Scotogenic dark matter and Dirac neutrinos from unbroken gauged $B − L$ symmetry

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We propose a simple extension of the standard model where neutrinos get naturally small "scotogenic" Dirac masses from an unbroken gauged $B − L$ symmetry, ensuring dark matter stability. The associated gauge boson gets mass through the Stueckelberg mechanism. Two scenarios are identified, and the resulting phenomenology briefly sketched.

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1. Introduction

Amongst the major drawbacks of the Standard Model (SM) is the absence of neutrino mass and the lack of a viable dark matter candidate. Amendments for both of these issues require new physics. Particle dark matter candidates should be stable, at least on cosmological time scales [1]. A simple way to ensure this is through the imposition of an adequate protecting symmetry whose nature is unknown. For example, dark matter stability can result from a residual $\mathbb{Z}_2$ matter-parity symmetry [2], from the R-parity symmetry in supersymmetric models [3] or from some $\mathbb{Z}_n$-like symmetry, such as quarticity [4–6]. Although these in general are *ad hoc* assumptions, it could be that they can follow naturally from the spontaneous breaking of an extended gauge symmetry [7–9].

A specially attractive possibility is that neutrino mass and dark matter have a common origin, i.e. the same physics being responsible for both. For example, dark matter could be mediator of neutrino mass generation [10–13]. Also the symmetry stabilising dark matter could be closely related to neutrinos. For example, it could be an unbroken subgroup of the flavour symmetry that helps understand the neutrino oscillation parameters [14–17]. In some cases this will lead to Dirac neutrinos, obtained as a consequence of flavour symmetry imposition [18,19].

Likewise, the dark matter stabilising symmetry could be a residual $\mathbb{Z}_n$ subgroup of lepton number symmetry or $B − L$. This may,

again, lead to Dirac neutrinos. Such a possibility was explored in [20], assuming that $B − L$ is spontaneously broken down to a $\mathbb{Z}_n$ symmetry stabilising dark matter, as well as in [21], where the residual dark matter stabilising symmetry follows from the soft breaking of $B − L$.

On the other hand, conservation of the full ungauged $B − L$ symmetry could stabilise dark matter. A scenario of this type has been suggested within a bound-state dark matter scenario [22]. Indeed, this is a justified hypothesis since, despite decade-long searches [23], there has been no experimental evidence of $B − L$ breakdown.

In this letter we consider the alternative case of gauged unbroken $B − L$ as the dark matter stabilisation symmetry. The promotion of the accidental $B − L$ global symmetry of the standard model to a local one stands out for its simplicity, since the inclusion of three right-handed neutrinos ν_{iR} is enough to make it anomaly free and hence consistent. Clearly $B − L$ preserving models are viable provided the associated Z' boson develops an adequate mass. Here we study a Stueckelberg [24] $B − L$ extension of the standard model with naturally small neutrino masses. These are achieved through the scotogenic approach, while the unbroken $B − L$ symmetry is responsible for both the Dirac nature of the neutrino mass and the stabilisation of a dark matter candidate.

The letter is organized a follow. In Sec. 2 we describe the theoretical setup, in Sec. 3 we discuss the scalar sector and in Sec. 4 we describe the Stueckelberg mechanism, while in Sec. 5 we give the mass generation mechanisms in the two alternative realisations of our scenario. Finally, in Sec. 6 we briefly comment on the phenomenology and summarize.

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Table 1
Standard Model fermions and scalars and their gauge transformation properties and the global $B - L$.

Fields	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_{B-L}$
L_{iL}	1	2	$-1/2$	-1
e_{iR}	1	1	-1	-1
Q_{iL}	3	2	$1/6$	$1/3$
u_{iR}	3	1	$2/3$	$1/3$
d_{iR}	3	1	$-1/3$	$1/3$
H	1	2	$1/2$	0

2. Two scenarios

We start from the basic setup provided by the standard model fermion and scalar sectors, as defined in Table 1. This has an automatic global "baryon number minus lepton number" symmetry, we call simply $U(1)_{B-L}$. This symmetry, however, cannot be directly promoted to a local one, as it exhibits non-vanishing $[U(1)_{B-L}]^3$ and $[Grav]^2 \times [U(1)_{B-L}]$ anomalies. Therefore, in order to "gauge" $U(1)_{B-L}$ consistently, we need to extend the standard model field content so as to ensure anomaly cancellation.[1]

A simple way to achieve this is by adding three right-handed fermions, ν_{iR}, singlets under the SM gauge group but with a $B - L$ charge of -1. Adding the ν_{iR} allows for a tree-level Dirac mass for neutrinos through the term $y\overline{L_L}\tilde{H}\nu_{iR}$. Viable neutrino masses require very tiny Yukawa couplings, $y \lesssim \mathcal{O}(10^{-11})$. On the other hand, if the ν_{iR} have a Majorana mass term, the $B - L$ symmetry is broken, and the active neutrinos get masses via the type-I seesaw mechanism, requiring the $B - L$ scale to be around the unification scale for Yukawa couplings of order 1.

Here we propose alternative ways of generating small neutrino masses, relying neither on unnaturally suppressed Yukawa couplings nor on inaccessibly large energy scales. At the same time, in our scenario dark matter candidates mediate neutrino mass generation in a scotogenic fashion, and are stabilised by the conserved $B - L$ symmetry. We suggest two standard model extensions with gauged $B - L$ and dark matter as the mediator of neutrino mass generation.

In Model A, we add a new doublet η and a singlet σ in the scalar sector, both charged under the $B - L$ symmetry. In the fermion sector, three right-handed neutrinos, ν_{iR}, as well as three gauge singlet fermions, S_{iR}, are introduced. The particle content and symmetry properties are given in Table 2. In addition to the gauge symmetries, a $\mathbb{Z}_2$ is imposed, under which all the standard fields in Table 1 transform trivially. Notice that, even though there is a Majorana mass term for the three 2-component SM singlet fermions S_R,[2] it conserves $B - L$. As a result, this is consistent with the Dirac nature of the light neutrinos.

We also propose a variant, Model B, shown in Table 3. This differs from Model A due to introduction of three extra fermion fields: S_L. In contrast to Model A, the charges of the new fields, except for those of ν_{iR}, are not fixed but defined by a single integer $n \neq 0$. Notice that, although S_L and S_R are SM singlets, they carry nonzero $B - L$ charges.

In both cases, the $\mathbb{Z}_2$ symmetry in Tables 2 and 3 prevents the appearance of $\overline{L_{iL}}\tilde{H}\nu_{iR}$ in the Yukawa sector. As a result there are no tree-level neutrino masses. Nevertheless, this symmetry is broken in the scalar sector, making it possible for neutrinos to get calculable Dirac masses at the one-loop level.

Table 2
SM extension A: new fields and their symmetry properties.

Fields	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_{B-L}$	$\mathbb{Z}_2$
ν_{iR}	1	1	0	-1	$-$
S_{iR}	1	1	0	0	$+$
η	1	2	$1/2$	1	$+$
σ	1	1	0	1	$-$

Table 3
SM extension B: new fields and their symmetry properties. Here, $n(\neq 0) \in \mathbb{Z}$.

Fields	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_{B-L}$	$\mathbb{Z}_2$
ν_{iR}	1	1	0	-1	$-$
S_{iL}	1	1	0	$2n$	$+$
S_{iR}	1	1	0	$2n$	$+$
η	1	2	$1/2$	$2n+1$	$+$
σ	1	1	0	$2n+1$	$-$

Notice that the $U(1)_{B-L}$ remains exactly conserved. This implies that the $\mathbb{Z}_2$ group, called matter parity, generated by $M_P = (-1)^{3(B-L)+2s}$, where s is the field's spin, is also exactly conserved. Under M_P, the SM fields and ν_{iR} transform trivially. On the other hand, S, η and σ, in Tables 2 and 3, are M_P-odd fields. Therefore, the lightest among them is stable by matter parity and can play the role of dark matter. In either Model A or B, the dark matter candidate can be scalar or fermionic. When fermionic, dark matter would be Majorana-type in Model A and Dirac-type in Model B.

3. Scalar spectrum

Taking into account the fields and symmetries in Tables 2 and 3, the simplest scalar potential, shared by the two schemes sketched above, can be written as

$$V = \sum_{s=H,\eta,\sigma} \left[\mu_s^2 (s^\dagger s) + \lambda_s (s^\dagger s)^2 \right] + \lambda_{H\eta}(H^\dagger H)(\eta^\dagger \eta)$$
$$+ \lambda'_{H\eta}(H^\dagger \eta)(\eta^\dagger H) + \lambda_{H\sigma}(H^\dagger H)(\sigma^* \sigma) + \lambda_{\eta\sigma}(\eta^\dagger \eta)(\sigma^* \sigma)$$
$$+ \frac{\mu_3}{\sqrt{2}}(\eta^\dagger H \sigma + h.c.), \tag{1}$$

where the last term breaks the $\mathbb{Z}_2$ symmetry softly, with the mass parameter μ_3 assumed to be real for simplicity.

Assuming that $B - L$ is not broken, it is easy to see that the matter parity of the new scalars is the opposite of that of the standard Higgs doublet H. As a result, M_P conservation requires the new scalars η and σ not to acquire any vev, and hence they do not mix with H. Therefore, similar to the standard model case, when the neutral component of H acquires a vev, the $SU(2)_L \otimes U(1)_Y$ group is broken down to a $U(1)_Q$ subgroup, generated by the conventional electric charge operator $Q = T_3 + Y$. The CP-even field in the neutral component of H becomes massive with $m_{h^0}^2 = 2\lambda_H v^2$ and is identified with the 125 GeV Higgs boson observed at the LHC in 2012. The remaining components of H are absorbed by the gauge sector, through the Higgs mechanism, making the $W^\pm$ and Z vector bosons massive.

The other scalars are in the M_P-odd or "dark sector" and do not acquire a vev. The first component of the scalar doublet η corresponds to a massive charged scalar field, $\eta^\pm$, whose mass is

$$m_{\eta^\pm}^2 = \frac{\lambda_{H\eta} v^2}{2} + \mu_\eta^2. \tag{2}$$

The spontaneous symmetry breaking through $\langle H \rangle$ induces a mixing between the second component of η, η^0, and the singlet σ, arising from $(\sigma, \eta^0) M_\varphi^2 (\sigma, \eta^0)^\dagger$, where

$$M_\varphi^2 = \frac{1}{2} \begin{pmatrix} 2\mu_\sigma^2 + \lambda_{H\sigma} v^2 & \mu_3 v \\ \mu_3 v & 2\mu_\eta^2 + (\lambda_{H\eta} + \lambda'_{H\eta})v^2 \end{pmatrix}. \qquad (3)$$

Upon diagonalising the mass matrix above, we find two complex neutral scalars in the spectrum

$$\begin{pmatrix} \varphi_1^0 \\ \varphi_2^0 \end{pmatrix} = \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} \sigma \\ \eta^0 \end{pmatrix}, \quad \text{with}$$

$$2\theta = \arctan(\epsilon) = \arctan\left[\frac{2\mu_3 v}{2(\mu_\sigma^2 - \mu_\eta^2) + (\lambda_{H\sigma} - \lambda_{H\eta} - \lambda'_{H\eta})v^2} \right]. \qquad (4)$$

Notice that when the $\mathbb{Z}_2$-soft-breaking term, μ_3, goes to zero, the mixing angle θ also vanishes. The mass eigenvalues associated with such states are

$$m_{\varphi_{(1,2)}^0}^2 = \frac{1}{4}\Bigg\{ 2(\mu_\eta^2 + \mu_\sigma^2) + v^2(\lambda_{H\eta} + \lambda'_{H\eta} + \lambda_{H\sigma})$$

$$\mp \mathcal{F}\sqrt{\left[2(\mu_\sigma^2 - \mu_\eta^2) + v^2(\lambda_{H\sigma} - \lambda_{H\eta} - \lambda'_{H\eta})\right]^2 + 4\mu_3 v^2} \Bigg\}, \qquad (5)$$

respectively, where $\mathcal{F} = 1$ for $(M_\varphi^2)_{22}/(M_\varphi^2)_{11} > 1$ and $\mathcal{F} = -1$ otherwise.

It is worth noticing that in Eq. (3), the real and imaginary parts of the scalar fields appear together, that is, real and imaginary parts are degenerate in mass. That means that, if dark matter is scalar, it is described by a complex field, in contrast to conventional scotogenic scenarios [10,13] in which they are nearly degenerate but not exactly so.

4. Stueckelberg mechanism

Neglecting kinetic mixing,[3] after electroweak spontaneous symmetry breaking the Z' boson associated with $U(1)_{B-L}$ remains unmixed with the standard EW gauge bosons. Its massive nature can be described by the kinetic Lagrangian [24]

$$\mathcal{L}_{\text{kin}}^{\text{St}} = -\frac{1}{4}Z'^{\mu\nu}Z'_{\mu\nu} + \frac{1}{2}(M_{Z'}Z'^\mu - \partial^\mu A)^2, \qquad (6)$$

which is invariant under the $U(1)_{B-L}$ gauge transformations

$$Z'^\mu \to Z'^\mu + \partial^\mu \Lambda,$$
$$A \to A + M_{Z'}\Lambda, \qquad (7)$$

where A is a scalar Stueckelberg compensator and $Z'^{\mu\nu} = \partial^\mu Z'^\nu - \partial^\nu Z'^\mu$. Upon gauge-fixing, implemented by the R_ξ gauge term

$$\mathcal{L}_{\text{fg}}^{\text{St}} = -\frac{1}{2\xi}(\partial_\mu Z'^\mu + M_{Z'}\xi A)^2, \qquad (8)$$

the Z' boson acquires mass $M_{Z'}$ and the auxiliary field A decouples, as

$$\mathcal{L}_{\text{kin}}^{\text{St}} + \mathcal{L}_{\text{fg}}^{\text{St}} = -\frac{1}{4}Z'^{\mu\nu}Z'_{\mu\nu} + \frac{1}{2}M_{Z'}^2 Z'^\mu Z'_\mu - \frac{1}{2\xi}(\partial_\mu Z'^\mu)^2$$

$$+ \frac{1}{2}\partial^\mu A \partial_\mu A - \frac{1}{2}M_{Z'}^2 \xi A^2, \qquad (9)$$

up to a total derivative. Here $M_{Z'}$ is a free parameter of the model, unrelated to any vev and disconnected from the neutrino mass generation mechanism.

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[3] Kinetic mixing has been discussed in Refs. [31–33].

The relevant Z' interactions are

$$\mathcal{L}_{\text{f}}^{Z'} = g'Z'_\mu \sum_{i=1}^{3} \Bigg[\frac{1}{3}(\bar{u}_i\gamma^\mu u_i + \bar{d}_i\gamma^\mu d_i) - \bar{e}_i\gamma^\mu e_i - \bar{\nu}_i\gamma^\mu \nu_i$$

$$+ 2n\bar{S}\gamma^\mu S \Bigg],$$

$$\mathcal{L}_{\text{s}}^{Z'} = ig'(2n+1)Z'_\mu \Bigg[\eta^- \partial^\mu \eta^+ - \eta^+ \partial^\mu \eta^-$$

$$+ \sum_{i=1}^{2} \left(\varphi_i^{0*}\partial^\mu \varphi_i^0 - \varphi_i^0 \partial^\mu \varphi_i^{0*} \right) \Bigg] \qquad (10)$$

$$+ g'^2(2n+1)^2 Z'^\mu Z'_\mu \left(\eta^-\eta^+ + \sum_{i=1}^{2} \varphi_i^{0*}\varphi_i^0 \right),$$

$$\mathcal{L}_{\text{s+g}}^{Z'} = 2eg'(2n+1)Z'^\mu \Big\{ \left[A_\mu + \cot(2\theta_W)Z_\mu \right] \eta^-\eta^+$$

$$- \csc(2\theta_W)Z_\mu \left| \varphi_1^0 \cos\theta - \varphi_2^0 \sin\theta \right|^2$$

$$+ \frac{\csc\theta_W}{\sqrt{2}} \left[W_\mu^+ \eta^-(\varphi_1^0 \cos\theta - \varphi_2^0 \sin\theta) + \text{h.c.} \right] \Big\},$$

with $n = 0$ in Model A and $n(\neq 0) \in \mathbb{Z}$ in Model B.

There are no gauge-mediated flavour-changing neutral currents and both $B - L$ and M_P remain unbroken to all orders in perturbation theory, preserving the Dirac nature of neutrinos.

5. Scotogenic neutrino masses

We now give the most general renormalisable Yukawa Lagrangians for our models. According to Tables 2 and 3 they can be written, respectively, as follows

$$-\mathcal{L}_Y^A = y^e \overline{L_L}H e_R + y^\nu \overline{L_L}\tilde{\eta}S_R + h\overline{(S_R)^c}\sigma \nu_R + \frac{1}{2}M_S^M \overline{(S_R)^c}S_R$$

$$+ \text{h.c.}, \qquad (11)$$

$$-\mathcal{L}_Y^B = y^e \overline{L_L}H e_R + y^\nu \overline{L_L}\tilde{\eta}S_R + h\overline{S_L}\sigma \nu_R + M_S^D \overline{S_L}S_R + \text{h.c.},$$

where the flavour indices have been omitted.

Due to the conservation of $B - L$ and M_P neutrino masses are not generated at the tree level, arising only as a calculable one-loop contribution via the diagrams in Fig. 1. In both cases the neutrino masses have the same form

$$(m_\nu)_{ij} = \frac{\sin(2\theta)}{32\pi^2} \sum_k y_{ik}^\nu h_{kj} m_{S_k}$$

$$\times \left[\frac{m_{\varphi_1}^2}{m_{\varphi_1}^2 - m_{S_k}^2} \ln \frac{m_{\varphi_1}^2}{m_{S_k}^2} - \frac{m_{\varphi_2}^2}{m_{\varphi_2}^2 - m_{S_k}^2} \ln \frac{m_{\varphi_2}^2}{m_{S_k}^2} \right], \qquad (12)$$

where, for model A, m_{S_k} are the eigenvalues of the Majorana mass matrix M_S^M, while, for model B, m_{S_k} are the Dirac mass matrix M_S^D eigenvalues. In the limit of a small $\mathbb{Z}_2$ soft-breaking term, i.e. $\mu_3 \ll 1$, we have that $\theta \ll 1$, then $\sin(2\theta) \simeq \epsilon$, with $\epsilon = \epsilon(\mu_3) \ll 1$, as defined in Eq. (4).

The internal fields in the loop are odd under matter parity, while the others are even. The lightest among the M_P-odd fields is stable and, if electrically neutral, can play the role of dark matter. Assuming that the charged component of the scalar doublet η is heavier than the other M_P-odd fields, the model can have either a complex neutral scalar or a fermion as dark matter. In the latter case it can either be a Majorana or a Dirac fermion.
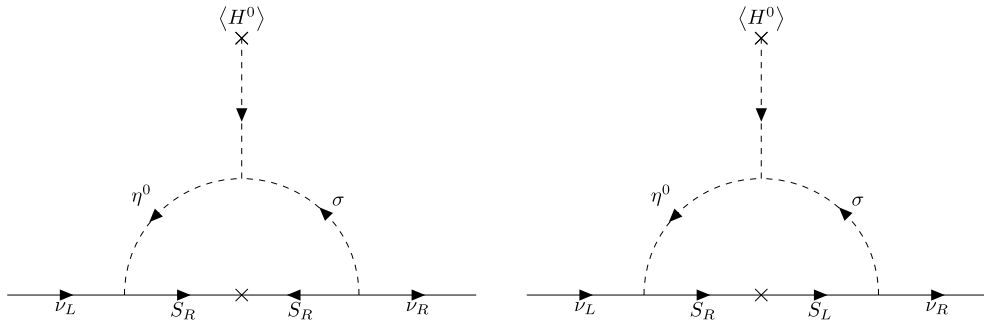
Fig. 1. One-loop Dirac neutrino masses for Model A and B, respectively.

6. Comments on phenomenology

The class of models suggested here has a broad range of phenomenological implications. Some of these are present in the minimal Stueckelberg $B - L$ extension of the standard model studied in Ref. [34]. For example, the existence of a Z' associated to the conserved $B - L$ symmetry implies new gauge couplings of the standard model fermions, as seen in Eq. (10). Hence its effects should be manifest at high energies, say, in electron-positron or proton-proton collisions. This implies that the ratio of the Z' mass and its corresponding gauge coupling is constrained by collider data. There are limits coming from Tevatron [35], LEPII and LHC [34]. The most stringent current value is

$$M_{Z'}/g' \geq 6.9 \text{ TeV at } 95\% \text{ C.L.} \qquad (13)$$

Besides the physics of the Z', our proposal harbours scotogenic dark matter [10], made stable by the conservation of $B - L$. The scenario differs from the proposal in Ref. [22] in that the dark matter here is elementary, and can be light. The pure dark matter phenomenology has similarities with the non-scotogenic scenario discussed in [36]. Let us first comment on the possibility of dark matter being a scalar candidate. In this case, the lightest among the M_P-odd complex scalars φ_1^0 and φ_2^0 will be stable and can play the role of dark matter. Consistency with direct detection experiments requires the coupling between the complex DM candidate and the Z-boson to be very small. This can be easily achieved here if the mixing between φ_1^0 and φ_2^0 is very small, and the lightest state is φ_1^0. In this case, the dark matter candidate φ_1^0 is mostly the scalar singlet σ, and couples to the Z-boson only through its suppressed mixing with η^0. Notice that since the mixing angle θ is governed by the $\mathbb{Z}_2$-soft-breaking parameter μ_3, it can be made naturally small since its absence is associated with an enhanced symmetry, and hence protected in 't Hooft sense.

The fate of mixed complex dark matter has been analysed in Ref. [37] in a simpler phenomenological setup with no Z' boson. In the case where only the Higgs and the Z-boson portals are available, the region of the parameter space compatible with the observed relic abundance and direct detection experiments is, in general, very constrained, unless co-annihilation takes place due to φ_1^0 and φ_2^0 being almost degenerate. The allowed region is considerably widened in the presence of a Majorana fermion, like S_R in our Model A, acting as a new channel for dark matter annihilation.

On the other hand, the dark matter candidate can be one of the neutral fermions, say S_1, if it is the lightest M_p-odd particle. For the case of Model A, the dominant process contributing to the thermal relic density of the Majorana fermion dark matter candidate S_1 is driven by the Yukawa couplings y^ν and h in Eq. (11). This scenario is analogous to the original scotogenic model [10], where fairly large Yukawas are required to produce

the correct dark matter abundance, in potential tension with experimental bounds from Lepton Flavour Violation processes like $\mu \rightarrow e\gamma$ [38]. Nevertheless, even in this case, certain regions of the parameter space are phenomenologically viable, as investigated in Ref. [39]. For the case of Model B with a Dirac fermion dark matter, there are new processes involved in setting the relic density, mediated by the $B - L$ gauge boson, according to the interactions shown in Eq. (10). Assuming the Z' to be the dominating channel, the correct relic density can be successfully reproduced around the resonance condition $M_{Z'} \approx 2M_{S_1}^D$ for any $n \neq 0$ in Table 3 [40]. Finally, it is worth emphasising that, although some of the dark matter scenarios are more constrained than others, all of them exhibit viable regions of the parameter space.

In short, we have proposed a simple extension of the standard model where neutrinos get naturally small "scotogenic" Dirac-type masses from an unbroken gauged $B - L$ symmetry. The associated gauge boson gets mass through the Stueckelberg mechanism. The conservation of $B - L$ and matter parity play a key role in ensuring dark matter stability. Dedicated studies of the resulting phenomenology will be presented elsewhere.

Declaration of competing interest

The authors declare that they have no known competing financial interests or personal relationships that could have appeared to influence the work reported in this paper.

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