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Dark Matter and Vector-like Leptons From Gauged Lepton Number

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We investigate a simple model where Lepton number is promoted to a local $U(1)_L$ gauge symmetry which is then spontaneously broken, leading to a viable thermal DM candidate and vector-like 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.

\section{I. INTRODUCTION}

With the recent discovery of a new resonance with standard model (SM) Higgs like properties \cite{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 beyond the standard (BSM) physics 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 \cite{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 $U(1)_L$ 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 \cite{4,5}. However, the first complete and consistent model of gauged lepton number (and baryon number) was not explored until more recently in \cite{6} with numerous variations following \cite{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 vector-like 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 $SU(3)_c \otimes SU(2)_W \otimes U(1)_Y \otimes U(1)_L$, where the SM leptons are charged under $U(1)_L$. The anomalous $U(1)_L$ requires us to add a new set of leptons with the appropriate quantum numbers to cancel anomalies. Typically, $U(1)_L$ is spontaneously broken by the vacuum expectation value 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 \cite{6}). Such constructions allow for a simple realization of the well known ‘see-saw’ mechanism of neutrino mass generation, but do not contain viable dark matter candidates without additional assumptions or particle content.

Here, motivated 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 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 vector-like (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 \cite{13–19} and was also recently shown to be useful for baryogenesis \cite{20,21}.

The organization of this papers 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 vector-like 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 [6–8]. All options include three generations of right-handed singlet neutrinos (\( \nu_{Ri} \)), 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 \( \nu_{Ri} \), 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 vector-like 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{align*}

\ell'_L &\equiv (\nu'_L, e'_L) \equiv (2, -1/2, L'), \\
\ell'_R &\equiv (1, -1, L'), \\
nu'_R &\equiv (1, 0, L').
\end{align*}
\]  

(1)

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

\[
\begin{align*}

\ell''_L &\equiv (\nu''_L, e''_L) \equiv (2, -1/2, L''), \\
\ell''_R &\equiv (1, -1, L''), \\
nu''_R &\equiv (1, 0, L'').
\end{align*}
\]  

(2)

where the condition,

\[
L' - L'' = -3
\]

is required by anomaly cancellation. The addition of two sets of chiral fermions carrying lepton number which together form a vector-like 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 [10] where baryon number is also gauged and one obtains a vector-like set of ‘leptoquarks’ 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{align*}

\mathcal{L} &\supset (D^\mu \Phi)^\dagger (D_\mu \Phi) + \frac{\epsilon}{2} Z^{\mu \nu}_{L L} B_{\mu \nu} \\
&+ \bar{\ell}_L D_\mu \gamma^\mu \ell_L + \bar{\nu}_R D_\mu \gamma^\mu \nu_R + i \bar{l} D_\mu \gamma^\mu l,
\end{align*}
\]  

(4)

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

The parameter \( \epsilon \) 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 \( \epsilon \) 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 \( \epsilon \) also leads to \( Z - Z_L \) mixing parameterized by [22],

\[
\tan 2\xi = \frac{2M_Z^2 s_W \epsilon \sqrt{1 - \epsilon^2}}{M_Z^2 - M^2_Z (1 - \epsilon^2) + M^2_Z s_W^2 \epsilon^2},
\]  

(5)

where \( \xi \) 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 \( \Phi \) allows for an expanded scalar potential,

\[
\begin{align*}

V(H, \Phi) &= -\mu^2_H H^\dagger H + \lambda_H |H^\dagger H|^2 \\
&- \mu_\phi^2 \Phi^\dagger \Phi + \lambda_\phi |\Phi|^2 + \lambda_{\phi H} \Phi^\dagger \Phi H^\dagger H,
\end{align*}
\]  

(6)

where \( H \equiv (2, -1/2, 0) \) is the SM Higgs doublet. Once lepton number is broken, the real component of \( \Phi \) 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 $v_\phi$ will be the 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 $\lambda_{h\phi}$ will generically lead to mixing between the real singlet components of $\Phi$ and $H$ parameterized by the mixing angle,

$$\tan 2\theta = \frac{\lambda_{h\phi}v_\phi v_\phi}{\lambda_H v_h^2 - \lambda_H v_h^2}.$$

This mixing leads to the mass eigenstates,

$$h = c_\theta h_o - s_\theta \phi_o,$$
$$\phi = s_\theta h_o + c_\theta \phi_o,$$

where $\phi_o$ and $h_o$ are the gauge eigenstates and $\phi, h$ are the mass eigenstates with masses,

$$m_{h,\phi}^2 = (\lambda_H v_h^2 + \lambda_\phi v_\phi^2) \pm \sqrt{\left(\lambda_H v_h^2 - \lambda_\phi v_\phi^2\right)^2 + \lambda_{h\phi}^2 v_h^2 v_\phi^2},$$

where we have assumed $\lambda_{h\phi} > m_h$ and defined $c_\theta = \cos \theta$, $s_\theta = \sin \theta$, etc. The coupling $\lambda_{h\phi}$ 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 $U(1)_L$ interactions. Yukawa interactions (assuming $L_\phi$ 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 can not 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 $\nu_R'$ and $\nu_R''$ 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 $L_\phi$ 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,

$$\mathcal{L} \supset -c_i \phi_i \bar{\nu}_R' \nu'_{L} - c_i \Phi \bar{\nu}_{L} e_R - c_i \Phi \bar{\nu}_{L}^c \nu'_{R} - y''_e \bar{H} \bar{\nu}'_{L} \nu'_{R} - y''_e \bar{H} \bar{\nu}'_{L} \nu''_{L} + h.c.$$  \hspace{0.5cm} (10)

In general these couplings are complex, containing phases which can lead to CP violation, but for simplicity we assume all couplings in Eq.(10) are real (but see [26, 27] for recent studies of CP violating effects on the diphoton rate coming from vector-like leptons). It is also clear from Eq.(10) that once $\Phi$ obtains a vev the couplings $\lambda_{c_\theta c_\phi}$, and $\lambda_{c_\theta s_\phi}$ will lead to vector-like (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 $\nu'_R$ and $\nu''_L$ 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 vector-like ‘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 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,

$$\mathcal{L} \supset y \Phi \bar{\nu}'_{L} \nu_{R_i} + h.c.$$  \hspace{0.5cm} (11)

Once $\Phi$ obtains a vev, this will lead to mixing between the SM right-handed neutrinos, $\nu_{R_i}$ and the exotic right handed neutrino, $\nu'_R$. 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 $(L', L'') \neq (0,3), (-3/2, -3/2), (-3,0)$. 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 $c_i \rightarrow 0$, 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 vector-like masses for the new leptons much smaller than $v_\phi$ are natural. We also note that small values of the $y_{\nu\ell}^{S,M}$ 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 [13, 15], but here arises from $U(1)_L$ gauge invariance and anomaly cancellation. Only one new scale ($v_\phi$) 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 insuring the existence of a stable DM particle are guaranteed by $U(1)_L$ 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 $\Phi$ 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

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

roughly independently of the value of $g'$.

This lower bound and the experimentally measured value of $m_h \approx 125$ GeV constrains the quartic couplings in the scalar potential of Eq.(6) through Eq.(7) and (9). Fixing $v_\phi = 1.7$ TeV and $m_h = 125$ GeV we can then examine the scalar mixing angle $\theta$, the Higgs quartic $\lambda_H$, and the heavy scalar mass eigenstate $m_\phi$ as functions of the scalar couplings $\lambda_{hp}$ and $\lambda_\Phi$. In Fig.1 we show contours of $\lambda_H (\lambda_\Phi, \lambda_{hp})$ (solid-orange), $\theta(\lambda_\Phi, \lambda_{hp})$ (dotted-red), and $m_\phi(\lambda_\Phi, \lambda_{hp})$ (solid-black) in the $\lambda_{hp} - \lambda_\Phi$ plane.

As can be seen, values of $\theta \lesssim 0.1 - 0.2$ can be obtained for quartic couplings of $O(1)$ 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 $m_\phi \gtrsim 1$ 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 $v_\phi$ 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 $\epsilon$ 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 $\epsilon \sim 0.002$. 

![Figure 1: Contours of Higgs mixing angle $\theta$ (red-dotted), Higgs quartic coupling $\lambda_H$ (orange-solid), and heavy scalar mass $m_\phi$ in GeV (black-solid) as defined in Eqs.(7), Eq.(9) as a function of scalar couplings $(\lambda_{hp}, \lambda_\Phi)$ in Eq.(6).](image1)

![Figure 2: Contours of $Z - Z_L$ mixing angle $\sin \xi (\times 10^4)$ in the $\epsilon - M_{Z_L}$ plane (see Eq.(5)).](image2)
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 \( v_{\phi} \) naturally also arises. One could imagine embedding this model in a supersymmetric version as was done in [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 for example [31].

Another possibility for generating natural values for not only \( v_{\phi} \), but also the electroweak scale \( (v_h) \) is through dimensional transmutation where \( v_{\phi} \) is generated radiatively [32]. This scale is then inherited by the SM through the ‘Higgs Portal’ as done recently in [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 \( v_{\phi} \) 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 \( \Lambda \) down to \( U(1)_{L} \). Below the scale \( \Lambda \), but above the lepton and electroweak breaking scales, loop corrections due to hyper-charged leptons vanish provided the leptons satisfy an orthogonality condition [34],

\[
\text{Tr} \left( L Y \right) = 0. \tag{13}
\]

Combined with the anomaly cancellation constraint in Eq.(3), this would determine the exotic lepton numbers to be \( L' = -3 \) and \( L'' = 0 \). Below \( v_{\phi} \) 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 \( \mathcal{L} \sim \frac{1}{\Lambda^2} \Phi \sigma q q q \ell \) (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, \( \frac{1}{\Lambda^4} (q q q q \ell)(\ell H)^2 \bar{\Phi} \), first occurs at dimension twelve while \( \Delta B = 2 \) operators with \( \Delta L = 0 \) are allowed at dimension 9 [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 is 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 WIMP 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 matter 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 \( \Phi \) and \( H \) obtain expectation values which gives,

\[
\mathcal{L} \supset -\frac{c_{\nu}v_{\phi}}{\sqrt{2}} (1 + \frac{\phi}{v_{\phi}}) \bar{\nu}_{L}^{'} \nu^{' \prime}_{L} - \frac{c_{\nu}v_{\phi}}{\sqrt{2}} (1 + \frac{\phi}{v_{\phi}}) \bar{\nu}_{L}^{''} \nu^{'' \prime}_{R} \tag{14}
\]

\[-y_{\nu}v_{h}(1 + \frac{h_{o}}{v_{h}}) \bar{\nu}_{R}^{''} \nu^{'' \prime}_{L} - y_{\nu}v_{h}(1 + \frac{h_{o}}{v_{h}}) \bar{\nu}_{L}^{'} \nu^{'' \prime}_{R} + h.c.,
\]

leading to the mass matrix,

\[
\mathcal{M}_{\nu} = \frac{1}{\sqrt{2}} \begin{pmatrix}
    c_{\nu}v_{\phi} y_{\nu}v_{h} \\
y_{\nu}v_{h} c_{\nu}v_{\phi}
\end{pmatrix}, \tag{15}
\]

which can be diagonalized using the singular value decomposition \( \mathcal{M}_{\nu D} = U^{\dagger}_{L} \mathcal{M}_{\nu} U_{R} \), where \( \mathcal{M}_{\nu D} \) is a diagonal mass matrix with positive mass eigenvalues \( m_{\nu_{x}} \) and \( m_{\nu_{4}} \).

While the Yukawa couplings to \( \Phi \) 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 \( y_{\nu}v_{h}, y_{\nu}v_{h} \ll c_{\nu}v_{\phi} \), the mass eigenvalues are approximately given by

\[
m_{\nu_{x}} \approx \frac{1}{\sqrt{2}} c_{\nu}v_{\phi}, \tag{16}
\]

\[
m_{\nu_{4}} \approx \frac{1}{\sqrt{2}} c_{\nu}v_{\phi}.
\]
In this limit, the eigenstate $\nu_4$ is mostly composed of the electroweak doublet neutrinos $\nu'_{\nu_R}$ and $\nu'_{\nu_L}$, while $\nu_X$ is a combination of the singlets $\nu'_{L}$ and $\nu'_{R}$ and with tiny couplings to the SM $W^\pm$ and $Z$ bosons. Since the doublet neutrino $\nu_4$ 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 $\nu_X$ is the DM candidate. Of course $\nu_X$ must be able to decay which means that at least one of the Yukawa couplings $y'_L, y''_L$ should be nonzero to allow $\nu_X$ to decay into a Higgs boson and $\nu_X$. Nonetheless, this requirement allows the $y'_L, y''_L$ to be small enough so as to be completely irrelevant in the discussion below.

### B. Annihilation Channels

In [13], annihilation through the interactions generated by $y'_\nu, y''_\nu$ was shown to give the correct relic abundance for DM with dominantly Majorana masses $\lesssim 100$ GeV. Here, because direct detection constraints require $y'_\nu, y''_\nu$ to be tiny, one would have to either rely on co-annihilation 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 [13], there are additional annihilation channels for $\nu_X$ into SM leptons. In particular, since $\nu_X$ 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 $\nu_X$ carry lepton number $L^\nu$ and $L'^\nu$, respectively, and $L^\nu - L'^\nu = -3$ implies a non-vanishing coupling of $\nu_X$ 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 $\nu_X$ 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 $\nu_X$ acquires a small doublet component through nonzero $y'_\nu, y''_\nu$ 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.

The dark matter also couples to the singlet scalar $\phi_L$, with a strength $c_{\nu} \approx \sqrt{2} m_{\nu_X} / v_\phi$. When the Higgs mixing angle $\theta$ is nonzero this will allow annihilation into SM particles through s-channel exchange of $h$ and $\phi$, 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 $c_{\nu}$.

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 $y'_\nu, y''_\nu$, 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 over-all picture. The relevant interactions come from Eq. (4) which before lepton number and electroweak symmetry breaking can be written as,

$$\mathcal{L} \supset g'Z_{L_R} \left( L'' \bar{\nu}_R \gamma^\mu \nu_R' + \bar{L'} \nu_L' \gamma^\mu \nu_L' + \bar{L'} \gamma^\mu l \right),$$

(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

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

(18)

where $P_R$ and $P_L$ are the right and left projection operators respectively and we have neglected any mixing.
between $\nu_X$ and $\nu_4$ generated by $y'_\nu$, $y''_\nu$. Using Eq. (18) a straight forward calculation of the diagram in Fig. 3 gives the annihilation cross section,

$$\sigma = \frac{g'^4((L'^2 + L''^2)(s - m_{\nu_X}^2) + 6L'L''m_{\nu_X}^2)}{8\pi(1 - 4m_{\nu_X}^2/s)^{1/2}((M_{Z_L}^2 - s)^2 + M_{Z_L}^2\Gamma_{Z_L}^2)}.$$  \tag{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 non-relativistic expansion, $s = 4m_{\nu_X}^2 + m_{\nu_X}^2 v^2$, and expanding the annihilation cross section in powers of $v$ to give $\langle \sigma v \rangle = a + b\langle v^2 \rangle + O(\langle v^4 \rangle)$ \cite{36}. Expanding Eq. (19) we obtain

$$a = \frac{3g'^4R_4(L' + L'')^2}{4\pi m_{\nu_X}^2(1 - 4R^2)^2} \tag{20}$$

for the velocity independent coefficient. Note, 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

$$b = \frac{g'^4R_4((L'^2 + L''^2)(11 + 4R^2) + L'L''(6 + 72R^2))}{32\pi m_{\nu_X}^2(1 - 4R^2)^3}. \tag{21}$$

Here we have defined $R = m_{\nu_X}/M_{Z_L}$ and neglected terms of order $\Gamma_{Z_L}/M_{Z_L}$. 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

$$a \approx \frac{3g'^4(L' + L'')^2R_4}{4\pi m_{\nu_X}^2} + O(R^6) \tag{22}$$

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 e.g. using the procedure presented in \cite{36}. We have opted instead to implement the model into the numerical code MICROMEGAS \cite{38}. Not only does this facilitate the exploration of regions of parameter space where the $O(\langle v^2 \rangle)$ expansion breaks down, but it also simplifies the computation of direct and indirect detection signals. The approximate calculation of the relic density following \cite{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 $m_{\nu_X}$ and $v_\phi$, for a few choices of $L'$, in Figure 4. The LEP II constraints on $v_\phi$ require dark matter masses greater than about 200 GeV, and (depending on $L'$), a thermal relic density enforces a tight correlation between $v_\phi$ and $m_{\nu_X}$.

In the limit $y'_\nu, y''_\nu \approx 0$, DM couples to $h$ and $\phi$ through $c_{\nu}$ and the Higgs mixing,

$$\mathcal{L} \supset \frac{c_\nu}{\sqrt{2}}(c_\nu \phi - s_\nu h)\bar{\nu}_X \nu_X,$$  \tag{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 $c_{\nu}$, 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 $y'_\nu, y''_\nu \to 0$ and 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 \cite{39}. 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.

Higgs exchange 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.
section due to the sizable coupling of the Z boson. In experiments such as LZ [40].

The dark matter relic density for different values of $L'$ is well in reach of second generation DM direct detection experiments such as LZ [40]. The green bands indicate regions with correct relic density for different values of $L'$.

$$\sin \theta$$

and present the results in Figure 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 limit from the XENON-100 experiment [29], while the dashed red line indicates the projected reach of the LZ experiment [40]. The green bands indicate regions with correct relic density for different values of $L'$.

Z-boson exchange induces a large DM-neutron cross section due to the sizable coupling of the Z to light quarks. We parameterize the coupling of the Z-boson to the DM as,

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

where $\epsilon'$ is either induced by $Z-Z'$ mixing or by nonzero neutrino Yukawa couplings $y_{\nu_L}, y_{\nu_R}'$. The upper bound on $\epsilon'$ 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 $y_{\nu_L}', y_{\nu_R}' \approx 0$, $\epsilon'$ 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 [40].

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 [41], Fermi [42], and AMS-02 [43]. At the same time, contributions to the anti-proton flux are very tiny, evading constraints from PAMELA [44].

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 [45], which are just short of being able to rule out thermal cross sections for dark matter masses around a few $10^3$s of GeV based on one sixth of the annihilations producing $\tau^+ \tau^-$. In the near future, such constraints are only relevant for $\nu_X$
dark matter which has been produced non-thermally.

Dark matter may also annihilate directly into $\gamma\gamma$ and/or $\gamma Z$ at loop level, providing mono-chromatic 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 [46], where it was found that $\gamma\gamma$, $\gamma Z$, 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 [47], 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 spino-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 [48].

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 $\Phi$ and $H$ obtain expectation values, the Lagrangian for the exotic charged lepton sector becomes,

$$\mathcal{L} = -\frac{c_{v\phi}}{\sqrt{2}} (1 + \frac{\phi_o}{v_o}) e_L'^\dagger e_L' - \frac{c_{v\phi}}{\sqrt{2}} (1 + \frac{\phi_o}{v_o}) e_R'^\dagger e_R'$$

$$- \frac{y_{e}'' v_h}{\sqrt{2}} (1 + \frac{h_o}{v_h}) e_R'^\dagger e_L' - \frac{y_{e}'' v_h}{\sqrt{2}} (1 + \frac{h_o}{v_h}) e_L'^\dagger e_R' + h.c.$$ (25)

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

$$M_e = \frac{1}{\sqrt{2}} \begin{pmatrix} c_{v\phi} & y_{e}'' v_h \\ y_{e}'' v_h & c_{v\phi} \end{pmatrix}.$$ (26)

Again we can diagonalize via $M_{LD} = U_L^\dagger M_L U_R$ to obtain the mass eigenvalues and eigenstates. The Lagrangian in Eq.(25) also leads to the interaction matrices for $\phi_o$ and $h_o$ given by,

$$N_{e}^{\phi} = \frac{v_h}{\sqrt{2}} \begin{pmatrix} 0 & y_{e}'' \\ y_{e}'' & 0 \end{pmatrix}, \ N_{e}^{h} = \frac{v_{\phi}}{\sqrt{2}} \begin{pmatrix} c_{v\phi} & 0 \\ 0 & c_{v\phi} \end{pmatrix},$$ (27)

which upon the rotation performed to diagonalize $M_e$ gives interaction matrices in the mass basis defined as $V_{\phi} = U_L^\dagger N_{e}^{\phi} U_R$ and $V_{h} = U_L^\dagger N_{e}^{h} U_R$. These matrices dictate the couplings of the exotic leptons to $\phi$ and $h$. We note also that Eq.(26) is the same mass matrix in the charged lepton sector considered in [13], with the difference being that in this model there are no explicit mass terms. In particular, when $v_{h}, v_{\phi} \rightarrow 0$ 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 di-photons, as we will see below.

A useful simplifying limit is $c_{v\phi} \approx c_{e} \equiv c_{e}$ 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,

$$m_{e_{1}} \approx \frac{1}{\sqrt{2}} (c_{e} v_{\phi} - y_{e} v_{h})$$

$$m_{e_{2}} \approx \frac{1}{\sqrt{2}} (c_{e} v_{\phi} + y_{e} v_{h}),$$ (28)

where we have assumed $c_{e} v_{\phi} > y_{e} v_{h}$. Thus we see that for fixed $y_{e}$ and $v_{\phi}$, 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 on Higgs mixing in Sec.II.B, we can write the modification of the SM Higgs partial width as,

$$\epsilon_{i} \equiv \frac{\Gamma_{h_{i}}}{\Gamma_{h_{SM}}} = \frac{|M(h \rightarrow i)|^2}{|M(h \rightarrow i_{SM})|^2}$$

$$= c_{\sigma}^2 \frac{|M(h_{0} \rightarrow i) - \frac{\Gamma_{h_{0}}}{\Gamma_{h_{SM}}} M(\phi_{0} \rightarrow i)|^2}{|M(h_{0} \rightarrow i_{SM})|^2},$$ (29)

where we have used Eq.(8) and $\Gamma_{h_{SM}}$ is the SM partial width to a final state $i$ and $\Gamma_{h_{i}}$ is the partial width for $h$ to decay into $i$. The rate expected at the LHC relative to the SM can be written as,

$$\mu_{i} = \frac{\sigma(j \rightarrow h)}{\sigma(j \rightarrow h_{SM})} \frac{\mathcal{B}(h \rightarrow i)}{\mathcal{B}(h_{SM} \rightarrow i)} = \epsilon_{i} \frac{\Gamma_{h_{SM}}}{\Gamma_{h}},$$ (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 $\Gamma_{h_{SM}}$ as the total SM Higgs width and $\Gamma_{h}$ 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 $\epsilon_{gg}$ are through Higgs mixing which gives $\epsilon_{gg} \approx c_{\sigma}^2$. 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 the only effect again comes from Higgs mixing, which gives \( \epsilon_{Z\gamma} = \epsilon_{WW} \sim c_{\theta}^2 \). 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, 49] 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 \( c_{\theta}^2 \) 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 \( \Gamma_{h\gamma\gamma}/\Gamma_h \approx c_{\theta}^{-2} \) which will cancel with the \( c_{\theta}^2 \) in the production channels \( \epsilon_{gg, VV} \). This gives finally for the relative rates \( \mu_i = c_{\theta}^2 \) for \( i \neq \gamma\gamma \) and for the final modified diphoton signal strength,

\[
\mu_{\gamma\gamma} = c_{\gamma\gamma}.
\]

Using the approach and conventions of [50], which examined the similar \( gg \rightarrow h \) process, we can go on to obtain the exotic charged lepton contributions to the \( h \rightarrow \gamma\gamma \) process, allowing for the modified diphoton signal strength, the amplitudes in Eq.(32) are evaluated at \( m_{h_*} = m_h \) and \( m_{\phi_*} = m_h \) where \( m_h \) is the physical scalar mass.

Using Eq.(29)-(32) we obtain,

\[
\mu_{\gamma\gamma} = \frac{c_{\theta}}{v_h} \left(F_{SM} + \sum_i (V_{h})_{ii} F_F(\tau_{e_i}) \right) \left( \sum_i (V_{\phi})_{ii} F_F(\tau_{e_i}) \right)^2
\]

\[
= c_{\theta}^2 \left( 1 + F_{SM} \sum_i (V_{h})_{ii} F_F(\tau_{e_i}) \right) - t_\theta \left( F_{SM}^{-1} \frac{v_h}{v_\phi} \sum_i (V_{\phi})_{ii} F_F(\tau_{e_i}) \right)^2,
\]

where \( F_{SM} \) 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 numerical value of \( \sim -6.5 \) for \( m_h = 125 \text{ GeV} \). Note only the diagonal entries in the interaction matrices \( (V_{h})_{ii} \) and \( (V_{\phi})_{ii} \) contribute in the \( h \rightarrow \gamma\gamma \) loop.

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

\[
\mu_{\gamma\gamma} \simeq c_{\theta}^2 \left[ 1 - \frac{v_h}{\sqrt{2} F_{SM}} \left[ y_e \left( \frac{F_F(\tau_{e_1})}{m_{e_1}} - \frac{F_F(\tau_{e_2})}{m_{e_2}} \right) + c_e t_\theta \left( \frac{F_F(\tau_{e_1})}{m_{e_1}} + \frac{F_F(\tau_{e_2})}{m_{e_2}} \right) \right] \right]^2,
\]

where \( m_{e_1,e_2} \) are given in Eq.(28) and satisfy \( m_{e_1} < m_{e_2} \). Remembering that \( F_{SM} < 0 \) we see in the limit \( t_\theta \rightarrow 0 \) we have an enhancement in the diphoton rate in the presence of mostly vector-like leptons entering through the \( h_* \) component of \( h \). This is, of course, expected from the low energy Higgs theorems (see e.g. [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 $\lambda_{hp}$ 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 $c_1$ and $c_2$ become almost purely vector-like and again the contribution only enters through Higgs mixing via the $\phi_o$ component of $h$. However as $v_\phi \to \infty$ one also has $t_\phi \to 0$ and the $\phi_o$ contribution eventually decouples from the $h \to \gamma\gamma$ amplitude as $v_\phi$ is taken large. Eq.(34) is in agreement with [15] for the case where their explicit mass term is put to zero.

To avoid the constraints discussed in Sec. IIE we choose $v_\phi = 1.7$ TeV and take the lightest charged lepton to have mass greater than $m_{min} \sim 100$ 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 [51]. Regardless this implies that these fermions must be mostly ‘vector-like’ 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,

$$m_{e_1} = \frac{c_e v_\phi - y_e v_h}{\sqrt{2}} \geq m_{min}, \quad (35)$$

which leads to a condition on the Yukawa coupling,

$$\frac{\sqrt{2} m_{min} + y_e v_h}{v_\phi} \leq c_e \lesssim 4\pi. \quad (36)$$

where we have also indicated $4\pi$ as the perturbative upper bound.

Since the mixing angle will affect all decay channels, we perform a fit to the full Higgs data [52, 53] 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, 3 $\sigma$ regions (purple) for the favored parameter space where the grey band shows the LEP II excluded region 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 $\lambda_{hp} < 0$, which can potentially lead to vacuum instabilities. On the other hand, positive values of $\theta \sim 0.5$ where $\lambda_{hp} > 0$ lead to no instability and as shown in [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 $\theta$ 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 $\sim 10 - 20\%$ 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.
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 $pp \rightarrow \ell^+ \ell^- Z_L$. The cross section for this process is calculated using the program CALCHEP [57] with the MRST2002 PDF set [58] and 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 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 there are three light $\nu_R$ 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 NLO corrections are moderate [59]. For this plot we have assumed that the lepton masses are given by Eq. (16) and Eq. (28). This leads to the following mass hierarchies for the exotic lepton sector,

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

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 $\nu_4$. Note also that for $y_e \sim 0.8$ and the $m_{e_1}$ range 100 GeV – 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 Eq.(16) and Eq.(28) the heavy charged state $e_2$ can have the following decay chain,

$$e_2 \rightarrow W\nu_4 \rightarrow WW e_1 \rightarrow WWW\nu_X. \tag{38}$$

Note that although we are neglecting mass mixing between $\nu_X$ and $\nu_4$ by assuming $y_{\nu_4} \ll 1$, it must be non-zero for the the heavy leptons to decay down to the DM.

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

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

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

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

The heavy neutrino state $\nu_4$ can decay via $Z$ and $h$ bosons through,

$$\nu_4 \rightarrow Z\nu_X, \ \nu_4 \rightarrow h\nu_X, \tag{41}$$

as well as $W$ bosons through,

$$\nu_4 \rightarrow W e_1 \rightarrow WW\nu_X. \tag{42}$$

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

$$pp \rightarrow e_1^+ e_1^- \rightarrow WW E_T \rightarrow l^+ l^- E_T. \tag{43}$$

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

$$pp \rightarrow e_1^+ \nu_4 \rightarrow WZ E_T \rightarrow l^+ l^- l^- E_T. \tag{44}$$

The signatures are similar to those from 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 light enough $\phi$ 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 $s_\phi^2$. Additionally, if kinematically allowed $\phi$ can also have the following decays to heavy leptons and dark matter,

$$\phi \rightarrow e_1 e_1, \ \phi \rightarrow e_2 e_2$$

$$\phi \rightarrow e_1 e_2, \ \phi \rightarrow \nu_4 \nu_X \tag{45}$$

It can of course also decay to Higgs pairs $\phi \rightarrow hh$ when kinematically allowed. As discussed in Sec.II.E, however, for $v_\phi \sim 1.7$ TeV we typically have $\phi$ 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 vector like 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 vector-like 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 vector like 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 GeV 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 \( h \rightarrow \gamma \gamma \), 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 \( O(1) \) 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 \( U(1)_L \) 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.

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