Canonical seesaw implication for two-component dark matter

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We show that the canonical seesaw mechanism implemented by the $U(1)_{B-L}$ gauge symmetry provides two-component dark matter naturally. The seesaw scale that breaks $B-L$ defines a residual gauge symmetry to be $Z_6 = Z_2 \otimes Z_3$, where $Z_2$ leads to the usual matter parity, while $Z_3$ is newly recognized, transforming quark fields nontrivially. The dark matter component that is odd under the matter parity has a mass ranging from keV to TeV. Another dark matter component that lies in a nontrivial representation of $Z_3$ can gain a mass in the range similar to the former component, in spite of the fact that it can be heavier than the light quarks $u, d$. This two-component dark matter setup can address the XENON1T anomaly recently observed.

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Motivation. Neutrino mass and dark matter are the two big questions in science, which require the new physics beyond the standard model [1].

It is well established that the canonical seesaw mechanism can generate appropriate small neutrino masses through the exchange of heavy Majorana right-handed neutrino singlets, $\nu_{aR}$ for $a = 1, 2, 3$, added to the standard model [2-10]. However, the canonical seesaw in its simple form does not naturally address the dark matter issue, unless some dark matter stability condition or parameter finetuning is ad hoc imposed.

The simplest gauge completion of the seesaw mechanism with $U(1)_{B-L}$ can provide a natural origin for the existence of the right-handed neutrinos and the right-handed neutrino mass scale [11-13]. This work shows that such theory manifestly yields a novel consequence of two-component dark matter, properly solving the recent XENON1T excess [14].

Description of the model. Indeed, the full gauge symmetry takes the form,

$$SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes U(1)_{B-L}. \tag{1}$$

Here the right-handed neutrino fields $\nu_{aR}$ transforming under the gauge symmetry as

$$\nu_{aR} \sim (1, 1, 0, -1) \tag{2}$$

are required in order to cancel the $[\text{Gravity}]^2 U(1)_{B-L}$ and $[U(1)_{B-L}]^3$ anomalies. Additionally, a scalar singlet transforming under the gauge symmetry as

$$\chi \sim (1, 1, 0, 2) \tag{3}$$

must be presented to break $U(1)_{B-L}$ for the model consistency, simultaneously generating the right-handed neutrino masses or the seesaw scale.

As usual, let us assign the standard model lepton, quark, and Higgs representations with respect to the new gauge symmetry to be,

$$l_{aL} = \frac{\nu_{aL}}{\epsilon_{aL}} \sim \left(1, 2, -\frac{1}{2}, -1\right), \tag{4}$$

$$e_{aR} \sim (1, 1, -1, -1), \tag{5}$$

$$q_{aL} = \frac{u_{aL}}{d_{aL}} \sim \left(3, 2, 1, \frac{1}{3}\right), \tag{6}$$

$$u_{aR} \sim \left(3, 1, 2, \frac{1}{3}\right), \tag{7}$$

$$d_{aR} \sim \left(3, 1, -1, \frac{1}{3}\right), \tag{8}$$

$$\phi = \left(\phi^+ \phi^0\right) \sim \left(1, 2, \frac{1}{2}, 0\right). \tag{9}$$

The scalar multiplets develop vacuum expectation values (VEVs), such as

$$\langle \chi \rangle = \frac{\Lambda}{\sqrt{2}}, \quad \langle \phi \rangle = \left(0, \frac{v}{\sqrt{2}}\right), \tag{10}$$

satisfying

$$\Lambda \gg v = 246 \text{ GeV}. \tag{11}$$

The Yukawa Lagrangian includes

$$\mathcal{L} \supset h_{ab} \bar{\nu}_{aL} \phi v_{bR} + \frac{1}{2} f^c_{ab} \nu^c_{aR} \chi v_{bR} + H.c.$$ \hspace{1cm} (12)

where

$$m_{ab} = -h_{ab} \frac{v}{\sqrt{2}}, \quad M_{ab} = -f^c_{ab} \frac{\Lambda}{\sqrt{2}}. \tag{13}$$

Hence, the canonical seesaw is naturally recognized in the $U(1)_{B-L}$ gauge completion given that $v \ll \Lambda$ or $m \ll M$, yielding the observed neutrino ($\sim \nu_{aL}$) masses to be

$$m_{\nu} = -m_{ab}^{-1} m_T = h_{\nu} (f_{\nu})^{-1} (h_{\nu})^T \frac{v^2}{\sqrt{2} \Lambda}. \tag{14}$$
while the heavy neutrinos ($\nu R$) obtain large masses at the $B - L$ breaking scale, $M \sim \Lambda$.

The neutrino oscillation data implies $m_{\nu} \sim 0.1$ eV [1], which leads to

$$\Lambda \sim [(h^\nu)^2 / f^\nu] 10^{14} \text{ GeV}. \quad (15)$$

The seesaw scale $\Lambda$ is close to the grand unification scale if $(h^\nu)^2 / f^\nu \sim 1$. If $(h^\nu)^2 / f^\nu$ is sufficiently small, say $f^\nu / h^\nu \sim 10^{-5.5} - 10^{-5}$ proportional to the electron Yukawa coupling, we derive $\Lambda \sim 1-10$ TeV, in agreement to the collider bounds [1].

The above results have been established in the literature. However, a proper realization of residual gauge symmetry of $B - L$ and its implication for dark matter have not emerged yet. Let us call the reader’s attention to previous works [15–22] relevant to this proposal.

Residual symmetry and dark matter. The symmetry breaking scheme is obtained as

$$SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes U(1)_{B-L} \downarrow \Lambda$$

$$SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes R \downarrow \nu$$

$$SU(3)_C \otimes U(1)_Q \otimes R$$

Here the electric charge is related to the isospin and hypercharge as $Q = T_3 + Y$. $R$ is a residual symmetry of $U(1)_{B-L}$ that conserves the $\chi$ vacuum, although this vacuum ($\chi$) = $\Lambda / \sqrt{2} \neq 0$ breaks $B - L$ by two unit. As being a $U(1)_{B-L}$ transformation, $R = e^{i \alpha (B-L)}$ where $\alpha$ is a transforming parameter. The vacuum conservation condition $R(\chi) = (\chi)$ leads to $e^{i \alpha (2)} = 1$, or equivalently $\alpha = k \pi$ for integer $k$. Hence, the residual symmetry is

$$R = e^{ik \pi (B-L)} = (e^{i \pi (B-L)})^k. \quad (16)$$

It is noted that the transformation with $k$ is conjugated to that with $-k$, i.e. $R^\dagger = (e^{i \pi (B-L)})^{-k} = R^{-1}$.

| Field $(\nu, e) (u, d) (\phi, \chi, A)$ | $(1,1,1)$ | $(1,-1,1)$ | $(1,1,1)$ |
|---------------------------------|----------|------------|----------|
| $(p^3) \sim -1 \ -1 \ -1$ | $Z_2$ | $Z_2$ | $Z_2$ |
| $(p^3) \sim w \ w \ w$ | $p^3$ | $p^3$ | $p^3$ |
| $(p^3) \sim w \ w \ w$ | $p^3$ | $p^3$ | $p^3$ |
| $(p^3) \sim w \ w \ w$ | $p^3$ | $p^3$ | $p^3$ |

TABLE II. Field representations under the residual symmetry $R = Z_2 \otimes Z_3$.

The field representations under $Z_2$ and $Z_3$ are computed in Table II where $w \equiv e^{i 2 \pi / 3}$ is the cube root of unity. Here note that $Z_2$ has two irreducible representations, $\{1\}$ according to $p^3 = 1$ and $\{1\}'$ according to $p^3 = -1$, whereas $Z_3$ has three irreducible representations, $\{1\}$ according to $(p^3, p^3) = (1,1)$ or $(1,-1)$, $\{1\}'$ according to $(p^2, p^3) = (w, w)$ or $(w,-w)$, and $\{1\}''$ according to $(p^2, p^2) = (w, w)$ or $(w^2, -w^2)$, which are independent of $p^3$ values, 1 or -1, that identify $Z_3$ elements in a coset of the quotient group, respectively. The representation $\{1\}''$ is not presented for the existing fields, but the antiquarks $(u', d')$ belong to $\{1\}''$ under $Z_3$.

For brevity, the quotient group can be defined as

$$Z_3 = \{[1], [p^2], [p^3]\}, \quad (21)$$

where each (coset) element $[x]$ consists of two elements of $Z_6$, the characteristic $x$ and the other $p^3 x$, as multiplied by $p^3$. Hence, $[1] = [p^2] = Z_2$, $[p] = [p^4] = \{p, p^4\}$, and $[p^2] = [p^3] = \{p^2, p^3\}$. Because of $[p^3] = [p^2]^3 = [p^3]^3$, the $Z_3$ group is completely generated by

$$[p^2] = [e^{i 2 \pi (B-L)}] = [w^3 (B-L)]. \quad (22)$$

That said, the $Z_3$ irreducible representations $\{1\}, \{1\}'$, and $\{1\}''$ are simply determined by $[p^2] = [1] \rightarrow 1$, $[p^3] = [w] \rightarrow w$, and $[p^2] = [w^2] \rightarrow w^2$, respectively. Here, the intermediate $Z_6$ representations $[r]$ consists of $r$ and $\pm r$ as multiplied by $p^3 = \pm 1$ respectively, which are homomorphic to that of $Z_3$, $[r] \rightarrow \{r, r\} \rightarrow r$.

Since the spin parity $P_S = (-1)^{2s}$ is always conserved by the Lorentz symmetry, we can conveniently multiply

\[ \text{(19)} \]

\[ \text{(20)} \]

1 The nontrivial representations of $Z_3$ obey $\{1\}' \otimes \{1\}'' = \{1\}$. $\{1\}'^3 = \{1\}^3 = \{1\}'$, $\{1\}''^3 = \{1\}^3 = \{1\}'^3$, and $\{1\}'^* = \{1\}^* = \{1\}''^*$. whereas that of $Z_2$ satisfies $\{1\}^2 = \{1\} = \{1\}'^2$, whereas that of $Z_2$ satisfies $\{1\}^2 = \{1\} = \{1\}'^2$.
the residual symmetry \( R = Z_2 \otimes Z_3 \) with spin-parity group \( S = \{ 1, P_3 \} \) to perform
\[
R \rightarrow R \otimes S = (Z_2 \otimes S) \otimes Z_3,
\]
where \( Z_3 \) is retained as the quotient group. The new invariant subgroup \( Z_2 \otimes S \) defines a matter parity
\[
P_M = p^3 \times P_S = (-1)^{3(B-L)+2s}, \tag{24}
\]
where the quotient group \( (Z_2 \otimes S)/P \) is a group of matter-parity symmetry by itself, which is an invariant subgroup of \( Z_2 \otimes S \). Therefore, we factorize
\[
R \otimes S = [(Z_2 \otimes S)/P] \otimes P \otimes Z_3. \tag{25}
\]
Here \( (Z_2 \otimes S)/P = \{ P, \{ p^3, P_S \} \} \) is conserved, if \( P_M \) is conserved. Therefore, instead of \( R \otimes S \), we can consider an alternative residual symmetry which is contained in
\[
R \otimes S \supset P \otimes Z_3, \tag{26}
\]
where the quotient group \( (Z_2 \otimes S)/P \) is neglected, since the theory automatically preserves it. Of course, the theory conserves both \( P \) and \( Z_3 \), under which the representations under these groups are given in Table III.

| Field \((u, e)(u, d)(\phi, \chi, A)\) | \( P_M \) | \( P \) | \( Z_3 \) |
|-----------------|------|------|------|
| \( P^0 \)       | 1    | 1    | 1    |
| \( P^+ \)       | 1    | 1    | 1    |
| \( |p^o| \)       | 1    | \( w \) | 1    |
| \( Z_3 \)       | 1    | \( 1' \) | 1    |

TABLE III. Field representations under the alternative residual symmetry \( P \otimes Z_3 \), where note that the matter parity group \( P \) is isomorphic to a \( Z_3 \) group generated by \( P_M \), while the quotient group \( Z_3 \) is generated by \(|p^o|\).

Hence, the model provides a natural stability mechanism for two-component dark matter, in which a dark matter component transforms nontrivially under the matter parity group \( P = Z_2 \) not confused with the beginning \( Z_2 \) in \( [19] \), i.e. in \( 1' \) of \( P \) characterized by \( P_M = -1 \), while the remaining dark matter component transforms nontrivially under the quotient group \( Z_3 \), i.e. in \( 1' \) or \( 1'' \) of \( Z_3 \) characterized by \(|p| = [w] \rightarrow w \) or \(|p^o| = [w^2] \rightarrow w^2 \), respectively. Upon \( P \otimes Z_3 \), let us assume the simplest dark matter candidates, as summarized in Table IV. Note that \( B-L \) charge of each dark field can deviate from the supplied value by an arbitrary even number that does not change the representations, because of the cyclic property of the residual symmetries.\(^2\) Here, \( F \) and \( \Phi \) mean fermion and scalar dark fields, respectively. Further, we assume the net mass of \( F_1 \) and \( F_2 \) is smaller than that of \( \Phi \).

The dark matter component stabilized by \( P \) (i.e., \( F_1 \)) can have an arbitrary mass. This is also valid for the dark matter component stabilized by \( Z_3 \) (i.e., \( F_2 \)), even though this component may be heavier than the light quarks \( u,d \), which all transform nontrivially under \( Z_3 \). Indeed, the \( Z_3 \) dark matter component must be color neutral, hence cannot decay to any colored final state, such as single quarks, because of the color conservation. This color conservation requires a color-neutral final state, if it comes from a dark matter decay. Obviously, the color-neutral final state if containing quarks must take only combinations of \( qf \) and/or \( qqq \). It follows that the final state is invariant (i.e. singlet) under \( Z_3 \) too, hence cannot be the product of any \( Z_3 \) dark matter decay, because of the \( Z_3 \) conservation. In other words, the \( SU(3)_C \) and \( Z_3 \) symmetries jointly suppress the decay of \( Z_3 \) dark matter component (i.e. stable), even if this component has a mass larger than that of quark.

| Field \([m, n, \nu, e(d), \Phi, \chi, A]\) | \( P_M \) | \( P \) | \( Z_3 \) |
|-----------------|------|------|------|
| \( F_1 \sim (1, 1, 0, 0) \) | \(-1\) | \( 1 \) | \( 1 \) |
| \( F_2 \sim (1, 1, 0, 1/3) \) | \( 1 \) | \( 1/2 \) | \( 1/2 \) |
| \( \Phi \sim (1, 1, 0, -1/3) \) | \(-1\) | \( 1' \) | \( 1' \) |

TABLE IV. Simplest dark matter candidates implied by the residual symmetry \( P \otimes Z_3 \), where the \( P \) and \( Z_3 \) representations are determined by the matter parity \( P_M \) and the quotient generator \(|p^o|\), respectively.

With this proposal, we have the novel, simplest model for two-component dark matter based upon \( F_1 \) and \( F_2 \) self-interacting through a heavier dark field \( \Phi \), which is of course implied by the residual symmetry \( P \otimes Z_3 \), thus the canonical seesaw. [We can have other scenarios for two-component dark matter, if more dark fields are introduced, but they are complicated and suppressed.] Notice that since \( F_{1,2} \) and \( \Phi \) are the standard model singlets, the \( U(1)_{B-L} \) dynamics is crucially/sufficiently governing the dark matter observables, besides the known consequences of neutrino mass and baryon asymmetry \([27]\).

See saw implication for the XENON1T excess. The XENON1T experiment has recently reported an excess in electronic recoil energy ranging from 1 keV to 7 keV, peaked around 2.4 keV, having a local statistical significance above 3σ \([14]\). Such signal of electron recoils seems to reveal the existence of a structured dark sector \([24, 50]\). Indeed, the dark matter component that scatters off electrons should be fast moving \( v \sim 0.03-0.25 \) for the dark matter mass \( m_2 \sim 0.1 \) MeV to 10 GeV, which exceeds the velocity of cold dark matter \( v \sim 10^{-3} \) (cf. \([25]\)).

This fast dark matter component \( (F_2) \) may be generated locally as a boosted dark matter from the annihilation or semi-annihilation of the heavier dark matter component \( (F_1) \), which is nicely implicated by our model. As a matter of the fact, the heavier dark matter component \( F_1 \sim (1, 1, 0, 0) \) which interacts with normal matter

\(^2\) The actual \( B-L \) charges of the dark fields would be phenomenologically determined.
only via gravity would dominate the cold dark matter, set
by its annihilation or co-annihilation to the lighter dark
matter component $F_2$. The lighter dark matter com-
ponent $F_2$ sub-dominates the dark matter abundance since
it strongly couples to normal matter via the $Z'$ portal.

The relevant Lagrangian terms are

$$\mathcal{L} \supset \bar{F}_1(i\gamma^\mu\partial_\mu - m_1)F_1 + \bar{F}_2(i\gamma^\mu D_\mu - m_2)F_2$$

$$+ \left[(D^\mu\Phi)(D_\mu\Phi) - m_2^2\Phi\Phi\right] + \bar{c}(i\gamma^\mu D_\mu - m_e)e + (h\bar{F}_1F_2\Phi + H.c.),$$

(27)

where $D_\mu = \partial_\mu + ig_{B-L}(B - L)Z'_\mu$ and the dark matter
masses obey $m_0 > m_1 + m_2$ and $m_1 > m_2$. Since the
$B - L$ charge of $F_1$ is fixed, the remaining dark fields can
possess more general $B - L$ charges,

$$F_2 \sim (1,1,0,1/3+2n), \quad \Phi \sim (1,1,0,-1/3-2n),$$

(28)

for $n = 0, \pm 1, \pm 2, \cdots$, as mentioned.

![FIG. 1. Annihilation (left) and co-annihilation (right) pro-
cesses of $F_1$ that set the cold dark matter density.](image)

The relic density of $F_1$ is governed by Feynman di-
agrams in Fig. 1. The co-annihilation process is only
enhanced when the masses of $F_1$ and $\Phi$ are highly degene-
rate. However, this work signifies $m_0 > m_1 + m_2 > m_1$
such that the co-annihilation contribution is negligible.
The dark matter abundance is given by the $F_1$ annihila-
tion in the left diagram.

Applying the Feynman rules, we obtain the thermal
average cross-section times relative velocity as

$$\langle \sigma v_{rel} \rangle \simeq \frac{|h|^4m_i^2}{8\pi(m_i^2 + m_0^2)} \left(1 - \frac{m_0^2}{m_i^2}\right)^{3/2},$$

(29)

which relates to the $F_1$ abundance, $\Omega h^2 \approx 0.1$ pb$/\langle \sigma v_{rel} \rangle$,
where $h$ is the reduced Hubble parameter without con-
fusion. Using the experimental data $\Omega h^2 \approx 0.12$ [1] and
the fact that $m_0 > m_1 > m_2 > m_e$, we get the constraint of
the dark matter self-coupling to be

$$|h| \simeq 0.015 \sqrt{\frac{m_0^2 + m_1^2}{m_1 \times 1 \text{ GeV}}} \left(1 - \frac{m_2^2}{m_1^2}\right)^{-3/8}$$

$$> 5 \times 10^{-4}.$$  

(30)

Of course, at present, $F_2$ is locally generated by the left
diagram in Fig. 1 which subsequently scatters off elec-
trons in the XENON1T experiment through the diagram in
Fig. 2. In the limit of mediator mass $m_{Z'}$, to be much

![FIG. 2. Scattering process of the boosted dark matter $F_2$
with electrons in the XENON1T experiment.](image)

larger than the momentum transfer, the $F_2$-electron scat-
tering cross-section can be written as [27]

$$\sigma_e = \frac{g_{B-L}^2(1/3 + 2n)^2m_e^2}{\pi m_{Z'}^2}.\tag{31}$$

This leads to the number of the signal events as related
to the scattering cross-section by [35, 58]

$$N_{\text{sig}} = \frac{1.6 \times \sigma_e}{3 \times 10^5 \text{ pb}} \left(\frac{1 \text{ GeV}}{m_1}\right)^2.\tag{32}$$

We require the number of the signal events about
100/ton/year in order to explain the XENON1T excess.
This yields the mass of dominant dark matter $F_1$ as re-
lated to the $U(1)_{B-L}$ breaking scale,

$$m_1 \simeq 0.001 \times |1 + 6n| \times (\text{GeV}^3/\Lambda^2).$$

(33)

Since the dominant dark matter $F_1$ is thermally pro-
duced, its mass should obey $m_1 > m_2 > m_e$ as given.
This suggests an upper bound on $\Lambda$ to be

$$\Lambda < 1.5\sqrt{|1 + 6n|} \text{ GeV}.\tag{34}$$

Additionally, the new physics scale must satisfy $\Lambda > O(1 \text{ TeV})$ in order for the seesaw mechanism properly
working. The free parameter $n$ that relates to the $B - L$
charges of $F_2$ and $\Phi$ obeys

$$|n| > 0.74 \times 10^5.$$  

(35)

Further, one demands a perturbative condition for the
$U(1)_{B-L}$ gauge interaction, i.e. $|1/3 + 2n|g_{B-L} < \sqrt{4\pi}$, which along with the above result implies

$$g_{B-L} < 2.4 \times 10^{-5},$$

(36)

corresponding to the $Z'$ mass bounded as

$$m_{Z'} < 50 \text{ MeV}.\tag{37}$$

It is verified that the $F_2$ relic density is negligible, where $F_2$ completely annihilates to the standard model
particles via the (s-channel) $Z'$ portal.

The small coupling and the mass of $Z'$ obviously sat-
sify the low energy constraints from the electron-positron
colliding experiment KLOE2 [59], NA64 experiment [60],
TEXONO experiment [61], or $(g - 2)_{\mu,e}$ [62].
Conclusion. We have discovered a seminal consequence of the canonical seesaw mechanism in addition to the known result of leptogenesis, such that this neutrino mass generation scheme with $B - L$ gauge completion resolves the long-standing hypothesis of structured dark matter stability. The seesaw scale has a nontrivial physical vacuum that preserves two residual $B - L$ symmetries related to the usual matter parity $P_M = (-1)^{(B - L) + 2}$ and the new $Z_3$ quotient generator $[p^z] = [a^{(B - L)}]$, respectively. This yields a novel scenario of two-component dark matter appropriate to the recent XENON1T experiment, where the cold dark matter $F_1$ has $B - L = 0$, while the boosted dark matter $F_2$ has a $B - L$ charge deviating from $1/3$ by five order of magnitude, which is allowed by the cyclic property of the residual generators. $F_{1,2}$ possess masses beyond the electron mass, while the $B - L$ gauge boson has the gauge coupling and mass limited below $2.4 \times 10^{-5}$ and 50 MeV, respectively. If the XENON1T anomaly is relaxed, this setup can provide a generic scenario of two-component dark matter weakly interacting with normal matter.

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