Dileptonic Scalar Dark Matter and Exotic Leptons

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Abstract

A simple scenario of long-lived dark matter is presented. Assuming lepton number \((L)\) conservation with Dirac neutrinos, the neutral component of an \(SU(2)_L\) scalar triplet with \(L = 2\) is a suitable candidate, which decays to two neutrinos. Exotic leptons with \(L = -1\) may also play a role, and act as anchors for small seesaw Dirac neutrino masses.
**Introduction:** Dark matter is a current topic of pervasive research. For recent reviews, see for example Refs.[1, 2, 3]. The conventional wisdom is that dark matter is a single particle which was in thermal equilibrium with matter in the early Universe and froze out as the temperature of the Universe fell below its mass, typically of order 100 GeV to a few TeV. Its relic abundance is now precisely known [4], i.e. $\Omega_c h^2 = 0.120 \pm 0.001$. A dark symmetry is also usually assumed to ensure that the lightest dark particle is stable.

The simplest model of dark matter is that of a stable real singlet scalar $S$ which is distinguished from matter by a $Z_2$ symmetry under which only $S$ is odd. For a recent comprehensive study, see Ref.[5]. The scalar potential simply consists of $S$ and the Higgs doublet $\Phi = (\phi^+, \phi^0)$ of the Standard Model (SM) of quarks and leptons. Assuming that small Majorana neutrino masses are obtained through the canonical seesaw mechanism with three heavy singlet right-handed neutrinos $N_R$, the terms $SN_R N_R$ are forbidden by the imposed $Z_2$. There is however a better way to understand this dark $Z_D^2$ parity. It is simply derivable [6] from lepton parity $Z_L^2 = (-1)^L$ as $Z_D^2 = Z_L^2 \times (-1)^{2j}$, where $j$ is the spin of the particle. This works by assigning $S$ as odd under $Z_L^2$, i.e. same as the leptons. In other words, $S$ is a leptonic scalar. Under $Z_D^2$, $S$ is then odd, whereas all the leptons (being fermions) are even. This simple connection applies to many other scenarios of Majorana neutrinos with dark matter, including the original 2006 scotogenic one-loop radiative model [7]. In the context of grand unification, this leptonic marker is most easily understood as $Q_\chi$ from the decomposition $SO(10) \rightarrow SU(5) \times U(1)_\chi$ [8].

In the continuing absence of experimental evidence for neutrinoless double beta decay [9], there has been a recent surge of theoretical interest in Dirac neutrinos, with the recent identification of lepton number with a corresponding dark symmetry [10], i.e. a generalization of dark parity $Z_D^2$ from lepton parity $Z_L^2$ to $Z_N^D$ from $Z_N^L$ with the factor $\omega^{-2j}$ where $\omega^N = 1$, as well as $U(1)_D$ from $U(1)_L$ with $D = L - 2j$. 


More generally, naturally small Dirac neutrino masses [11] may be obtained in analogy to Majorana neutrino masses [12]. In any such framework, dark matter may also have something to do with lepton number. In this paper, a simple scenario is presented by the addition of an SU(2)\textsubscript{L} scalar triplet \( \rho = (\rho^+, \rho^0, \rho^-) \) with \( L = 2 \).

**Dileptonic scalar SU(2)\textsubscript{L} triplet**: Consider first \( \rho \) with \( L = 0 \). This would have a trilinear interaction with the SM Higgs doublet \( \Phi \) through \( \rho(\Phi^\dagger \Phi)^3 \), where

\[
(\Phi^\dagger \Phi)^3 = [\bar{\phi}^0 \phi^+ + (\phi^- \phi^0 - \bar{\phi}^0 \phi^0) / \sqrt{2}, \phi^- \phi^0].
\]

(1)

This means that \( \rho^0 \) will develop a nonzero vacuum expectation value and contributes to the \( W^\pm \) mass. It is clearly not a candidate for dark matter.

Consider next \( \rho \) with \( L = 1 \). This means that \( \rho^- \) is not the complex conjugate of \( \rho^+ \) and that \( \rho^0 \) is complex, not real. It also means that it does not couple to \( (\Phi^\dagger \Phi)_3 \) or any linear combination of two leptons which must have \( L = 0 \) or \( L = 2 \). In other words, the designation of \( L = 1 \) is arbitrary for \( \rho \). This symmetry is in fact new and has nothing to with \( L \).

Things are different if \( \rho \) has \( L = 2 \). It may now couple to two leptons through the following three dimension-six operators:

\[
\mathcal{L}_1 = \Lambda_1^{-2} \rho^\dagger (\Phi^\dagger \Phi)_3 (LL)_3, \\
\mathcal{L}_2 = \Lambda_2^{-2} \rho^\dagger (\Phi^\dagger \Phi)_3 \nu_R \nu_R, \\
\mathcal{L}_3 = \Lambda_3^{-2} \partial^\mu (\rho^\dagger \Phi)_{2} \overline{\nu_R} \gamma_\mu L,
\]

(2) \hspace{1cm} (3) \hspace{1cm} (4)

where \( L = (\nu, e)_L \),

\[
(\Phi^\dagger \Phi)_3 = (\phi^+ \phi^+, \sqrt{2} \phi^+ \phi^0, \phi^0 \phi^0),
\]

(5)

\[
(LL)_3 = [\nu_L \nu_L, (\nu_L e_L + e_L \nu_L) / \sqrt{2}, e_L e_L],
\]

(6)

\[
(\rho^\dagger \Phi)_3 = \left[ -\sqrt{\frac{2}{3}} \phi^- \phi^0, -\sqrt{\frac{1}{3}} \phi^+ \phi^- \phi^0, \frac{1}{3} \phi^0 \phi^0 + \sqrt{\frac{2}{3}} \phi^- \phi^+ \right].
\]

(7)
The structure of $\mathcal{L}_3$ is such that it is suppressed by lepton masses compared to $v = \langle \phi^0 \rangle$, so it may safely be ignored. Now the only two-body decay of $\rho^0$ is to two neutrinos, as in an earlier proposal with $Z_3$ lepton symmetry [14]. In $\mathcal{L}_1$, $\rho \rightarrow \nu_L \nu_L$ has the amplitude $-\sqrt{2}v^2/\Lambda_1^2$. In $\mathcal{L}_2$, $\rho \rightarrow \nu_R \nu_R$ has the amplitude $-\sqrt{2}v^2/\Lambda_2^2$. The resulting decay rate is then

$$\Gamma_{1,2} = \frac{m_\rho v^4}{16\pi\Lambda_{1,2}^4},$$  

(8)

assuming the dominant decay to only one Dirac neutrino. Since $\rho^0$ has $L = 2$ and lepton number is conserved, other two-body decay modes such as $e^+e^-$ and $\gamma\gamma$ are strictly forbidden. However, three-body modes such Higgs + $\nu\nu$ and $W + \nu e$ are possible even though they are subdominant. They would then disrupt the Cosmic Microwave Background (CMB) and push the limit of the dark matter’s lifetime to beyond $10^{25}$ seconds [13]. As a result,

$$\frac{m_\rho}{2 \text{ TeV}} < \left( \frac{\Lambda_{1,2}}{8.63 \times 10^{14} \text{ GeV}} \right)^4.$$  

(9)

The components of the scalar triplet $\rho$ are split by their one-loop gauge interactions, with the neutral component lower in mass than the charged ones [15] [16] by about 166 MeV. Hence the latter would decay to the former by emitting a virtual $W^\pm$ boson which converts to $\pi^\pm$ or $e^\pm, \mu^\pm$ and neutrinos.

As for relic abundance, $\rho$ behaves analogously to fermion triplet dark matter [17] as well as the wino in supersymmetry [18]. Their annihilation to electroweak gauge bosons has the correct cross section for $m_\rho$ in a range near 2 TeV.

Since $\rho^0$ does not couple to $Z$ at tree level, it may only interact with quarks in one loop through $W^\pm$ exchange, or through the SM Higgs boson directly. These interactions allow it to be detected in underground experiments. With $m_\rho \sim 2$ TeV, a possible discovery is on the horizon [16] of current experiments.

**Ultraviolet completions**: To obtain $\mathcal{L}_{1,2}$ in a renormalizable theory, new heavy particles are required. Consider first a scalar triplet $\xi = (\xi^{++}, \xi^+, \xi^0)$ with $L = -2$. This enables the
well-known interaction
\[ \xi^0 \nu_L \nu_L - \xi^+ (\nu_L e_L + e_L \nu_L) / \sqrt{2} + \xi^{++} e_L e_L. \]  
(10)

The usual next step is to allow \( L \) to be broken spontaneously \[19\] or explicitly by the soft term \[20, 21\]
\[ \xi^0 \bar{\phi}^0 + \sqrt{2} \xi^+ \phi^- \bar{\phi}^0 + \xi^{++} \phi^- \phi^0, \]  
(11)
in which case the well-known Type II seesaw \[12\] is realized for obtaining small Majorana neutrino masses. Here \( L \) is strictly conserved, so (10) is allowed but not (11). With \( \rho \) and \( \xi \), the triple product \( \rho^\dagger \times (\Phi \Phi)_3 \times \xi^\dagger \) is allowed, thereby enabling \( \mathcal{L}_1 \) as shown in Fig. 1. After integrating out \( \xi \), the dimension-six operator \( \mathcal{L}_1 \) is obtained.

\[ \rho \Phi \Phi \xi \]
\[ \Phi \]
\[ \rho \]
\[ \xi \]
\[ L \]
\[ L \]

Figure 1: Dimension-six operator for \( \rho^0 \rightarrow \nu_L \nu_L \).

Consider next a heavy neutral scalar singlet \( \zeta \) with \( L = -2 \), then \( \zeta \nu_R \nu_R \) is allowed, as well as the quartic scalar interaction \( \rho^\dagger (\Phi \Phi)_3 \zeta^\dagger \). Together they generate Fig. 2, resulting in \( \mathcal{L}_2 \).

Note that in a previous proposal \[22\], \( \zeta \) is assumed to be very light and acts as the dilepton mediator for self-interacting leptonic dark matter. In that case, it must decay very quickly to avoid disrupting the standard cosmological scenario of the early Universe. In the singlet-triplet Majoron model with soft breaking of \( L \), the pseudo-Majoron \[23\] decays instead to \( \nu_L \nu_L \).
**Exotic leptons**: With the new insight of having particles with nonzero lepton numbers, an interesting addition to the SM would be the following vectorlike doublet \( E_{L,R} = (E_0^0, E^-)_{L,R} \) and singlet \( N_{L,R} \) with \( L = -1 \). They have invariant mass terms \((E_L^0 E_R^0 + E_L^- E_R^-)\) and \( \bar{N}_L N_R \), as well as terms mixing \( E \) and \( N \) through \( \Phi \), and those linking them to the SM leptons \( L = (\nu, e)_L, e_R, \nu_R \) with \( L = 1 \), because \( \bar{N}_L \) and \( \nu_R \) transform identically. The two sectors mix, through the terms

\[
N_R \nu_R, \quad (\nu_L \phi^0 - e_L \phi^+) N_L, \quad (E_R^0 \phi^0 - E_R^- \phi^+) \nu_R.
\]

(12)

This mixing is of course assumed to be very small. The neutral Dirac fermion matrix linking \((\bar{\nu}_L, N_R, E_R^0)\) to \((\nu_R, \bar{N}_L, E_L^0)\) is of the form

\[
M = \begin{pmatrix}
m_\nu & m_{\nu N} & 0 \\
m_{N \nu} & M_N & m_{N E} \\
m_{E \nu} & m_{E N} & M_E
\end{pmatrix},
\]

(13)

where \( m_{N \nu}, m_{E \nu}, \) and \( m_{\nu N} \) mix the \( L = \pm 1 \) sectors.

There is however another important coupling, i.e. \( \rho^i(\bar{E}_R L)_3 \), where

\[
(\bar{E}_R L)_3 = [\bar{E}_R^- \nu_L, (\bar{E}_R^0 \nu_L - \bar{E}_R^- e_L) / \sqrt{2}, \bar{E}_R^0 e_L].
\]

(14)

Using this, the dimension-six operators \( \mathcal{L}_{1,3} \) are derived as shown in Figs. 3 and 4. Whereas it appears that only \( E_R \) and \( N_L \) are needed for these operators, the presence of \( E_L \) and \( N_R \) allows them to have heavy invariant masses greater than \( m_\rho \) for the validity of the interpretation of \( \mathcal{L}_{1,3} \) as effective operators.
Seesaw Dirac neutrino masses: The presence of $N_{L,R}$ with $L = -1$ is also useful for obtaining small seesaw Dirac neutrino masses. A $Z_2$ symmetry is imposed \cite{11} under which $\nu_R$ (with $L = 1$) is odd and all other fields even. All dimension-four terms must respect this $Z_2$. This means that $m_\nu$ and $m_{E\nu}$ in Eq. (13) are zero. However, $Z_2$ is allowed to be broken softly, hence $m_{N\nu}$ is nonzero. Note that $L$ is still strictly conserved. Now $\mathcal{M}$ of Eq. (13) becomes

$$\mathcal{M} = \begin{pmatrix} 0 & m_{\nu N} & 0 \\ m_{N\nu} & M_N & m_{N E} \\ 0 & m_{E N} & M_E \end{pmatrix}, \quad (15)$$

which induces a small Dirac neutrino mass $m_{N\nu}m_{\nu N}/M_N$, assuming negligible mixing of $N$ with $E$.

Conclusion: If neutrinos are Dirac fermions with strictly conserved lepton number $L$, the existence of dark matter may be related to $L$. A simple scenario is presented where an electroweak scalar triplet $(\rho^+, \rho^0, \rho^-)$ with $L = 2$ is a possible candidate. The lightest
component, i.e. $\rho^0$, decays to two neutrinos with a lifetime exceeding that of the Universe. A possible connection to exotic leptons with $L = -1$ is also discussed, with the natural appearance of small seesaw Dirac neutrino masses.

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References

[1] G. Bertone and D. Hooper, Rev. Mod. Phys. 90, 045002 (2018).

[2] M. Lisanti, TASI 2015 Lectures, arXiv:1603.03797

[3] B.-L. Young, Front. Phys. 12, 121201, 121202 (2017).

[4] Planck Collaboration: N. Aghanim et al., Astron. Astrophys. 641, A6 (2020).

[5] The GAMBIT Collaboration: P. Athron et al., Eur. Phys. J. C77, 568 (2017).

[6] E. Ma, Phys. Rev. Lett. 115, 011801 (2015).

[7] E. Ma, Phys. Rev. D73, 077301 (2006).

[8] E. Ma, Phys. Rev. D98, 091701 (2018).

[9] A. Giuliani et al., arXiv:1910.04688 [hep-ex].

[10] E. Ma, Phys. Lett. B809, 135736 (2020).

[11] E. Ma and O. Popov, Phys. Lett. B764, 142 (2017).

[12] E. Ma, Phys. Rev. Lett. 81, 1171 (1998).

[13] T. R. Slatyer and C. L. Wu, Phys. Rev. D95, 023010 (2017).
[14] E. Ma, N. Pollard, R. Srivastava, and M. Zakeri, Phys. Lett. B750, 135 (2015).

[15] M. Sher, Phys. Rev. D52, 3136 (1995).

[16] M. Cirelli, N. Fornego, and A. Strumia, Nucl. Phys. B753, 178 (2006).

[17] E. Ma and D. Suematsu, Mod. Phys. Lett. A24, 583 (2009).

[18] M. Beneke et al., JHEP 1603, 119 (2016).

[19] G. B. Gelmini and M. Roncadelli, Phys. Lett. B99, 411 (1981).

[20] J. Schechter and J. W. F. Valle, Phys. Rev. D22, 2227 (1980).

[21] E. Ma and U. Sarkar, Phys. Rev. Lett. 80, 5716 (1998).

[22] E. Ma, Mod. Phys. Lett. A33, 1850226 (2018).

[23] E. Ma and M. Maniatis, JHEP 1707, 140 (2017).