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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 ν_L couples to n_R only through ϕ_1^0 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 Φ_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.

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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 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

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

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

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

This basic structure was already known many years ago [5,6] 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:

$$\Phi = \begin{pmatrix} \phi_1^0 & \phi_2^+ \\ \phi_1^- & \phi_2^0 \end{pmatrix}, \qquad \Phi_L = \begin{pmatrix} \phi_L^+ \\ \phi_L^0 \end{pmatrix}, \qquad \Phi_R = \begin{pmatrix} \phi_R^+ \\ \phi_R^0 \end{pmatrix}, \tag{4}$$

as well as two triplets for making ν and n massive separately:

$$\Delta_L = \begin{pmatrix} \Delta_L^+/\sqrt{2} & \Delta_L^{++} \\ \Delta_L^0 & -\Delta_L^+/\sqrt{2} \end{pmatrix}, \qquad \Delta_R = \begin{pmatrix} \Delta_R^+/\sqrt{2} & \Delta_R^{++} \\ \Delta_R^0 & -\Delta_R^+/\sqrt{2} \end{pmatrix}. \tag{5}$$

Their assignments under S are listed in Table I.

The Yukawa terms allowed by S are then $\bar{\psi}_L \Phi \psi_R$, $\bar{Q}_L \tilde{\Phi} Q_R$, $\bar{Q}_L \Phi_L d_R$, $\bar{Q}_R \Phi_R h_L$, $\psi_L \psi_L \Delta_L$, and $\psi_R \psi_R \Delta_R$,

TABLE I. Scalar content of the proposed model.

Scalar	$SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)$	S
Φ	(1, 2, 2, 0)	1/2
$\tilde{\Phi} = \sigma_2 \Phi^* \sigma_2$	(1, 2, 2, 0)	−1/2
Φ_L	(1, 2, 1, 1/2)	0
Φ_R	(1, 1, 2, 1/2)	−1/2
Δ_L	(1, 3, 1, 1)	−2
Δ_R	(1, 1, 3, 1)	−1

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whereas $\bar{\psi}_L \tilde{\Phi} \psi_R$, $\bar{Q}_L \Phi Q_R$, and $\bar{h}_L d_R$ are forbidden. Hence m_e, m_u come from $v_2 = \langle \phi_2^0 \rangle$, m_d comes from $v_3 = \langle \phi_L^0 \rangle$, m_h comes from $v_4 = \langle \phi_R^0 \rangle$, m_ν 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 ν 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. $(-)^{3B+L+2j}$. The known quarks and leptons have even R, but n, h, $W_R^\pm$, $\phi_R^\pm$, $\Delta_R^\pm$, $\phi_1^\pm$, $\mathrm{Re}(\phi_1^0)$, and $\mathrm{Im}(\phi_1^0)$ 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 *nonsupersymmetric* model.

III. $SU(2)_R$ HIGGS STRUCTURE

There exists an experimental bound [1] on $M_{Z'}$ of 850 GeV from the 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 the 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 Φ_R and Δ_R:

$$\begin{aligned} V_R = {} & m_4^2 \Phi_R^\dagger \Phi_R + m_6^2 \,\mathrm{Tr}(\Delta_R^\dagger \Delta_R) + \tfrac{1}{2}\lambda_1 (\Phi_R^\dagger \Phi_R)^2 \\ & + \tfrac{1}{2}\lambda_2 [\mathrm{Tr}(\Delta_R^\dagger \Delta_R)]^2 + \tfrac{1}{4}\lambda_3 \,\mathrm{Tr}(\Delta_R^\dagger \Delta_R - \Delta_R \Delta_R^\dagger)^2 \\ & + f_1 (\Phi_R^\dagger \Phi_R) \,\mathrm{Tr}(\Delta_R^\dagger \Delta_R) + f_2 \Phi_R^\dagger (\Delta_R^\dagger \Delta_R - \Delta_R \Delta_R^\dagger) \Phi_R \\ & + \mu (\Phi_R^\dagger \Delta_R \tilde{\Phi}_R + \tilde{\Phi}_R^\dagger \Delta_R^\dagger \Phi_R), \end{aligned} \tag{6}$$

where

$$\Phi_R^\dagger \Phi_R = \phi_R^- \phi_R^+ + \bar{\phi}_R^0 \phi_R^0, \tag{7}$$

$$\mathrm{Tr}\,(\Delta_R^\dagger \Delta_R) = \Delta_R^{--} \Delta_R^{++} + \Delta_R^- \Delta_R^+ + \bar{\Delta}_R^0 \Delta_R^0, \tag{8}$$

$$\Delta_R^\dagger \Delta_R - \Delta_R \Delta_R^\dagger = \begin{pmatrix} \bar{\Delta}_R^0 \Delta_R^0 - \Delta_R^{--} \Delta_R^{++} & \sqrt{2}(\Delta_R^- \Delta_R^{++} - \bar{\Delta}_R^0 \Delta_R^+) \\ \sqrt{2}(\Delta_R^{--} \Delta_R^+ - \Delta_R^- \Delta_R^0) & -\bar{\Delta}_R^0 \Delta_R^0 + \Delta_R^{--} \Delta_R^{++} \end{pmatrix}, \tag{9}$$

$$\tilde{\Phi}_R^\dagger \Delta_R^\dagger \Phi_R = \phi_R^0 \phi_R^0 \bar{\Delta}_R^0 + \sqrt{2} \phi_R^0 \phi_R^+ \Delta_R^- - \phi_R^+ \phi_R^+ \Delta_R^{--}. \tag{10}$$

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

$$V_0 = m_4^2 v_4^2 + m_6^2 v_6^2 + \tfrac{1}{2}\lambda_1 v_4^4 + \tfrac{1}{2}\lambda_2 v_6^4 + \tfrac{1}{2}\lambda_3 v_6^4 + f_1 v_4^2 v_6^2 - f_2 v_4^2 v_6^2 + 2\mu v_4^2 v_6, \tag{11}$$

where $v_{4,6}$ are determined by

$$\frac{\partial V_0}{\partial v_4} = 2v_4 [m_4^2 + \lambda_1 v_4^2 + (f_1 - f_2) v_6^2 + 2\mu v_6] = 0, \tag{12}$$

$$\frac{\partial V_0}{\partial v_6} = 2v_6 [m_6^2 + (\lambda_2 + \lambda_3) v_6^2 + (f_1 - f_2) v_4^2] + 2\mu v_4^2 = 0. \tag{13}$$

The physical mass-squared matrices are given by

$$\mathcal{M}^2(\mathrm{Re}\phi_R^0, \mathrm{Re}\Delta_R^0) = \begin{pmatrix} 2\lambda_1 v_4^2 & 2(f_1 - f_2) v_4 v_6 + 2\mu v_4 \\ 2(f_1 - f_2) v_4 v_6 + 2\mu v_4 & 2(\lambda_2 + \lambda_3) v_6^2 - \mu v_4^2 / v_6 \end{pmatrix}, \tag{14}$$

$$\mathcal{M}^2(\mathrm{Im}\phi_R^0, \mathrm{Im}\Delta_R^0) = \begin{pmatrix} -4\mu v_6 & 2\mu v_4 \\ 2\mu v_4 & -\mu v_4^2 / v_6 \end{pmatrix}, \tag{15}$$

$$\mathcal{M}^2(\phi_R^\pm, \Delta_R^\pm) = \begin{pmatrix} 2v_6 (f_2 v_6 - \mu) & -\sqrt{2} v_4 (f_2 v_6 - \mu) \\ -\sqrt{2} v_4 (f_2 v_6 - \mu) & v_4^2 / v_6 (f_2 v_6 - \mu) \end{pmatrix}, \tag{16}$$

$$\mathcal{M}^2(\Delta_R^{\pm\pm}) = 2 f_2 v_4^2 - 2\lambda_3 v_6^2 - \mu v_4^2 / v_6. \tag{17}$$

As expected, the linear combinations

$$\frac{(v_4\,\mathrm{Im}\phi_R^0 + 2v_6\,\mathrm{Im}\Delta_R^0)}{\sqrt{v_4^2+4v_6^2}}, \qquad \frac{(v_4\phi_R^\pm + \sqrt{2}v_6\Delta_R^\pm)}{\sqrt{v_4^2+2v_6^2}}$$

have zero mass, corresponding to the longitudinal components of Z' and $W_{\mathrm{R}}^\pm$. Their orthogonal combinations

$$A_{\mathrm{R}} = \frac{\sqrt{2}(v_4\,\mathrm{Im}\Delta_{\mathrm{R}}^0 - 2v_6\,\mathrm{Im}\phi_{\mathrm{R}}^0)}{\sqrt{v_4^2+4v_6^2}},$$

$$\xi_{\mathrm{R}}^\pm = \frac{(v_4\Delta_{\mathrm{R}}^\pm - \sqrt{2}v_6\phi_{\mathrm{R}}^\pm)}{\sqrt{v_4^2+2v_6^2}}$$

have mass-squares $-\mu(v_4^2+4v_6^2)/v_6$ and $(f_2-\mu/v_6)\times(v_4^2+2v_6^2)$, respectively. Since n_{R} couples to $\Delta_{\mathrm{R}}^\pm$, but not to $\phi_{\mathrm{R}}^\pm$, the discussion on dark-matter relic abundance from nn annihilation to lepton pairs through the $\Delta_{\mathrm{R}}^\pm$ exchange in Ref. [1] applies only if $v_6^2 \ll v_4^2$. 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

$$\frac{f_n}{\sqrt{2}}(\Delta_{\mathrm{R}}^0 n_{\mathrm{R}} n_{\mathrm{R}} + \sqrt{2}\delta_{\mathrm{R}}^+ n_{\mathrm{R}} e_{\mathrm{R}} + \Delta_{\mathrm{R}}^{++} e_r e_{\mathrm{R}}), \tag{18}$$

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 $M_{W_{\mathrm{R}}} = 832$ GeV, where $v_2 = 95$ GeV and $v_3 = 146$ GeV have been used to ensure zero $Z - Z'$ mixing at tree level (see the next section).

The physical charged scalar ξ_{R}^+ is now 99.4% Δ_{R}^+ and its mass is given by

$$m_{\xi_{\mathrm{R}}^+}^2 = (f_2 - \mu/v_6)(v_4^2+2v_6^2) = [220\ \mathrm{GeV}]^2. \tag{19}$$

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

$$\frac{m_{A_{\mathrm{R}}}^2}{1+4v_6^2/v_4^2} + m_{\Delta_{\mathrm{R}}^{++}}^2 = \frac{2m_{\xi_{\mathrm{R}}^+}^2}{1+2v_6^2/v_4^2} - 2\lambda_3 v_6^2. \tag{20}$$

This means that both $m_{A_{\mathrm{R}}}$ and $m_{\Delta_{\mathrm{R}}^{++}}$ are bounded from above as a function of λ_3 which should not be larger than about one in magnitude. We plot in Fig. 1 $m_{\Delta_{\mathrm{R}}^{++}}$ versus $m_{A_{\mathrm{R}}}$ for various values of λ_3.

We now come to a very important conclusion. To satisfy the dark-matter relic density in this model, m_n and $m_{\xi_{\mathrm{R}}^+}$ have to be of order 200 GeV. This in turn implies that $m_{A_{\mathrm{R}}}$ and $m_{\Delta_{\mathrm{R}}^{++}}$ are bounded in such a way that the decays $A_{\mathrm{R}} \to nn$ and $\Delta_{\mathrm{R}}^{++} \to \xi_{\mathrm{R}}^+\xi_{\mathrm{R}}^+$ are kinematically forbidden. This means that the dominant decay of Δ_{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 $A_{\mathrm{R}} \to \Delta_{\mathrm{R}}^{++}\Delta_{\mathrm{R}}^{--}$. Note that the present experimental bound on $m_{\Delta^{++}}$ is 110 GeV [10].

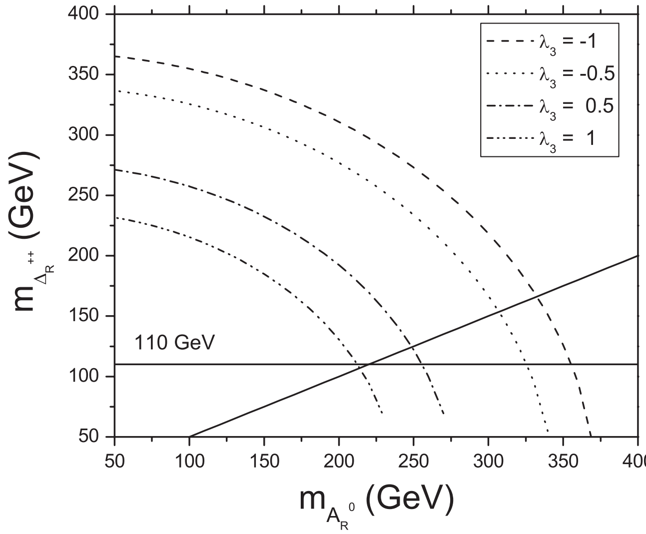

FIG. 1. Plot of $M_{A_{\mathrm{R}}^0}$ versus $M_{\Delta_{\mathrm{R}}^{++}}$ for different values of λ_3 with $v_2 = 95$ GeV, $v_3 = 146$ GeV, $v_4 = 1851$ GeV, $v_6 = 141$ GeV, $M_{\xi_{\mathrm{R}}^+} = 220$ GeV, and $M_{Z'} = 1$ TeV, $M_{W_{\mathrm{R}}^+} = 832$ GeV. The lower bound of 110 GeV comes from the Tevatron data. The line $m_{A_{\mathrm{R}}} = 2m_{\Delta_{\mathrm{R}}^{++}}$ is also shown.

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

$$m_{H_{\mathrm{R1}}^0}^2 = \frac{m_{A_{\mathrm{R}}}^2}{1+4v_6^2/v_4^2} + \frac{2v_6^2}{\lambda_1}[\lambda_1(\lambda_2+\lambda_3) - (f_1 - 0.014)^2]. \tag{21}$$

Finally, we need to consider the scalar bidoublet and the scalar left doublet. In this model, (ϕ_1^0, ϕ_1^-) will be heavy, and the two doublets (ϕ_2^+, ϕ_2^0), $(\phi_{\mathrm{L}}^+, \phi_{\mathrm{L}}^0)$ are similar to the usual two Higgs doublets considered in the SM. The linear combination $(v_3\Phi_2 - v_2\Phi_{\mathrm{L}})/\sqrt{v_2^2+v_3^2} = [H_{\mathrm{L}}^+, (H_{\mathrm{L}}^0 + iA_{\mathrm{L}}^0)/\sqrt{2}]$ is physical and light.

IV. GAUGE SECTOR

Since e has $L = 1$ and n has $L = 0$, the W_{R}^+ of this model must have $L = S - T_{3\mathrm{R}} = 0 - 1 = -1$. This also means that unlike the conventional LRM, $W_{\mathrm{R}}^\pm$ does not mix with the $W_{\mathrm{L}}^\pm$ of the SM at all. This important property allows the $SU(2)_{\mathrm{R}}$ breaking scale to be much lower than it would be otherwise, as explained already 23 years ago [2,3]. Assuming that $g_{\mathrm{L}} = g_{\mathrm{R}}$ and let $x \equiv \sin^2\theta_W$, then the neutral gauge bosons of the DLRM (as well as the ALRM) are given by

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

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

$$J_{Z'} = xJ_{3L} + (1-x)J_{3R} - xJ_{em} \tag{23}$$

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

$$M_{W_L}^2 = \frac{e^2}{2x}(v_2^2 + v_3^2), \qquad M_Z^2 = \frac{M_{W_L}^2}{1-x}, \qquad M_{W_R}^2 = \frac{e^2}{2x}(v_2^2 + v_4^2 + 2v_6^2), \tag{24}$$

$$M_{Z'}^2 = \frac{e^2(1-x)}{2x(1-2x)}(v_2^2 + v_4^2 + 4v_6^2) - \frac{x^2 M_{W_L}^2}{(1-x)(1-2x)}, \tag{25}$$

where zero $Z - Z'$ mixing has been assumed, using the tree-level condition [3] $v_2^2/(v_2^2 + v_3^2) = x/(1-x)$, which is subject to radiative corrections. Although $Z - Z'$ mixing is common to all models with extended gauge groups with nontrivial Higgs sectors, the present precision electroweak measurements are such that it must somehow be made very small [11]. Note that in the ALRM, Δ_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. In Ref. [1], it has already been shown that with an integrated luminosity of $L = 10\ \text{fb}^{-1}$, up to $M_{Z'} \sim 2.4$ TeV may be discovered. Once $M_{Z'}$ is determined, then the DLRM predicts the existence of $W_R^\pm$ with a mass in the range

$$\frac{(1-2x)}{2(1-x)}M_{Z'}^2 + \frac{x}{2(1-x)^2}M_{W_L}^2 < M_{W_R}^2 < \frac{(1-2x)}{(1-x)}M_{Z'}^2 + \frac{x^2}{(1-x)^2}M_{W_L}^2. \tag{26}$$

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 $0.2\ \text{fb}^{-1}$ is required [1]. Its discovery channel is presumably $\mu^+\mu^-$, but it will also have 4 charged muons in the final state from $\Delta_R^{++}\Delta_R^{--}$, and perhaps even 8 charged muons, as shown below.

In addition to all SM particles, Z' also decays into $n\bar{n}$, $\Delta_R^{++}\Delta_R^{--}$, $\xi_R^+\xi_R^-$, $A_R^0 H_{R1}^0$, $H_L^+ H_L^-$, and $A_L^0 H_L^0$. In particular, the subsequent decay $\Delta_R^{\pm\pm} \to \mu^\pm\mu^\pm$ will be a unique signature where the like-sign dimuons have identical invariant masses.[1]

The interactions of Z' with fermions come from

$$\mathcal{L} = -g' Z'_\mu J^\mu_{Z'}, \tag{27}$$

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

$$\Gamma(Z' \to \bar{f}f) = \frac{(g')^2 M_{Z'}}{24\pi}[c_L^2 + c_R^2], \tag{28}$$

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

$$u_L = -\frac{x}{6}, \qquad u_R = \frac{1}{2} - \frac{7x}{6}, \qquad d_L = -\frac{x}{6}, \qquad d_R = \frac{x}{3}, \tag{29}$$

$$\nu_L = \frac{x}{2}, \qquad n_R = \frac{1-x}{2}, \qquad e_L = \frac{x}{2}, \qquad e_R = -\frac{1}{2} + \frac{3x}{2}. \tag{30}$$

Here we need to consider 3 families for u, d, ν, e but only one for n.

The decay of $Z' \to A_R^0 H_{R1}^0$ to scalars comes from

$$\mathcal{L} = -g'(1-x)Z'_\mu[(\partial^\mu H_{R1}^0)A_R^0 - (\partial^\mu A_R^0)H_{R1}^0], \tag{31}$$

with the partial decay width

$$\Gamma(Z' \to A_R^0 H_{R1}^0) = \frac{(g')^2 M_{Z'}(1-x)^2}{48\pi}, \tag{32}$$

where $(1-x)$ comes from $I_{3L} = 0$, $I_{3R} = -1$, $Q = 0$. For $Z' \to \xi_R^+\xi_R^-$, 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 (ϕ_2^+, ϕ_2^0) and (ϕ_L^+, ϕ_L^0). The physical

[1]Not all models involving doubly charged scalars have this decay, see, for example, [12].

TABLE II. Z' decay widths and branching fractions.

Final state	Partial width in Γ_0	Branching fraction (%)
$\bar{u}u$	$(9/2) - 21x + 25x^2 = 0.9925$	39.4
$\bar{d}d$	$5x^2/2 = 0.132\,25$	5.3
$\bar{\nu}\nu$	$3x^2/2 = 0.079\,35$	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
$\bar{n}n$	$(1-x)^2/2 = 0.296\,45$	11.8
$A_R^0 H_{R1}^0$	$(1-x)^2 = 0.5929$	23.6
$\xi_R^+ \xi_R^-$	$x^2 = 0.0529$	2.1
$\Delta_R^{++}\Delta_R^{--}$	$(1-3x)^2 = 0.0961$	3.8
$H_L^+ H_L^-$	$(1-3x)^2/4 = 0.024\,025$	0.9
$A_L^0 H_L^0$	$(1-3x)^2/4 = 0.024\,025$	0.9
All	2.514 05	100.0

linear combination is $(v_3\Phi_2 - v_2\Phi_L)/\sqrt{v_2^2 + v_3^2}$. 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 $Z' \to A_L^0 H_L^0$ is $(1-3x)/2$.

Let $\Gamma_0 = (g')^2 M_{Z'}/48\pi$, then the partial decay widths in units of Γ_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 $m_{H_{R1}^0} > 2m_{\Delta_R^{++}}$ as well, we will have the spectacular decay chain $Z' \to A_R^0 H_{R1}^0 \to \Delta_R^{++}\Delta_R^{--} + \Delta_R^{++}\Delta_R^{--}$, resulting in 8 charged muons as shown in Fig. 2. This branching fraction is of order 20%, given the fact that both A_R^0 and H_{R1}^0 decay predominantly into $\Delta_R^{++}\Delta_R^{--}$, and the dominant decay mode of $\Delta_R^{\pm\pm}$ is into two charged muons. In other parts of the parameter space, the decay $A_R^0 \to \Delta_R^{++}\Delta_R^{--}$ 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 this case, for $M_{Z'} = 1$ TeV and an integrated luminosity of 10 fb^{-1} at $\sqrt{s} = 14$ TeV, LHC will produce about 330 Z' events. Thus, with a branching fraction for $\mathrm{BR}(Z' \to \Delta_R^{++}\Delta_R^{--}) \sim 0.04$, and $\Delta_R^{++}(\Delta_R^{--})$ decaying into $\mu^+\mu^+(\mu^-\mu^-)$ with branching fraction ~ 1, there will be $13\mu^+\mu^+\mu^-\mu^-$ events. The requirement that each same-sign dimuon pair should form identical invariant masses ought to make this observation essentially background free.

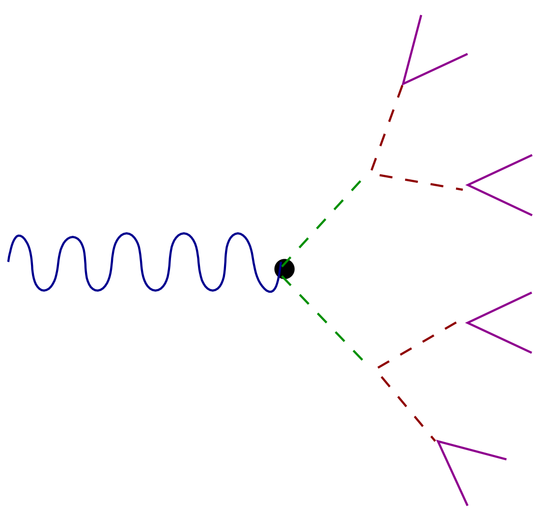

FIG. 2 (color online). Diagram for the process $Z' \to A_R^0 H_R^0 \to \Delta_R^{++}\Delta_R^{--} + \Delta_R^{++}\Delta_R^{--} \to 2(\mu^+\mu^+) + 2(\mu^-\mu^-)$.

In the above, we have assumed that the Δ_R scalar triplet couples only to muons. This means that the corresponding scotino n_μ is part of the $SU(2)_R$ doublet $(n_\mu, \mu)_R$ with $m_{n_\mu} = 200$ GeV. If Δ_R couples to electrons, then $e^+e^- \to e^+e^-$ scattering through the $\Delta_R^{\pm\pm}$ exchange would be much too big to be consistent with the known data. We also assume no flavor mixing, i.e. Δ_R does not couple to μe, for example, or lepton flavor violating processes such as $\mu \to eee$ and $\mu \to e\gamma$ would be too big. However, Δ_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 χ of $S = -1$. The interaction $(\bar{N}\nu_\mu + \bar{E}\mu)\chi$ will contribute positively and compensate for Δ_R. One final complication is that n_e should have a mass greater than n_μ in order that n_μ is dark matter. Since it cannot come from Δ_R, n_e must have a Dirac mass partner, i.e. a n_L singlet. Of course, we can avoid all constraints by considering n_τ instead as the scotino, in which case Δ_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 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_R^0 and H_{R1}^0 will decay into $\Delta_R^{++}\Delta_R^{--}$ as well, so that $Z' \to A_R^0 H_{R1}^0$ 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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