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# Variants of the dark left-right gauge model: Neutrinos and scotinos

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In the recently proposed dark left-right gauge model (DLRM) of particle interactions, the usual left-handed lepton doublet $(\nu, e)_L$ transforming under $SU(2)_L$ is accompanied by the *unusual* right-handed fermion doublet $(n, e)_R$ transforming under $SU(2)_R$, where $n_R$ is *not* the Dirac mass partner of $\nu_L$. In this scenario, whereas $\nu_L$ is certainly a neutrino, $n_R$ should be considered a *scotino*, i.e. a dark-matter fermion. Variants of this basic idea are discussed, including its minimal *scotogenic* realization.

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## I. INTRODUCTION

The gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$ of the standard model (SM) of particle interactions treats left-handed and right-handed fermions differently, with the electric charge given by $Q = T_{3L} + Y$. To restore left-right symmetry, it is often proposed that the extension $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ be considered, where $Q = T_{3L} + T_{3R} + (B - L)/2$. In that case, the fermion content of the SM gains one extra particle, i.e. $\nu_R$ in the right-handed lepton doublet $(\nu, l)_R$. Connecting this with the usual left-handed lepton doublet $(\nu, l)_L$ through a Higgs bidoublet, $\nu_R$ pairs with $\nu_L$ to obtain a Dirac mass, just as $l_R$ does with $l_L$. Assuming $SU(2)_R \times U(1)_{B-L}$ is broken to $U(1)_Y$ through a Higgs triplet transforming as (1, 1, 3, 1), $\nu_R$ gets a large Majorana mass, thereby inducing a small seesaw mass for $\nu_L$. The above is a well-known scenario for what the addition of $\nu_R$ would do for understanding the existence of tiny neutrino masses. For a more general discussion of the $SU(2)_R$ breaking scale, see Ref. [1].

Suppose the mass connection between $\nu_R$ and $\nu_L$ is severed without affecting $l_R$ and $l_L$, then $\nu_L$ and $\nu_R$ can be different particles, with their own interactions. Whereas $\nu_L$ is clearly still the well-known neutrino, $\nu_R$ may become something else entirely. As shown in Ref. [2], it may in fact be a *scotino*, i.e. a dark-matter fermion, and to avoid confusion, it is renamed $n_R$. This is accomplished in a nonsupersymmetric $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)$ model with the imposition of a global $U(1)$ symmetry $S$, such that the breaking of $SU(2)_R \times S$ will leave the generalized lepton number $L = S - T_{3R}$ unbroken. It is called the dark left-right model (DLRM), to distinguish it from the alternative left-right model (ALRM) proposed 22 years ago [3,4] which has the same crucial property that $n_R$ is *not* the mass partner of $\nu_L$.

## II. FERMION CONTENT

The fermion structure of the DLRM under $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1) \times S$ is given by [2]

$$\psi_L = \begin{pmatrix} \nu \\ e \end{pmatrix}_L \sim (1, 2, 1, -1/2; 1), \qquad \psi_R = \begin{pmatrix} n \\ e \end{pmatrix}_R \sim (1, 1, 2, -1/2; 1/2), \tag{1}$$

$$Q_L = \begin{pmatrix} u \\ d \end{pmatrix}_L \sim (3, 2, 1, 1/6; 0), \qquad d_R \sim (3, 1, 1, -1/3; 0), \tag{2}$$

$$Q_R = \begin{pmatrix} u \\ h \end{pmatrix}_R \sim (3, 1, 2, 1/6; 1/2), \qquad h_L \sim (3, 1, 1, -1/3; 1), \tag{3}$$

where $h$ is a new heavy quark of charge $-1/3$. The above fermionic content was first studied in Ref. [5] and also in Ref. [6], without the identification of $n_R$ as a scotino.

To allow $e_L$ to pair with $e_R$ to form a Dirac fermion, the Higgs bidoublet

$$\Phi = \begin{pmatrix} \phi_1^0 & \phi_2^+ \\ \phi_1^- & \phi_2^0 \end{pmatrix} \sim (1, 2, 2, 0; 1/2) \tag{4}$$

is added so that $m_e$ is obtained from $v_2 = \langle \phi_2^0 \rangle$. At the same time, $\nu_L$ is connected to $n_R$ through $\phi_1^0$. However, $\langle \phi_1^0 \rangle = 0$ will be maintained because $\phi_1^0$ has $S - T_{3R} = 1$, whereas that of $\phi_2^0$ is zero. As shown in Ref. [2], the spontaneous breaking of $SU(2)_R \times S$ leaves the residual symmetry $L = S - T_{3R}$ unbroken, where $L$ is the conventional lepton number assigned to $\nu$ and $e$. Here $n$ has $L = S - T_{3R} = 0$, whereas $W_R^\pm$ has $L = S - T_{3R} = \mp 1$, and it does not mix with $W_L^\pm$, in contrast to the case of the conventional left-right model, where such mixing is unavoidable. Further, the bidoublet

$$\tilde{\Phi} = \sigma_2 \Phi^* \sigma_2 = \begin{pmatrix} \bar{\phi}_2^0 & -\phi_1^+ \\ -\phi_2^- & \bar{\phi}_1^0 \end{pmatrix} \sim (1, 2, 2, 0; -1/2) \tag{5}$$

is prevented by $S$ from coupling $\psi_L$ to $\psi_R$, thereby ensuring the absence of tree-level flavor-changing neutral currents, which was not possible in the conventional non-supersymmetric left-right model.

In the quark sector, $Q_L$ couples to $Q_R$ through $\tilde{\Phi}$, but not $\Phi$. Hence $m_u$ is obtained from $v_2$, and there is no mixing between $d$ and $h$. The former has $L = 0$, but the latter has $L = S - T_{3R} = 1$. For $d_L$ to pair with $d_R$, and $h_R$ to pair with $h_L$, the Higgs doublets

$$\Phi_L = \begin{pmatrix} \phi_L^+ \\ \phi_L^0 \end{pmatrix} \sim (1, 2, 1, 1/2; 0), \quad \Phi_R = \begin{pmatrix} \phi_R^+ \\ \phi_R^0 \end{pmatrix} \sim (1, 1, 2, 1/2; -1/2) \tag{6}$$

are needed. Note that $v_4 = \langle \phi_R^0 \rangle$ will break $SU(2)_R \times U(1)$ to $U(1)_Y$ as desired, and the leptoquark $h$ gets a heavy mass of order $v_4$.

## III. EXOTIC VARIANTS

The fermion sector may be more exotic. For example, Eq. (1) may be replaced by

$$\psi_L = \begin{pmatrix} \nu \\ e \end{pmatrix}_L \sim (1, 2, 1, -1/2; 1), \quad e_R \sim (1, 1, 1, -1; 1), \tag{7}$$

$$\psi_R = \begin{pmatrix} n \\ E \end{pmatrix}_R \sim (1, 1, 2, -1/2; -1/2), \quad E_L \sim (1, 1, 1, -1; 0). \tag{8}$$

In this case, $E$ has $L = 0$, $n$ has $L = -1$, $m_e$ comes from $v_3 = \langle \phi_L^0 \rangle$, $m_E$ from $v_4$, and neither $\Phi$ nor $\tilde{\Phi}$ couples to $\bar{\psi}_L \psi_R$.

As another example, Eqs. (2) and (3) may be replaced by

$$Q_L = \begin{pmatrix} u \\ d \end{pmatrix}_L \sim (3, 2, 1, 1/6; 0), \qquad u_R \sim (3, 1, 1, 2/3; 0), \tag{9}$$

$$Q_R = \begin{pmatrix} f \\ d \end{pmatrix}_L \sim (3, 1, 2, 1/6; -1/2), \quad f_L \sim (3, 1, 1, 2/3; -1). \tag{10}$$

In this case, $f$ is an exotic quark of charge 2/3 and $L = -1$. Here, $Q_L$ couples to $Q_R$ through $\Phi$, but not $\tilde{\Phi}$, $m_d$ comes from $v_2$, $m_u$ from $v_3$, and $m_f$ from $v_4$.

## IV. MASSES FOR $\nu_L$ AND $n_R$

With the above Higgs content, $\nu_L$ and $n_R$ remain massless. Consider now the various ways that they acquire masses:

(1) In the DLRM [2], Higgs triplets under $SU(2)_L$ and $SU(2)_R$ are used separately for $\nu_L$ and $n_R$ masses.

(2) In Ref. [5], they are massless.

(3) In Ref. [6], they acquire radiative masses separately from the addition of two charged scalar singlets.

(4) In the ALRM [3], the usual lepton doublet is actually part of a bidoublet:

$$\begin{pmatrix} \nu_e & E^c \\ e & N_E^c \end{pmatrix}_L \sim (1, 2, 2, 0), \tag{11}$$

which means that $\nu_e$ and $e$ have $SU(2)_R$ interactions. In the original proposal, $\nu_L$ and $n_R$ are massless, but they can acquire seesaw masses separately through the many other fields available in the $\underline{27}$ representation of $E_6$, as explained in Ref. [7]. One of the three $n_R$ copies in this supersymmetric model pairs with the neutral gaugino from the breaking of $SU(2)_R \times U(1) \to U(1)_Y$ to form a Dirac fermion. The other two are light and considered as sterile neutrinos which mix with $\nu_L$ through the soft term $\nu_e N_E^c - e E^c$ which breaks $R$ parity.

(5) A simple variation of the DLRM also allows neutrino masses to be radiatively generated by dark matter (i.e. *scotogenic*) in one loop [8–23]. Instead of $\Delta_L$, a scalar singlet $\chi \sim (1, 1, 1, 0; -1)$ is added, then the trilinear scalar term $\mathrm{Tr}(\Phi \tilde{\Phi}^\dagger)\chi$ is allowed. Using the soft term $\chi^2$ to break $L$ to $(-)^L$, a scotogenic neutrino mass is obtained as shown in Fig. 1. It is also possible to do this in two loops [24,25] and three loops [26–28].

(6) Since the scotogenic mechanism of Fig. 1 does not care how $n_R$ acquires a Majorana mass, it may be accomplished with three neutral singlet fermions $n_L$ with $S = 0$ instead of the Higgs triplet $\Delta_R$. Now the Yukawa coupling $\bar{n}_L(n_R \phi_R^0 - e_R \phi_R^+)$ is allowed, as well as a Majorana mass for $n_L$. Hence $n_R$ gets an induced Majorana mass which is essential for Fig. 1. Note that $n_L$ does not couple to $(\nu_L \phi_L^0 - e_L \phi_L^+)$ because of $S$.

In this minimal variant, the $Z'$ mass comes entirely from the $\Phi_R$ doublet as in the ALRM, hence the prediction

$$M_{W_R}^2 = \frac{(1 - 2x)}{(1 - x)} M_{Z'}^2 + \frac{x^2}{(1 - x)^2} M_{W_L}^2, \tag{12}$$

where $x \equiv \sin^2\theta_W$ and zero $Z - Z'$ mixing has been as-

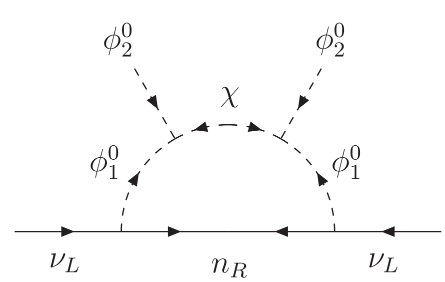


FIG. 1. One-loop scotogenic neutrino mass.

sumed [4], i.e. $v_2^2/(v_2^2 + v_3^2) = x/(1-x)$. Currently, the experimental bound on $M_{Z'}$ is 850 GeV [2].

The diagram of Fig. 1 is exactly calculable [8]. The $\bar{n}_R \nu_L \bar{\phi}_1^0$ coupling is given by $(m_\alpha/v_2)U_{\alpha i}$, where $\alpha = e, \mu, \tau$ and $i = 1$ to 6 refer to the mass eigenstates of the $6 \times 6$ $(n_R, n_L)$ mass matrix

$$\mathcal{M}_n = \begin{pmatrix} 0 & m_D \\ m_D & m_n \end{pmatrix}. \tag{13}$$

In the bases $(\mathrm{Re}\phi_1^0, \mathrm{Re}\chi)$, $(\mathrm{Im}\phi_1^0, \mathrm{Im}\chi)$, the respective mass-squared matrices are

$$\begin{pmatrix} m_\phi^2 & \mu v_2 \\ \mu v_2 & m_\chi^2 + \mu_\chi^2 \end{pmatrix}, \qquad \begin{pmatrix} m_\phi^2 & -\mu v_2 \\ -\mu v_2 & m_\chi^2 - \mu_\chi^2 \end{pmatrix}. \tag{14}$$

Let their mass eigenstates and mixing angles be $(m_{R1}^2, m_{R2}^2, \theta_R)$ and $(m_{I1}^2, m_{I2}^2, \theta_I)$, then

$$(\mathcal{M}_\nu)_{\alpha\beta} = \sum_i \frac{m_\alpha m_\beta U_{\alpha_i} U_{\beta i} M_i}{16\pi^2 v_2^2}\Bigg[\cos^2\theta_R \frac{m_{R1}^2}{m_{R1}^2 - M_i^2} \ln\frac{m_{R1}^2}{M_i^2} - \cos^2\theta_I \frac{m_{I1}^2}{m_{I1}^2 - M_i^2} \ln\frac{m_{I1}^2}{M_i^2} + \sin^2\theta_R \frac{m_{R2}^2}{m_{R2}^2 - M_i^2} \times \ln\frac{m_{R2}^2}{M_i^2} - \sin^2\theta_I \frac{m_{I2}^2}{m_{I2}^2 - M_i^2} \ln\frac{m_{I2}^2}{M_i^2}\Bigg]. \tag{15}$$

In the limit $\mu_\chi^2 = 0$, $\mathcal{M}_\nu$ vanishes because $m_{R1} = m_{I1}$, $m_{R2} = m_{I2}$, and $\theta_R = -\theta_I$. In the limit $\mu = 0$, $\mathcal{M}_\nu$ also vanishes because $m_{R1} = m_{I1}$ and $\theta_R = \theta_I = 0$. Furthermore, if $m_n = 0$ in Eq. (13), $\mathcal{M}_\nu$ is again zero because $n$ is a Dirac particle. This latter is an example of the inverse seesaw mechanism [29–33]. Hence neutrino masses are suppressed in this scenario by three possible limits, and the scale of $SU(2)_R$ breaking (associated with the masses of the dark-matter particles of Fig. 1) may well be as low as 1 TeV, as advocated.

## V. CONCLUSION

The neutral component $n$ of the $SU(2)_R$ doublet $(n, e)_R$ is proposed as a dark-matter fermion (scotino). Variants of this basic idea, the dark left-right gauge model, are discussed. A minimal version is considered, where neutrino masses are radiatively generated by dark matter (scotogenic) and naturally suppressed, allowing the $SU(2)_R$ breaking scale to be as low as 1 TeV.

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