has revealed that neutrinos have nonzero masses and that lepton flavors are mixed [1]. In addition, precise cosmological observations have confirmed the existence of nonbaryonic cold dark matter: $\Omega_D h^2 = 0.1123 \pm 0.0035$ [2]. Together with the cosmic baryon asymmetry, these important discoveries cannot be accommodated in the minimal SM without introducing extra ingredients.

Perhaps the most attractive approach towards understanding the origin of small neutrino masses is using the dimension-5 Weinberg operator [3],

$$\frac{1}{4} \kappa_{gf} \overline{\ell}_{Lc}^g \epsilon_{cd} \phi_d \ell_{Lb}^f \epsilon_{ba} \phi_a + \text{H.c.}, \quad (1)$$

which comes from integrating out some new superheavy particles. A simple way to obtain the operator in Eq. (1) is through the type-I seesaw mechanism [4], in which three right-handed neutrinos N_R having large Majorana masses but no SM charges are introduced. Through Yukawa interactions of the N_R with the SM leptons, the three active neutrinos then acquire tiny Majorana masses as given by

the type-I seesaw formula: i.e., the mass matrix of light neutrinos is given by

$$M_\nu = -M_D M_R^{-1} M_D^T \quad (2)$$

where M_ν is the mass matrix for the light neutrinos, M_D is the Dirac mass matrix linking the left-handed light neutrinos to the N_R , and M_R is the mass matrix for the N_R . Variations of this idea, including the type-II [5] and type-III [6] seesaw mechanisms have been widely discussed in the literature.

Seesaw mechanisms are among the most natural ways to generate tiny neutrino masses, and they can be embedded into more fundamental frameworks such as grand unified theories or string theory. A salient feature of the seesaw mechanism is that leptogenesis [7] can work well to account for the matter-antimatter asymmetry of the Universe. A lepton asymmetry is dynamically generated by the CP -violating and out-of-equilibrium decays of right-handed neutrinos and then converted into a baryon asymmetry via $(B + L)$ -violating sphaleron interactions [8] that exist in the SM. However, seesaw mechanisms typically lose direct testability on the experimental side. A direct test of the seesaw mechanism would involve the detection of these heavy seesaw particles at a collider or in other neutrino experiments, as well as the measurement of their Yukawa couplings with the electroweak doublets. In the canonical seesaw mechanism, heavy seesaw particles turn out to be too heavy, i.e., $\sim 10^{14} - 10^{16}$ GeV, to be experimentally accessible.

In view of the prospects for probing new TeV-scale physics, it is interesting to ask how one might lower the conventional seesaw scale to an experimentally accessible one. Indeed, there exist many approaches to lowering the

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seesaw scale [9–17], but these may seem unnatural. In this paper we propose a new TeV-scale seesaw mechanism. A guiding principle in constructing this framework is to invoke as few new degrees of freedom as possible within the context of the seesaw mechanism while addressing the cosmic matter content and yielding some degree of testability. In doing so, we seek to explore features that might be generically present in UV-complete theories.

To this end, we work in the framework of a two-Higgs-doublet model (2HDM) with a $U(1)'$ gauge symmetry. The second Higgs doublet H_n rather naturally gets a vacuum expectation value (VEV) through its interaction with the SM Higgs boson in the presence of an additional SM singlet scalar Φ . We first show that its VEV can be relatively small. Through an appropriate choice of $U(1)'$ charges for the SM fields as well as the H_n , Φ , and N_R , we obtain neutrino masses through the standard seesaw mechanism but allow $\mathcal{O}(1)$ Yukawa couplings between the light sector and TeV-scale N_R . We show that properties of the massive Z' associated with the spontaneously broken $U(1)'$ are compatible with present phenomenological constraints, while the new SM gauge singlets required for anomaly cancellation may provide a viable cold dark matter candidate. Constraints on the model from lepton flavor violation decays are studied. The model may also generate the matter-antimatter asymmetry of the Universe through flavor-dependent leptogenesis.

In principle, this scenario contains several ingredients that should lead to experimental signatures: new TeV-scale degrees of freedom, including the N_R , Z' , and dark matter fields; $\mathcal{O}(1)$ Yukawa couplings in the neutrino sector, implying the possibility of significant charged lepton flavor violation; and several new scalar degrees of freedom that can mix with the SM Higgs. We find, however, that in the absence of additional lepton flavor symmetries, most of the dynamics associated with neutrino mass generation, leptogenesis, and production of the dark matter relic density will be generally difficult to test experimentally. The magnitude of the neutrino Yukawa couplings directly relevant to the leptogenesis lepton asymmetry must be smaller than necessary for observable signatures in the next generation of charged lepton flavor-violation searches, assuming an anarchical Yukawa matrix structure. The coupling of the Z' boson to SM particles is suppressed by a tiny $Z - Z'$ mixing angle. The mixing of the new scalar singlet with the SM Higgs, which governs the dark matter annihilation cross section, must be sufficiently small to ensure a dark matter relic density consistent with observation, making its impact on Higgs phenomenology at the Large Hadron Collider unlikely to be observable. In the absence of additional flavor symmetries in the lepton sector, the only possible experimental signature of this scenario would be a spin-independent signal in dark matter direct detection searches. Even this signal would not in itself distinguish this scenario from others that lead to a

similar signature. In short, it is possible that new TeV-scale physics may be responsible for three of the strongest motivations for extending the SM yet may remain hidden from view. Moreover, unlike many other TeV-scale BSM scenarios that require additional symmetries¹ in order to suppress experimental signatures, the present scenario would require the adoption of new symmetries or structure in order to allow for effects at the observable level.

Our discussion of the model and these features is arranged as follows: In Sec. II we give a brief introduction to the model. We study neutrino masses, lepton flavor violation and dark matter phenomenology in Secs. III and IV. Section V is devoted to flavor-dependent leptogenesis. We summarize in Sec. VI.

II. THE MODEL

Our purpose is to explain the tiny but nonzero neutrino masses, dark matter and the cosmic baryon asymmetry via TeV-scale new physics. To do so in a minimal manner, we extend the SM with three right-handed Majorana neutrinos N_R , one extra Higgs doublet H_n , and one scalar singlet Φ . From the standpoint of the seesaw mechanism, sufficiently small Dirac masses M_D for TeV-scale M_R and $\mathcal{O}(1)$ Yukawa couplings can arise if the VEV of H_n is small and solely responsible for a nonvanishing M_D . Ensuring that the SM Higgs doublet H does not contribute to M_D in this case requires imposition of an additional symmetry. Such a symmetry can be a global $U(1)'$ symmetry that is broken both spontaneously and explicitly (softly) [20–23], a discrete flavor symmetry [24–28], or a local $U(1)'$ symmetry. In Table I we classify variations of the possible $U(1)'$ extensions, including the relevant particle content and $U(1)'$ charges. We distinguish two possibilities for the global $U(1)'$, corresponding to whether or not the additional doublet is charged under this group. In the case of the local $U(1)'$ we denote the additional SM gauge singlet fermions required for anomaly cancellation as Ψ .

As we discuss below, the local $U(1)'$ and type-II global $U(1)'$ also admit a dark matter candidate. In the latter instance, the pseudo-Goldstone boson associated with spontaneous symmetry breaking is the dark matter particle, while explicit, soft breaking is needed to provide a dark matter mass. A simpler version that includes only the SM doublet H and singlet Φ and that does not consider neutrino masses has been studied in Ref. [20]. In the present context, the generation of the neutrino Dirac mass in the type-II global $U(1)'$ entails an additional, explicit breaking of the symmetry, as one may see from the set of possible Yukawa interactions:

$$-\mathcal{L}_Y = \bar{\ell}_L Y_{eij} H E_R + \bar{\ell}_L^i Y_{vij} \tilde{H} N_R^j + \bar{\ell}_L^i \lambda_{vij} \tilde{H}_n N_R^j + \text{H.c.} \quad (3)$$

¹E.g., the assumption of minimal lepton flavor violation [18], Peccei-Quinn [19] symmetry, etc.

TABLE I. Quantum numbers of fields under local, type-I global, and type-II global $U(1)'$ symmetry. SB denotes an explicit symmetry-breaking term, and DM denotes a possible dark matter candidate in this scenario.

	Local $U(1)'$	Type-I global $U(1)'$	Type-II global $U(1)'$
H	0	0	0
H_n	1	1	0
Φ	1	1	1
N_R	1	1	1
Ψ	1	\times	\times
Z_2	\checkmark	\times	\times
SB	\times	\checkmark	\checkmark
DM	Ψ	\times	δ

Given the $U(1)'$ charges of the N_R and H , the second term explicitly breaks the symmetry, as does the third term if H_n is uncharged. For the type-II global $U(1)'$, for which one must include an explicit symmetry-breaking term in the potential to generate the dark matter mass, it would be unreasonable not to include both of the symmetry-breaking terms in Eq. (3) as well. In this case, one would have to explain the disparity between a tiny coupling Y_ν associated with the SM Higgs vacuum expectation value (VEV) and an $\mathcal{O}(1)$ λ in the presence of the tiny VEV of the neutral component of H_n . Consequently, we do not consider this scenario further.

In contrast, the type-I global $U(1)'$ allows one to forbid the second term in Eq. (3), while the third term needed for the TeV-scale seesaw mechanism is allowed. However, the most general scalar potential also includes operators that are invariant under the new $U(1)'$ but also provide a mechanism for decay of the pseudo-Goldstone boson, thereby precluding its viability as a dark matter particle. Consequently, we focus our attention on the local $U(1)'$ scenario, for which both a viable dark matter particle and the TeV-scale seesaw mechanism appear possible without the introduction of explicit symmetry-breaking terms.

For the local $U(1)'$ case, the neutral component H_n^0 gets a small vacuum expectation value (VEV) through its coupling with the SM Higgs and is solely responsible for the origin of the Dirac neutrino mass matrix. Three new SM gauge singlets ψ_L are introduced to cancel the anomalies. The most general scalar potential is then given by

$$\begin{aligned}
V = & -m_1^2 H^\dagger H + m_2^2 H_n^\dagger H_n - m_0^2 \Phi^\dagger \Phi + \lambda_0 (\Phi^\dagger \Phi)^2 \\
& + \lambda_1 (H^\dagger H)^2 + \lambda_2 (H_n^\dagger H_n)^2 + \lambda_3 (H^\dagger H)(H_n^\dagger H_n) \\
& + \lambda_4 (H^\dagger H_n)(H_n^\dagger H) + \lambda_5 (\Phi^\dagger \Phi)(H^\dagger H) \\
& + \lambda_6 (\Phi^\dagger \Phi)(H_n^\dagger H_n) + [\lambda_n \Phi(H_n^\dagger H) + \text{H.c.}]. \quad (4)
\end{aligned}$$

Here we assume λ_n is a real parameter with a positive mass dimension.² In the case of the type-II global $U(1)'$ scenario,

²Note that this operator is multiplicatively renormalized logarithmically and is, therefore, technically natural.

the corresponding operator breaks a Z_2 symmetry that would otherwise guarantee stability of the possible dark matter candidate. It should also be noted that the interaction term, $(H^\dagger H_n)(H_n^\dagger H) + \text{H.c.}$, is forbidden by the $U(1)'$ symmetry. In what follows, we define $H = (h^+, (h + iA + v_1)/\sqrt{2})^T$, $H_n = (\rho^+, (\eta + i\delta + v_2)/\sqrt{2})^T$ and $\Phi = (\phi + i\varphi + v_0)/\sqrt{2}$.

After imposing the conditions for the VEVs v_j to yield an extremum, one has

$$\begin{aligned}
v_0^2 & \approx \frac{4m_0^2 \lambda_1 - 2m_1^2 \lambda_5}{4\lambda_0 \lambda_1 - \lambda_5^2}, \\
v_1^2 & \approx \frac{4m_1^2 \lambda_0 - 2m_0^2 \lambda_5}{4\lambda_1 \lambda_0 - \lambda_5^2}, \\
v_2 & \approx -\frac{\lambda_n v_0 v_1}{\sqrt{2} m_2^2}. \quad (5)
\end{aligned}$$

It is clear that v_2 is suppressed by m_2^2 . By setting m_2^2 to be sufficiently large or λ_n to be tiny, v_2 can be of the order of several MeV—the salient feature of our model. As a typical example, we obtain $v_2 \sim 2$ MeV by setting $\lambda_n \sim 1$ MeV, $m_2 \sim 150$ GeV and $v_0 \sim 250$ GeV. We also note that a small nonvanishing λ_n can be viewed as a slight breaking of a Z_2 symmetry,³ under which $\Phi \leftrightarrow -\Phi$, $H \leftrightarrow H$, $H_n \leftrightarrow H_n$; hence, the smallness of the λ_n is natural, in the 't Hooft sense [29], even through there is no dynamical understanding of this smallness.

In the basis (h, η, ϕ) we then derive the mass-squared matrix for the CP -even scalars

$$M_{CP\text{even}}^2 = \begin{pmatrix} m_{hh}^2 & v_1 v_2 (\lambda_3 + \lambda_4) + \frac{\lambda_n v_0}{\sqrt{2}} & v_1 v_0 \lambda_5 + \frac{\lambda_n v_0}{\sqrt{2}} \\ \clubsuit & m_{\eta\eta}^2 & v_2 v_0 \lambda_6 + \frac{\lambda_n v_1}{\sqrt{2}} \\ \clubsuit & \clubsuit & m_{\phi\phi}^2 \end{pmatrix} \quad (6)$$

where

$$\begin{aligned}
m_{hh}^2 & = 2v_1^2 \lambda_1 - \frac{\lambda_n v_0 v_2}{\sqrt{2} v_1}, \\
m_{\eta\eta}^2 & = 2v_2^2 \lambda_2 - \frac{\lambda_n v_0 v_1}{\sqrt{2} v_2}, \\
m_{\phi\phi}^2 & = 2v_0^2 \lambda_0 - \frac{\lambda_n v_1 v_2}{\sqrt{2} v_0}, \quad (7)
\end{aligned}$$

and where the “ \clubsuit ” symbol indicates that an entry $(M^2)_{jk}$ is equal to $(M^2)_{kj}$. We also derive the mass matrix for the CP -odd scalars in the basis (A, δ, ϕ) :

³We have estimated the lifetimes of ρ , η , and δ , which give $\tau_\rho \ll \tau_\eta \sim 10^{-11}$ s and $\tau_\delta \sim 10^{-9}$ s by assuming $m_2 \sim 150$ MeV. In short, there exist no long-lived scalars charged under Z_2 that would potentially interfere with big bang nucleosynthesis.

$$M_{CP\text{ odd}}^2 = \begin{pmatrix} m_{AA}^2 & -\frac{v_0\lambda_n}{\sqrt{2}} & -\frac{v_2\lambda_n}{\sqrt{2}} \\ \clubsuit & m_{\delta\delta}^2 & +\frac{v_2\lambda_n}{\sqrt{2}} \\ \clubsuit & \clubsuit & m_{\varphi\varphi}^2 \end{pmatrix}, \quad (8)$$

where

$$m_{AA}^2 = \frac{\lambda_n v_0 v_2}{\sqrt{2} v_1}, \quad m_{\delta\delta}^2 = \frac{\lambda_n v_0 v_1}{\sqrt{2} v_2}, \quad m_{\varphi\varphi}^2 = \frac{\lambda_n v_1 v_2}{\sqrt{2} v_0}. \quad (9)$$

It is straightforward to see that there is only one massive CP -odd scalar with a squared-mass eigenvalue: $-\lambda_n(v_0^2 v_1^2 + v_1^2 v_2^2 + v_0^2 v_2^2)(v_0 v_1 v_2)^{-1}/\sqrt{2}$. The other two are would-be Goldstone bosons eaten by the gauge fields Z and Z' , respectively.

Before proceeding, we make additional comparisons with the global $U(1)'$ scenarios. As indicated earlier, the spontaneous breaking of a global $U(1)'$ symmetry would lead to a massless Goldstone boson, a possibility that is severely constrained by big bang nucleosynthesis [30] and observations of Bullet Cluster galaxies [31]. Hence, such global $U(1)'$ symmetry must be explicitly broken, which may be accomplished by adding the following terms to the Higgs potential [20]:

$$V_\Delta = \Delta_1 \Phi^2 + \Delta_2 \bar{\Phi} + \text{H.c.} \quad (10)$$

Note that the operators in V_Δ close under renormalization. Though Δ_1 and Δ_2 are in general complex parameters, we set them to be real for simplicity. Adding Eq. (10) to Eq. (4) and imposing the minimization conditions, one has

$$\begin{aligned} M_{CP\text{ even}}'^2 &= M_{CP\text{ even}}^2 + \Delta M^2, \\ M_{CP\text{ odd}}'^2 &= M_{CP\text{ odd}}^2 - \Delta M^2, \end{aligned} \quad (11)$$

where

$$\Delta M^2 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 2\Delta_1 \end{pmatrix} \quad (12)$$

The VEVs of Higgs fields in this case are quite similar to those in Eq. (5), only up to the replacement $m_0^2 \rightarrow m_0^2 + \Delta_1$.

Compared with the local $U(1)'$ case, one does not require the extra fermion singlets to cancel anomalies. However, one does need to add explicit $U(1)'$ symmetry-breaking terms to the Higgs potential by hand. As discussed above, there then exists no reason not to include the explicit $U(1)'$ -breaking terms in the \mathcal{L}_Y as well. For the local $U(1)'$ scenario, the extra fermion singlet can be a dark matter candidate, as we discuss in Sec. IV. For the type-I global $U(1)'$ scenario, there is no dark matter candidate. For the type-II global $U(1)'$ scenario, the imaginary part of the

scalar singlet can be a dark matter candidate [20], but the $U(1)'$ symmetry in this case is only responsible for the stability of ϕ_a , and it has nothing to do with the VEV of the new Higgs doublet.

As with any local $U(1)'$ scenario, one must consider $Z - Z'$ mixing and constraints from electroweak precision observables. In the present instance, the kinetic term $(D_\mu H_n)^\dagger (D^\mu H_n)$ contributes to the $Z - Z'$ mixing at both the tree level and the one-loop level. Effects of the latter are characterized by the mixing tensor $\Pi_{ZZ'}^T$ given in the Appendix. Taking both contributions into account, the mass eigenvalues of Z and Z' are then

$$\begin{aligned} M_Z^2 &= \frac{1}{4} c^2 (g^2 + g'^2) (v_1^2 + v_2^2) + s^2 g'^2 (v_0^2 + v_2^2) \\ &\quad - cs (g'' \sqrt{g^2 + g'^2} v_2^2 + 2\hat{\Pi}_{ZZ'}^T), \end{aligned} \quad (13)$$

$$\begin{aligned} M_{Z'}^2 &= \frac{1}{4} s^2 (g^2 + g'^2) (v_1^2 + v_2^2) + c^2 g'^2 (v_0^2 + v_2^2) \\ &\quad + cs (g'' \sqrt{g^2 + g'^2} v_2^2 + 2\hat{\Pi}_{ZZ'}^T), \end{aligned} \quad (14)$$

where $c = \cos \theta$ and $s = \sin \theta$, with

$$\tan 2\theta = \frac{4v_2^2 g'' \sqrt{g^2 + g'^2} + 8\hat{\Pi}_{ZZ'}^T}{4g'^2 (v_0^2 + v_2^2) - (g^2 + g'^2) (v_1^2 + v_2^2)}. \quad (15)$$

The mixing matrix that appears in the change of the basis form $(A_{3\mu}, B_\mu, B'_\mu)$ to (Z_μ, A_μ, Z'_μ) is given by

$$\begin{pmatrix} Z_\mu \\ A_\mu \\ Z'_\mu \end{pmatrix} = \begin{pmatrix} c & 0 & -s \\ 0 & 1 & 0 \\ s & 0 & c \end{pmatrix} \begin{pmatrix} c_w & -s_w & 0 \\ s_w & c_w & 0 \\ 0 & 0 & 1 \end{pmatrix} \begin{pmatrix} A_{3\mu} \\ B_\mu \\ B'_\mu \end{pmatrix}, \quad (16)$$

where $(c_w, s_w) \equiv (\cos \theta_w, \sin \theta_w)$, with θ_w being the Standard Model weak mixing angle. Phenomenological constraints typically require the $Z - Z'$ mixing angle θ to be less than $\sim 1 - 2 \times 10^{-3}$ [32] and the mass of the extra neutral gauge boson to be heavier than $865 \sim 910$ GeV [33]. A suitable mass hierarchy and mixing between Z and Z' are maintained by setting v_0 to be relatively large (i.e., $v_0 > 1$ TeV), with $v_2 \sim 1$ MeV and $m_{\eta,\delta} \sim 150$ GeV. However, it should be mentioned that Z' in our model only couples to the BSM fields such that the constraints given above could be relaxed. The masses of W^\pm are $1/2g\sqrt{v^2 + v_2^2}$. It is straightforward to determine that for this choice of parameters, the constraint on the model from the ρ parameter, $\rho \equiv \frac{M_W}{M_Z \cos \theta_w} \approx 1 \pm 0.001$, is maintained as the constraint on $Z - Z'$ mixing is fulfilled.

We further study the constraints implied by the second Higgs-doublet loop contributions to the electroweak precision observables (EWPO), which can be characterized

by the leading effects of these corrections in terms of the oblique parameters [34]. The most important constraint comes from the T parameter. The resulting contribution from the second scalar doublet is

$$T \approx \frac{1}{2\pi\hat{s}^2} \frac{(m_\rho^2 - m_\eta^2)^2}{M_W^2(m_\rho^2 + m_\eta^2)}, \quad (17)$$

where \hat{s} is the sine function of the weak mixing angle in the $\bar{M}S$ scheme. The latest global fit to the EWPO yields $T = -0.08 \pm 0.11(+0.09)$ [1]. A straightforward calculation gives us a 2σ upper bound of the mass splitting between the neutral and charged components of H_n , about 12 GeV, by assuming $m_\eta \sim 150$ GeV. Since v_2 and λ_n are small, the only term that can generate viable squared-mass splitting between ρ and η is $\lambda_4(H^\dagger H_n)(H_n^\dagger H)$, as was shown in Eq. (4). Assuming $m_\eta \sim 150$ GeV, the T parameter puts a constraint on $|\lambda_4|$, which should be roughly less than $\mathcal{O}(0.1)$.⁴

III. NEUTRINO MASSES AND LEPTON FLAVOR VIOLATIONS

Tiny Majorana neutrino masses arise naturally in this framework, as one may observe from Eqs. (3) and (5). We now set $Y_\nu = 0$ as required by the $U(1)'$ symmetry and charge assignments. The entries in the Dirac neutrino mass matrix are proportional to $\lambda_{ij}v_2$, such that they may be at the MeV scale while keeping $\lambda \sim \mathcal{O}(1)$. In this case we only need TeV-scale right-handed Majorana neutrinos to suppress the active neutrino masses to the eV scale.

The mass term for N_R requires introduction of the dimension-5 operator

$$\frac{1}{\Lambda} \Phi^2 \bar{N}_R N_R^C + \text{H.c.}, \quad (18)$$

where the scale Λ is presumably associated with integrating out additional heavy fields. For example, one may make the theory renormalizable by introducing three extra fields S_L that are gauge singlets and uncharged under the $U(1)'$ symmetry. Their mass term, as well as Yukawa interactions with the right-handed neutrinos, can be written as

$$\bar{S}_L Y_N \Phi^\dagger N_R + \bar{S}_L^C M_S S_L + \text{H.c.} \quad (19)$$

Integrating out heavy fields S_L gives $1/\Lambda \sim Y_N^T Y_N / M_S$, yielding the right-handed neutrino mass matrix written as

$$M_R = -v_0^2 Y_N M_S^{-1} Y_N^T. \quad (20)$$

Tiny, nonvanishing Majorana masses of active neutrinos can be obtained by integrating out heavy Majorana neutrinos:

⁴The authors thank the referee for pointing out possible constraints of EWPO on the parameter space of our model.

$$M_\nu = v_2^2 \lambda M_R^{-1} \lambda^T = v_2^2 v_0^{-2} \lambda (Y_N^T)^{-1} M_S Y_N^{-1} \lambda^T. \quad (21)$$

Setting $\mathcal{O}(v_2) \sim 1$ MeV, $\mathcal{O}(v_0) \sim 1$ TeV, and $\mathcal{O}(M_S) \sim 100$ GeV, eV-scale neutrino masses only require λ and Y_N to be $\mathcal{O}(1)$.

The structure of the Yukawa matrix λ is, of course, constrained by the results of neutrino oscillation studies. Below we provide an illustrative example that is consistent with these constraints. In addition, the Yukawa interactions in Eq. (3) lead to nonconservation of charged lepton flavor at the one-loop level. Given the TeV-scale masses for the N_R and $\mathcal{O}(1)$ Yukawa couplings λ , one might anticipate observable charged lepton flavor violation (CLFV). At present, the most stringent constraint arises from the nonobservation of the decay $\mu \rightarrow e + \gamma$. The associated branching ratio is given by

$$BR(\mu \rightarrow e + \gamma) = \frac{3e^2}{64\pi^2 G_F^2} |\mathcal{A}|^2 \left(1 - \frac{m_e^2}{m_\mu^2}\right)^3, \quad (22)$$

where the amplitude \mathcal{A} arises at one-loop order and depends on the mass of the η scalar and masses M_i of the TeV-scale N_R as

$$\mathcal{A} = \frac{\lambda_{ei} \lambda_{\mu i}^*}{12(M_i^2 - m_\eta^2)} \left\{ 2 + \frac{9m_\eta^2}{M_i^2 - m_\eta^2} + 6 \left(\frac{m_\eta^2}{M_i^2 - m_\eta^2} \right)^2 - \frac{6M_i^4 m_\eta^2}{(M_i^2 - m_\eta^2)^3} \ln \left(\frac{M_i^2}{m_\eta^2} \right) \right\}. \quad (23)$$

Here, a sum over i is assumed.

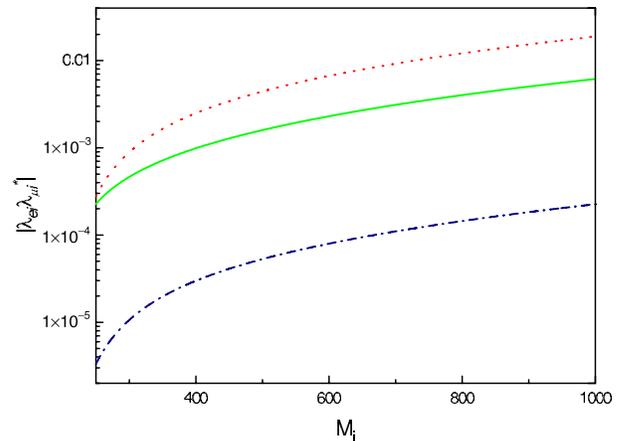


FIG. 1 (color online). $|\lambda_{\mu i}^* \lambda_{ei}|$ as a function of M_i (assumed here to be degenerate) constrained by charged lepton flavor-violating processes. The solid line is the constraint from the null result for $\mu \rightarrow e\gamma$ [37]. The dotted line is the constraint of the current null result for μ - e conversion in 197Au. The dot-dashed line gives the reach of future μ - e conversion searches, assuming a sensitivity to the conversion branching ratio of 10^{-18} [38].

Looking to the future, experiments searching for the $\mu - e$ conversion in nuclei with competitive sensitivity are planned at Fermilab and J-PARC. For this process, the dominant contribution arises from the exchange of a virtual photon that couples to a loop-induced $\mu - e$ charge radius operator. Nuclear coherence enhances this contribution by $\sim Z^2$ over that associated with the magnetic dipole operator, while the Z^0 -exchange dipole and charge-radius conversion amplitude is suppressed by a factor of m_μ^2/m_Z^2 . Retaining only the leading contribution yields the branching ratio [35,36]

$$BR_{\mu \rightarrow e}^A = R_{\mu \rightarrow e}^0(A) \left| 1 + \frac{\tilde{g}_{LV}^p V^p(A)}{A_R D(A)} + \frac{\tilde{g}_{LV}^n V^n(A)}{A_R D(A)} \right|^2 \times BR(\mu \rightarrow e\gamma), \quad (24)$$

with

$$R_{\mu \rightarrow e}^0(A) = \frac{G_F^2 m_\mu^5}{192 \pi^2 \Gamma_{\text{capt}}^A} |D(A)|^2 \quad (25)$$

and

$$\begin{aligned} \tilde{g}_{LV}^p &= 2g_{LV}^u + g_{LV}^d, & \tilde{g}_{LV}^n &= g_{LV}^u + 2g_{LV}^d, & A_R &= \frac{\sqrt{2}}{8} \frac{\mathcal{A}}{G_F m_\mu}, \\ g_{LV}^q &= -\frac{s_w^2}{72\pi^2} \frac{m_W^2}{M_i^2 - m_\eta^2} Q_q \lambda_{ei} \lambda_{\mu i}^* \left[2 + \frac{3M_i^2}{M_i^2 - m_\eta^2} + 6 \left(\frac{M_i^2}{M_i^2 - m_\eta^2} \right)^2 + 6 \left(\frac{M_i^2}{M_i^2 - m_\eta^2} \right)^3 \ln \frac{M_i^2}{m_\eta^2} \right], \end{aligned}$$

where $D(A)$, $V^p(A)$, and $V^n(A)$ are overlap integrals as a function of atomic number [36].

The current experimental upper bounds for the $BR(\mu \rightarrow e\gamma)$ and $\mu - e$ conversion in ^{197}Au are 1.2×10^{-11} and 7.0×10^{-13} [36–38], respectively. To illustrate the impact of these present and prospective limits, we assume that the heavy neutrinos are degenerate and plot in Fig. 1 constraints on $|\lambda_{\mu i}^* \lambda_{ei}|$ as a function of M_i implied by the $\mu \rightarrow e\gamma$ (solid line) and $\mu - e$ conversion (dotted line) processes. The dot-dashed line is the possible constraint from the future $\mu - e$ conversion experiments [38]. By assuming $M_i \sim 300$ GeV, $m_\rho \sim 150$ GeV, we obtain an upper bound for $|\lambda_{ei} \lambda_{\mu i}^*|$ of roughly 4×10^{-4} . As a result, $\lambda \sim \mathcal{O}(10^{-2})$. It can be found from the figure that the current bounds on $\mu - e$ conversion do not provide significant constraints on the parameter space. On the other hand, the sensitivities of prospective future $\mu - e$ conversion searches [38] exceed that of the MEG experiment.

IV. DARK MATTER

The fact that about 23% of the Universe is made of dark matter has been firmly established, while the nature of this dark matter still eludes us. A weakly interacting massive

particle (WIMP) is a promising dark matter candidate, since the WIMP relic density can be naturally near the experimental observed value for a WIMP mass around the electroweak scale. In our model, the fields ψ_L , introduced to cancel anomalies of the $U(1)'$ gauge symmetry (one for each N_R), provide a WIMP candidate. The ψ_L mass can be generated through a dimension-5 effective operator, in a similar way to that for the right-handed neutrinos:

$$\frac{1}{\Lambda} \tilde{Y}_{ij} \Phi^2 \bar{\psi}_{Li}^c \psi_{Lj}, \quad (26)$$

where Λ is a cutoff scale. After the spontaneous breaking of the $U(1)'$, ψ gets a nonzero mass:

$$(M_\psi)_{ij} = \frac{v_0^2}{\Lambda} \tilde{Y}_{ij}. \quad (27)$$

Due to the Z_2 discrete flavor symmetry, the lightest ψ_L is a massive stable particle, and thus can be the cold dark matter candidate.

There are three annihilation channels for the dark matter illustrated in Fig. 2: (1) $\psi\bar{\psi}$ annihilates into the SM fields through the mixing between Z and Z' ; (2) $\psi\bar{\psi}$ annihilates into the SM particles directly through the mixing between

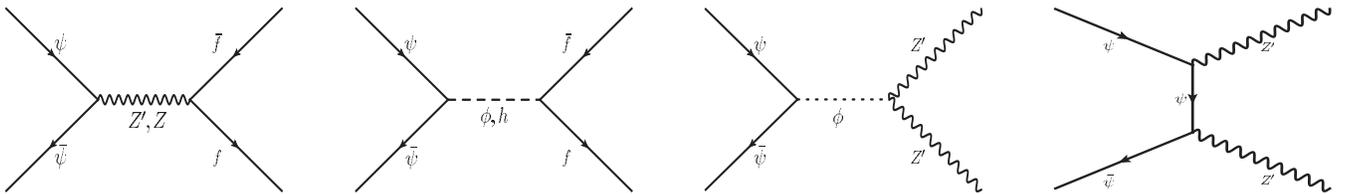


FIG. 2. Feynman diagrams for the annihilation of dark matter.

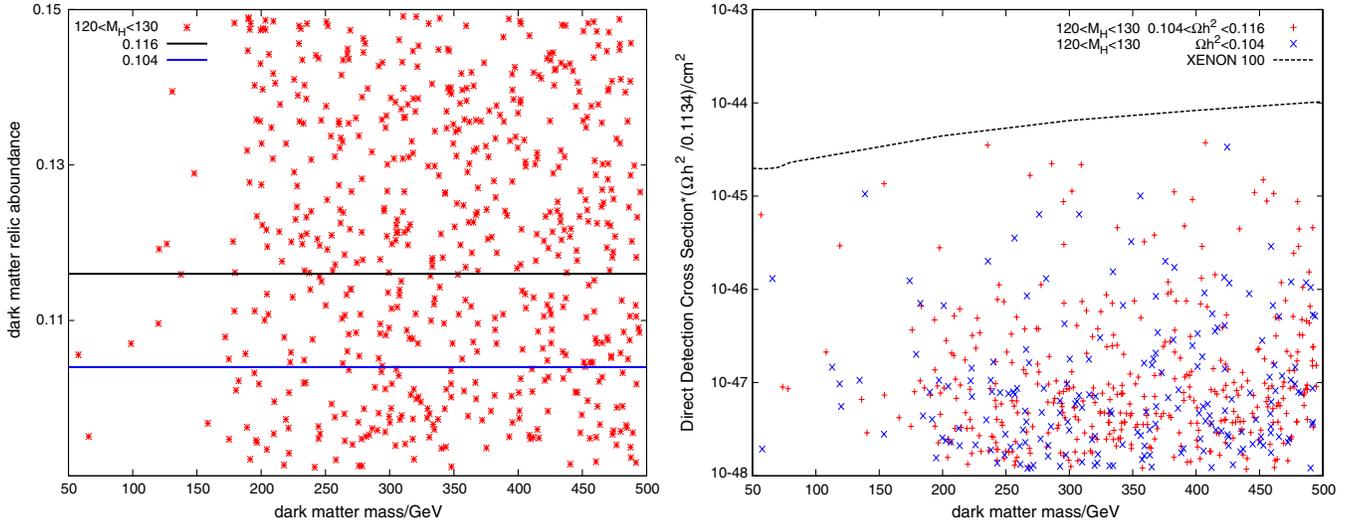


FIG. 3 (color online). Left panel: Dark matter relic abundance as a function of the dark matter mass. The horizontal band indicates the region consistent with the current relic density measurement from WMAP [2]. Right panel: Scaled dark matter direct-detection cross section as a function of the dark matter mass. The dashed line indicates the current bound from XENON 100 [42].

the SM Higgs and the scalar singlet; and (3) $\psi\bar{\psi}$ annihilates into the $Z'Z'$, which decays subsequently into the SM fields through mixing with the Z boson. We calculate the resulting relic density and the related direct-detection cross section using Micromegas [39,40], which solves the Boltzman equation numerically and utilizes CALCHEP [41] to calculate the relevant cross sections.

We plot in Fig. 3 (left panel) the dark matter relic abundance as a function of the dark matter mass while varying the parameters in the scalar potential over the following ranges: $m_0 \in [1 \times 10^2 \text{ GeV}, 1 \times 10^4 \text{ GeV}]$, $m_1 \in [1 \times 10^2 \text{ GeV}, 1 \times 10^3 \text{ GeV}]$, $\lambda_0 \in [0.1, 10]$, λ_1 and $\lambda_5 \in [0.01, 1]$, and $M_{Z'} \in [200 \text{ GeV}, 500 \text{ GeV}]$. The horizontal band represents the region consistent with the current relic density measurement from WMAP, $0.104 < \Omega h^2 < 0.116$ [2]. It is evident from the figure that one can achieve the observed relic density with a dark matter mass in the range $50 \sim 500 \text{ GeV}$.

We plot in Fig. 3 (right panel) the dark matter direct-detection cross section, scaled by the fraction of the relic density produced, as a function of the dark matter mass. The black dashed line gives the latest experimental constraint from XENON 100 [42]. We conclude that the expected spin-independent direct-detection scattering cross section is below the current experimental bound. However, one could expect a signal in the future XENON1T [43], which aims to probe the cross sections of order $\sigma \sim 2 \times 10^{-47} \text{ cm}^2$ within two years of operation.

It is interesting to analyze the dependence of the relic density and direct detection cross section on the $\phi-h$ mixing angle θ , which is the 13 mixing angle of the mass matrix given in Eq. (6). To that end, we plot in Fig. 4 (left panel) the mixing angle as a function of the dark matter mass, assuming saturation of the relic density. We observe

that in general the magnitude of the mixing angle is less than 15° , making the impact on Higgs boson production at the LHC marginal. In the right panel, we show the scaled direct-detection cross section as a function of the mixing angle and dark matter mass. We see that a nonvanishing signal in future searches with a sensitivity of $\sigma_{\text{SI-scaled}} \sim 10^{-47}$ would require a mixing angle larger than 1° in magnitude. In principle, it is conceivable that a future high-precision Higgs factory could provide a complementary probe of doublet-singlet mixing at this level.

V. LEPTOGENESIS

In addition to providing a viable dark matter candidate, the $U(1)'$ scenario allows for a viable baryogenesis via leptogenesis scenario. As the Majorana masses are in the hundreds of GeV range, there exists at least some possibility of testing or constraining this leptogenesis mechanism experimentally. To explore this possibility in an illustrative case, we assume that right-handed Majorana neutrinos are hierarchical, $M_1 \ll M_{2,3}$, so that studying the evolution of the number density of N_1 is sufficient. In the mass scale regime of interest here, the interactions mediated by all charged lepton Yukawas are in equilibrium. Consequently, we should consider flavor-dependent leptogenesis. The CP -violating asymmetries generated by N_1 decays are [44,45]

$$\varepsilon_{\alpha\alpha} = \frac{1}{8\pi} \frac{1}{(\lambda^\dagger \lambda)_{11}} \sum_j \text{Im}\{\lambda_{\alpha 1}^* (\lambda^\dagger \lambda)_{1j} \lambda_{\alpha j}\} g\left(\frac{m_j^2}{m_1^2}\right), \quad (28)$$

where the wave-function plus vertex contributions are included in [44–46]

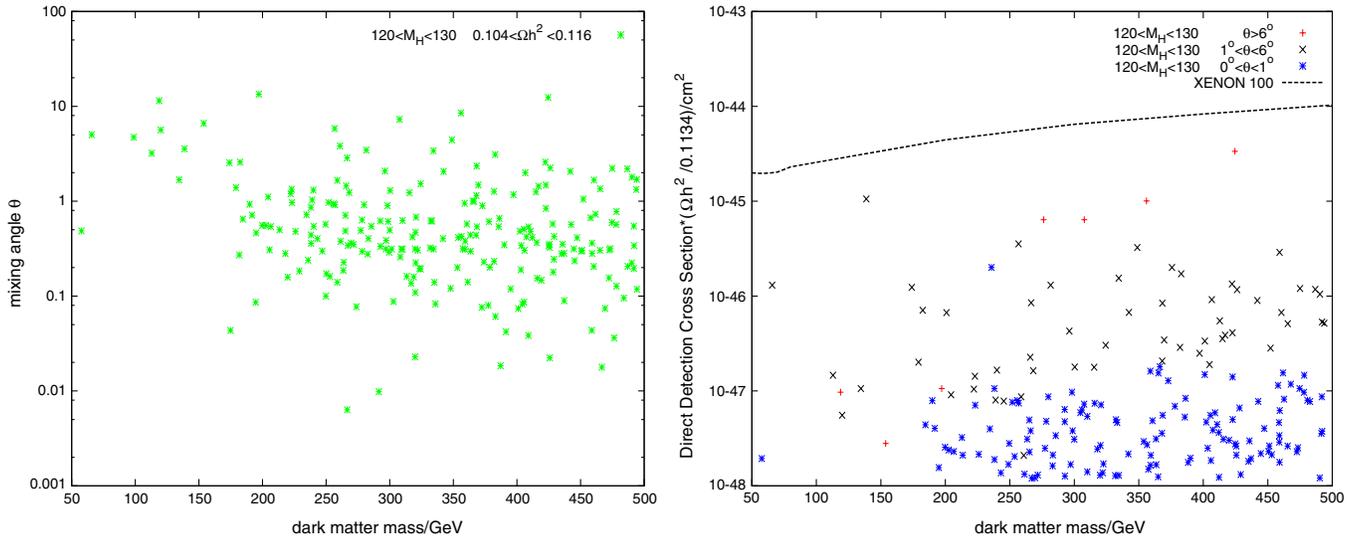


FIG. 4 (color online). Left panel: $h - \phi$ mixing angle θ (in degrees) as a function of the dark matter mass. Right panel: Scaled dark matter direct-detection cross section as a function of the dark matter mass, assuming a relic density consistent with the WMAP determination. The dashed line shows the current bound from XENON 100 [42].

$$g(x) = \sqrt{x} \left[\frac{1}{1-x} + 1 - (1+x) \ln \left(\frac{1+x}{x} \right) \right] \rightarrow -\frac{3}{2}x^{-1/2} - \frac{5}{6}x^{-3/2} \quad (1 \ll x). \quad (29)$$

In addition to $\varepsilon_{\alpha\alpha}$, the final baryon asymmetry depends on washout parameters:

$$K_{\alpha\alpha} \equiv \frac{\Gamma(N_1 \rightarrow \eta \ell_\alpha)}{H(M_1)} = \frac{|\lambda_{\alpha 1}|^2 M_1 / 4\pi}{\sqrt{g_*} M_1^2 / M_{\text{pl}}} \equiv \frac{\tilde{m}_{\alpha\alpha}}{\tilde{m}^*}, \quad (30)$$

where $H(M_1)$ denotes the value of the Hubble rate evaluated at a temperature $T = M_1$, M_{pl} is the Planck mass, M_1 is the mass of the lightest right-handed neutrino, $\tilde{m}^* = 3 \times 10^{-3}$ eV, and $\tilde{m} = 3 \times 10^{-3}$. Note that we have expressed $K_{\alpha\alpha}$ in terms of a scale associated with the neutrino mass, $\tilde{m}_{\alpha\alpha}$, and translated the remaining dimensional factors into \tilde{m}^* .

The washout factor $K_{\alpha\alpha}$ should be smaller than some maximum, K_{Max} ; otherwise, the washout effect would be too strong to generate the proper matter-antimatter asymmetry. For conventional thermal leptogenesis, $K_{\text{Max}} \sim 1$, while it may be as large as ~ 1000 for resonant leptogenesis [47]. As we show below, $K_{\alpha\alpha}$ can be of $\mathcal{O}(100)$ of the present instance, since $\varepsilon_{\alpha\alpha}$ in this model can be much larger than that of the conventional thermal leptogenesis case. The lepton asymmetry for the flavor α can be expressed approximately as [46]

$$Y_{\alpha\alpha} \approx \frac{\varepsilon_{\alpha\alpha}}{g_*} \left[\left(\frac{\tilde{m}_{\alpha\alpha}}{8.25 \times 10^{-3} \text{ eV}} \right)^{-1} + \left(\frac{0.2 \times 10^{-3} \text{ eV}}{\tilde{m}_{\alpha\alpha}} \right)^{-1.16} \right]^{-1}. \quad (31)$$

Nonperturbative sphaleron interactions partially convert this lepton asymmetry into a net baryon number asymmetry. Taking into account the flavor effects, the final baryon asymmetry is given by [46]

$$Y_B \approx -\frac{8}{25} \left(\frac{40}{13} Y_{ee} + \frac{51}{12} Y_{\mu\mu} + \frac{51}{13} Y_{\tau\tau} \right). \quad (32)$$

It should be noted that the analytical solutions given in Eqs. (31) and (32) only work for the case in which the $\Delta L = 1$ washout effect is much larger than the $\Delta L = 2$ washout effect, which can be expressed as [48]

$$K_{\Delta L=2} \sim \sum_{ij} |(\lambda^\dagger \lambda)_{ij}^2| \frac{1}{M_i M_j} \times \frac{T M_{\text{pl}}}{32\pi^3 \zeta(3)} \sqrt{\frac{90}{8\pi^2 g_*}} \quad (33)$$

for $T < M_1$. If $K_{\Delta L=2}$ is comparable with the $K_{\alpha\alpha}$, one must resort to a numerical, rather than analytic, solution to the Boltzmann equations.

We emphasize that the neutrino Yukawa couplings enter both $\varepsilon_{\alpha\alpha}$ and the washout factor, translated into $\tilde{m}_{\alpha\alpha}$ in Eq. (31). As a result, the magnitudes of the $\lambda_{\alpha 1}$ are strongly bounded: $|\lambda_{\alpha 1}| \lesssim 10^{-6}$. If the structure of the neutrino Yukawa matrix is anarchical, we would then expect the magnitudes of the couplings that enter CLFV observables to be $|\lambda_{\mu i}^* \lambda_{ei}| \lesssim 10^{-12}$, far below a level that could be probed in the next generation of CLFV searches (see Fig. 1). Consequently, observable CLFV could arise only in the presence of a strongly hierarchical Yukawa matrix, with $|\lambda_{\mu i}^* \lambda_{ei}|$ for $i = 2, 3$ being 6 to 7 orders of magnitude larger than for $i = 1$. Such a strong deviation from the assumption of anarchy would suggest the presence of an additional lepton flavor symmetry, as we now discuss.

In general, the neutrino Yukawa coupling that is compatible with low-energy data, written in the basis where the charged lepton Yukawa coupling and right-handed neutrino mass matrix are diagonal, is given by $Y_\nu = V_{PMNS}/v_n\sqrt{m_\nu}R\sqrt{M_N}$ [49], where V_{PMNS} is the neutrino mixing matrix and R is a 3×3 orthogonal matrix that parameterizes the information lost in decoupling the three right-handed neutrinos. In our model, we need a hierarchical neutrino Yukawa coupling Y_ν , which might be obtained by choosing a special but complicated structure for R . Here we instead use a flavor symmetry to generate a hierarchical Yukawa coupling Y_ν directly. To be concrete, we assume the neutrino sector admits a $U(1)_\ell$ symmetry [50]. The $U(1)_\ell$ charges of the fields are given by $Q(\ell_L) = Q(\ell_R) = 1$, $Q(1/\sqrt{2}(e^{i\alpha}N_{R2} + e^{i\beta}N_{R3})) = -Q(1/\sqrt{2}(e^{-i\alpha}N_{R2} - e^{-i\beta}N_{R3})) = 1$, and $Q(N_{R1}) = 0$. We assume that the $U(1)_\ell$ symmetry is explicitly broken by the GUT or Planck-scale physics, such that Yukawa interaction $\varepsilon\ell_L H_n N_{R1}$ emerges. In this case, the Yukawa interaction matrix can be given by

$$\lambda = \begin{pmatrix} \varepsilon & ae^{i\alpha} & ae^{i\beta} \\ \varepsilon & be^{i\alpha} & be^{i\beta} \\ \varepsilon & ce^{i\alpha} & ce^{i\beta} \end{pmatrix}, \quad (34)$$

where ε arises from the lepton number-violating Yukawa interaction term and is thus relatively small; a , b , and c are arbitrary complex parameters, the scale of which is restricted by lepton flavor-violating decays. Following from Eq. (34), the mass matrix of the right-handed neutrinos is given by

$$M_R = \begin{pmatrix} M_1 & 0 & 0 \\ 0 & M_2 & M_2 c_{\alpha-\beta} \\ 0 & M_2 c_{\alpha-\beta} & M_2 \end{pmatrix} + \Delta M, \quad (35)$$

where $c_{\alpha-\beta} = \cos(\alpha - \beta)$. Given the neutrino Yukawa coupling in Eq. (34) and the heavy neutrino mass matrix in Eq. (35), the structure of the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix implied by neutrino oscillation data cannot arise solely from the neutrino sector. We recall that the PMNS matrix follows from the mismatch between the diagonalizations of the neutrino mass matrix and the charged lepton mass matrix, i.e., $V_{PMNS} = V_e^\dagger V_\nu$, where V_e and V_ν , respectively, rotate the left-handed charged and neutral lepton flavor eigenstates to the mass eigenstates. The correct PMNS matrix may then emerge from an appropriately chosen structure for the charged lepton Yukawa matrix.

Notice that any lepton asymmetry produced after the $(B + L)$ -violating interactions drop out of thermal equilibrium will not be converted into a baryon asymmetry. It was found [50] that when the mass of the right-handed neutrino is very close to the critical temperature T_C for the

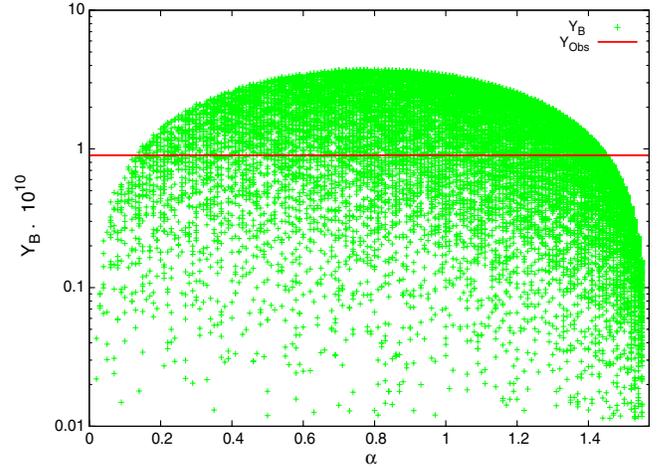


FIG. 5 (color online). Y_B as a function of α . The horizontal line is the center value of the experimental observation.

electroweak phase transition ($T_C \approx 129$ GeV by taking the SM Higgs boson mass of 125 GeV), freeze-out effects from sphaleron processes dropping out of equilibrium need to be considered. Nevertheless, the correct baryon asymmetry can be generated in this low-mass region because the lepton asymmetry can be sufficiently large in both resonant leptogenesis and our case. In Fig. 5, we plot Y_B as a function of α by setting $\varepsilon \sim 10^{-6}$, $\beta = 0$, and $|a| = |b| = |c| \times 10^{-2} = 10^{-4}$, with their phases varying arbitrarily in the range $[0, \pi]$. We also set $M_1 = 300$ GeV and $M_2 = 500$ GeV. In this case, the $\Delta L = 2$ washout factor is $K_{\Delta L=2}^{\text{Max}} \sim 0.1$, which can be neglected compared with the $\Delta L = 1$ washout factor, $K_{\alpha\alpha} \sim 50$. In this case, we may safely take Eq. (31) as the analytical solution to the Boltzmann equations. The horizontal line in the figure represents the current experimental value of Y_B . It is clear that the matter-antimatter asymmetry of the Universe can be generated. Notice that $K_{\alpha\alpha} \gg 1$ in our case, which is a salient feature of this scenario compared with the conventional thermal leptogenesis case but not one requiring near degeneracies in the heavy neutrino spectrum. Note that for the choice of parameters given here, $|\lambda_{\mu i}^* \lambda_{ei}| < 10^{-8}$ for $i = 2, 3$, implying an unobservable signal in the next generation of CLFV searches.⁵

VI. CONCLUSION

The search for new physics at the TeV scale is motivated in part by naturalness considerations and in part by the possibility that new TeV-scale dynamics may account for the origin of the visible and dark matter of the Universe. On the other hand, the observation of neutrino oscillations and the tiny scale of neutrino masses point to new physics

⁵More generally, we have estimated that for $|\lambda_{ai}^* \lambda_{\beta i}| \lesssim 10^{-6}$ ($\alpha, \beta = e, \mu, \tau$ and $i = 2, 3$), $\Delta L = 1$ washout processes dominate over $\Delta L = 2$ processes.

at much higher scales, as suggested by the conventional seesaw paradigm. It is interesting to ask whether nature may have generated both neutrino mass and the matter content of the Universe at the TeV scale. If so, then one would anticipate signatures in experiments sensitive to BSM physics at this scale.

In this paper, we have analyzed a simple BSM scenario that can account for dark matter, baryogenesis, and neutrino mass with new TeV-scale degrees of freedom and shown that, nonetheless, the experimental signatures are likely to be sparse at best. We have made no attempt to alleviate the Higgs-mass fine-tuning problem, though it is possible that an embedding of this scenario in a UV-complete model may do so. In this minimal scenario, a neutrinophilic 2HDM with a local $U(1)'$ symmetry, the second Higgs doublet is entirely responsible for neutrino mass. Its VEV can be naturally small, allowing for $\mathcal{O}(1)$ Yukawa couplings, TeV-scale right-handed neutrinos, and the possibility of observable CLFV. The associated flavor-dependent low-scale (nonresonant) leptogenesis can account for the cosmic baryon asymmetry, while the fermions needed for anomaly cancellation provide a suitable dark matter candidate.

Despite having the ingredients for a plethora of experimental signatures, we find that simultaneously solving the neutrino mass, dark matter, and baryon asymmetry problems implies that most distinctive features of this scenario (apart from a nondistinctive spin-independent direct-detection signal) would be out of reach in the foreseeable future in the absence of additional lepton flavor symmetries. This situation contrasts with a variety of other BSM scenarios, for which new symmetries must be imposed in order to suppress otherwise large deviations from the SM that are inconsistent with observation. In short, nature's solutions to some of the key problems at the interface of particle physics and cosmology may lie at the TeV scale, yet either remain hidden from view or point to an even more complex flavor problem.

ACKNOWLEDGMENTS

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APPENDIX: Z – Z' MIXING AT THE ONE-LOOP LEVEL

The one-loop contributions to the $\Pi_{ZZ'}^T$ are given by

$$\begin{aligned} \Pi_{ZZ'}^{T1} = & + \frac{gg''(c_w^2 - s_w^2)}{(4\pi)^2 c_w} \left\{ \left(-\frac{1}{6}q^2 + m_\rho^2 \right) (\alpha_\varepsilon + 1) \right. \\ & + q^2 F(m_\rho^2, m_\rho^2, q^2) - 2m_\rho^2 F_1(m_\rho^2, m_\rho^2, q^2) \left. \right\} \\ & + \frac{gg''}{(4\pi)^2 c_w} \left\{ \left(-\frac{1}{6}q^2 + \frac{1}{2}(m_\eta^2 + m_\delta^2) \right) (\alpha_\varepsilon + 1) \right. \\ & + q^2 F(m_\eta^2, m_\delta^2, q^2) - m_\eta^2 F_1(m_\delta^2, m_\eta^2, q^2) \\ & \left. - m_\delta^2 F_1(m_\eta^2, m_\delta^2, q^2) \right\}, \end{aligned} \quad (\text{A1})$$

$$\begin{aligned} \Pi_{ZZ'}^{T2} = & - \frac{c_w gg' g^3 v_2^2}{2(4\pi)^2} [\alpha_\varepsilon - F_0(m_W^2, m_\rho^2, q^2)] \\ & + \frac{g^3 g' v_2^2}{2c_w^3 (4\pi)^2} [\alpha_\varepsilon - F_0(m_Z^2, m_\eta^2, q^2)] \\ & + \frac{gg'^3 v_2^2}{2c_w (4\pi)^2} [\alpha_\varepsilon - F_0(m_{Z'}^2, m_\eta^2, q^2)], \end{aligned} \quad (\text{A2})$$

$$\begin{aligned} \Pi_{ZZ'}^{T3} = & - \frac{gg''}{2c_w (4\pi)^2} \{ (m_\eta^2 + m_\delta^2) (\alpha_\varepsilon + 1) \\ & - m_\eta^2 \ln m_\eta^2 - m_\delta^2 \ln m_\delta^2 \} \\ & - \frac{gg''(c_w^2 - s_w^2)m_\rho^2}{c_w (4\pi)^2} (\alpha_\varepsilon + 1 - \ln m_\rho^2), \end{aligned} \quad (\text{A3})$$

where $\alpha_\varepsilon = 1/\varepsilon - \gamma_E + \ln 4\pi + \ln \mu^2$ and g'' is the coupling constant of the new $U(1)$ gauge symmetry. $\Pi_{ZZ'}^{T1}$ comes from the vertex with two Higgs bosons and one gauge field, and $\Pi_{ZZ'}^{T2}$ comes from the vertex with two gauge fields and one Higgs boson, while $\Pi_{ZZ'}^{T3}$ comes from the vertex with two gauge fields and two Higgs bosons. We define $\hat{\Pi}_{ZZ'}^T \equiv \Pi_{ZZ'}^{T1} + \Pi_{ZZ'}^{T2} + \Pi_{ZZ'}^{T3}$.

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