PHYSICAL REVIEW D **88,** 035001 (2013)

Dark matter and vectorlike leptons from gauged lepton number

Pedro Schwaller,[1,2] Tim M. P. Tait,[3] and Roberto Vega-Morales[4,5]

[1]*Argonne National Laboratory, HEP Division, Argonne, Illinois 60439, USA*
[2]*Physics Department, University of Illinois-Chicago, Chicago, Illinois 60607, USA*
[3]*Physics Department, University of California-Irvine, Irvine, California 92697, USA*
[4]*Fermi National Accelerator Laboratory (FNAL), Batavia, Illinois 60510, USA*
[5]*Department of Physics and Astronomy, Northwestern University, Evanston, Illinois 60208, USA*
(Received 19 May 2013; published 1 August 2013)

We investigate a simple model where lepton number is promoted to a local $\mathrm{U}(1)_L$ gauge symmetry which is then spontaneously broken, leading to a viable thermal dark matter (DM) candidate and vectorlike leptons as a byproduct. The dark matter arises as part of the exotic lepton sector required by the need to satisfy anomaly cancellation and is a Dirac electroweak (mostly) singlet neutrino. It is stabilized by an accidental global symmetry of the renormalizable Lagrangian which is preserved even after the gauged lepton number is spontaneously broken and can annihilate efficiently to give the correct thermal relic abundance. We examine the ability of this model to give a viable DM candidate and discuss both direct and indirect detection implications. We also examine some of the LHC phenomenology of the associated exotic lepton sector and in particular its effects on Higgs decays.

DOI: 10.1103/PhysRevD.88.035001 PACS numbers: 12.60.Cn, 95.35.+d

I. INTRODUCTION

With the recent discovery of a new resonance with standard model (SM) Higgs-like properties [1,2] the final piece of the SM appears to be in place. It is well known, however, that there are questions for which the SM has no answer and physics beyond the standard model is needed. Chief among these questions is the nature of dark matter (DM) and the mechanism which makes it stable. It is also well known that the renormalizable SM Lagrangian possesses an (anomalous) accidental global symmetry associated with the conservation of overall lepton number. If one allows for higher-dimensional operators, lepton-violating interactions can occur at dimension five, but to date no such processes (with the possible ambiguous exception of neutrino masses) have been observed experimentally [3]. This is perhaps an indication that lepton number is a more fundamental symmetry which prevents the generation of SM lepton-number-violating operators. In this work, we connect the apparent lack of lepton number violation to the stability of thermal relic dark matter, by deriving both from a $\mathrm{U}(1)_L$ gauge symmetry associated with lepton number.

Gauging lepton number is attractive for both phenomenological as well as theoretical reasons and the possibility of lepton number (and also baryon number) as a local gauge symmetry was first explored in Refs. [4,5]. However, the first complete and consistent model of gauged lepton number (and baryon number) was not explored until more recently in Ref. [6] with numerous variations following [7–12]. Here we explore a particular realization where the DM arises as part of the exotic lepton sector required by gauging lepton number and the attendant need to cancel anomalies.

The DM candidate is a Dirac electroweak (mostly) singlet neutrino stabilized by an accidental global symmetry of the renormalizable Lagrangian which is preserved even after lepton number is spontaneously broken. As we will see, as a byproduct of the lepton-breaking mechanism and the requirement of a viable DM candidate, one also obtains a set of vectorlike leptons which can have interesting phenomenology at the LHC through either direct production or through modifications of Higgs decays to SM particles.

We extend the SM gauge group to $\mathrm{SU}(3)_c \otimes \mathrm{SU}(2)_W \otimes \mathrm{U}(1)_Y \otimes \mathrm{U}(1)_L$ where the SM leptons are charged under $\mathrm{U}(1)_L$. The anomalous $\mathrm{U}(1)_L$ requires us to add a new set of leptons with the appropriate quantum numbers to cancel anomalies. Typically, $\mathrm{U}(1)_L$ is spontaneously broken by the vacuum expectation value (VEV) of a SM singlet scalar in such a way that Majorana masses can be generated for the right-handed neutrinos (whose presence is required by anomaly cancellation [6]). Such constructions allow for a simple realization of the well known "seesaw" mechanism of neutrino mass generation, but do not contain viable dark matter candidates without additional assumptions or particle content.

Here, motived by the desire for a thermal DM candidate, we choose to break lepton number with a SM singlet scalar carrying $L = 3$. This leads to a remnant global U(1) symmetry preventing the decay of the lightest new lepton which stabilizes the DM candidate. This global symmetry is a consequence of the gauge symmetry and particle content of the model and does not need to be additionally imposed. It also ensures that the model is safe from dangerous flavor-violating processes which are highly constrained by experiment. An automatic consequence of this construction is that one also obtains a new generation of vectorlike (with respect to the SM) leptons after the spontaneous breaking of lepton number. This type of lepton spectrum has garnered recent interest in the context of modifications to the Higgs decay into diphotons [13–19] and was also recently shown to be useful for baryogenesis [20,21].

The organization of this paper is as follows. In Sec. II we briefly review the gauging of lepton number and cancellation of anomalies. We also discuss the details of the lepton-breaking mechanism as well as the particle content and Lagrangian. In Sec. III we discuss the DM candidate and stability and obtain the relic abundance for a range of DM masses. We also examine the direct and indirect detection prospects. In Sec. IV we discuss constraints as well as LHC phenomenology and examine the effect of the vectorlike leptons on the Higgs-to-diphoton rate. We present our conclusions and an overview of possible future work in Sec. V.

II. THE MODEL

The SM gauge group is extended to $SU(3)_c \otimes SU(2)_W \otimes U(1)_Y \otimes U(1)_L$ where L represents the lepton charge. We restrict ourselves to the minimal particle content consisting of a set of anomaly-canceling exotic leptons, plus the new gauge field and a SM singlet scalar which breaks lepton number spontaneously. In principle, this theory is UV complete up to large energies, and we restrict ourselves to considering renormalizable interactions. We discuss each of these ingredients, including the interactions, below.

A. Anomaly cancellation

The anomalies introduced when gauging lepton number and various ways to cancel them with the addition of new fermions are discussed in detail in Refs. [6–8]. All options include three generations of right-handed singlet neutrinos (ν_{Ri}, considered as part of the SM) with quantum numbers $\nu_{Ri} \equiv (1, 0, 1)$ under $(SU(2)_W, U(1)_Y, U(1)_L)$ and $i = e, \mu, \tau$. We define all SM leptons to have $L = 1$. In addition to ν_{Ri}, one must add new electroweak doublet and singlet leptons to cancel the gauge anomalies. There are several options; here we focus on a simple construction making use of two exotic generations of chiral fermions which together form a vectorlike set under the SM gauge group [8], insuring that anomaly cancellation in the SM gauge factors is preserved. The first set of new fermions is a sequential fourth generation of leptons carrying lepton number $L = L'$,

$$\begin{aligned} \ell'_L &\equiv (\nu'_L e'_L) \equiv (2, -1/2, L'), \\ e'_R &\equiv (1, -1, L'), \\ \nu'_R &\equiv (1, 0, L'). \end{aligned} \tag{1}$$

The second is a mirror set of opposite chirality with lepton number $L = L'' = L' + 3$,

$$\begin{aligned} \ell''_R &\equiv (\nu''_R e''_R) \equiv (2, -1/2, L''), \\ e''_L &\equiv (1, -1, L''), \\ \nu''_L &\equiv (1, 0, L''), \end{aligned} \tag{2}$$

where the condition

$$L' - L'' = -3 \tag{3}$$

is required by anomaly cancellation. The addition of two sets of chiral fermions carrying lepton number which together form a vectorlike set under the SM also avoids the need to add new quarks to cancel anomalies, although scenarios with exotic quarks are also interesting and have been explored in the context of gauged baryon number [6–8]. The particle content in Eqs. (1) and (2) is similar to that obtained in Ref. [10] where baryon number is also gauged and one obtains a vectorlike set of "lepto-quarks" as well as a potential DM candidate. Here we focus on only gauging lepton number which requires a simpler scalar sector and fewer new particles.

B. Gauge and Higgs sector

The gauging of lepton number will introduce a new spin-1 vector boson which we label Z_L. In addition to the usual Abelian vector-field kinetic terms, the $U(1)_L$ gauge field will have interactions,

$$\begin{aligned} \mathcal{L} \supset (D^\mu \Phi)^\dagger (D_\mu \Phi) + \frac{\epsilon}{2} Z_L^{\mu\nu} B_{\mu\nu} + \bar{l}'_L D_\mu \gamma^\mu l'_L \\ + \bar{l}''_R D_\mu \gamma^\mu l''_R + \bar{l}_i D_\mu \gamma^\mu l_i, \end{aligned} \tag{4}$$

where $D^\mu = \partial^\mu + ig' L Z_L^\mu$ with L the lepton number assignment for a particular field. $\Phi \equiv (1, 0, L_\Phi)$ is the SM singlet scalar carrying lepton number whose VEV (v_ϕ) breaks the $U(1)_L$ spontaneously. The index $i = e, \mu, \tau$ runs over all SM leptons while $l = \ell, e, \nu$ where ℓ is an SU(2) doublet and e, ν are singlets. Note that there is no $\delta M^2 Z_{L\mu} Z^\mu$ term since Φ is not charged under the SM and the Higgs does not carry L.

The parameter ϵ encapsulates the degree of kinetic mixing between $U(1)_L$ and $U(1)_Y$. One can in principle impose $\epsilon = 0$ at tree level through symmetries, but in general it is a free parameter of the theory and is additively renormalized by loops of leptons. While any value of ϵ at the weak scale can be engineered, the loop-induced piece is typically of order 10^{-3}, small enough to be consistent with experimental constraints without undue fine-tuning.

After lepton and electroweak symmetry breaking ϵ also leads to $Z - Z_L$ mixing parametrized by [22]

$$\tan 2\xi = \frac{2M_Z^2 s_W \epsilon \sqrt{1-\epsilon^2}}{M_{Z_L}^2 - M_Z^2(1-\epsilon^2) + M_Z^2 s_W^2 \epsilon^2}, \tag{5}$$

where ξ is the $Z_L - Z$ mixing angle and M_Z, M_{Z_L} are the masses. In the absence of mixing, $M_{Z_L} = L_\Phi g' v_\phi$. As we will see, since this mixing is constrained to be small by direct searches for dark matter (with weaker constraints from precision measurements [22–24]) we take M_Z, M_{Z_L} as the physical masses as well.

In the Higgs sector the existence of Φ allows for an expanded scalar potential,

$$V(H, \Phi) = -\mu_H^2 H^\dagger H + \lambda_H |H^\dagger H|^2 - \mu_\Phi^2 \Phi^\dagger \Phi + \lambda_\Phi |\Phi^\dagger \Phi|^2 + \lambda_{\rm hp} \Phi^\dagger \Phi H^\dagger H, \tag{6}$$

where $H \equiv (2, -1/2, 0)$ is the SM Higgs doublet. Once lepton number is broken, the real component of Φ obtains a vacuum expectation value $\langle \Phi \rangle = v_\phi/\sqrt{2}$, while the Higgs boson H obtains its own VEV, $\langle H \rangle = (0, v_h/\sqrt{2})$, to break the electroweak symmetry. The scale v_ϕ will be the only new dimensional scale introduced, with all of the other parameters being dimensionless couplings. We will see below in Sec. II D that $L_\Phi = 3$ is preferred.

The presence of the "Higgs portal" coupling $\lambda_{\rm hp}$ will generically lead to mixing between the real singlet components of Φ and H parametrized by the mixing angle,

$$\tan 2\theta = \frac{\lambda_{\rm hp} v_h v_\phi}{\lambda_\Phi v_\phi^2 - \lambda_H v_h^2}. \tag{7}$$

This mixing leads to the mass eigenstates,

$$h = c_\theta h_o - s_\theta \phi_o, \qquad \phi = s_\theta h_o + c_\theta \phi_o, \tag{8}$$

where ϕ_o and h_o are the gauge eigenstates and ϕ, h are the mass eigenstates with masses

$$m_{h,\phi}^2 = (\lambda_H v_h^2 + \lambda_\Phi v_\phi^2) \mp \sqrt{(\lambda_\Phi v_\phi^2 - \lambda_H v_h^2)^2 + \lambda_{\rm hp}^2 v_h^2 v_\phi^2}, \tag{9}$$

where we have assumed $m_\phi > m_h$ and defined $c_\theta = \cos\theta$, $s_\theta = \sin\theta$, etc. The coupling $\lambda_{\rm hp}$ will also lead to a tree-level shift in the Higgs quartic coupling [25], which provides a mechanism for stabilizing the vacuum in the presence of the exotic charged leptons with large Yukawa couplings to the SM Higgs. It was shown to be a particularly efficient stabilization mechanism when $m_\phi \gg m_h$, even for small mixing angles [15].

C. Global symmetries and breaking *L*

The two new sets of leptons along with the SM lepton sector comprise three separate sectors labeled by their lepton number $L = 1$, L', L'' for which global U(1) symmetries can be associated. These global symmetries are each separately conserved by the SM and $\mathrm{U}(1)_L$ interactions. Yukawa interactions (assuming L_Φ permits them) will break these symmetries in realistic models, as discussed below. A combination of precision electroweak, collider, and direct detection constraints prohibit a stable lepton which carries electroweak charge. Thus, couplings to the Higgs must not be too large and the DM cannot receive its mass solely from the SM Higgs, leading to the need to generate an additional contribution to the DM mass which does not come from electroweak symmetry breaking.

From these considerations one concludes that the SM singlets ν_R' and ν_L'' or some combination must compose the majority of the DM. Majorana masses can be generated by choosing the lepton-breaking scalar to carry $L_\Phi = 2L'$ or $L_\Phi = 2L''$. However, this choice still leaves either L' or L'' unbroken meaning that the lightest lepton of the corresponding sector will be stable and only receive its mass from its couplings to the Higgs, which as discussed is ruled out by experiment. It is clear that in order to avoid a heavy stable lepton with unacceptably large couplings to the Z or Higgs boson one must choose L_Φ such that it generates an interaction between the L' and L'' sectors. The anomaly cancellation condition of Eq. (3) ensures that the only possibility is $L_\Phi = 3$.

D. Yukawa sector

Given $L_\Phi = 3$, the Lagrangian for the Yukawa sector of the new leptons can be written as

$$\mathcal{L} \supset -c_\ell \Phi \bar{\ell}_R'' \ell_L' - c_e \Phi \bar{e}_L'' e_R' - c_\nu \Phi \bar{\nu}_L'' \nu_R' - y_e' H \bar{\ell}_L' e_R' - y_e'' H \bar{\ell}_R'' e_L'' - y_\nu' \tilde{H} \bar{\ell}_L' \nu_R' - y_\nu'' \tilde{H} \bar{\ell}_R'' \nu_L'' + \mathrm{H.c.} \tag{10}$$

In general these couplings are complex, containing phases which can lead to CP violation, but for simplicity we assume that all couplings in Eq. (10) are real (but see Refs. [26,27] for recent studies of CP-violating effects on the diphoton rate coming from vectorlike leptons). It is also clear from Eq. (10) that once Φ obtains a VEV the couplings c_ℓ, c_e, and c_ν will lead to vectorlike (with respect to the SM) masses for the exotic leptons. The new leptons will also receive mass contributions from electroweak symmetry breaking through the $y_{\nu,e}'$, $y_{\nu,e}''$ couplings. Note also that unless L', $L'' = 0$, explicit Majorana masses for ν_R' and ν_L'' are not allowed nor will they be generated after lepton number breaking unless $L' = -L'' = -3/2$. (This case was considered explicitly in the context of gauged lepton and baryon number with vectorlike "lepto-quarks" [10].) We avoid these choices in what follows.

In principle there may still be couplings between the exotic and SM leptons. Since we have taken the SM lepton number to be $L = 1$, this implies that L', $L'' \neq 1$ in order to avoid mixing with SM leptons which can lead to dangerous flavor-changing neutral currents as well as the decay of the DM. If we choose $L' = -4$, which fixes $L'' = -1$, then, in addition to those in Eq. (10), one can also generate interactions between the SM and the new lepton sector given by

$$\mathcal{L} \supset y \Phi \bar{\nu}_R^{c\prime} \nu_{Ri} + \mathrm{H.c.} \tag{11}$$

Once Φ obtains a VEV, this will lead to mixing between the SM right-handed neutrinos, ν_{Ri} and the exotic right-handed neutrino, ν_R'. This also implies that the exotic lepton sector can decay to the SM, thus eliminating this scenario as an explanation for dark matter. To summarize, in order to avoid mixing with the SM and ensure a stable DM candidate, we take $(L', L'') \neq (1, 4)$, $(-4, -1)$, $(-2, 1)$. Furthermore, to avoid Majorana mass terms we also assume $(L', L'') \neq (0, 3)$, $(-\frac{3}{2}, \frac{3}{2})$, $(-3, 0)$. Thus our complete Yukawa sector Lagrangian is given by Eq. (10) and L' can otherwise be any real number satisfying $L' = -3 + L''$.

In the limit that the Yukawa couplings $c_i \to 0$, one recovers the global symmetries which separately preserve L', L'' and L_{SM}. As a result, $c_i \ll 1$ are technically natural, implying that vectorlike masses for the new leptons much smaller than v_ϕ are natural. We also note that small values of the $y'_{\nu,e}$, $y''_{\nu,e}$, and $y^{\mathrm{SM}}_{\nu i}$ Higgs Yukawa couplings are technically natural.

It is worth noting that Eq. (10) is very similar to the Yukawa sectors proposed in a generic framework in Refs. [13,15], but here they arise from $\mathrm{U}(1)_L$ gauge invariance and anomaly cancellation. Only one new scale (v_ϕ) is introduced, with the masses of the new fermions following from dimensionless couplings. Furthermore, the global symmetries needed to protect against dangerous mixing with SM leptons and to ensure the existence of a stable DM particle are guaranteed by $\mathrm{U}(1)_L$ gauge invariance as opposed to being imposed by hand.

E. Experimental constraints

Low-energy experiments place a limit on the parameters which describe the Z_L sector. Since the SM Higgs does not carry lepton number and Φ is a SM singlet, there is no mass mixing between Z_L and the SM electroweak interaction at tree level. Furthermore, since Z_L does not couple to quarks, direct search limits from the LHC are rather weak, and the strongest limits are obtained from constraints on four-lepton operators derived from LEP II data [28]; these require

$$v_\phi \gtrsim 1.7 \text{ TeV}, \tag{12}$$

roughly independently of the value of g'.

This lower bound and the experimentally measured value of $m_h \simeq 125$ GeV constrains the quartic couplings in the scalar potential of Eq. (6) through Eqs. (7) and (9). By fixing $v_\phi = 1.7$ TeV and $m_h = 125$ GeV we can then examine the scalar mixing angle θ, the Higgs quartic λ_H, and the heavy scalar mass eigenstate m_ϕ as functions of the scalar couplings λ_{hp} and λ_Φ. In Fig. 1 we show contours of $\lambda_H(\lambda_\Phi, \lambda_{\mathrm{hp}})$ (solid-orange), $\theta(\lambda_\Phi, \lambda_{\mathrm{hp}})$ (dotted-red), and $m_\phi(\lambda_\Phi, \lambda_{\mathrm{hp}})$ (solid-black) in the $\lambda_{\mathrm{hp}} - \lambda_\Phi$ plane. As can be seen, values of $\theta \lesssim 0.1$–0.2 can be obtained for quartic couplings of $\mathcal{O}(1)$ and heavy scalar masses $\sim$2.5 TeV. To obtain mixings as large as $\theta \sim 0.4$ requires $\lambda_H \sim 3$ and small $\lambda_\Phi \lesssim 0.5$ with $m_\phi \sim 1.5$ TeV. In general we find

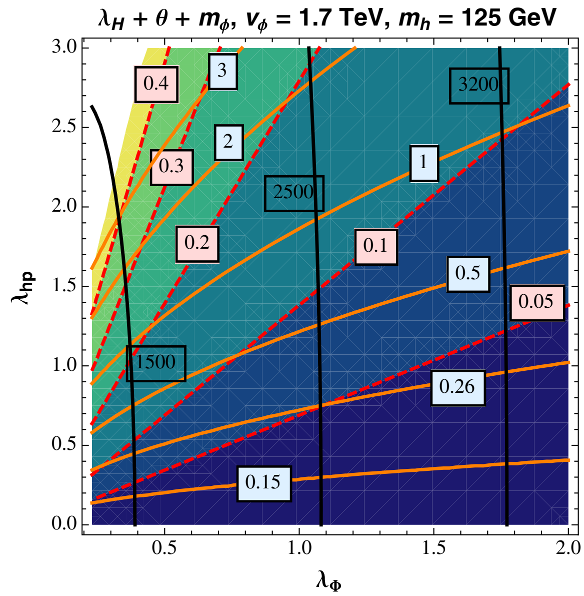

FIG. 1 (color online). Contours of Higgs mixing angle θ (red-dotted), Higgs quartic coupling λ_H (orange-solid), and heavy scalar mass m_ϕ in GeV (black-solid) as defined in Eqs. (7) and (9) as a function of the scalar couplings $(\lambda_{\mathrm{hp}}, \lambda_\Phi)$ in Eq. (6).

$m_\phi \gtrsim 1$ TeV for $v_\phi = 1.7$ TeV, possibly within reach of the LHC, but more likely too heavy to be produced directly.

Precision measurements on the Z-pole also constrain the degree of $Z_L - Z$ mass mixing. Since this occurs at loop level (through loops of the SM and exotic leptons as well as scalars), it will typically be small enough ($\lesssim 10^{-3}$) for any v_ϕ consistent with the LEP II bound. There are also constraints [via $\sin\xi$ in Eq. (5)] on the kinetic mixing parameter from direct detection [29], which are comparable to the expected size induced by loops of leptons. Using Eq. (5) we examine the $\epsilon - M_{Z_L}$ parameter space for typically allowed values of $\sin\xi \lesssim 10^{-4}$ over a range of Z_L masses. In Fig. 2 we present contours of $\sin\xi \times 10^4$ in the $\epsilon - M_{Z_L}$ plane for small values of the kinetic mixing parameter ϵ as would be favored in theories where $\epsilon = 0$ at tree level, as discussed in Sec. II B. We can see that for $M_{Z_L} \sim 1$ TeV one can obtain a $Z - Z_L$ mixing angle of $\sin\xi \sim 0.1 \times 10^{-4}$ with a kinetic mixing of $\epsilon \sim 0.002$.

F. Possible extensions

There are a number of possibilities for how one could extend this model or embed it into a more complete theory. For instance,with the need to break lepton number spontaneously, the question as to how one obtains v_ϕ naturally also arises. One could imagine embedding this model in a supersymmetric version as was done in Refs. [8,9,30] for other gauged lepton number constructions. Another possibility is to have the scalar sector of this model arise as part of a set of Goldstone bosons resulting from a strongly broken global symmetry, as in, for example, Ref. [31].

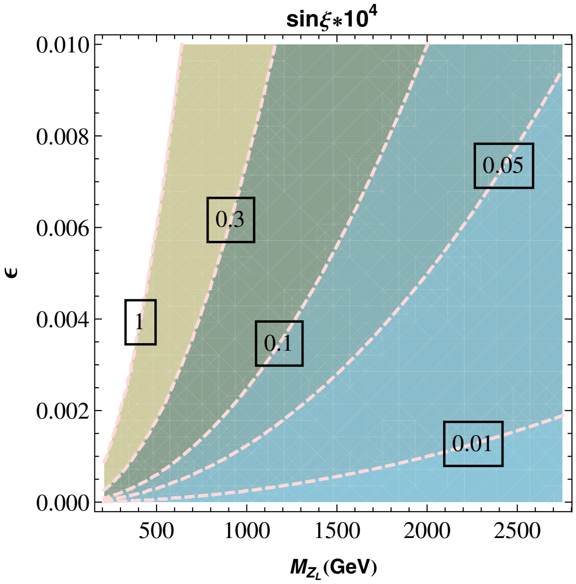

FIG. 2 (color online). Contours of the $Z - Z_L$ mixing angle $\sin\xi$ ($\times\, 10^4$) in the $\epsilon - M_{Z_L}$ plane [see Eq. (5)].

Another possibility for generating natural values for not only v_ϕ but also the electroweak scale (v_h) is through dimensional transmutation where v_ϕ is generated radiatively [32]. This scale is then inherited by the SM through the "Higgs portal," as was done recently in Ref. [33] for a hidden U(1) gauge extension of the SM, but we leave it to a future study to explore this possibility. For the remainder of this study we simply set v_ϕ to its lower bound of $v_\phi = 1.7$ TeV.

One can also extend the theory to obtain $\epsilon = 0$ at tree level in Eq. (4) by positing that the $U(1)_L$ gauge symmetry arises out of a larger non-Abelian gauge symmetry which forbids $\epsilon \neq 0$ [34] and is broken at some high scale Λ down to $U(1)_L$. Below the scale Λ, but above the lepton- and electroweak-breaking scales, loop corrections due to hypercharged leptons vanish provided the leptons satisfy an orthogonality condition [34],

$$\mathrm{Tr}(LY) = 0. \tag{13}$$

Combined with the anomaly cancellation constraint in Eq. (3), this would determine the exotic lepton numbers to be $L' = -3$ and $L'' = 0$. Below v_ϕ and v_h there will be loop-induced (from both leptons and scalars) corrections which generate a kinetic mixing, but typically $\epsilon \ll 1$.

Note that although we have only gauged lepton number, this is enough to prevent the dimension-six operators of the form $\mathcal{L} \sim \frac{1}{\Lambda^2} qqq\ell$ (for appropriate lepton-number assignment to the lepton-breaking scalar) which might lead to proton decay. However, while baryon-number-violating operators at dimension six are forbidden, higher-order operators are still allowed since baryon number is not protected by a gauge symmetry. The leading operator that might mediate proton decay, $\frac{1}{\Lambda^8}(qqq\ell)(\ell H)^2\Phi^\dagger$, first occurs at dimension 12 while $\Delta B = 2$ operators with $\Delta L = 0$ are allowed at dimension nine [35], as in the SM. For scales $\Lambda \gtrsim \mathcal{O}(100)$ TeV the model considered here should be reasonably safe from the effects of these potentially dangerous operators. Of course one can extend this model to include gauged baryon number as well to prevent these operators [10].

Finally, it is worth mentioning that this model possesses many ingredients which may be helpful for explaining the baryon asymmetry of the universe. The current construction automatically contains new massive states as well as new interactions containing CP-violating phases. It would be interesting to explore whether or not it is capable of explaining this asymmetry as well as dark matter. Since the weakly interacting massive particle in this theory is a Dirac fermion, there is potential to realize a theory with asymmetric dark matter. We leave it to future studies to explore these possibilities.

III. DARK MATTER

Here we examine the DM candidate in this model. We first discuss the stability which results from an accidental global symmetry of the Lagrangian and identify the DM as a heavy mostly singlet neutrino. This global symmetry is a consequence of the particle content and underlying lepton gauge symmetry, much in the same way that lepton number is an accidental global symmetry in the SM. We then discuss the various annihilation channels and calculate the relic abundance of the DM candidate to establish the allowed masses. We also discuss various other phenomenological features.

A. DM candidate and stability

We begin by examining the neutrino sector once Φ and H obtain expectation values, which gives

$$\begin{aligned}\mathcal{L} \supset & -\frac{c_\ell v_\phi}{\sqrt{2}}\left(1+\frac{\phi_o}{v_\phi}\right)\bar{\nu}''_R\nu'_L - \frac{c_\nu v_\phi}{\sqrt{2}}\left(1+\frac{\phi_o}{v_\phi}\right)\bar{\nu}''_L\nu'_R \\ & -\frac{y''_\nu v_h}{\sqrt{2}}\left(1+\frac{h_o}{v_h}\right)\bar{\nu}''_R\nu''_L - \frac{y'_\nu v_h}{\sqrt{2}}\left(1+\frac{h_o}{v_h}\right)\bar{\nu}'_L\nu'_R + \mathrm{H.c.},\end{aligned} \tag{14}$$

leading to the mass matrix

$$\mathcal{M}_\nu = \frac{1}{\sqrt{2}}\begin{pmatrix} c_\ell v_\phi & y'_\nu v_h \\ y''_\nu v_h & c_\nu v_\phi \end{pmatrix}, \tag{15}$$

which can be diagonalized using the singular value decomposition $\mathcal{M}_{\nu D} = U_L^\dagger \mathcal{M}_\nu U_R$, where $\mathcal{M}_{\nu D}$ is a diagonal mass matrix with positive mass eigenvalues m_{ν_X} and m_{ν_4}.

While the Yukawa couplings to Φ and H break the global U(1) symmetries associated with L' and L'' explicitly, there is a residual Z_2 symmetry under which all heavy leptons are odd and all SM leptons are even, which is preserved after spontaneous breaking of the lepton number and electroweak gauge symmetries. Assuming that the new charged leptons are heavier, this residual global symmetry guarantees the stability of the lighter of the two neutrino mass eigenstates, opening up the possibility for dark matter.

In the limit where $y'_\nu v_h$, $y''_\nu v_h \ll c_{\ell,\nu} v_\phi$, the mass eigenvalues are approximately given by

$$m_{\nu_X} \approx \frac{1}{\sqrt{2}} c_\nu v_\phi, \qquad m_{\nu_4} \approx \frac{1}{\sqrt{2}} c_\ell v_\phi. \tag{16}$$

In this limit, the eigenstate ν_4 is mostly composed of the electroweak doublet neutrinos ν''_R and ν'_L, while ν_X is a combination of the singlets ν''_L and ν'_R and with tiny couplings to the SM $W^\pm$ and Z bosons. Since the doublet neutrino ν_4 couples directly to the Z boson, direct detection experiments render it unacceptable as a DM candidate. Therefore we require $c_\nu < c_\ell$, such that ν_X is the DM candidate. Of course ν_4 must be able to decay, which means that at least one of the Yukawa couplings y'_ν, y''_ν should be nonzero to allow ν_4 to decay into a Higgs boson and ν_X. Nonetheless, this requirement allows the y_ν's to be small enough so as to be completely irrelevant in the discussion below.

B. Annihilation channels

In Ref. [13], annihilation through the interactions generated by y'_ν, y''_ν was shown to give the correct relic abundance for DM with dominantly Majorana masses $\lesssim 100$ GeV. Here, because direct detection constraints require y'_ν, y''_ν to be tiny, one would have to either rely on coannihilation with one of the charged leptons or annihilation through a nearly on-shell Higgs. We instead will assume in the following that these couplings are too tiny to affect the DM phenomenology directly, although they do play a role in direct and indirect searches as well as LHC phenomenology, to be described below.

Compared to Ref. [13], there are additional annihilation channels for ν_X into SM leptons. In particular, since ν_X is a Dirac fermion, annihilation through a vector boson is s-wave and unsuppressed, in contrast to the case of Majorana DM. Indeed, the left- and right-handed components of ν_X carry lepton number L'' and L', respectively, and $L' - L'' = -3$ implies a nonvanishing coupling of ν_X to Z_L, allowing $\nu_X \bar{\nu}_X$ to annihilate into SM leptons through s-channel Z_L exchange, shown in the top diagram of Fig. 3. There are additional annihilation channels which arise through mixing in the neutrino as well as in the Higgs sectors. We discuss the various annihilation modes in more detail below, assuming that ν_X is mostly singlet with at most a small doublet component, i.e. $y'_\nu v_h$, $y''_\nu v_h \ll c_\nu v_\phi$.

If ν_X acquires a small doublet component through nonzero y'_ν, y''_ν couplings, annihilation into SM particles through Z or h exchange becomes possible, but again we will assume that these couplings are sufficiently small such that these annihilation channels can be neglected. This is also required since otherwise a large direct detection cross section through Z-boson exchange would be induced. At the same time this suppresses annihilation into W^+W^- through a heavy charged lepton exchanged in the t channel.

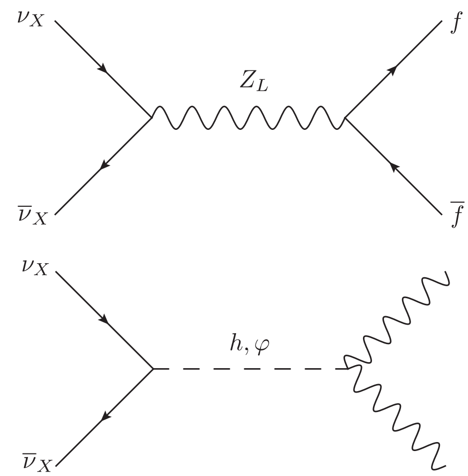

FIG. 3. Diagrams leading to s-channel $\nu_X \bar{\nu}_X$ annihilation into SM states through the exchange of Z_L, h, or ϕ.

The dark matter also couples to the singlet scalar ϕ_o with a strength $c_\nu \approx \sqrt{2} m_{\nu_X}/v_\phi$. When the Higgs mixing angle θ is nonzero this will allow annihilation into SM particles through s-channel exchange of h and ϕ, shown in the bottom diagram of Fig. 3. While not generally negligible, the contribution of these annihilation channels turns out to be suppressed compared to the Z_L channel in the regime of interest where $v_\phi \sim 1.7$ TeV and DM $m_{\nu_X} \sim v_h$, leading to somewhat small values for c_ν. Furthermore, the Z_L channel leads to unsuppressed annihilation into all SM leptons, while most of the scalar channels are suppressed by the small Yukawa couplings of the SM quarks and leptons. We thus expect annihilation through Z_L to be the dominant contribution to the relic abundance in this regime. Note also that in this regime we have $m_{\nu_X} \ll M_{Z_L}$, which as we will see leads to a relic abundance which is largely independent of the lepton gauge coupling g' [see Eq. (22)].

C. Relic abundance

Motivated by the requirement for small y'_ν, y''_ν, we first consider the dominant annihilation through the Z_L into SM lepton pairs, and then demonstrate that scalar exchange is unlikely to change the overall picture. The relevant interactions come from Eq. (4), which before lepton number and electroweak symmetry breaking can be written as

$$\mathcal{L} \supset g' Z_{L\mu} (L'' \bar{\nu}''_R \gamma^\mu \nu''_R + L' \bar{\nu}'_L \gamma^\mu \nu'_L + \bar{l} \gamma^\mu l), \tag{17}$$

where l runs over SM leptons, all of which have $L = 1$, which implies that the left- and right-handed couplings of the SM leptons to Z_L are equal. This is in contrast to the case for the exotic leptons since $L' \neq L''$. After lepton number breaking and rotating to the mass basis Eq. (17) becomes

$$\mathcal{L} \supset g' Z_{L\mu}(\bar{\nu}_X \gamma^\mu (L'' P_R + L' P_L)\nu_X + \bar{l}\gamma^\mu l), \quad (18)$$

where P_R and P_L are the right and left projection operators, respectively, and we have neglected any mixing between ν_X and ν_4 generated by y'_ν, y''_ν. Using Eq. (18), a straightforward calculation of the diagram in Fig. 3 gives the annihilation cross section,

$$\sigma = \frac{g'^4((L'^2 + L''^2)(s - m_{\nu_X}^2) + 6L'L'' m_{\nu_X}^2)}{8\pi(1 - 4m_{\nu_X}^2/s)^{1/2}((M_{Z_L}^2 - s)^2 + M_{Z_L}^2\Gamma_{Z_L}^2)}, \quad (19)$$

where an overall factor of 6 is implicit for the three generations of charged leptons and neutrinos in the SM. As is well known, the annihilation cross section $\langle \sigma v \rangle$ is well approximated by a nonrelativistic expansion, $s = 4m_{\nu_X}^2 + m_{\nu_X}^2 v^2$, and by expanding the annihilation cross section in powers of v to give $\langle \sigma v \rangle = a + b\langle v^2 \rangle + \mathcal{O}(\langle v^4 \rangle)$ [36]. Expanding Eq. (19), we obtain

$$a = \frac{3g'^4 R^4 (L' + L'')^2}{4\pi m_{\nu_X}^2 (1 - 4R^2)^2} \quad (20)$$

for the velocity-independent coefficient. Note that this is in contrast to the case of Majorana dark matter annihilating through a gauge boson, in which case $a = 0$ up to corrections that are suppressed by the final-state fermion masses. For the $\langle v^2 \rangle$ coefficient we have

$$b = \frac{g'^4 R^4 ((L'^2 + L''^2)(11 + 4R^2) + L'L''(6 + 72R^2))}{32\pi m_{\nu_X}^2 (1 - 4R^2)^3}. \quad (21)$$

Here we have defined $R = m_{\nu_X}/M_{Z_L}$ and neglected terms of order Γ_{Z_L}/M_{Z_L}. In general the contribution from a will dominate since the contribution from b is suppressed by the relatively small value of v^2 at freeze-out. It is useful to consider the limit of heavy Z_L mass compared to the DM mass, or $R \ll 1$. Keeping only the leading term after expanding in powers of R, we have

$$a \approx \frac{3g'^4 (L' + L'')^2 R^4}{4\pi m_{\nu_X}^2} + \mathcal{O}(R^6). \quad (22)$$

Since $M_{Z_L} = 3g' v_\phi$, the dependence on the gauge coupling g' cancels in the leading term, as is usual for the contact interaction that describes vector exchange at low energies. For a fixed choice of the quantum numbers L' and L'', the annihilation rate is therefore largely determined by the ratio $m_{\nu_X}^2/v_\phi^4$.

From these results a good approximation for the relic density can be obtained by e.g. using the procedure presented in Ref. [36]. We have opted instead to implement the model into the numerical code MICROMEGAS [37]. Not only does this facilitate the exploration of regions of parameter space where the $\mathcal{O}(v^2)$ expansion breaks down, but it also simplifies the computation of direct and indirect detection signals. The approximate calculation of the relic density following

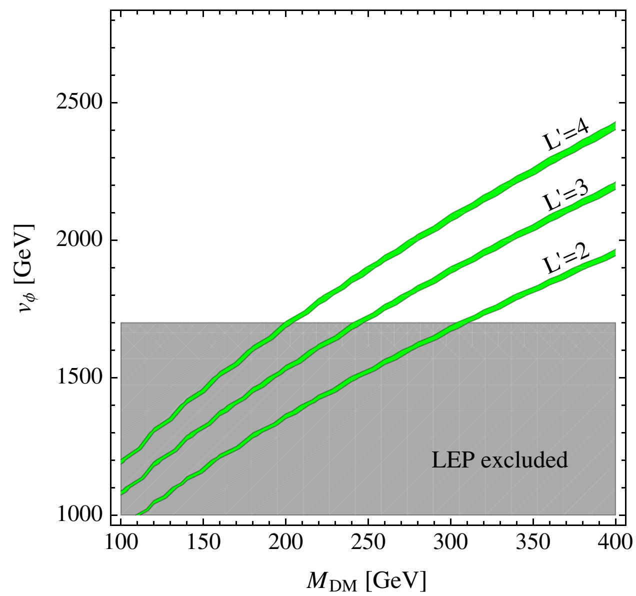

FIG. 4 (color online). Relic density as a function of DM mass and VEV v_ϕ, in the absence of mixing and taking $M_{\mathrm{DM}} = m_{\nu_X}$. The green bands indicate regions in agreement with the measured value of $\Omega h^2 = 0.120 \pm 0.003$ [56] for different choices of L', as indicated in the figure.

Ref. [36] was used as validation of the MICROMEGAS implementation of the model. The resulting relic density (including all sub-leading effects) is shown as a function of m_{ν_X} and v_ϕ, for a few choices of L', in Fig. 4. The LEP II constraints on v_ϕ require dark matter masses greater than about 200 GeV, and (depending on L') a thermal relic density enforces a tight correlation between v_ϕ and m_{ν_X}.

In the limit y'_ν, $y''_\nu \approx 0$, DM couples to h and ϕ through c_ν and the Higgs mixing,

$$\mathcal{L} \supset \frac{c_\nu}{\sqrt{2}}(c_\theta \phi - s_\theta h)\bar{\nu}_X \nu_X, \quad (23)$$

where we have used Eq. (8). These couplings allow the DM to annihilate through the bottom diagram shown in Fig. 3. Since dark matter masses of order the weak scale require a relatively small c_ν, annihilation through Higgs exchange only has a small effect on the relic density. On the other hand it is crucial for direct detection, which will be discussed in the next section.

D. Direct and indirect detection

In the limit y'_ν, $y''_\nu \to 0$ and with negligible mixing in the Higgs sector, the dark matter couples to SM leptons through Z_L, but has no tree-level interactions with quarks. This is a challenging situation for dark matter direct detection experiments, because of the wave-function suppression to scatter off of atomic electrons or loop suppression of the induced dark matter dipole moment [38]. Consequently, even a small amount of $Z - Z_L$ or $H - \Phi$ mixing can dominate the rate, which effectively disconnects the expectations at direct detection experiments from the relic density.

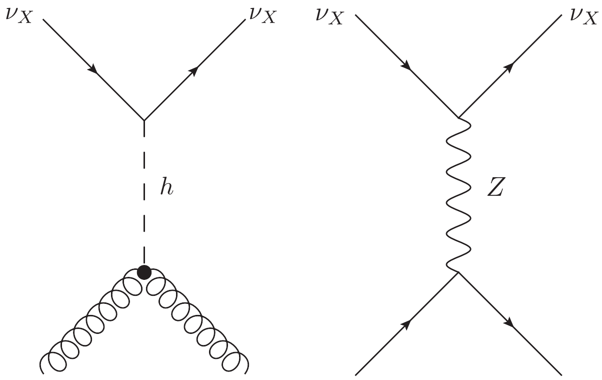

FIG. 5. Diagrams leading to scattering with nucleons mediated by the exchange of a Higgs or Z boson.

Higgs exchange, as shown in Fig. 5, leads to spin-independent scattering with nuclei. We compute the rate as a function of the DM mass and Higgs mixing angle $\sin\theta$ using MICROMEGAS and present the results in Fig. 6 for DM masses 100–400 GeV. For moderate Higgs mixing, The DM-nucleon cross section lies about one order of magnitude below the current best limit from the XENON100 experiment, but is well within reach of second generation DM direct detection experiments such as LZ [39].

Z-boson exchange, as shown in Fig. 5, induces a large DM-neutron cross section due to the sizable coupling of the Z to light quarks. We parametrize the coupling of the Z boson to the DM as

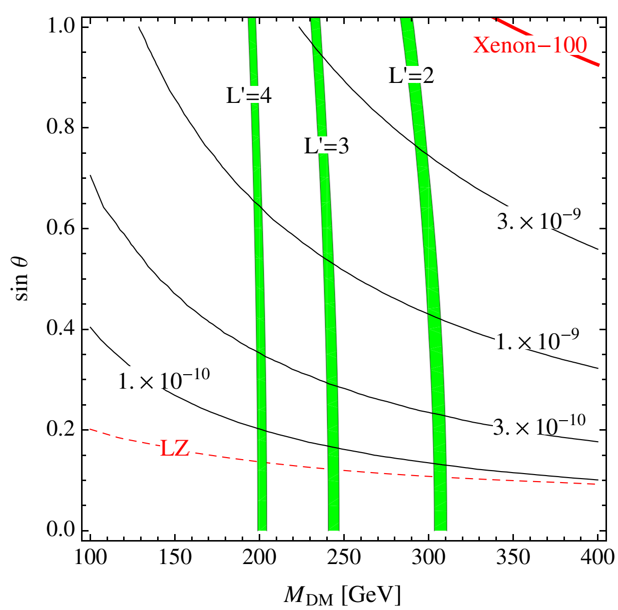

FIG. 6 (color online). DM-nucleon cross section in pb, as a function of the Higgs mixing angle $\sin\theta$ and of the DM mass, for $v_\phi = 1.7$ TeV and $M_{\rm DM} = m_{\nu_X}$. The solid red line indicates the current limit from the XENON100 experiment [29], while the dashed red line indicates the projected reach of the LZ experiment [39]. The green bands indicate regions with correct relic density for different values of L'.

$$\mathcal{L} \supset \epsilon' g' Z_\mu \bar{\nu}_X \gamma^\mu (L'' P_R + L' P_L)\nu_X, \tag{24}$$

where ϵ' is either induced by $Z - Z'$ mixing or by nonzero neutrino Yukawa couplings y'_ν, y''_ν. The upper bound on ϵ' from direct detection for $L' = 2$ is shown in Fig. 7, for DM masses 100–400 GeV. One can see that for $g' = 0.5$ and $v_\phi = 1.7$ TeV, direct detection requires roughly $\epsilon' \lesssim 1 - 2 \times 10^{-4}$ depending on the DM mass. In the limit y'_ν, $y''_\nu \approx 0$, ϵ' is due solely to $Z - Z_L$ mixing and gives $\epsilon' = \sin\xi$ as defined in Eq. (5). Since $M_{Z_L} = 3g'v_\phi = 2.55$ TeV, Eq. (5) and Fig. 2 together imply that direct detection signals roughly 20 times below the current bound can be obtained for a gauge kinetic mixing parameter [see Eq. (4)] of $\epsilon \sim 7 \times 10^{-3}$, within range of future direct detection experiments [39].

Dark matter can also be observed indirectly, by searching for the products of DM annihilation. Here, the dark matter annihilates predominantly into charged leptons or neutrinos. While there is a large rate into positrons, it is characterized by roughly the thermal relic cross section and is thus quite a bit too small to account for the anomalous positron fraction observed by PAMELA [40], Fermi [41], and AMS-02 [42]. At the same time, contributions to the antiproton flux are very tiny, evading constraints from PAMELA [43].

Annihilation into charged leptons will also produce gamma rays as secondaries. Currently, the tightest constraints on such production are from the Fermi LAT null observations of dwarf spheroidal galaxies [44], which are just short of being able to rule out thermal cross sections

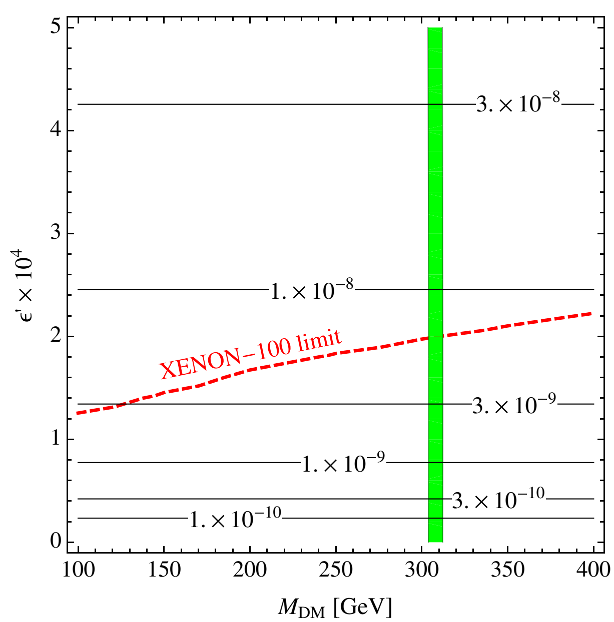

FIG. 7 (color online). DM-nucleon cross section in pb, as a function of the DM-Z coupling parameter ϵ' and of the DM mass (where $M_{\rm DM} = m_{\nu_X}$), for $v_\phi = 1.7$ TeV, $L' = 2$ and $g' = 0.5$, which implies $M_{Z_L} = 2.55$ TeV. The red dashed line indicates the current limit from the XENON100 experiment.

for dark matter masses around a few 10's of GeV based on one sixth of the annihilations producing $\tau^+\tau^-$. In the near future, such constraints are only relevant for ν_X dark matter which has been produced nonthermally.

Dark matter may also annihilate directly into $\gamma\gamma$ and/or γZ at loop level, providing monochromatic gamma-ray lines, whose distinctive energy profile can help compensate for a tiny rate. Predictions for the class of models including $U(1)_L$ were studied in Ref. [45], where it was found that $\gamma\gamma$, γZ, and $\gamma\phi$ (if kinematically accessible) final states can be generated. The largest signal is likely to be $\gamma\phi$, which is expected to be at least an order of magnitude below the current Fermi bounds [46], but may be visible to future experiments.

The rate for dark matter to be captured in the Sun or Earth and then annihilate into high-energy neutrinos is controlled by the spin-dependent cross section which in turn is controlled by the degree of $Z - Z_L$ mixing. Thus, despite a large annihilation fraction into SM neutrinos, the precision constraints render it difficult to imagine an observable rate at ICECUBE in the near future [47].

IV. LHC PHENOMENOLOGY AND CONSTRAINTS

The presence of new particles required by the $U(1)_L$ gauge symmetry leads to a variety of potentially interesting LHC phenomenology. In this section we discuss various aspects of the phenomenology of this model as well as the relevant constraints coming from the LHC. We also examine in more detail the charged lepton sector and its effects on the Higgs decays.

A. Exotic charged lepton sector

Once Φ and H obtain expectation values, the Lagrangian for the exotic charged lepton sector becomes

$$\mathcal{L} \supset -\frac{c_\ell v_\phi}{\sqrt{2}}\left(1+\frac{\phi_o}{v_\phi}\right)\bar{e}''_R e'_L - \frac{c_e v_\phi}{\sqrt{2}}\left(1+\frac{\phi_o}{v_\phi}\right)\bar{e}''_L e'_R - \frac{y''_e v_h}{\sqrt{2}}\left(1+\frac{h_o}{v_h}\right)\bar{e}''_R e''_L - \frac{y'_e v_h}{\sqrt{2}}\left(1+\frac{h_o}{v_h}\right)\bar{e}'_L e'_R + \text{H.c.}, \tag{25}$$

which gives a mass matrix of the same form as that found in the neutrino sector,

$$\mathcal{M}_e = \frac{1}{\sqrt{2}}\begin{pmatrix} c_\ell v_\phi & y''_e v_h \\ y'_e v_h & c_e v_\phi \end{pmatrix}. \tag{26}$$

Again, we can diagonalize via $\mathcal{M}_{eD} = U_L^\dagger \mathcal{M}_e U_R$ to obtain the mass eigenvalues and eigenstates. The Lagrangian in Eq. (25) also leads to the interaction matrices for ϕ_o and h_o, given by

$$\mathcal{N}_e^h = \frac{v_h}{\sqrt{2}}\begin{pmatrix} 0 & y''_e \\ y'_e & 0 \end{pmatrix}, \qquad \mathcal{N}_e^\phi = \frac{v_\phi}{\sqrt{2}}\begin{pmatrix} c_\ell & 0 \\ 0 & c_e \end{pmatrix}, \tag{27}$$

which upon the rotation performed to diagonalize $\mathcal{M}_e$ gives interaction matrices in the mass basis defined as $\mathcal{V}_\phi = U_L^\dagger \mathcal{N}_e^\phi U_R$ and $\mathcal{V}_h = U_L^\dagger \mathcal{N}_e^h U_R$. These matrices dictate the couplings of the exotic leptons to ϕ and h. We note also that Eq. (26) is the same mass matrix in the charged lepton sector considered in Ref. [13], with the difference being that in this model there are no explicit mass terms. In particular, when $v_h, v_\phi \to 0$ all masses go to zero, which makes the gauged-lepton-number model more constrained and relates the electroweak- and lepton-breaking scales to the rate of Higgs decay to diphotons, as we will see below.

A useful simplifying limit is $c_\ell \approx c_e \equiv c_e$ and $y'_e \approx y''_e \equiv y_e$, in which case the charged leptons are maximally mixed and one obtains the simple relations for the mass eigenvalues

$$m_{e_1} \approx \frac{1}{\sqrt{2}}(c_e v_\phi - y_e v_h), \qquad m_{e_2} \approx \frac{1}{\sqrt{2}}(c_e v_\phi + y_e v_h), \tag{28}$$

where we have assumed $c_e v_\phi > y_e v_h$. Thus we see that for fixed y_e and v_ϕ, the mass of the charged leptons is controlled by c_e. Along with the scalar mixing discussed in Sec. II B we now have the pieces necessary for examining the modification to Higgs decays.

B. Modifications of Higgs decays

Assuming that the Higgs cannot decay directly into new particles, the primary effect of the new lepton sector on Higgs decays will be through loop effects. From the discussion of Higgs mixing in Sec. II B, we can write the modification of the SM Higgs partial width as

$$\epsilon_i \equiv \frac{\Gamma_{hi}}{\Gamma^{\rm SM}_{h_o i}} = \frac{|\mathcal{M}(h\to i)|^2}{|\mathcal{M}(h_o \underset{\rm SM}{\to} i)|^2} = \frac{c_\theta^2|\mathcal{M}(h_o\to i) - t_\theta \mathcal{M}(\phi_o\to i)|^2}{|\mathcal{M}(h_o \underset{\rm SM}{\to} i)|^2}, \tag{29}$$

where we have used Eq. (8) and $\Gamma^{\rm SM}_{h_o i}$ is the SM partial width to a final state i and Γ_{hi} is the partial width for h to decay into i. The rate expected at the LHC relative to the SM can be written as

$$\mu_i = \frac{\sigma(j\to h)}{\sigma(j \underset{\rm SM}{\to} h_o)}\frac{\mathcal{B}(h\to i)}{\mathcal{B}(h_o \underset{\rm SM}{\to} i)} = \epsilon_j \frac{\Gamma^{\rm SM}_{h_o}}{\Gamma_h}\epsilon_i, \tag{30}$$

where we have made use of the narrow-width approximation, $\mathcal{B}$ signifies the branching fraction, and the production channels are labeled $j = VV, gg$. We also define $\Gamma^{\rm SM}_{h_o}$ as the total SM Higgs width and Γ_h as the total decay width for the mass eigenstate h. Since this model does not contain any new colored particles the only new effects entering ϵ_{gg} are through Higgs mixing, which gives $\epsilon_{gg} \approx c_\theta^2$. Since ZZ

and WW already occur at tree level in the SM, we assume the loop corrections due to the new leptons are negligible, which implies that the only effect again comes from Higgs mixing, which gives $\epsilon_{ZZ} = \epsilon_{WW} \sim c_\theta^2$. Similarly for the SM Higgs Yukawa interactions we have $\epsilon_Y \sim c_\theta^2$.

This leaves the $Z\gamma$ and $\gamma\gamma$ channels, which first occur at one loop in the SM, as the most promising possibilities for these effects to manifest themselves. However, in Refs. [13,48] the modification to $Z\gamma$ was shown to be only $\sim 5\%$ for a corresponding $\gamma\gamma$ enhancement of $\sim 50\%$, and to good approximation $\epsilon_{Z\gamma} \sim c_\theta^2$. Thus, in addition to the universal c_θ^2 suppression from Higgs mixing, the only additional modifications to the total decay width comes from the $\gamma\gamma$ channel through loops of exotic charged leptons. Since for the modifications we are interested in $\Gamma_{h\gamma\gamma} \ll \Gamma_h$, this implies $\Gamma_{h_o}^{\rm SM}/\Gamma_h \approx c_\theta^{-2}$ which will cancel with the c_θ^2 in the production channels $\epsilon_{gg,VV}$. Finally, this gives for the relative rates $\mu_i = c_\theta^2$ for $i \neq \gamma\gamma$ and for the final modified diphoton signal strength

$$\mu_{\gamma\gamma} = \epsilon_{\gamma\gamma}. \tag{31}$$

Using the approach and conventions of Ref. [49], which examined the similar $gg \to h$ process, we can go on to obtain the exotic charged lepton contributions to the $h \to \gamma\gamma$ amplitudes (omitting photon polarization vectors),

$$\begin{aligned}
&\mathcal{M}^{\mu\nu}(h_o \to \gamma\gamma) \\
&\quad = \left(\frac{\alpha}{2\pi v_h}\right)\sum_i \frac{(\mathcal{V}_h)_{ii} F_F(\tau_{e_i})}{m_{e_i}}\left(p_1^\nu p_2^\mu - \frac{m_h^2}{2} g^{\mu\nu}\right), \\
&\mathcal{M}^{\mu\nu}(\phi_o \to \gamma\gamma) \\
&\quad = \left(\frac{\alpha}{2\pi v_\phi}\right)\sum_i \frac{(\mathcal{V}_\phi)_{ii} F_F(\tau_{e_i})}{m_{e_i}}\left(p_1^\nu p_2^\mu - \frac{m_h^2}{2} g^{\mu\nu}\right),
\end{aligned} \tag{32}$$

where the index $i = 1, 2$ runs over the exotic charged lepton mass eigenstates found after diagonalizing the mass matrix in Eq. (26), and F_F are the fermonic loop functions with $\tau_{e_i} = m_h^2/4m_{e_i}^2$ as defined in Ref. [49]. Note that the amplitudes in Eq. (32) are evaluated at $m_{h_o} = m_h$ and $m_{\phi_o} = m_h$, where m_h is the physical scalar mass.

Using Eqs. (29)–(32), we obtain

$$\begin{aligned}
\mu_{\gamma\gamma} &= \frac{\left|\frac{c_\theta}{v_h}\left(F_{\rm SM} + \sum_i \frac{(\mathcal{V}_h)_{ii}}{m_{e_i}} F_F(\tau_{e_i})\right) - \frac{s_\theta}{v_\phi}\left(\sum_i \frac{(\mathcal{V}_\phi)_{ii}}{m_{e_i}} F_F(\tau_{e_i})\right)\right|^2}{|F_{\rm SM}/v_h|^2} \\
&= c_\theta^2 \left|\left(1 + F_{\rm SM}^{-1} \sum_i \frac{(\mathcal{V}_h)_{ii}}{m_{e_i}} F_F(\tau_{e_i})\right)\right. \\
&\quad \left. - t_\theta\left(F_{\rm SM}^{-1} \frac{v_h}{v_\phi} \sum_i \frac{(\mathcal{V}_\phi)_{ii}}{m_{e_i}} F_F(\tau_{e_i})\right)\right|^2,
\end{aligned} \tag{33}$$

where $F_{\rm SM}$ is the SM loop function which includes the dominant and negative $W^\pm$-boson contribution as well as the smaller and positive t quark, which sum to give a numerical value of ~ -6.5 for $m_h = 125$ GeV. Note that only the diagonal entries in the interaction matrices $(\mathcal{V}_h)_{ii}$ and $(\mathcal{V}_\phi)_{ii}$ contribute in the $h \to \gamma\gamma$ loop.

After the approximations leading to the masses in Eq. (28), which give $(\mathcal{V}_\phi)_{11} = (\mathcal{V}_\phi)_{22} \approx c_e v_\phi/\sqrt{2}$ and $(\mathcal{V}_h)_{11} = -(\mathcal{V}_h)_{22} \approx -y_e v_h/\sqrt{2}$, we obtain (approximately) for the modified signal strength

$$\begin{aligned}
\mu_{\gamma\gamma} \simeq c_\theta^2 &\left| 1 - \frac{v_h}{\sqrt{2}F_{\rm SM}}\left[y_e\left(\frac{F_F(\tau_{e_1})}{m_{e_1}} - \frac{F_F(\tau_{e_2})}{m_{e_2}}\right)\right.\right. \\
&\left.\left. + c_e t_\theta\left(\frac{F_F(\tau_{e_1})}{m_{e_1}} + \frac{F_F(\tau_{e_2})}{m_{e_2}}\right)\right]\right|^2,
\end{aligned} \tag{34}$$

where m_{e_1,e_2} are given in Eq. (28) and satisfy $m_{e_1} < m_{e_2}$. Remembering that $F_{\rm SM} < 0$ we see in the limit $t_\theta \to 0$ that we have an enhancement in the diphoton rate in the presence of mostly vectorlike leptons entering through the h_o component of h. This is, of course, expected from the low-energy Higgs theorems (see e.g. Ref. [13]). We see also that the contribution from Higgs mixing is constructive for $t_\theta > 0$ and destructive for $t_\theta < 0$, which also corresponds to the sign of the coupling $\lambda_{\rm hp}$ in Eq. (6). In the limit $y_e \to 0$ the enhancement enters entirely through Higgs mixing and thus requires large mixing angles and Yukawa coupling c_e. In the realistic limit $v_\phi \gg v_h$, the e_1 and e_2 become almost purely vectorlike and again the contribution only enters through Higgs mixing via the ϕ_o component of h. However, as $v_\phi \to \infty$ one also has $t_\theta \to 0$ and the ϕ_o contribution eventually decouples from the $h \to \gamma\gamma$ amplitude as v_ϕ is taken large. Equation (34) is in agreement with Ref. [15] for the case where their explicit mass term is put to zero.

To avoid the constraints discussed in Sec. II E we choose $v_\phi = 1.7$ TeV and take the lightest charged lepton to have mass greater than $m_{\rm min} \sim 100$ GeV. Measurements of the Higgs decays at the LHC indicate rates consistent with the SM with the possibility of a slight, though not significant, enhancement in the diphoton channel [50]. Regardless, this implies that these fermions must be mostly "vectorlike" since otherwise their effects would lead to destructive interference [13] with the SM contribution, giving a reduced rate, which is disfavored. This allows us to write

$$m_{e_1} = \frac{c_e v_\phi - y_e v_h}{\sqrt{2}} \gtrsim m_{\rm min}, \tag{35}$$

which leads to a condition on the Yukawa coupling,

$$\frac{\sqrt{2}m_{\rm min} + y_e v_h}{v_\phi} \lesssim c_e \lesssim 4\pi, \tag{36}$$

where we have also indicated 4π as the perturbative upper bound.

Since the mixing angle will affect all decay channels, we perform a fit to the full Higgs data set in the $c_e - \theta$ plane for fixed $y_e = 0.8$ and $v_\phi = 1.7$ TeV. We show in Fig. 8 the 1, 2, and 3σ regions (purple) for the favored parameter

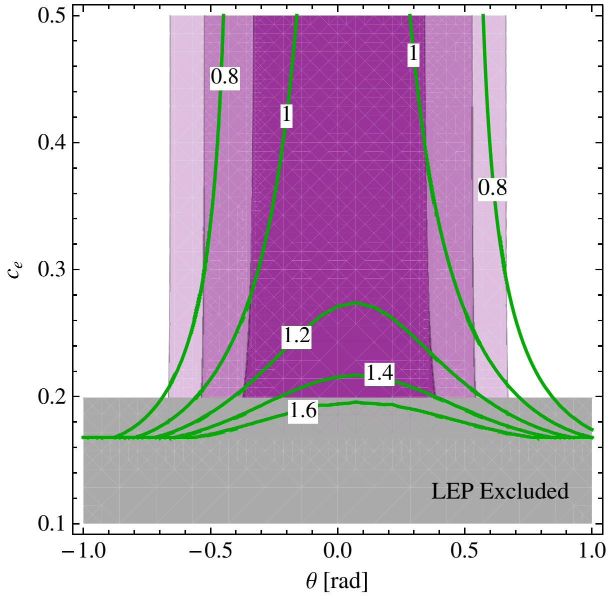

FIG. 8 (color online). Fits to the full Higgs data set in the $c_e - \theta$ plane for $y_e = 0.8$ and $v_\phi = 1.7$ TeV. Here the purple contours show the 1, 2, and 3σ regions while the grey band shows the LEP excluded region and the green lines are contours of constant $\mu_{\gamma\gamma}$. Details on the fitting procedure can be found in Ref. [57].

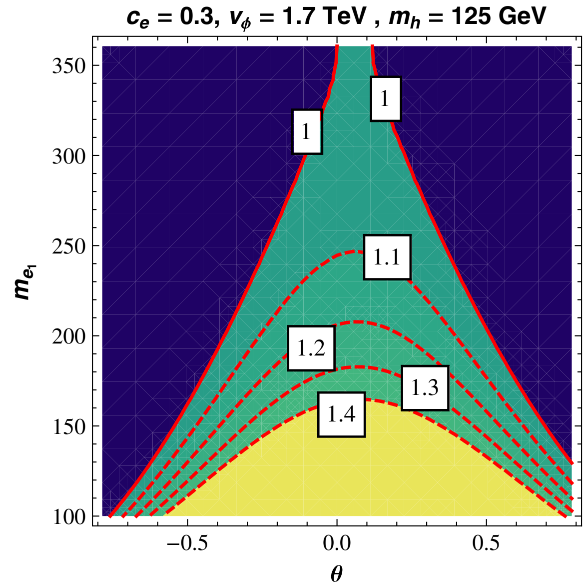

FIG. 9 (color online). Contours of the relative diphoton rate as a function of the Higgs mixing angle θ and lightest exotic charged lepton mass m_{e_1}.

space, where the grey band shows the excluded region by LEP II for which $m_{e_1} < 100$ GeV. Values as large as $\theta \sim \pm 0.5$ give a good fit to the Higgs data, while larger values are disfavored due to the $\cos\theta$ suppression of the signal rates. We also show contours of the relative diphoton rate shown in the green curves, though it is also worth noting that with the current data, the diphoton rate has no significant impact on the quality of the fit. Negative values of the mixing angle correspond to $\lambda_{\rm hp} < 0$, which can potentially lead to vacuum instabilities. On the other hand, positive values of $\theta \sim 0.5$ where $\lambda_{\rm hp} > 0$ lead to no instability and, as shown in Ref. [15], can be made consistent with constraints coming from the S and T parameters.

Choosing instead to fix $c_e = 0.3$ and trading in y_e for the lightest charged lepton mass, we can examine contours of $\mu_{\gamma\gamma}$ as a function of m_{e_1} and θ, as seen in Fig. 9. Since the DM mass serves as a lower bound on the charged lepton mass we see for the DM masses $\gtrsim 200$ GeV found in Sec. III that modifications up to ~10–20% can be obtained for $\theta \sim 0.3$–0.4 and $m_{e_1} \gtrsim 200$ GeV. Of course one can lower this bound by considering larger values of L', as can be seen in Fig. 4, or by tuning the Z_L mass such that the DM annihilation is resonantly enhanced.

Allowing c_e and y_e to vary instead while fixing $\theta = 0.4$ and $v_\phi = 1.7$ TeV, we show $\mu_{\gamma\gamma}$ contours in the $c_e - y_e$ plane in Fig. 10. As can be seen, observable modifications can be obtained for $\mathcal{O}(1)$ values of the Yukawa couplings for which vacuum stability issues can be avoided [15]. For these ranges of Yukawa couplings, m_{e_1} lies in the range 100–500 GeV, such that the exotic leptons can be produced at the LHC. We will discuss possible collider signatures below.

If one is willing to push the Yukawa couplings as large as the perturbative limit $\sim 4\pi$, one can realize large deviations in $\mu_{\gamma\gamma}$ even for multi-TeV masses. In Fig. 11, we show the deviation in the plane of $m_{e_1} - m_{e_2}$ for fixed $v_\phi = 1.7$ TeV, right above the LEP II limit. Even for a lightest exotic charged lepton with mass $m_{e_1} \sim 2$–3 TeV, one can obtain appreciable modifications to the Higgs diphoton rate, reflecting the fact that the fermion masses

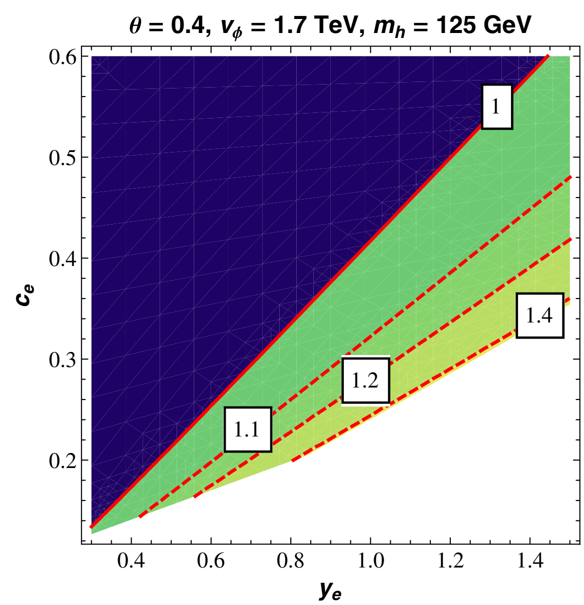

FIG. 10 (color online). Contours of the relative diphoton rate as a function of exotic charged lepton Yukawa couplings.

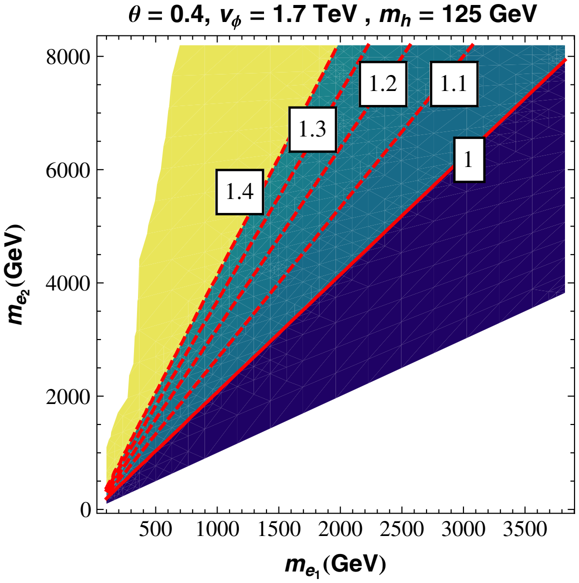

FIG. 11 (color online). Contours of the relative diphoton rate as a function of exotic charged lepton masses. Here we allow the masses to be as large as allowed by perturbativity and $v_\phi = 1.7$ TeV.

here are purely the result of Yukawa couplings, and thus do not exhibit decoupling [51]. Of course, all exotic contributions to the $h \to \gamma\gamma$ amplitude decouple in the limit of $v_\phi \to \infty$. It should also be noted that the required large Yukawa couplings can induce vacuum instabilities in the Higgs potential at scales close to the masses of the exotic leptons. Additional structures like supersymmetry would be required to restore vacuum stability. Some work in this direction recently appeared in Refs. [18,19,52].

C. Other potential LHC signatures

Since the LHC is a hadron machine, weakly coupled extensions of the SM such as the model presented here are not heavily constrained by the current LHC data. Currently, constraints on the masses of the new leptons and of Z_L mostly derive from the LEP experiments. Exotic charged leptons must be heavier than about 100 GeV for consistency with direct search limits. The Z_L mass should be larger than the LEP-2 center-of-mass energy of 209 GeV, and furthermore its coupling is subject to the constraint $M_{Z_L} = 3g' v_\phi$, where $v_\phi \geq 1.7$ TeV (and we have neglected any kinetic mixing with the Z boson).

One of the defining features of our model is Z_L, the gauge boson of the lepton number symmetry. Since it does not couple to quarks, it is difficult to produce at the LHC. The most promising option is to radiate a Z_L from a pair of Drell-Yan-produced leptons, in the process $pp \to \ell^+\ell^- Z_L$. The cross section for this process is calculated using the program CALCHEP [53] with the MRST2002 PDF set [54] and is shown in Fig. 12, where one can see it is at most of order 10^{-2} fb at the 14 TeV LHC. As long as the

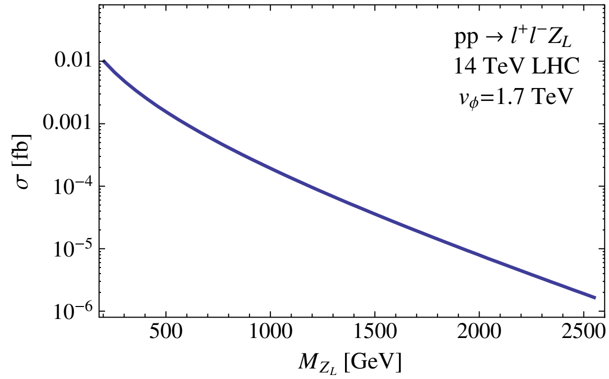

FIG. 12 (color online). Cross section for the process $pp \to \ell^+\ell^- Z_L$ at the 14 TeV LHC, for $v_\phi = 1.7$ TeV, and summed over SM leptons, $\ell^\pm = e^\pm, \mu^\pm, \tau^\pm$.

new leptons are heavier than half the Z_L mass, the gauge boson will decay into charged SM leptons with a branching ratio of 50%, while the other 50% are into neutrinos (recalling that there are three light ν_{Ri} in this model). The final state with four charged leptons, two of which reconstruct the Z_L mass, is essentially background free. Nevertheless even at a possible high-luminosity upgrade of the LHC with 3 ab^{-1} it will be difficult to probe Z_L masses above 500 GeV.

Pairs of charged and neutral leptons can be pair produced at the LHC in the Drell-Yan process. The cross sections for the different processes at the 14 TeV LHC are shown in Fig. 13, and were again obtained using CALCHEP. The processes are similar to chargino/neutralino pair production, for which next-to-leading-order corrections are moderate [55]. For this plot we have assumed that the lepton masses are given by Eq. (16) and (28). This leads to the following mass hierarchies for the exotic lepton sector:

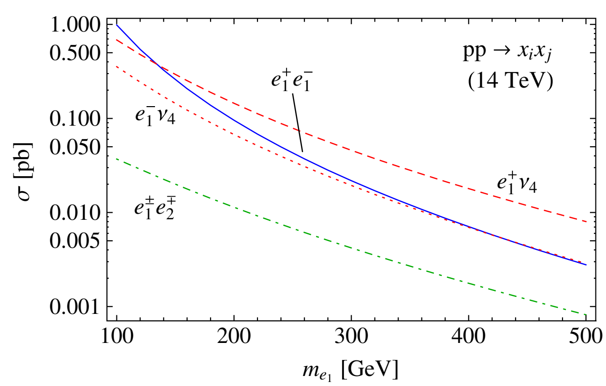

FIG. 13 (color online). Cross sections for the pair production of exotic leptons at the 14 TeV LHC, as a function of the lightest charged lepton mass m_{e_1} in the limit leading to Eqs. (16) and (28). For the processes involving $e_2^\pm$ and ν_4 we have assumed that $m_{e_2} = m_{e_1} + 280$ GeV, which implies $m_{\nu_4} = m_{e_1} + 140$ GeV.

$$m_{e_2} > m_{\nu_4} > m_{e_1} > m_{\nu_X}. \tag{37}$$

In this limit the mass splitting between e_1 and e_2 is given by $m_{e_2} - m_{e_1} = \sqrt{2} y_e v_h$, while $m_{\nu_4} - m_{e_1} = \frac{1}{\sqrt{2}} y_e v_h$. For $y_e \sim 0.8$ this gives a mass splitting of $\sim$280 GeV between the charged leptons and a splitting of $\sim$140 GeV between e_1 and ν_4. Note also that for $y_e \sim 0.8$ and the m_{e_1} range 100–500 GeV shown in Fig. 13 one also has $0.2 \lesssim c_e \lesssim 0.53$. The cross sections can be as large as one pb for particle masses close to the LEP limits, and up to 50 fb for particle masses in the several hundred GeV range.

The decays of the exotic leptons will lead to a number of signatures at the LHC via their decays to electroweak gauge and Higgs bosons as well as DM. In the limits leading to Eqs. (16) and (28) the heavy charged state e_2 can have the following decay chain:

$$e_2 \to W\nu_4 \to WWe_1 \to WWW\nu_X. \tag{38}$$

Note that although we are neglecting mass mixing between ν_X and ν_4 by assuming $y_\nu \ll 1$, it must be nonzero for the heavy leptons to decay down to the DM.

One can also have the heavy charged state decaying to DM more directly via

$$e_2 \to Wh\nu_X, \qquad e_2 \to WZ\nu_X, \qquad e_2 \to W\nu_X, \tag{39}$$

while the light charged state only has one tree-level decay,

$$e_1 \to W\nu_X. \tag{40}$$

The heavy neutrino state ν_4 can decay via Z and h bosons through

$$\nu_4 \to Z\nu_X, \qquad \nu_4 \to h\nu_X, \tag{41}$$

as well as W bosons through

$$\nu_4 \to We_1 \to WW\nu_X. \tag{42}$$

Thanks to the large mass differences between the particles, all intermediate gauge bosons are on-shell, such that their final states can easily be reconstructed at the LHC. These decay patterns can change in more general lepton-mixing scenarios, but should offer promising channels at the LHC.

For low masses, we see from Fig. 13 that $e_1^+ e_1^-$ has the largest production rate. Assuming leptonic decays of the W bosons, this leads to a signature

$$pp \to e_1^+ e_1^- \to WW\not{E}_T \to l^+ l^- \not{E}_T. \tag{43}$$

For larger masses the $e_1^+ \nu_4$ channel becomes dominant, and can give rise to a striking trilepton signature through

$$pp \to e_1^+ \nu_4 \to WZ\not{E}_T \to l^+ l^+ l^- \not{E}_T. \tag{44}$$

The signatures are similar to those from the production of weakly charged supersymmetric particles at the LHC. While limits can be obtained in special cases from the 8 TeV run of the LHC, we expect that at least 100 fb^{-1} at the 14 TeV LHC are needed to probe the exotic lepton sector at the LHC.

For a light enough ϕ there is also the potential to produce it resonantly at the LHC through Higgs mixing. This scalar would inherit the SM Higgs decays, but be suppressed by s_θ^2. Additionally, a kinematically allowed ϕ can also have the following decays to heavy leptons and dark matter:

$$\begin{aligned} \phi &\to e_1 e_1, & \phi &\to e_2 e_2, \\ \phi &\to e_1 e_2, & \phi &\to \nu_4 \nu_X. \end{aligned} \tag{45}$$

If so, it can of course also decay to Higgs pairs $\phi \to hh$ when kinematically allowed. As discussed in Sec. II E, however, for $v_\phi \sim 1.7$ TeV we typically have ϕ in the TeV range (see Fig. 1), making it phenomenologically irrelevant for much of the parameter space.

V. CONCLUSIONS/OUTLOOK

We have constructed a theory based on the gauging of lepton number, and found that for many choices of the parameters, the exotic leptons required to cancel gauge anomalies contain a dark matter candidate whose thermal relic density naturally saturates the requirements of cosmological observation. The dark matter is a Dirac (mostly singlet) neutrino and we find that masses $\gtrsim$200 GeV give the correct thermal relic abundance via annihilation through the massive vector boson associated with the gauged lepton number. Higgs scalar mixing as well as gauge kinetic mixing, which are found in this model, also allow for a direct detection signal and give reasonably good prospects for detection in near future experiments.

The theory introduces only one new scale, the vacuum expectation value of a SM singlet scalar which breaks the lepton number and is constrained by experiment to be $\gtrsim$1.7 TeV. The global symmetry which stabilizes the dark matter is a consequence of the gauge structure and particle content of the theory and does not need to be additionally imposed. Furthermore, as a consequence of the lepton number breaking, the dark matter is also accompanied by a set of vectorlike leptons charged under the SM gauge group with couplings to the SM Higgs. The same global symmetry which stabilizes the dark matter also prevents any dangerous flavor-changing neutral currents or mass mixing with SM leptons. For a lepton-breaking scale $\sim$1.7 TeV phenomenologically viable dark matter and exotic vectorlike leptons can be obtained.

The model contains a variety of potential LHC signals, though rates will be challenging. Some of the signatures, such as a four-lepton final state with a Z_L resonance in two of the leptons, are fairly novel and specific, but otherwise most LHC phenomenology resembles other vectorlike lepton constructions along with singlet scalar phenomenology.

The 14 TeV run of the LHC should be able to probe some of the parameter space in the exotic lepton sector, although an e^+e^- collider with center-of-mass energies between 250 and 500 GeV is more suitable for this task. Unless the Z_L is very light, direct production is unlikely to be observable at the LHC. The indirect effect on four-lepton interactions can however be probed at a linear collider, vastly extending the reach of the LEP experiments.

The exotic charged leptons can also lead to observable modifications of the Higgs decays and in particular to $h \to \gamma\gamma$, which is also affected by Higgs mixing. We have examined these effects for a range of model parameters and lepton masses which can potentially be produced at the LHC. Potential vacuum stability issues due to the presence of charged leptons with $\mathcal{O}(1)$ couplings to the Higgs can be alleviated with the presence of the gauge and scalar sector of this model, but one can also easily embed it into a more fundamental UV completion which would presumably solve such problems.

While $\mathrm{U}(1)_L$ is an attractive gauge symmetry, which may contribute to the answer as to how dark matter can be massive and yet remain stable, many open questions remain in the current construction. For example, the hierarchy problem remains unaddressed, and almost certainly would require more structure and would lead to new phenomena. The current construction automatically contains new massive states as well as new interactions potentially containing CP-violating phases, which may be useful for explaining the baryon asymmetry of the Universe. One can also easily imagine embedding this model into a supersymmetric version or some other construction which solves the hierarchy problem or generates the lepton-breaking scale naturally, but we leave these possibilities to a future study.

ACKNOWLEDGMENTS

The authors thank André de Gouvêa, Bogdan Dobrescu, Patrick Fox, Roni Harnik, Carlos Wagner, and Felix Yu for useful conversations. The research of T. M. P. T. is supported in part by NSF Grant No. PHY-0970171 and by the University of California, Irvine through a Chancellor's fellowship. R. V.-M. is supported by the Fermilab Graduate Student Fellowship program. This research is also partially supported by Fermi Research Alliance, LLC under Contract No. De-AC02-07CH11359 with the U.S. Department of Energy. The work of P. S. is supported in part by the U.S. Department of Energy, Division of High Energy Physics, under Grants No. DE-AC02-06CH11357 and No. DE-FG02-12ER41811.

[1] G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **716**, 1 (2012).
[2] S. Chatrchyan *et al.* (CMS Collaboration), Phys. Lett. B **716**, 30 (2012).
[3] J. Beringer *et al.* (Particle Data Group), Phys. Rev. D **86**, 010001 (2012).
[4] R. Foot, G. C. Joshi, and H. Lew, Phys. Rev. D **40**, 2487 (1989).
[5] S. Rajpoot, Int. J. Theor. Phys. **27**, 689 (1988).
[6] P. Fileviez Perez and M. B. Wise, Phys. Rev. D **82**, 011901 (2010).
[7] T. R. Dulaney, P. Fileviez Perez, and M. B. Wise, Phys. Rev. D **83**, 023520 (2011).
[8] P. Fileviez Perez and M. B. Wise, J. High Energy Phys. 08 (2011) 068.
[9] P. Fileviez Perez and M. B. Wise, Phys. Rev. D **84**, 055015 (2011).
[10] M. Duerr, P. Fileviez Perez, and M. B. Wise, Phys. Rev. Lett. **110**, 231801 (2013).
[11] W. Chao, Phys. Lett. B **695**, 157 (2011).
[12] P. Dong and H. Long, arXiv:1010.3818.
[13] A. Joglekar, P. Schwaller, and C. E. Wagner, J. High Energy Phys. 12 (2012) 064.
[14] N. Arkani-Hamed, K. Blum, R. T. D'Agnolo, and J. Fan, J. High Energy Phys. 01 (2013) 149.
[15] B. Batell, S. Jung, and H. M. Lee, J. High Energy Phys. 01 (2013) 135.
[16] C. Arina, R. N. Mohapatra, and N. Sahu, Phys. Lett. B **720**, 130 (2013).
[17] G. Moreau, Phys. Rev. D **87**, 015027 (2013).
[18] W.-Z. Feng and P. Nath, Phys. Rev. D **87**, 075018 (2013).
[19] A. Joglekar, P. Schwaller, and C. E. M. Wagner, J. High Energy Phys. 07 (2013) 046.
[20] H. Davoudiasl, I. Lewis, and E. Ponton, Phys. Rev. D **87**, 093001 (2013).
[21] P. Fileviez Perez and M. B. Wise, J. High Energy Phys. 05 (2013) 094.
[22] K. Babu, C. F. Kolda, and J. March-Russell, Phys. Rev. D **57**, 6788 (1998).
[23] K. R. Dienes, C. F. Kolda, and J. March-Russell, Nucl. Phys. **B492**, 104 (1997).
[24] A. Hook, E. Izaguirre, and J. G. Wacker, Adv. High Energy Phys. **2011**, 859762 (2011).
[25] J. Elias-Miro, J. R. Espinosa, G. F. Giudice, H. M. Lee, and A. Strumia, J. High Energy Phys. 06 (2012) 031.
[26] M. Voloshin, Phys. Rev. D **86**, 093016 (2012).
[27] D. McKeen, M. Pospelov, and A. Ritz, Phys. Rev. D **86**, 113004 (2012).
[28] M. S. Carena, A. Daleo, B. A. Dobrescu, and T. M. Tait, Phys. Rev. D **70**, 093009 (2004).
[29] E. Aprile *et al.* (XENON100 Collaboration), Phys. Rev. Lett. **109**, 181301 (2012).
[30] P. Fileviez Perez, Phys. Lett. B **711**, 353 (2012).

[31] B. Gripaios, A. Pomarol, F. Riva, and J. Serra, J. High Energy Phys. 04 (2009) 070.
[32] S. R. Coleman and E. J. Weinberg, Phys. Rev. D **7**, 1888 (1973).
[33] C. Englert, J. Jaeckel, V. V. Khoze, and M. Spannowsky, J. High Energy Phys. 04 (2013) 060.
[34] C. D. Carone and H. Murayama, Phys. Rev. D **52**, 484 (1995).
[35] R. Mohapatra, J. Phys. G **36**, 104006 (2009).
[36] G. Servant and T. M. Tait, Nucl. Phys. **B650**, 391 (2003).
[37] G. Belanger, F. Boudjema, P. Brun, A. Pukhov, S. Rosier-Lees, P. Salati, and A. Semenov, Comput. Phys. Commun. **182**, 842 (2011).
[38] P. J. Fox, R. Harnik, J. Kopp, and Y. Tsai, Phys. Rev. D **84**, 014028 (2011).
[39] D. Malling *et al.*, arXiv:1110.0103.
[40] O. Adriani *et al.* (PAMELA Collaboration), Nature (London) **458**, 607 (2009).
[41] M. Ackermann *et al.* (Fermi LAT Collaboration), Phys. Rev. Lett. **108**, 011103 (2012).
[42] M. Aguilar *et al.* (AMS Collaboration), Phys. Rev. Lett. **110**, 141102 (2013).
[43] O. Adriani *et al.* (PAMELA Collaboration), Phys. Rev. Lett. **105**, 121101 (2010).
[44] M. Ackermann *et al.* (Fermi-LAT collaboration), Phys. Rev. Lett. **107**, 241302 (2011).
[45] C. Jackson, G. Servant, G. Shaughnessy, T. M. P. Tait, and M. Taoso, arXiv:1302.1802.
[46] M. Ackermann *et al.* (LAT Collaboration), Phys. Rev. D **86**, 022002 (2012).
[47] M. Aartsen *et al.* (IceCube Collaboration), Phys. Rev. Lett. **110**, 131302 (2013).
[48] M. Carena, I. Low, and C. E. Wagner, J. High Energy Phys. 08 (2012) 060.
[49] K. Kumar, R. Vega-Morales, and F. Yu, Phys. Rev. D **86**, 113002 (2012).
[50] A. Falkowski, F. Riva, and A. Urbano, arXiv:1303.1812.
[51] T. Appelquist and J. Carazzone, Phys. Rev. D **11**, 2856 (1975).
[52] B. Kyae and C. S. Shin, J. High Energy Phys. 06 (2013) 102.
[53] A. Belyaev, N. D. Christensen, and A. Pukhov, Comput. Phys. Commun. **184**, 1729 (2013).
[54] A. Martin, R. Roberts, W. Stirling, and R. Thorne, Eur. Phys. J. C **28**, 455 (2003).
[55] W. Beenakker, R. Hopker, and M. Spira, arXiv:hep-ph/9611232.
[56] P. Ade *et al.* (Planck Collaboration), arXiv:1303.5076.
[57] A. Freitas and P. Schwaller, Phys. Rev. D **87**, 055014 (2013).