Physics Letters B 750 (2015) 135–138

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Gauge B–L model with residual Z_3 symmetry

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ARTICLE INFO

Article history:
Received 22 July 2015
Received in revised form 28 August 2015
Accepted 3 September 2015
Available online 7 September 2015
Editor: J. Hisano

ABSTRACT

We study a gauge B–L extension of the standard model of quarks and leptons with unconventional charges for the singlet right-handed neutrinos, and extra singlet scalars, such that a residual Z_3 symmetry remains after the spontaneous breaking of B–L. We discuss the phenomenological consequences of this scenario, including the possibility of long-lived self-interacting dark matter and Z' collider signatures.

Lepton number L is a familiar concept. It is usually defined as a global $U(1)$ symmetry, under which the leptons of the standard model (SM), i.e. e, μ, τ together with their neutrinos ν_e, ν_μ, ν_τ have $L = 1$, and all other SM particles have $L = 0$. In the case of nonzero Majorana neutrino masses, this continuous symmetry is broken to a discrete Z_2 symmetry, i.e. $(-1)^L$ or lepton parity. In this paper, we consider a gauge B–L extension of the SM, such that a residual Z_3 symmetry remains after the spontaneous breaking of B–L. This is then a realization of the unusual notion of Z_3 lepton symmetry. It has specific phenomenological consequences, including the possibility of a long-lived particle as a dark-matter candidate.

The conventional treatment of gauge B–L has three right-handed singlet neutrinos $\nu_{R1}, \nu_{R2}, \nu_{R3}$ transforming as $-1, -1, -1$ under B–L. It is well known that this assignment satisfies all the anomaly-free conditions for $U(1)_{B-L}$. However, another assignment [1]

$$\nu_{R1}, \nu_{R2}, \nu_{R3} \sim 5, -4, -4 \tag{1}$$

works as well, because

$$5 - 4 - 4 = -3, \quad (5)^3 - (4)^3 - (4)^3 = -3. \tag{2}$$

To obtain a realistic model with this assignment, it was recently proposed [2] that three additional neutral singlet Dirac fermions $N_{1,2,3}$ be added with B–$L = -1$, together with a singlet scalar χ_3 with B–$L = 3$. Consequently, the tree-level Yukawa couplings $\bar{\nu}_L N_R \bar{\phi}^0$ and $\bar{N}_L \nu_{R2} \chi_3$, $\bar{N}_L \nu_{R3} \chi_3$ are allowed, where $\Phi = (\phi^+, \phi^0)$ is the one Higgs doublet of the SM. Together with the invariant $\bar{N}_L N_R$ mass terms, the 6×5 neutrino mass matrix linking $(\bar{\nu}_L, \bar{N}_L)$ to (ν_R, N_R) is of the form

$$\mathcal{M}_{\nu N} = \begin{pmatrix} 0 & \mathcal{M}_0 \\ \mathcal{M}_3 & \mathcal{M}_N \end{pmatrix}, \tag{3}$$

where $\mathcal{M}_0$ and $\mathcal{M}_N$ are 3×3 mass matrices and $\mathcal{M}_3$ is 3×2 because ν_{R1} has no tree-level Yukawa coupling. This means that one linear combination of ν_L is massless. Of course, if the dimension-five term $\bar{\nu}_{R1} N_L \chi_3^2$ also exists, then $\mathcal{M}_3$ is 3×3 and $\mathcal{M}_{\nu N}$ is 6×6.

The form of $\mathcal{M}_{\nu N}$ allows nonzero seesaw Dirac neutrino masses for ν [3], i.e.

$$\mathcal{M}_\nu \simeq \mathcal{M}_0 \mathcal{M}_N^{-1} \mathcal{M}_3. \tag{4}$$

Without the implementation of a flavor symmetry, any 3×3 $\mathcal{M}_\nu$ is possible. Although the gauge B–L is broken, a residual global L symmetry remains in this model with ν, l, N all having $L = 1$. Because the pairing of any two neutral fermions of the same chirality always results in a nonzero B–L charge not divisible by 3 units in this model, it is impossible to construct an operator of any dimension for a Majorana mass term which violates B–L. Hence the neutrinos are indeed exactly Dirac.

We now add two more scalar singlets: χ_2 with B–$L = 2$ and χ_6 with B–$L = -6$. The important new terms in the Lagrangian are

$$\bar{N}_L \nu_{R1} \chi_6, \quad \chi_2 N_L N_L, \quad \chi_2 N_R N_R, \quad \chi_2^3 \chi_6, \quad \chi_3^2 \chi_6. \tag{5}$$

Now B–L is broken by $\langle \chi_3 \rangle = u_3$ as well as $\langle \chi_6 \rangle = u_6$, and all neutrinos become massive. If χ_2 is absent, then again a residual global L symmetry exists with $L = 1$ for ν, l, N and $L = 0$ for $\chi_{3,6}$. However, the existence of χ_2 shows that the residual symmetry is then Z_3, such that χ_2 and all leptons transform as $\omega = \exp(2\pi i/3)$ under Z_3 with $\chi_{3,6} \sim 1$. This is thus the first example of a lepton

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http://dx.doi.org/10.1016/j.physletb.2015.09.010

symmetry which is not Z_2 (for Majorana neutrinos), nor $U(1)$ or Z_4 [4,5] (for Dirac neutrinos). Note that Z_3 is also sufficient to guarantee that all the neutrinos remain Dirac.

Although there is no stabilizing symmetry here for dark matter, χ_2 has very small couplings to two neutrinos through the Yukawa terms of Eq. (5) from the mixing implied by Eq. (3). This means that χ_2 may have a long enough lifetime to be suitable for dark matter, as shown below.

Consider for simplicity the coupling of χ_2 to just one N, with the interaction

$$\mathcal{L}_{int} = \frac{1}{2} f_L \chi_2 N_L N_L + \frac{1}{2} f_R \chi_2 N_R N_R + \text{H.c.} \tag{6}$$

Let the ν_L–N_L mixing be $\zeta_0 = m_0/m_N$ and ν_R–N_R mixing be $\zeta_3 = m_3/m_N$, then the decay rate of χ_2 is

$$\Gamma(\chi_2 \to \bar{\nu}\bar{\nu}) = \frac{m_\chi}{32\pi}(f_L^2 \zeta_0^4 + f_R^2 \zeta_3^4). \tag{7}$$

If we set this equal to the age of the Universe (13.75×10^9 yr), and assuming $m_\chi = 100$ GeV, $f_L = f_R$ and $\zeta_0 = \zeta_3$, then $f\zeta^2 = 8.75 \times 10^{-22}$. Hence

$$\sqrt{f}\zeta << 3 \times 10^{-11} \tag{8}$$

would guarantee the stability of χ_2 to the present day, and allow it to be a dark-matter candidate. This sets the scale of m_N at about 10^{13} GeV, which is also the usual mass scale for the heavy Majorana singlet neutrino in the canonical seesaw mechanism.

In this model, there is of course a gauge boson Z' which couples to B–L. Its production at the Large Hadron Collider (LHC) is due to its couplings to quarks. Once produced, it decays into quarks and leptons. In the conventional B–L assignment for ν_R, its branching fractions to quarks, charged leptons, and neutrinos are 1/4, 3/8, and 3/8 respectively. In this model, the ν_R charges are $(5, -4, -4)$, hence their resulting partial widths are very large. Assuming that Z' decays also into χ_2, the respective branching fractions into quarks, charged leptons, neutrinos, and χ_2 as dark matter are then 1/18, 1/12, 5/6, and 1/36. This means Z' has an 86% invisible width. Using the production of Z' via $u\bar{u}$ and $d\bar{d}$ initial states at the LHC and its decay into e^-e^+ or $\mu^-\mu^+$ as signature, the current bound on $m_{Z'}$ assuming $g' = g$, i.e. the $SU(2)_L$ gauge coupling of the SM, is about 3 TeV, based on recent LHC data [6,7]. However, because the branching fraction into l^-l^+ is reduced by a factor of 2/9 in our B–L model, this bound is reduced to about 2.5 TeV, again for $g' = g$. There is also a similar bound [8] from precision $e^-e^+ \to e^-e^+$ measurements at the Large Electron Positron Collider (LEP), i.e. $m_{Z'}/g' >$ a few TeV.

Since χ_2 interacts with nuclei through Z', there is also a significant constraint from dark-matter direct-search experiments. The cross section per nucleon is given by

$$\sigma_0 = \frac{1}{\pi}\left(\frac{m_\chi m_n}{m_\chi + A m_n}\right)^2 \left(\frac{2g'^2}{m_{Z'}^2}\right)^2, \tag{9}$$

where A is the number of nucleons in the target and m_n is the nucleon mass. Consider for example $m_\chi = 100$ GeV, then $\sigma_0 < 1.25 \times 10^{-45}$ cm^2 from the recent LUX data [9]. This implies $m_{Z'}/g' > 16.2$ TeV, as shown in Fig. 1. If $g' = g$, then $m_{Z'} > 10.6$ TeV. This limit is thus much more severe than the LHC bound of 2.5 TeV. If $g' < g$, then both the LHC and LUX bounds on $m_{Z'}$ are relaxed. However, it also means that it is unlikely that Z' would be discovered at the LHC even with the 14 TeV run.

Consider now the annihilation cross section of $\chi_2\chi_2^*$ for obtaining its thermal relic abundance. The process $\chi_2\chi_2^* \to Z' \to$ SM

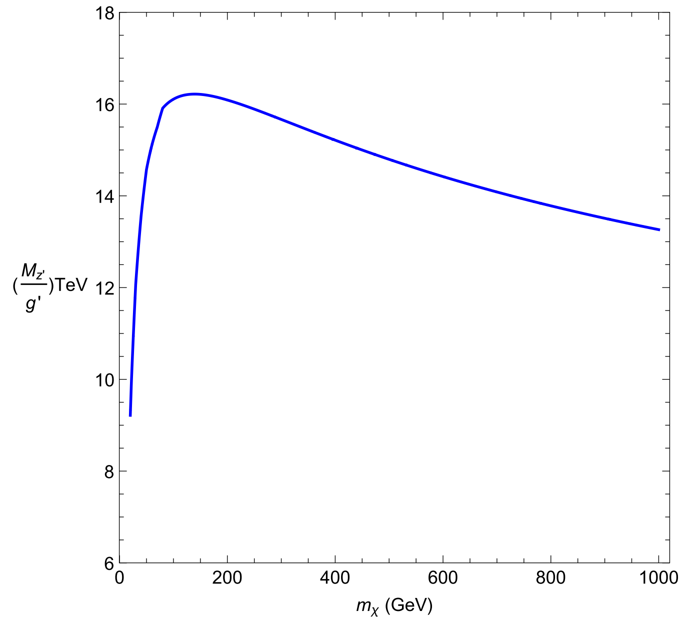

Fig. 1. Lower bound on $m_{Z'}/g'$ versus m_χ from LUX data.

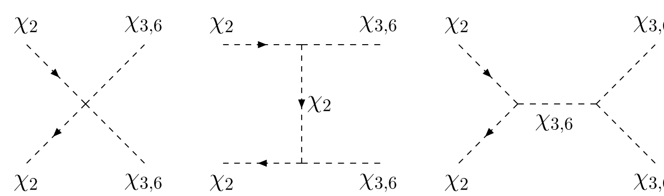

Fig. 2. $\chi_2\chi_2^\dagger$ annihilation to $\chi_{3,6}$ final states.

particles is p-wave suppressed and is unlikely to be strong enough for this purpose. We may then consider the well-studied process $\chi_2\chi_2^* \to h \to$ SM particles, where h is the SM Higgs boson. If this is assumed to account for all of the dark-matter relic abundance of the Universe, then it has recently been shown [10] that the required strength of this interaction is in conflict with LUX data except for a small region near $m_\chi = m_h/2$. On the other hand, another analysis [11] claims that a region with m_χ somewhat greater than m_h is still allowed.

In this paper, we will consider the following alternative scenario. We assume that the $h\chi_2\chi_2^*$ interaction is negligible, so that neither Higgs nor Z' exchange is important for $\chi_2\chi_2^*$ annihilation. Instead we invoke the new interactions of Fig. 2. Since $\chi_{3,6}$ may interact freely with h, thermal equilibrium is maintained with the other SM particles. This scenario requires of course that m_χ to be greater than at least one physical mass eigenvalue in the $\chi_{3,6}$ sector.

To summarize, $\chi_2 \sim \omega$ under Z_3 and decays into two antineutrinos, but its lifetime is much longer than the age of the Universe. It is thus an example of Z_3 dark matter [12–16]. It is also different from previous Z_2 proposals [17,18] based on Ref. [1]. It has significant elastic interactions with nuclei through Z' and Higgs exchange and may be discovered in direct-search experiments. On the other hand, its relic abundance is determined not by Z' or Higgs interactions, but by its annihilation to other scalars of this model which maintain thermal equilibrium with the SM particles through the SM Higgs boson. Note that this is also the mechanism used in a recently proposed model of vector dark matter [19].

We now discuss the details of the scalar sector of this model. Consider the scalar potential

$$\begin{aligned} V = & -\mu_0^2(\Phi^\dagger\Phi) + m_2^2(\chi_2^*\chi_2) - \mu_3^2(\chi_3^*\chi_3) - \mu_6^2(\chi_6^*\chi_6) \\ & + \frac{1}{2}\lambda_0(\Phi^\dagger\Phi)^2 + \frac{1}{2}\lambda_2(\chi_2^*\chi_2)^2 + \frac{1}{2}\lambda_3(\chi_3^*\chi_3)^2 \\ & + \frac{1}{2}\lambda_6(\chi_6^*\chi_6)^2 + \lambda_{02}(\chi_2^*\chi_2)(\Phi^\dagger\Phi) \\ & + \lambda_{03}(\chi_3^*\chi_3)(\Phi^\dagger\Phi) + \lambda_{06}(\chi_6^*\chi_6)(\Phi^\dagger\Phi) \\ & + \lambda_{23}(\chi_2^*\chi_2)(\chi_3^*\chi_3) + \lambda_{26}(\chi_2^*\chi_2)(\chi_6^*\chi_6) \\ & + \lambda_{36}(\chi_3^*\chi_3)(\chi_6^*\chi_6) \\ & + [\frac{1}{2}f_{36}(\chi_3^2\chi_6) + \text{H.c.}] + [\frac{1}{6}\lambda'_{26}(\chi_2^3\chi_6) + \text{H.c.}]. \end{aligned} \quad (10)$$

Let $\langle\phi^0\rangle = v$, $\langle\chi_3\rangle = u_3$, $\langle\chi_6\rangle = u_6$, then the minimum of V is determined by

$$\mu_0^2 = \lambda_0 v^2 + \lambda_{03}u_3^2 + \lambda_{06}u_6^2, \quad (11)$$

$$\mu_3^2 = \lambda_3 u_3^2 + \lambda_{03}v^2 + \lambda_{36}u_6^2 + f_{36}u_6, \quad (12)$$

$$\mu_6^2 = \lambda_6 u_6^2 + \lambda_{06}v^2 + \lambda_{36}u_3^2 + \frac{f_{36}u_3^2}{2u_6}. \quad (13)$$

There is one dark-matter scalar boson χ_2 with mass given by

$$m_\chi^2 = m_2^2 + \lambda_{02}v^2 + \lambda_{23}u_3^2 + \lambda_{26}u_6^2. \quad (14)$$

There is one physical pseudoscalar boson

$$A = \sqrt{2}\text{Im}(2u_6\chi_3 + u_3\chi_6)/\sqrt{u_3^2 + 4u_6^2} \quad (15)$$

with mass given by

$$m_A^2 = -f_{36}(u_3^2 + 4u_6^2)/2u_6. \quad (16)$$

There are three physical scalar bosons spanning the basis $[h, \sqrt{2}\text{Re}(\chi_3), \sqrt{2}\text{Re}(\chi_6)]$, with 3×3 mass-squared matrix given by

$$M^2 = \begin{pmatrix} 2\lambda_0 v^2 & 2\lambda_{03}u_3 v & 2\lambda_{06}u_6 v \\ 2\lambda_{03}u_3 v & 2\lambda_3 u_3^2 & 2\lambda_{36}u_3u_6 + f_{36}u_3 \\ 2\lambda_{06}u_6 v & 2\lambda_{36}u_3u_6 + f_{36}u_3 & 2\lambda_6 u_6^2 - f_{36}u_3^2/2u_6 \end{pmatrix}. \quad (17)$$

For illustration, we consider the special case $\lambda_{03} = \lambda_{06} = 0$, so that h decouples from $\chi_{3,6}$. It then becomes identical to that of the SM, and may be identified with the 125 GeV particle discovered [20,21] at the LHC. We now look for a solution with

$$S = \sqrt{2}\text{Re}(-u_3\chi_3 + 2u_6\chi_6)/\sqrt{u_3^2 + 4u_6^2}, \quad (18)$$

$$S' = \sqrt{2}\text{Re}(2u_6\chi_3 + u_3\chi_6)/\sqrt{u_3^2 + 4u_6^2}, \quad (19)$$

as mass eigenstates. This is easily accomplished for example with

$$u_3 = 2u_6, \quad 4\lambda_3 = \lambda_6 - f_{36}/u_6. \quad (20)$$

In this case,

$$S = -\text{Re}\chi_3 + \text{Re}\chi_6, \; m_S^2 = 2\lambda_6 u_6^2 - 4\lambda_{36}u_6^2 - 4f_{36}u_6, \quad (21)$$

$$S' = \text{Re}\chi_3 + \text{Re}\chi_6, \; m_{S'}^2 = 2\lambda_6 u_6^2 + 4\lambda_{36}u_6^2, \quad (22)$$

$$A = \text{Im}\chi_3 + \text{Im}\chi_6, \; m_A^2 = -4f_{36}u_6, \quad (23)$$

$$m_{Z'} = 12g'u_6. \quad (24)$$

The couplings of $\chi_2\chi_2^*$ to S and S' are given by

$$\chi_2\chi_2^*[u_6(\lambda_{26} - 2\lambda_{23})S + u_6(\lambda_{26} + 2\lambda_{23})S']. \quad (25)$$

Since S plays the same role in breaking B–L as the Higgs boson h does in breaking $SU(2)_L \times U(1)_Y$, it is expected to be massive of order $\sqrt{u_3^2 + 4u_6^2} = 2\sqrt{2}u_6$. This allows $m_{S'}$ to be adjusted to be very small, then it may serve as a light scalar mediator for χ_2 as self-interacting dark matter [22]. This is not a necessary assumption of our model and requires fine tuning of scalar parameters to achieve. We merely want to demonstrate that such a possible scenario exists within our model. For $m_{S'} \simeq 0$, we need $\lambda_{36} = -\lambda_6/2$. In that case, using Eq. (20), we find

$$m_S^2 = 16\lambda_3 u_6^2, \quad m_A^2 = m_S^2 - 4\lambda_6 u_6^2. \quad (26)$$

We assume that the relic density of χ_2 is dominated by the $\chi_2\chi_2^*$ annihilation to $S'S'$. This may have to be revised if the semi-annihilation $\chi_2\chi_2^* \to \chi_2 S'$ is sizeable. Here we simply assume that λ'_{26} is small. For illustration, we set to zero the $\chi_2\chi_2^*S'S'$ coupling, i.e. $\lambda_{23} + \lambda_{26} = 0$, as well as the $SS'S'$ coupling, i.e. $-12\lambda_3 + 6\lambda_6 + 2\lambda_{36} - f_{36}/u_6 = 0$. This implies $\lambda_3 = \lambda_6/2$ so that the $S'S'S'$ coupling is also zero and $m_A^2 = m_S^2/2$. This choice of parameters means that only the middle diagram of Fig. 2 contributes to the $\chi_2\chi_2^*$ annihilation cross section with

$$\sigma \times v_{rel} = \frac{1}{64\pi m_\chi^2}\left|\frac{\lambda_{26}^2 u_6^2}{m_\chi^2}\right|^2. \quad (27)$$

Equating this to the optimal value [23] of 4.4×10^{-26} cm^3 s^{-1} for the correct dark-matter relic density of the Universe, we find for $m_\chi = 100$ GeV

$$\lambda_{26} = 0.0295\left(\frac{1\text{ TeV}}{u_6}\right). \quad (28)$$

We assume of course that $m_A > 2m_\chi$.

For S' to be in thermal equilibrium with the SM particles, we consider nonzero values of λ_{03} and λ_{06}. This is possible in our chosen parameter space if $2\lambda_{03} + \lambda_{06} \simeq 0$, so that the $S'h$ mixing is very small and yet the $S'S'h$ coupling $\lambda_{06}v/4\sqrt{2}$ and $S'S'hh$ coupling $\lambda_{06}/16$ may be significant. Note that the Sh mixing is now fixed at $(\lambda_{06}/\lambda_6)(v/2\sqrt{2}u_6)$ which may yet be suitably suppressed for h to be essentially the one Higgs boson of the SM. Even if $\lambda_{03,06}$ are negligible, the gauge interaction $S'AZ'$ may also be sufficient to maintain thermal equilibrium. This may also affect the magnitude of the self-interacting $\chi_2\chi_2^*$ cross section.

The $h \to S'S'$ decay width is given by

$$\Gamma(h \to S'S') = \frac{\lambda_{06}^2 v^2}{256\pi m_h} = \left(\frac{\lambda_{06}}{0.04}\right)^2 0.5\text{ MeV}. \quad (29)$$

It is invisible at the LHC because S' decays slowly to e^-e^+ only through its mixing with h, if $m_{S'} \sim 10$ MeV for S' as a light mediator for the self-interacting dark matter χ_2.

In conclusion, we have considered the unusual case of a gauge B–L symmetry which is spontaneously broken to Z_3 lepton number. Neutrinos are Dirac fermions transforming as $\omega = \exp(2\pi i/3)$ under Z_3. A complex neutral scalar χ_2 exists which also transforms as ω. It is not absolutely stable, but decays to two antineutrinos with a lifetime much greater than that of the Universe. It is thus an example of Z_3 dark matter. In addition to the one Higgs boson h of the SM, there are three neutral scalars S, S', A and one heavy vector gauge boson Z'. From direct-search experiments, $m_{Z'}/g'$ is constrained to be very large, thus making it impossible to discover Z' at the LHC even with the current run. The relic abundance of χ_2 is determined by its annihilation into S' which is a candidate for the light mediator by which χ_2 obtains its long-range self-interaction.

Acknowledgement

This work is supported in part by the U.S. Department of Energy under Grant No. DE-SC0008541.

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