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Bound on Z' mass from CDMS II in the dark left-right gauge model II

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With the recent possible signal of dark matter from the CDMS II experiment, the Z' mass of a new version of the dark left-right gauge model (DLRM II) is predicted to be at around a TeV. As such, it has an excellent discovery prognosis at the operating Large Hadron Collider.

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I. INTRODUCTION

One year ago, we proposed [1] that dark-matter fermions (scotinos) were naturally present in an unconventional left-right gauge extension of the standard model (SM) of particle interactions, which we call the dark left-right model (DLRM). It is a nonsupersymmeric variation of the alternative left-right model (ALRM) discussed already 23 years ago [2,3]. One important difference of both the DLRM and the ALRM with the conventional left-right model (LRM) [4] is the fact that tree-level flavor-changing neutral currents [5] are naturally absent so that the $SU(2)_R$ breaking scale may easily be at around a TeV, allowing both the charged $W_R^\pm$ and the extra neutral Z' gauge bosons to be observable at the LHC. Interesting phenomenology of Z' decay into scalar bosons in the DLRM has just recently been discussed [6].

In this paper, we propose a new variant of this extension which we call DLRM II. (Other more exotic variants are also possible [7].) Instead of having Majorana scotinos as dark matter, we now have Dirac scotinos. Their interactions with nuclei through the Z' are thus relevant for understanding the recent result of the dark-matter direct-search experiment CDMS II [8]. It will be shown that the Z' mass may indeed be around a TeV, and its discovery prognosis at the LHC is excellent.

II. MODEL

Consider the gauge group $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)$. The conventional leptonic assignments are $\psi_L = (\nu, e)_L \sim (1, 2, 1, -1/2)$ and $\psi_R = (\nu, e)_R \sim (1, 1, 2, -1/2)$. Hence ν and e obtain Dirac masses through the Yukawa terms $\bar{\psi}_L \Phi \psi_R$ and $\bar{\psi}_L \tilde{\Phi} \psi_R$, where $\Phi = (\phi_1^0, \phi_1^-; \phi_2^+, \phi_2^0) \sim (1, 2, 2, 0)$ is a Higgs bidoublet and $\tilde{\Phi} = \sigma_2 \Phi^* \sigma_2 = (\bar{\phi}_2^0, -\phi_2^-; -\phi_1^+, \bar{\phi}_1^0)$ transforms in the same way. Both $\langle \phi_1^0 \rangle$ and $\langle \phi_2^0 \rangle$ contribute to m_ν and m_e, and similarly m_u and m_d in the quark sector, resulting thus in the appearance of tree-level flavor-changing neutral currents.

Suppose the term $\bar{\psi}_L \tilde{\Phi} \psi_R$ is forbidden by a symmetry, then the same symmetry may be used to maintain $\langle \phi_1^0 \rangle = 0$ and only e gets a mass through $\langle \phi_2^0 \rangle \neq 0$. At the same time, ν_L and ν_R are not Dirac mass partners, so they could in fact be completely different particles with independent masses of their own. Whereas ν_L is clearly the neutrino we observe in the usual weak interactions, ν_R can now be something else entirely. Here we rename ν_R as n_R and show that it may in fact be a scotino, i.e. a fermionic dark-matter candidate.

In our previous proposal [1], we imposed a new global $U(1)$ symmetry S in such a way that the spontaneous breaking of $SU(2)_R \times S$ will leave the combination $L = S - T_{3R}$ unbroken. We then showed that L is a generalized lepton number, with $L = 1$ for the known leptons, and $L = 0$ for all known particles which are not leptons. Here we consider instead the case $L = S + T_{3R}$. Our model is nonsupersymmetric, but it may be rendered supersymmetric by the usual procedure which takes the SM to the minimal supersymmetric standard model. Under $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1) \times S$, the fermions transform as shown in Table I. Note the necessary appearance of the exotic quark h, which will turn out to carry a lepton number as well.

The scalar sector consists of one bidoublet and two doublets:

$$\Phi = \begin{pmatrix} \phi_1^0 & \phi_2^+ \\ \phi_1^- & \phi_2^0 \end{pmatrix}, \qquad \Phi_L = \begin{pmatrix} \phi_L^+ \\ \phi_L^0 \end{pmatrix}, \qquad \Phi_R = \begin{pmatrix} \phi_R^+ \\ \phi_R^0 \end{pmatrix}. \tag{1}$$

Their assignments under S are listed in Table II.

TABLE I. Fermion content of the proposed model.

Fermion	$SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)$	S
$\psi_L = (\nu, e)_L$	(1, 2, 1, −1/2)	1
$\psi_R = (n, e)_R$	(1, 1, 2, −1/2)	3/2
ν_R	(1, 1, 1, 0)	1
n_L	(1, 1, 1, 0)	2
$Q_L = (u, d)_L$	(3, 2, 1, 1/6)	0
$Q_R = (u, h)_R$	(3, 1, 2, 1/6)	−1/2
d_R	(3, 1, 1, −1/3)	0
h_L	(3, 1, 1, −1/3)	−1

TABLE II. 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)	$1/2$

The Yukawa terms allowed by S are then $\bar{\psi}_L \Phi \psi_R$, $\bar{\psi}_L \tilde{\Phi}_L \nu_R$, $\bar{\psi}_R \tilde{\Phi}_R n_L$, $\bar{Q}_L \tilde{\Phi} Q_R$, $\bar{Q}_L \Phi_L d_R$, and $\bar{Q}_R \Phi_R h_L$, whereas $\bar{\psi}_L \tilde{\Phi} \psi_R$, $\bar{n}_L \nu_R$, $\bar{Q}_L \Phi Q_R$, and $\bar{h}_L d_R$ are forbidden. Hence m_e, m_u come from $v_2 = \langle \phi_2^0 \rangle$; m_ν, m_d come from $v_3 = \langle \phi_L^0 \rangle$; and m_n, m_h come from $v_4 = \langle \phi_R^0 \rangle$. This structure shows clearly that flavor-changing neutral currents are guaranteed to be absent at tree level.

As it stands, both the neutrino ν and the scotino n are Dirac fermions, and lepton number L is conserved. If we now introduce a soft term $\nu_R \nu_R$ which breaks L by two units, then ν_L gets a Majorana mass through the canonical seesaw mechanism, as is usually assumed. As for n, it remains a Dirac fermion, being protected by a residual global $U(1)$ symmetry, under which n, W_R^+ transform as 1, and h, $\phi_1^{0,-}$ transform as -1.

III. GAUGE SECTOR

Since e has $L = 1$ and n has $L = 2$, the W_R^+ 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 23 years ago [2,3]. Let $e/g_L = s_L = \sin\theta_W$ and $s_R = e/g_R$, with $c_{L,R} = \sqrt{1 - s_{L,R}^2}$, then $g_B = e/\sqrt{c_L^2 - s_R^2}$ and the neutral gauge bosons of the DLRM (as well as the ALRM) are given by

$$\begin{pmatrix} A \\ Z \\ Z' \end{pmatrix} = \begin{pmatrix} s_L & s_R & \sqrt{c_L^2 - s_R^2} \\ c_L & -s_L s_R/c_L & -s_L\sqrt{c_L^2 - s_R^2}/c_L \\ 0 & \sqrt{c_L^2 - s_R^2}/c_L & -s_R/c_L \end{pmatrix} \times \begin{pmatrix} W_L^0 \\ W_R^0 \\ B \end{pmatrix}. \quad (2)$$

Whereas Z couples to the current $J_{3L} - s_L^2 J_{\rm em}$ with coupling $e/s_L c_L$ as in the SM, Z' couples to the current

$$J_{Z'} = s_R^2 J_{3L} + c_L^2 J_{3R} - s_R^2 J_{\rm em} \quad (3)$$

with coupling $g_{Z'} = e/s_R c_L \sqrt{c_L^2 - s_R^2}$.

The masses of the gauge bosons are given by

$$M_{W_L}^2 = \frac{e^2}{2s_L^2}(v_2^2 + v_3^2), \qquad M_Z^2 = \frac{M_{W_L}^2}{c_L^2}, \qquad M_{W_R}^2 = \frac{e^2}{2s_R^2}(v_4^2 + v_2^2), \quad (4)$$

$$M_{Z'}^2 = \frac{e^2 c_L^2}{2s_R^2(c_L^2 - s_R^2)}(v_4^2 + v_2^2) - \frac{s_L^2 s_R^2 M_{W_L}^2}{c_L^2(c_L^2 - s_R^2)}, \quad (5)$$

where zero $Z - Z'$ mixing has been assumed, using the condition [3] $v_2^2/(v_2^2 + v_3^2) = s_R^2/c_L^2$.

IV. DIRECT-SEARCH CONSTRAINT FROM CDMS II

The Z' couplings to u, d, n (in units of $g_{Z'}$) are given by

$$u_L = -\tfrac{1}{6}s_R^2, \qquad u_R = \tfrac{1}{2}c_L^2 - \tfrac{2}{3}s_R^2, \qquad u_V = \tfrac{1}{4}c_L^2 - \tfrac{5}{12}s_R^2, \quad (6)$$

$$d_L = -\tfrac{1}{6}s_R^2, \qquad d_R = \tfrac{1}{3}s_R^2, \qquad d_V = \tfrac{1}{12}s_R^2, \quad (7)$$

$$n_L = 0, \qquad n_R = \tfrac{1}{2}c_L^2, \qquad n_V = \tfrac{1}{4}c_L^2. \quad (8)$$

The effective Lagrangian for elastic scattering of the scotino n off quarks is then given by

$$\mathcal{L} = \frac{g_{Z'}^2 n_V}{M_{Z'}^2}(\bar{n}\gamma_\mu n)(u_V \bar{u}\gamma^\mu u + d_V \bar{d}\gamma^\mu d). \quad (9)$$

In the original DLRM [1], n is a Majorana scotino, so it does not contribute to the s-wave elastic spin-independent scattering cross section in the nonrelativistic limit. Here n is a Dirac scotino, so it will contribute. Let

$$f_P = g_{Z'}^2 n_V(2u_V + d_V)/M_{Z'}^2, \qquad f_N = g_{Z'}^2 n_V(u_V + 2d_V)/M_{Z'}^2, \quad (10)$$

then its elastic cross section per nucleon is given by [9]

$$\sigma_0 = \frac{4m_r^2}{\pi}\frac{[Zf_P + (A - Z)f_N]^2}{A^2}, \quad (11)$$

where Z and A are the atomic and mass numbers of the target nucleus, and $m_r = m_n m_P/(m_n + m_P) \simeq m_P$. The CDMS II Collaboration [8] observed two possible signal events with an expected background of 0.6 ± 0.1. Using ^{73}Ge, i.e. $Z = 32$ and $A - Z = 41$, as a representative estimate of σ_0, this result could also be considered as an upper bound, i.e.

$$\sigma_0 = \frac{\pi\alpha^2 m_P^2(105c_L^2 - 137s_R^2)^2}{(146)^2 s_R^4(c_L^2 - s_R^2)^2 M_{Z'}^4} < 3.8 \times 10^{-8}\ \text{pb}, \quad (12)$$

which occurs at $m_n = 70$ GeV.

V. PHENOMENOLOGICAL ANALYSIS

We consider the range $e^2 < s_R^2 < c_L^2 - e^2$, where the lower bound corresponds to $g_R = 1$ and the upper bound to $g_B = 1$. The values of $g_{Z'}$ and $\Gamma_{Z'}/M_{Z'}$ are plotted in Figs. 1(a) and 1(b) where Z' is assumed to decay only into SM fermions.

We compute the production and decay of Z' to e^+e^- at the Tevatron as a function of $M_{Z'}$ for various values of s_R^2 and compare it to data [10] at $E_{\rm cm} = 1.96$ TeV and an integrated luminosity of 2.5 fb^{-1} in Fig. 2(a). We then plot the exclusion limits on $M_{Z'}$ from both the new CDMS II data and the Tevatron as a function of s_R^2 in Fig. 2(b). Note that the CDMS II bound is stronger than the Tevatron bound for $s_R^2 < 0.5$. Note also that due to the accidental cancellation in the numerator of σ_0 in Eq. (12), the observed events at CDMS II cannot be interpreted as signals of dark matter in this model if $s_R^2 > 0.5$, because they would be excluded by the Tevatron data.

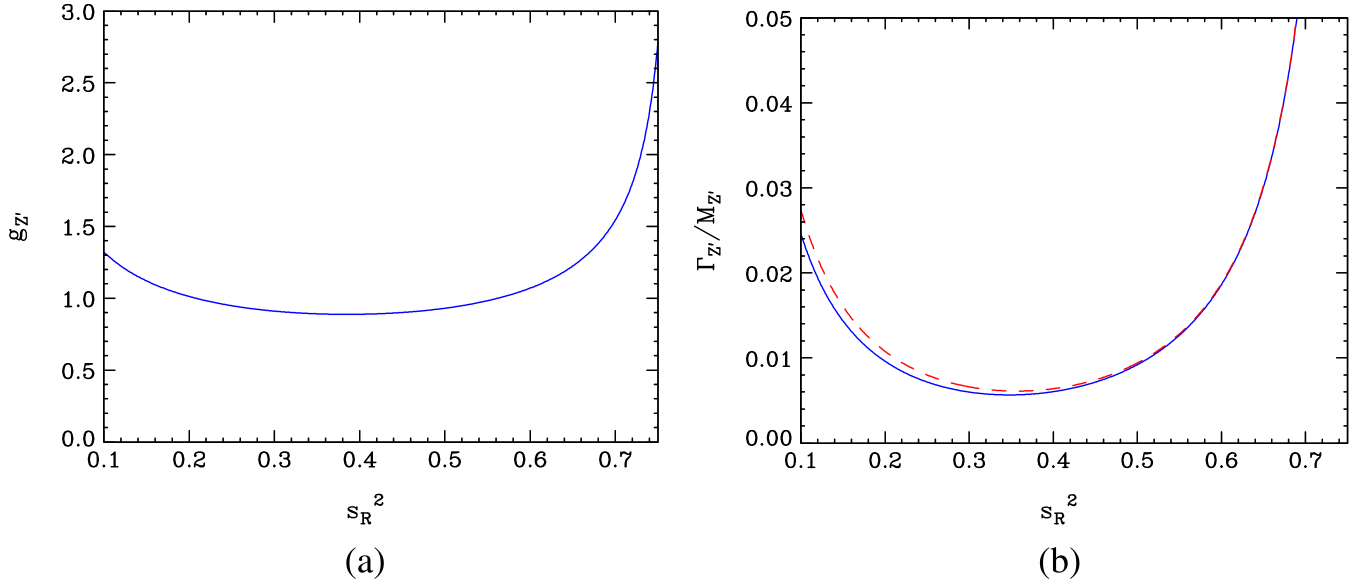

FIG. 1 (color online). (a) $g_{Z'}$ vs s_R^2. (b) $\Gamma_{Z'}/M_{Z'}$ vs s_R^2 for SM fermions decay products only in the cases $M_{Z'} = 500$ GeV (solid blue line) and $M_{Z'} \to \infty$ (dashed red line).

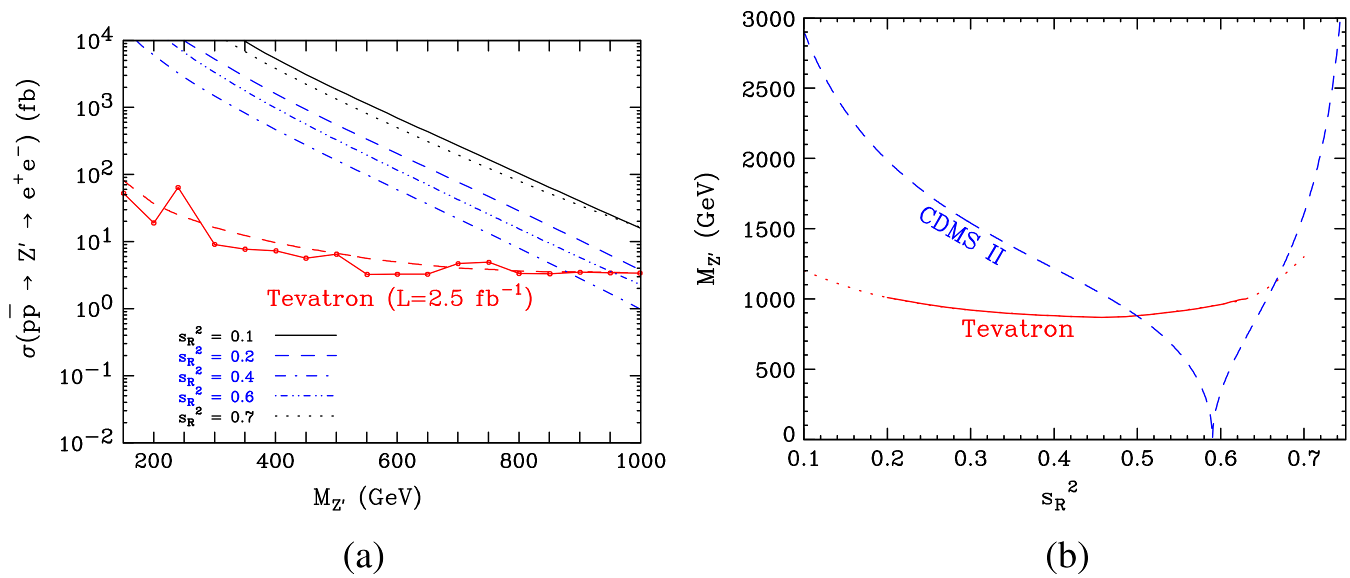

FIG. 2 (color online). (a) $\sigma(p\bar{p} \to Z' \to e^+e^-)$ vs $M_{Z'}$ in this model compared against the Tevatron dielectron search. (b) Lower bounds on $M_{Z'}$ vs s_R^2 from the Tevatron search (solid red line) and from the CDMS II search at $m_n = 70$ GeV (dashed blue line). The dotted segments assume a simple extrapolation of the Tevatron data.

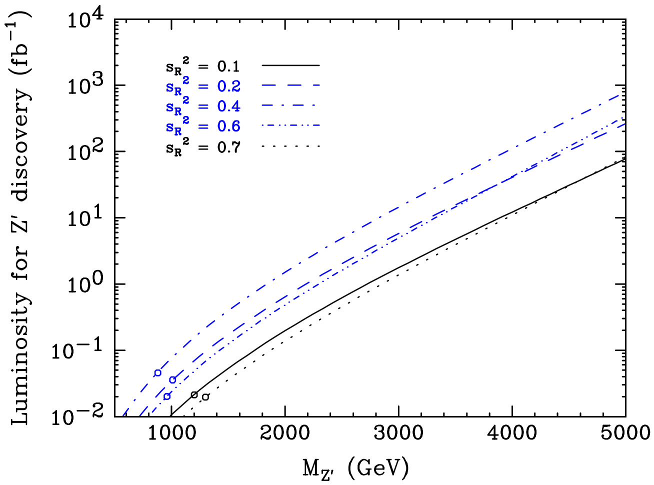

FIG. 3 (color online). Luminosity for Z' discovery by 10 dielectron events at LHC. Small circles are Tevatron limits.

Given that $M_{Z'}$ is allowed to be in the TeV range, its discovery prognosis is excellent at the LHC. We show in Fig. 3 its discovery reach (assuming $E_{\rm cm} = 14$ TeV) by 10 dilepton events (either dielectron or dimuon) which satisfy the following basic cuts on their transverse momenta, rapidities, and invariant mass: $p_T > 20$ GeV (each lepton), $|\eta| < 2.4$ (each lepton), and $|M_{\ell\bar{\ell}} - M_{Z'}| < 3\Gamma_{Z'}$. Using these cuts, the dominant SM background from γ/Z (Drell-Yan) is negligible. With an integrated luminosity of 1 fb^{-1}, the Z' of DLRM II with $M_{Z'} \sim 2$ TeV may then be discovered at the LHC.

VI. DARK-MATTER RELIC ABUNDANCE

In this model, the dark-matter relic abundance is presumably determined by the annihilation $n\bar{n} \to Z' \to$ SM fermions. The thermally averaged cross section multiplied by relative velocity is approximately given by

$$\langle \sigma v_{\rm rel} \rangle_{Z'} = \frac{g_{Z'}^4 c_L^4 m_n^2 \sum_f (f_L^2 + f_R^2)}{32\pi(4m_n^2 - M_{Z'}^2)^2} = \frac{\pi\alpha^2(3 - 9r + 10r^2)m_n^2}{2c_L^4 r^2(1-r)^2(4m_n^2 - M_{Z'}^2)^2}, \quad (13)$$

where $r = s_R^2/c_L^2$. Fixing the above at 1 pb as a typical value to satisfy the requirement of dark-matter relic abundance, it can easily be shown that for $m_n = 70$ GeV, the required $M_{Z'}$ is very much below the CDMS II bound. (For example, for $s_R^2 = 0.4$, $M_{Z'} = 267$ GeV would be required.) In other words, the $n\bar{n} \to Z'$ annihilation cross section would be too small to account for the observed dark-matter relic abundance. To remedy this situation, the mechanism proposed in the original DLRM may be invoked, i.e. $n\bar{n} \to l^- l^+$ through Δ_R^+ exchange. However, this requires adding the $SU(2)_R$ scalar triplet $(\Delta_R^{++}, \Delta_R^+, \Delta_R^0)$, which is not necessary in our present version and thus not so motivated. The alternative is to consider a larger value of m_n.

The CDMS II bound on σ_0 is very well approximated in the range $0.3 < m_n < 1.0$ TeV by the expression

$$\sigma_0 < 2.2 \times 10^{-7} \text{ pb } (m_n/1 \text{ TeV})^{0.86}. \quad (14)$$

Using this on the right-hand side of Eq. (12), we plot in Figs. 4(a) and 4(b) the $M_{Z'}$ bounds for $m_n = 400$ and 600 GeV, as well as the solutions of $M_{Z'}$ (with $M_{Z'} > 2m_n$) to Eq. (13) for 1 pb. We see that there are indeed consistent solutions (where the solid line is higher than the dashed line) for a range of s_R^2 in each case. If m_n falls below 300 GeV, then there is no solution because $M_{Z'}$ would then

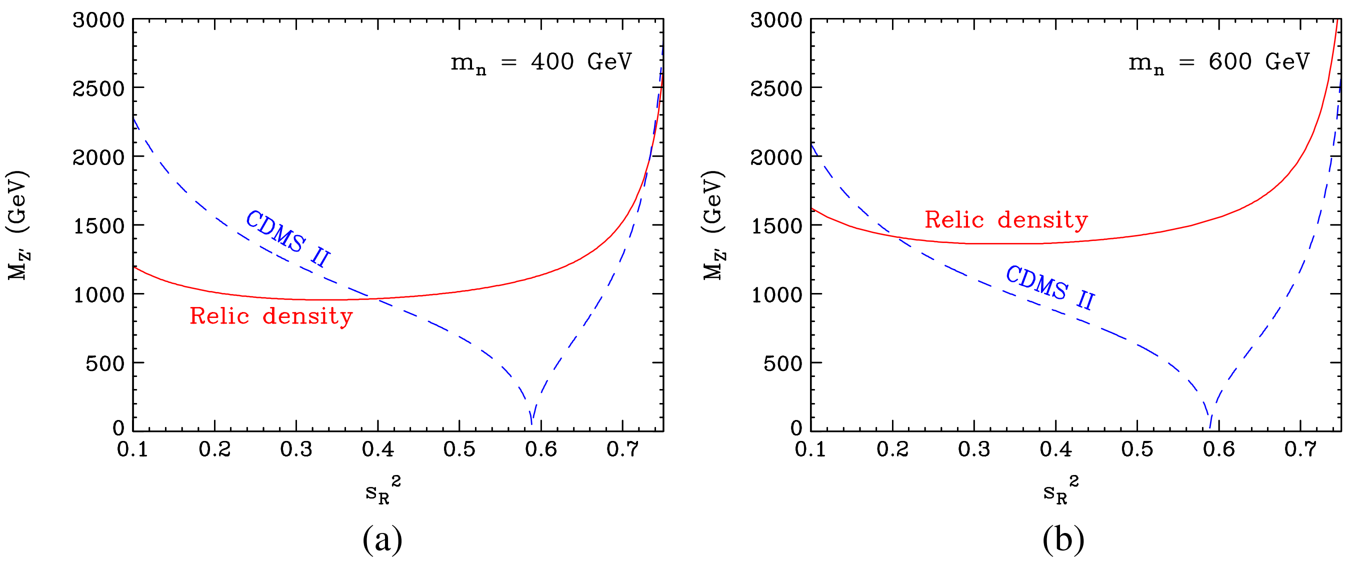

FIG. 4 (color online). (a) For $m_n = 400$ GeV, the CDMS II bound on $M_{Z'}$ (dashed blue line) and its value (solid red line) from $\langle \sigma v_{\rm rel} \rangle_{Z'} = 1$ pb vs s_R^2. (b) Same as in (a) for $m_n = 600$ GeV.

be excluded by the Tevatron bound. If the CDMS II events are considered as true dark-matter events, then the intersection of the two lines is the prediction of $M_{Z'}$ in this model. We note also that only a modest resonance enhancement is needed from the denominator of Eq. (13). The $n\bar{n}$ annihilation to l^+l^- through W_R exchange also contributes to dark-matter relic abundance, but it is an order of magnitude less:

$$\langle \sigma v_{\rm rel} \rangle_{W_R} = \frac{3 g_R^4 m_n^2}{64\pi (m_n^2 + M_{W_R}^2)^2}. \qquad (15)$$

VII. LEPTON FLAVOR VIOLATION

Unlike the original DLRM, where a scalar triplet $(\Delta_R^{++}, \Delta_R^{+}, \Delta_R^{0})$ may mediate lepton flavor violating processes such as $\mu \to eee$ at tree level, and must be forbidden by hand, the DLRM II is safe because it has no such interactions. Nevertheless, lepton (as well as quark) flavor violation occurs in one loop in the $SU(2)_R$ sector, in complete analogy to that of the SM in the $SU(2)_L$ sector. The branching fraction of $\mu \to e\gamma$ is then

$$B(\mu \to e\gamma) = \frac{3\alpha |\delta_R|^2}{64\pi} \left(\frac{s_L^2 M_{W_L}^2}{s_R^2 M_{W_R}^2} \right)^2 < 1.2 \times 10^{-11}, \quad (16)$$

where the experimental upper bound has also been displayed, and δ_R is the analog of the well-known suppression factor $\delta_L = \sum_i U_{ei}^* U_{\mu i} (m_{\nu_i}^2 / M_{W_L}^2)$ in the SM. Assuming $s_R^2 = s_L^2$ and $M_{W_R} = 1.5$ TeV, we find $|\delta_R| < 0.116$. Since the flavor structure of scotino mixing and their mass-squared differences are unknown, this upper bound could be saturated, and the observation of $\mu \to e\gamma$ may be imminent. The same holds for other lepton flavor violating processes such as $\mu - e$ conversion in nuclei. Note that the contribution to the muon anomalous magnetic moment here is about 10^{-10}, well below the experimental sensitivity. A more comprehensive study, including $D^0 - \bar{D}^0$ mixing [11], will be given elsewhere.

VIII. CONNECTING THE Z' AND DARK-MATTER SEARCHES

As the LHC begins its operation, one of its first possible discoveries could be a Z' through the process $q\bar{q} \to Z' \to l^+l^-$. There are many Z' models, and some of them could also be invoked [12] to explain the CDMS II results. However, the coupling of the dark matter to the Z' in these models is in general not related to the Z' leptonic couplings. Here they are intimately connected and predicted as a function of only s_R^2. In fact, if we assume $s_R^2 = s_L^2$ (i.e. left-right symmetry), then there is no free parameter. Our numerical analysis in this paper is only a rough estimate for illustration, but it points to the important assertion that the Z' interactions in this model are fixed with respect to direct dark-matter search and the detection of Z' itself at an accelerator. In these exciting times of having both the functioning LHC and ongoing dark-matter search experiments, the dark-matter mystery in astroparticle physics may be near a solution.

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