Dark Left-Right Gauge Model: SU(2)$_R$ Phenomenology

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In the recently proposed dark left-right gauge model of particle interactions, the left-handed fermion doublet $(\nu, e)_L$ is connected to its right-handed counterpart $(n, e)_R$ through a scalar bidoublet, but $\nu_L$ couples to $n_R$ only through $\phi^0_1$ which has no vacuum expectation value. The usual $R$ parity, i.e. $R = (-)^{3B+L+2j}$, can be defined for this nonsupersymmetric model so that both $n$ and $\Phi_1$ are odd together with $W_{R}^{\pm}$. The lightest $n$ is thus a viable dark-matter candidate (scotino). Here we explore the phenomenology associated with the $SU(2)_R$ gauge group of this model, which allows it to appear at the TeV energy scale. The exciting possibility of $Z' \to 8$ charged leptons is discussed.

I. INTRODUCTION

The nonsupersymmetric dark left-right model (DLRM) proposed recently [1] is a variant of a supersymmetric left-right extension of the standard model (SM) of particle interactions based on $E_6$ and inspired by string theory some 23 years ago [2, 3]. It has a number of

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desirable properties, the chief of which is the absence of tree-level flavor-changing neutral currents, thus allowing the $SU(2)_R$ breaking scale to be as low as experimentally allowed by collider data. This became known in the literature as the alternative left-right model (ALRM) \[4\]. Here we explore further consequences of the DLRM, coming from the $SU(2)_R$ sector.

II. MODEL

Consider the gauge group $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1) \times S$, where $S$ is a global symmetry such that the breaking of $SU(2)_R \times S$ will leave the combination $L = S - T_{3R}$ unbroken. This allows $L$ to be a generalized lepton number which is conserved \[1\] in all interactions except those which are responsible for Majorana neutrino masses. The fermion content of the DLRM is given by

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

(1)

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

(2)

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

(3)

This basic structure was already known many years ago \[2, 6\] but without realizing that $n$ is a scotino, i.e. a dark-matter fermion.

The scalar sector of the DLRM consists of one bidoublet and two doublets:

$$
\Phi = \begin{pmatrix} \phi_1^0 & \phi_2^+ \\ \phi_1^- & \phi_2^0 \end{pmatrix}, \quad \Phi_L = \begin{pmatrix} \phi_L^+ \\ \phi_L^0 \end{pmatrix}, \quad \Phi_R = \begin{pmatrix} \phi_R^+ \\ \phi_R^0 \end{pmatrix},
$$

(4)

as well as two triplets for making $\nu$ and $n$ massive separately:

$$
\Delta_L = \begin{pmatrix} \Delta_L^+ / \sqrt{2} & \Delta_L^{++} \\ \Delta_L^- / \sqrt{2} & -\Delta_L^- / \sqrt{2} \end{pmatrix}, \quad \Delta_R = \begin{pmatrix} \Delta_R^+ / \sqrt{2} & \Delta_R^{++} \\ \Delta_R^- / \sqrt{2} & -\Delta_R^- / \sqrt{2} \end{pmatrix}.
$$

(5)

Their assignments under $S$ are listed in Table I.

The Yukawa terms allowed by $S$ are then $\bar{\psi}_L \Phi \psi_R$, $\bar{Q}_L \Phi Q_R$, $\bar{Q}_L \Phi_L d_R$, $\bar{Q}_R \Phi_R h_L$, $\psi_L \psi_L \Delta_L$, and $\psi_R \psi_R \Delta_R$, whereas $\bar{\psi}_L \Phi \psi_R$, $\bar{Q}_L \Phi Q_R$, and $\bar{t}_L d_R$ are forbidden. Hence $m_e, m_u$ come from
TABLE I: Scalar content of proposed model.

| Scalar       | $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)$ | $S$  |
|--------------|-----------------------------------------------|------|
| $\Phi$       | (1, 2, 2, 0)                                    | 1/2  |
| $\tilde{\Phi} = \sigma_2 \Phi^* \sigma_2$ | (1, 2, 2, 0)                                   | −1/2 |
| $\Phi_L$     | (1, 2, 1, 1/2)                                  | 0    |
| $\Phi_R$     | (1, 1, 2, 1/2)                                  | −1/2 |
| $\Delta_L$   | (1, 3, 1, 1)                                    | −2   |
| $\Delta_R$   | (1, 1, 3, 1)                                    | −1   |

$v_2 = \langle \phi_2^0 \rangle$, $m_d$ comes from $v_3 = \langle \phi_L^0 \rangle$, $m_h$ comes from $v_4 = \langle \phi_R^0 \rangle$, $m_\nu$ comes from $v_5 = \langle \Delta_L^0 \rangle$, and $m_n$ comes from $v_6 = \langle \Delta_R^0 \rangle$. This structure shows clearly that flavor-changing neutral currents are guaranteed to be absent at tree level [7].

The generalized lepton number $L = S - T_{3R}$ remains 1 for $\nu$ and $e$, and 0 for $u$ and $d$, but the new particle $n$ has $L = 0$ and $h$ has $L = 1$, whereas $W_R^\pm$ has $L = \mp 1$ and $Z'$ has $L = 0$, etc. As neutrinos acquire Majorana masses, $L$ is broken to $(-)^L$. The generalized $R$ parity is then defined in the usual way, i.e. $(-)^{3B+L+2}$. The known quarks and leptons have even $R$, but $n$, $h$, $W_R^\pm$, $\phi_R^\pm$, $\Delta_R^\pm$, $\phi_T^\pm$, $\text{Re}(\phi_1^0)$, and $\text{Im}(\phi_1^0)$ have odd $R$. Hence the lightest $n$ can be a viable dark-matter candidate if it is also the lightest among all the particles having odd $R$. Note that $R$ parity has now been implemented in a non-supersymmetric model.

III. $SU(2)_R$ HIGGS STRUCTURE

There exists an experimental bound [1] on $M_{Z'}$ of 850 GeV from Tevatron data [8]. As for the recent CDMS-II results [9], they impose no additional constraint because $n$ is Majorana and does not contribute to the s-wave elastic spin-independent scattering cross section through $Z'$ exchange in the nonrelativistic limit. Assuming thus that $M_{Z'} > 850$ GeV only, we study the $SU(2)_R$ Higgs structure of this model and identify those new particles which may be relatively light and be observable at the Large Hadron Collider (LHC). Consider then the most general Higgs potential consisting of $\Phi_R$ and $\Delta_R$: 
\[ V_R = m_4^2 \Phi_R^\dagger \Phi_R + m_6^2 Tr(\Delta_R^\dagger \Delta_R) + \frac{1}{2} \lambda_1 (\Phi_R^\dagger \Phi_R)^2 + \frac{1}{2} \lambda_2 [Tr(\Delta_R^\dagger \Delta_R)]^2 \]
\[ + \frac{1}{4} \lambda_3 Tr(\Delta_R^\dagger \Delta_R - \Delta_R \Delta_R^\dagger)^2 + f_1 (\Phi_R^\dagger \Phi_R) Tr(\Delta_R^\dagger \Delta_R) \]
\[ + f_2 \Phi_R^\dagger (\Delta_R^\dagger \Delta_R - \Delta_R \Delta_R^\dagger) \Phi_R + \mu (\Phi_R^\dagger \Delta_R \Phi_R + \Phi_R^\dagger \Delta_R^\dagger \Phi_R), \]

where

\[ \Phi_R^\dagger \Phi_R = \phi_R^+ \phi_R^- + \overline{\phi}_R^0 \phi_R^0, \]
\[ Tr(\Delta_R^\dagger \Delta_R) = \Delta_R^- \Delta_R^+ + \Delta_R^- \Delta_R^- + \Delta_R^0 \Delta_R^0, \]
\[ \Delta_R^\dagger \Delta_R - \Delta_R \Delta_R^\dagger = \begin{pmatrix} \Delta_R^0 \Delta_R^0 - \Delta_R^- \Delta_R^+ & \sqrt{2}(\Delta_R^- \Delta_R^+ - \Delta_R^0 \Delta_R^0) \\ \sqrt{2}(\Delta_R^- \Delta_R^- - \Delta_R^0 \Delta_R^0) & -\Delta_R^0 \Delta_R^0 + \Delta_R^- \Delta_R^+ \end{pmatrix}, \]
\[ \Phi_R^\dagger \Delta_R^\dagger \Phi_R = \phi_R^0 \phi_R^0 \Delta_R^0 + \sqrt{2} \phi_R^0 \phi_R^0 \Delta_R^- - \phi_R^0 \phi_R^0 \Delta_R^-. \]

Let \( \langle \phi_R^0 \rangle = v_4 \) and \( \langle \Delta_R^0 \rangle = v_6 \), as already noted, then the minimum of \( V_R \) is given by

\[ V_0 = m_4^2 v_4^2 + m_6^2 v_6^2 + \frac{1}{2} \lambda_1 v_4^4 + \frac{1}{2} \lambda_2 v_6^4 + \frac{1}{2} \lambda_3 v_6^4 + f_1 v_4^2 v_6^2 - f_2 v_4^2 v_6^2 + 2 \mu v_4^2 v_6, \]

where \( v_{4,6} \) are determined by

\[ \frac{\partial V_0}{\partial v_4} = 2 v_4 [m_4^2 + \lambda_1 v_4^2 + (f_1 - f_2) v_6^2 + 2 \mu v_6] = 0, \]
\[ \frac{\partial V_0}{\partial v_6} = 2 v_6 [m_6^2 + (\lambda_2 + \lambda_3) v_6^2 + (f_1 - f_2) v_4^2] + 2 \mu v_4^2 = 0. \]

The physical mass-squared matrices are given by

\[ \mathcal{M}^2(Re\phi_R^0, Re\Delta_R^0) = \begin{pmatrix} 2\lambda_1 v_4^2 & 2(f_1 - f_2) v_4 v_6 + 2 \mu v_4 \\ 2(f_1 - f_2) v_4 v_6 + 2 \mu v_4 & 2(\lambda_2 + \lambda_3) v_6^2 - \mu v_4^2 / v_6 \end{pmatrix}, \]
\[ \mathcal{M}^2(Im\phi_R^0, Im\Delta_R^0) = \begin{pmatrix} -4 \mu v_6 & 2 \mu v_4 \\ 2 \mu v_4 & -\mu v_4^2 / v_6 \end{pmatrix}, \]
\[ \mathcal{M}^2(\phi_R^\dagger, \Delta_R^\dagger) = \begin{pmatrix} 2 v_6 (f_2 v_6 - \mu) & -\sqrt{2} v_4 (f_2 v_6 - \mu) \\ -\sqrt{2} v_4 (f_2 v_6 - \mu) & v_4^2 / v_6 (f_2 v_6 - \mu) \end{pmatrix}, \]
\[ \mathcal{M}^2(\Delta_R^{\pm \pm}) = 2 f_2 v_4^2 - 2 \lambda_3 v_6^2 - \mu v_4^2 / v_6. \]

As expected, the linear combinations

\[ \frac{(v_4 Im\phi_R^0 + 2 v_6 Im\Delta_R^0)}{\sqrt{v_4^2 + 4 v_6^2}}, \quad \frac{(v_4 \phi_R^\dagger + \sqrt{2} v_6 \Delta_R^\dagger)}{\sqrt{v_4^2 + 2 v_6^2}} \]
have zero mass, corresponding to the longitudinal components of $Z'$ and $W^\pm_R$. Their orthogonal combinations

$$A_R = \frac{\sqrt{2}(v_4 I m\Delta_R^0 - 2v_6 I m\phi_R^0)}{\sqrt{v_4^2 + 4v_6^2}}, \quad \xi_R^\pm = \frac{(v_4 \Delta_R^\pm - \sqrt{2}v_6 \phi_R^\pm)}{\sqrt{v_4^2 + 2v_6^2}}$$

have mass-squares $-\mu (v_4^2 + 4v_6^2)/v_6$ and $(f_2 - \mu/v_6)(v_4^2 + 2v_6^2)$ respectively. Since $n_R$ couples to $\Delta_R^\pm$, but not to $\phi_R^\pm$, the discussion on dark-matter relic abundance from $nn$ annihilation to lepton pairs through $\Delta_R^\pm$ exchange in Ref. [1] applies only if $v_6^2 << v_4^2$. This turns out to be exactly what the model requires because $m_n$ comes from $v_6$ and $m_n$ of order 200 GeV is needed for dark-matter relic abundance.

To be specific, we will assume in fact that $m_n = 200$ GeV. If this value is changed, some details in the following will be changed, but all the qualitative features of this model will remain. The first thing to notice is that for $m_n = 200$ GeV, Fig. 3 of Ref. [1] requires $m_{\Delta_R^+} = 220$ GeV. From the Yukawa coupling

$$\frac{f_n}{\sqrt{2}}(\Delta_R^0 n_R n_R + \sqrt{2}\delta_R R n_R e_R + \Delta_R^{++} e_R e_R), \quad (18)$$

we get $m_n = \sqrt{2} f_n v_6$. Since $f_n = 1$ is assumed in computing the relic abundance in Ref. [1], we obtain $v_6 = 141$ GeV. Let us now assume $M_{Z'} = 1$ TeV for illustration. Then $v_4 = 1851$ GeV and $M_{W_R} = 832$ GeV, where $v_2 = 95$ GeV and $v_3 = 146$ GeV have been used to ensure zero $Z - Z'$ mixing at tree level (see next section).

The physical charged scalar $\xi_R^+$ is now 99.4% $\Delta_R^+$ and its mass is given by

$$m_{\xi_R^+}^2 = (f_2 - \mu/v_6)(v_4^2 + 2v_6^2) = [220 \text{ GeV}]^2. \quad (19)$$

This implies that $f_2 - \mu/v_6 = 0.014$. We now note that the $\Delta_R$ scalar triplet masses satisfy the important sum rule

$$\frac{m_{\phi_R^+}^2}{1 + 4v_6^2/v_4^2} + m_{\Delta_R^+}^2 = \frac{2m_{\xi_R^+}^2}{1 + 2v_6^2/v_4^2} - 2\lambda_3 v_6^2. \quad (20)$$

This means that both $m_{\phi_R^+}$ and $m_{\Delta_R^+}$ are bounded from above as a function of $\lambda_3$ which should not be larger than about one in magnitude. We plot in Fig. $m_{\Delta_R^+}$ versus $m_{\phi_R^+}$ for various values of $\lambda_3$.

We now come to a very important conclusion. To satisfy the dark-matter relic density in this model, $m_n$ and $m_{\xi_R^+}$ have to be of order 200 GeV. This in turn implies that $m_{\phi_R^+}$
FIG. 1: Plot of $M_{A_R^0}$ versus $M_{A_R^{++}}$ for different values of $\lambda_3$ with $v_2 = 95$ GeV, $v_3 = 146$ GeV, $v_4 = 1851$ GeV, $v_6 = 141$ GeV, $M_{\xi_R^+} = 220$ GeV, and $M_{Z'} = 1$ TeV, $M_{W_R^{\pm}} = 832$ GeV. The lower bound of 110 GeV comes from Tevatron data. The line $m_{A_R} = 2m_{A_R^{++}}$ is also shown.

and $m_{A_R^{++}}$ are bounded in such a way that the decays $A_R \rightarrow nn$ and $\Delta_R^{++} \rightarrow \xi_R^+\xi_R^+$ are kinematically forbidden. This means that the dominant decay of $\Delta_R^{++}$ is into two like-sign leptons, which is a great experimental signature. There is also an allowed region in parameter space which enables the decay $A_R \rightarrow \Delta_R^{++}\Delta_R^{--}$. Note that the present experimental bound on $m_{A^{++}}$ is 110 GeV [10].

As for the remaining two scalar masses from diagonalizing Eq. (14), $H_{R2}^0$ will be heavy with $m_{H_{R2}^0}^2 = 2\lambda_1 v_4^2$, whereas $H_{R1}^0$ will be light with mass given by

$$m_{H_{R1}^0}^2 = \frac{m_{A_R}^2}{1 + 4v_6^2/v_4^2} + \frac{2v_6^2}{\lambda_1}[\lambda_1(\lambda_2 + \lambda_3) - (f_1 - 0.014)^2].$$ (21)
Finally, we need to consider the scalar bidoublet and the scalar left doublet. In this model, \((\phi_0^0, \phi_1^1)\) will be heavy, and the two doublets \((\phi_2^+, \phi_2^0), (\phi_L^+, \phi_L^0)\) are similar to the usual two Higgs doublets considered in the SM. The linear combination \((v_3 \Phi_2 - v_2 \Phi_L)/\sqrt{v_2^2 + v_3^2} = [H_L^+, (H_L^0 + iA_L^0)/\sqrt{2}]\) is physical and light.

### IV. GAUGE SECTOR

Since \(e\) has \(L = 1\) and \(n\) has \(L = 0\), the \(W_R^\pm\) of this model must have \(L = S - T_{3R} = 0 - 1 = -1\). This also means that unlike the conventional LRM, \(W_R^\pm\) does not mix with the \(W_L^\pm\) of the SM at all. This important property allows the \(SU(2)_R\) breaking scale to be much lower than it would be otherwise, as explained already 22 years ago \([2, 3]\). Assuming that \(g_L = g_R\) and let \(x \equiv \sin^2 \theta_W\), then the neutral gauge bosons of the DLRM (as well as the ALRM) are given by

\[
\begin{pmatrix}
A \\
Z \\
Z'
\end{pmatrix}
= \begin{pmatrix}
\sqrt{x} & \sqrt{x} & \sqrt{1-2x} \\
\sqrt{1-x} & -x/\sqrt{1-x} & -\sqrt{x(1-2x)/(1-x)} \\
0 & \sqrt{(1-2x)/(1-x)} & -x/(1-x)
\end{pmatrix}
\begin{pmatrix}
W_L^0 \\
W_R^0 \\
B
\end{pmatrix}. \quad (22)
\]

Whereas \(Z\) couples to the current \(J_{3L} - x J_{em}\) with coupling \(e/\sqrt{x(1-x)}\) as in the SM, \(Z'\) couples to the current

\[
J_{Z'} = xJ_{3L} + (1-x)J_{3R} - xJ_{em}
\]  

with the coupling \(e/\sqrt{x(1-x)(1-2x)}\). The masses of the gauge bosons are given by

\[
M_{W_L}^2 = \frac{e^2}{2x}(v_2^2 + v_3^2), \quad M_Z^2 = \frac{M_{W_L}^2}{1-x}, \quad M_{W_R}^2 = \frac{e^2}{2x}(v_2^2 + v_4^2 + 2v_5^2), \\
M_{Z'}^2 = \frac{e^2(1-x)}{2x(1-2x)}(v_2^2 + v_4^2 + 4v_5^2) - \frac{x^2 M_{W_L}^2}{(1-x)(1-2x)}, \quad (24)
\]

where zero \(Z - Z'\) mixing has been assumed, using the condition \[v_5^2/(v_2^2 + v_3^2) = x/(1-x)\]. Note that in the ALRM, \(\Delta_R\) is absent, hence \(v_6 = 0\) in the above. Also, the assignment of \((\nu, e)_L\) there is different, hence the \(Z'\) of the DLRM is not identical to that of the ALRM. At the LHC, if a new \(Z'\) exists which couples to both quarks and leptons, it will be discovered with relative ease. Once \(M_{Z'}\) is determined, then the DLRM predicts the existence of \(W_R^\pm\) with a mass in the range

\[
\frac{(1-2x)}{2(1-x)} M_{Z'}^2 + \frac{x}{2(1-x)^2} M_{W_L}^2 < M_{W_R}^2 < \frac{(1-2x)}{(1-x)} M_{Z'}^2 + \frac{x^2}{(1-x)^2} M_{W_L}^2. \quad (26)
\]
In the ALRM, since \( v_6 = 0 \), \( M_{W_R} \) takes the value of the upper limit of this range. The prediction of \( W_R^\pm \) in addition to \( Z' \) distinguishes these two models from the multitude of other proposals with an extra \( U(1)' \) gauge symmetry.

V. \( Z' \) Decay

Consider the possible discovery of \( Z' \) at the LHC. For \( M_{Z'} = 1 \) TeV, only an integrated luminosity of 0.2 fb\(^{-1} \) is required \([1]\). Its discovery channel is presumably \( \mu^+\mu^- \), but it will also have 4 charged muons in the final state from \( \Delta_R^\pm \Delta_R'^- \), and perhaps even 8 charged muons, as shown below.

In addition to all SM particles, \( Z' \) also decays into \( n\bar{n}, \Delta_R^\pm \Delta_R'^-, \xi_R^\pm \xi_R'^-, A_R^0 H^0_L, H^+_L H^-_L, \) and \( A_R^0 H^0_L \). In particular, the subsequent decay \( \Delta_R^\pm \rightarrow \mu^\pm \mu^\pm \) will be a unique signature where the like-sign dimuons have identical invariant masses \(^1\).

The interactions of \( Z' \) with fermions come from

\[
\mathcal{L} = -g' Z'_\mu J^\mu_{Z'},
\]

where \( g' = e/\sqrt{x(1-x)(1-2x)} \). Ignoring fermion masses, each fermionic partial width is given by

\[
\Gamma(Z' \rightarrow \bar{f} f) = \left(\frac{g'}{24\pi}\right)^2 M_{Z'}^2 \left|c_L^2 + c_R^2\right|,
\]

where \( c_{L,R} \) are the coefficients from \( J_{Z'} = x J_{3L} + (1-x) J_{3R} - x J_{em} \), and a color factor of 3 should be added for each quark. In the DLRM, we have

\[
u_L = \frac{x}{2}, \quad \nu_R = \frac{1-x}{2}, \quad e_L = \frac{x}{2}, \quad e_R = -\frac{1}{2} + \frac{3x}{2}.
\]

Here we need to consider 3 families for \( u, d, \nu, e \) but only one for \( n \).

The decay of \( Z' \rightarrow A_R^0 H^0_{R1} \) to scalars come from

\[
\mathcal{L} = -g'(1-x) Z'_\mu [(\partial^\mu H^0_{R1}) A_R^0 - (\partial^\mu A_R^0) H^0_{R1}],
\]

with the partial decay width

\[
\Gamma(Z' \rightarrow A_R^0 H^0_{R1}) = \frac{(g')^2 M_{Z'}(1-x)^2}{48\pi},
\]

\(^1\) Not all models involving doubly charged scalars have this decay, see for example \([11]\).
where \((1 - x)\) comes from \(I_{3L} = 0, I_{3R} = -1, Q = 0\). For \(Z' \to \xi^+\xi^-\), the factor is \(x\), coming from \(I_{3L} = 0, I_{3R} = 0, Q = 1\). For \(Z' \to \Delta_R^{++}\Delta_R^{-}\), the factor is \((1 - 3x)\), coming from \(I_{3L} = 0, I_{3R} = 1, Q = 2\).

The decay of \(Z'\) to the physical Higgs bosons of the effective two-doublet electroweak sector should also be considered. They are \((\phi_2^+, \phi_2^0)\) and \((\phi_L^+, \phi_L^0)\). The physical linear combination is \((v_3\Phi - v_2\Phi_L)/\sqrt{v_2^2 + v_3^2}\). Since \(v_2^2/v_3^2 = x/(1 - 2x)\), the \(Z'\) couplings are completely determined. The resulting factor for both \(Z' \to H_L^+H_L^-\) and \(Z' \to A_R^0H_R^0\) is \((1 - 3x)/2\).

Let \(\Gamma_0 = (g')^2M_{Z'}/48\pi\), then the partial decay widths in units of \(\Gamma_0\) and their respective branching fractions (%) are given in Table II. In the special case where \(m_{A_R} > 2m_{\Delta_R^{++}}\), which is allowed in part of the parameter space shown in Fig. 1 and assuming that \(m_{H_R^0} > 2m_{\Delta_R^{++}}\) as well, we will have the spectacular decay chain \(Z' \to A_R^0H_R^0 \to \Delta_R^{++}\Delta_R^{-}\), resulting in 8 charged muons as shown in Fig. 2. This branching fraction is of order 20 percent, given the fact that both \(A_R^0\) and \(H_R^0\) decay predominantly into \(\Delta_R^{++}\Delta_R^{-}\), and the

| final state | partial width in \(\Gamma_0\) | branching fraction (%) |
|-------------|--------------------------|-----------------------|
| \(\bar{u}u\) | \((9/2) - 21x + 25x^2 = 0.9925\) | 39.4 |
| \(\bar{d}d\) | \(5x^2/2 = 0.13225\) | 5.3 |
| \(\bar{\nu}\nu\) | \(3x^2/2 = 0.07935\) | 3.2 |
| \(\bar{e}e\) | \((1/2) - 3x + 5x^2 = 0.0745\) | 3.0 |
| \(\bar{\mu}\mu\) | \((1/2) - 3x + 5x^2 = 0.0745\) | 3.0 |
| \(\bar{\tau}\tau\) | \((1/2) - 3x + 5x^2 = 0.0745\) | 3.0 |
| \(\bar{n}n\) | \((1 - x)^2/2 = 0.29645\) | 11.8 |
| \(A_R^0H_R^0\) | \((1 - x)^2 = 0.5929\) | 23.6 |
| \(\xi_R^+\xi_R^-\) | \(x^2 = 0.0529\) | 2.1 |
| \(\Delta_R^{++}\Delta_R^{-}\) | \((1 - 3x)^2 = 0.0961\) | 3.8 |
| \(H_L^+H_L^-\) | \((1 - 3x)^2/4 = 0.024025\) | 0.9 |
| \(A_R^0H_R^0\) | \((1 - 3x)^2/4 = 0.024025\) | 0.9 |
| all | 2.51405 | 100.0 |
The dominant decay mode of $\Delta^\pm_R$ is into two charged muons. In other parts of the parameter space, the decay $A^0_R \to \Delta^{++}_R \Delta^{--}_R$ is kinematically forbidden, but the branching fraction for $Z' \to \Delta^{++}_R \Delta^{--}_R$ is still substantial, yielding 4 muons in the final state.

In the above, we have assumed that the $\Delta_R$ scalar triplet couples only to muons. This means that the corresponding scotino $n_\mu$ is part of the $SU(2)_R$ doublet $(n_\mu, \mu)_R$ with $m_{n_\mu} = 200$ GeV. If $\Delta_R$ couples to electrons, then $e^+e^- \to e^+e^-$ scattering through $\Delta_R^{\pm\pm}$ exchange would be much too big to be consistent with known data. We also assume no flavor mixing, i.e. $\Delta_R$ does not couple to $\mu e$ for example, or lepton flavor violating processes such as $\mu \to eee$ and $\mu \to e\gamma$ would be too big. However, $\Delta_R$ still contributes to the muon anomalous magnetic moment which turns out to have the magnitude of the experimental discrepancy but of the wrong sign. To remedy this situation, one possibility is to add $SU(2)_L$ fermion doublets $(N,E)_{L,R}$ with $S = 0$ and a neutral scalar singlet $\chi$ of $S = -1$. The interaction $(\bar{N}\nu_\mu + \bar{E}\mu)\chi$ will contribute positively and compensate for $\Delta_R$. One final complication is that $n_e$ should have a mass greater than $n_\mu$ in order that $n_\mu$ is dark matter. Since it cannot come from $\Delta_R$, $n_e$ must have a Dirac mass partner, i.e. an $n_L$ singlet. Of course, we can avoid all constraints by considering $n_\tau$ instead as the scotino, in which case $\Delta^{++}_R$ will decay into $\tau^+\tau^+$. The resulting experimental signature would then be much more difficult to pick out.
VI. CONCLUSION

We have explored the possible phenomenology of an unconventional $SU(2)_R$ model at the TeV scale called the Dark Left-Right Model (DLRM) [1]. The scalar sector associated with the $SU(2)_R$ gauge group has been studied in detail, including its mass spectrum and its most relevant signature, namely, the decay of the doubly charged scalar into same-sign dileptons: $\Delta_R^{\pm\pm} \to l^\pm l^\pm$. From the requirement of dark-matter relic abundance that the $SU(2)_R$ scalar triplet must be relatively light, we find that the $Z'$ of this model should decay into them with large branching fractions. In particular, $Z' \to \Delta_R^{++} \Delta_R^{--}$ will yield 4 charged muons, with 1.3 times the event rate of $Z' \to \mu^+ \mu^-$ directly. More spectacularly, if kinematically allowed, $A^0_R$ and $H^0_{R1}$ will decay into $\Delta_R^{++} \Delta_R^{--}$ as well, so that $Z' \to A^0_R H^0_{R1}$ will yield 8 charged muons, with 7.9 times the event rate of $Z' \to \mu^+ \mu^-$. Since a modest luminosity of 0.2 fb$^{-1}$ at the LHC will produce 10 dimuon events from this $Z'$ with $M_{Z'} = 1$ TeV, the predicted events with 4 muons and 8 muons will be clear signals of our proposal.

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[1] S. Khalil, H.-S. Lee, and E. Ma, Phys. Rev. D79, 041701(R) (2009).
[2] E. Ma, Phys. Rev. D36, 274 (1987).
[3] K. S. Babu, X.-G. He, and E. Ma, Phys. Rev. D36, 878 (1987).
[4] See for example J. L. Hewett and T. G. Rizzo, Phys. Rept. 183, 193 (1989).
[5] P. Ramond and D. B. Reiss, Phys. Lett. 80B, 87 (1978).
[6] K. S. Babu and V. S. Mathur, Phys. Rev. D38, 3550 (1988).
[7] S. L. Glashow and S. Weinberg, Phys. Rev. D15, 1958 (1977).
[8] T. Aaltonen et al. [CDF Collaboration], Phys. Rev. Lett. 102, 031801, 091805 (2009).
[9] Z. Ahmed et al. [CDMS Collaboration], arXiv:0912.3592 [astro-ph].
[10] V. M. Abazov et al. [D0 Collaboration], Phys. Rev. Lett. 93, 141801 (2004); D. E. Accosta et al. [CDF Collaboration], Phys. Rev. Lett. 93, 221802 (2004).
[11] A. Aranda, J. Hernandez-Sanchez, and P. Q. Hung, JHEP 0811, 092 (2008).