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Dirac neutrino mass generation from dark matter

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In 2006, a simple extension of the standard model was proposed in which neutrinos obtain radiative Majorana masses at one-loop level from their couplings with dark matter, hence the term "scotogenic," from the Greek "scotos" meaning darkness. Here an analogous mechanism for Dirac neutrino masses is discussed in a minimal model. In different ranges of the parameter space, various candidates for dark matter are possible. In particular, the lightest Dirac fermion which appears in the loop diagram generating neutrino mass can be a viable dark-matter candidate. Such a possibility does not exist for the Majorana case. Realistic neutrino mixing in the context of A_4 is discussed. A possible supersymmetric extension is also briefly discussed.

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Dirac neutrino masses have not received much attention in the literature mainly because of two reasons: (1) In the standard model (SM) of particle interactions, there are left-handed lepton doublets $(\nu, l)_L$ and right-handed charged-lepton singlets l_R but no ν_R because it transforms trivially under the $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge symmetry and there is no need for its existence. If it is added in by hand, the neutrino can then obtain a Dirac mass m_D in the same way as all the other fermions (quarks and charged leptons), i.e., from the vacuum expectation value of the scalar Higgs doublet of the SM. However, since ν_R is a neutral singlet, there is no symmetry which prevents it from having a large Majorana mass M. As a result, ν_L obtains an effective small Majorana mass from the seesaw mechanism [1], i.e., $m_\nu \simeq -m_D^2/M$. (2) If a symmetry is imposed in such a way that the lepton number is conserved, the Majorana mass term for ν_R will be forbidden. In that case, because neutrino masses are known to be of order 1 eV or less, the corresponding Yukawa couplings must be of order 10^{-11} or smaller. Such a small value is considered by many to be intrinsically unacceptable.

Nevertheless, up to now, there is not any indisputable evidence for the Majorana nature of the neutrinos from the searches for the neutrinoless double beta decay. Thus, the possibility of Dirac neutrino masses cannot be discounted. To overcome the above theoretical objections, it is proposed in this paper that neutrinos are Dirac fermions, with two important properties. (1) They are protected from becoming Majorana fermions by a $U(1)_{B-L}$ global or gauge symmetry. (2) They are protected from having a tree-level mass by a Z_2 symmetry which is identifiable with that of dark matter, as well as another Z_2 symmetry which sets them apart from other Dirac fermions. The latter symmetry is broken explicitly by soft terms. It may also be replaced by supersymmetry, but that would require a much larger Higgs content. As a result, neutrinos acquire one-loop radiative masses through their couplings with dark matter, hence the term "scotogenic," from the Greek "scotos" meaning darkness. These Dirac neutrino masses can be highly suppressed in the same way that the usual seesaw Majorana neutrino masses are highly suppressed. Their smallness can be also explained by the smallness of the soft terms breaking the Z_2 symmetry.

In 2006, it was proposed [2,3] that neutrinos are massive only because of their couplings with dark matter. This idea connects two of the most important issues in the particle physics and astrophysics. The idea was easily implemented [2] in a simple extension of the SM by adding a second scalar doublet (η^+, η^0) and three neutral singlet fermions N_i which are odd under an extra exactly conserved discrete Z_2 symmetry [4], in analogy to the R parity in supersymmetry. As a result, either $\eta_R = \sqrt{2}\,\mathrm{Re}(\eta^0)$ or the lightest N may be considered a candidate for dark matter. In particular η has been called the "inert" scalar doublet in a model proposed [5] after Ref. [2] and studied by many authors since then [6]. Variations of the original idea also abound and have become an active area of research [7–13].

In almost all previous such applications, neutrino masses have always been assumed to be of Majorana type. Suppose they are exactly Dirac. Is the connection between neutrino mass and dark matter still possible? If so, what are the necessary theoretical ingredients for it to happen, and what are the phenomenological consequences? In Ref. [14], using scalar singlets, a radiative Dirac neutrino mass is obtained; however, in this mechanism, the dark-matter fields do not propagate in the loop. Employing an idea similar to that proposed in Refs. [2,15] suggests a model both for a dark-matter candidate and generation of radiative Dirac neutrino mass. As indicated below, this model shares some features with the model introduced in the present paper. In Ref. [13], a model is introduced in which neutrinos obtain a Dirac mass via a one-loop diagram similar to that in Ref. [15] and a Majorana mass via two-loop diagrams after spontaneous breaking of the lepton number symmetry. In our model described below, the neutrino mass is purely of the Dirac type.

Consider first the imposition of a conserved additive lepton number to protect the neutrino mass from becoming Majorana. We choose to do so by extending the SM to include $B-L$ as either a global or gauged $U(1)$ symmetry. The latter has long been known to be a well-motivated anomaly-free extension which requires the existence of three singlet right-handed neutrinos. Of course, in breaking the gauged $U(1)_{B-L}$, we have to be sure that the global $U(1)_{B-L}$ symmetry of the sector relevant to the present study remains intact. This can be done easily by a scalar field transforming under $U(1)_{B-L}$ but not coupling to other fields with nonzero $B-L$. The second step is to forbid a tree-level Dirac neutrino mass m_ν, and yet allow a tree-level charged-lepton mass m_l. To do this, the simplest way is to impose a $Z_2^{(A)}$ symmetry such that ν^c is odd but all other fermions are even. There is therefore no connection between ν and ν^c at the tree level. To make them connect in one loop, new particles are postulated which are odd under an exactly conserved $Z_2^{(B)}$, and the previous $Z_2^{(A)}$ is allowed to be broken by soft terms. Another way is to make the model supersymmetric as well so that m_l comes from $\Phi_1 = (\phi_1^0, \phi_1^-)$ but m_ν is forbidden to couple to $\Phi_2 = (\phi_2^+, \phi_2^0)$ which is assumed odd under $Z_2^{(B)}$. In either case, we need to add heavy neutral singlet Dirac fermions (N_i, N_i^c) of odd $Z_2^{(B)}$ transforming under $U(1)_{B-L}$ and a neutral singlet scalar χ^0 of odd $Z_2^{(B)}$ which is trivial under $U(1)_{B-L}$.

First, let us consider the minimal nonsupersymmetric model. It is a simple extension of the SM in the same spirit of Ref. [2]. Its particle content is listed in Table I. In addition to the usual particles of the SM, we have added three copies of the Weyl spinors ν^c, three copies of the Dirac spinor pairs (N, N^c), one extra scalar doublet $\eta = (\eta^+, \eta^0)$ and one real scalar χ^0. The $B-L$ symmetry prevents N, N^c as well as ν^c from having a Majorana

TABLE I. Assignments of the particles of the minimal model under $B-L$, $Z_2^{(A)}$ and $Z_2^{(B)}$.

particles	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_{B-L}$	$Z_2^{(A)}$	$Z_2^{(B)}$
(u, d)	3	2	1/6	1/3	+	+
u^c	3*	1	−2/3	−1/3	+	+
d^c	3*	1	1/3	−1/3	+	+
(ν, e)	1	2	−1/2	−1	+	+
e^c	1	1	1	1	+	+
ν^c	1	1	0	1	−	+
(ϕ^+, ϕ^0)	1	2	1/2	0	+	+
(η^+, η^0)	1	2	1/2	0	+	−
χ^0	1	1	0	0	−	−
N	1	1	0	−1	+	−
N^c	1	1	0	1	+	−

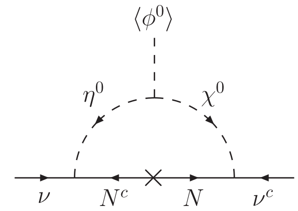

FIG. 1. One-loop generation of Dirac neutrino mass in the minimal model.

mass. Note that $Z_2^{(A)}$ is broken softly by the trilinear term $A\chi\Phi^\dagger\eta$, whereas $Z_2^{(B)}$ remains unbroken. ($\Phi = (\phi^+, \phi^0)$ is the SM Higgs doublet.) The one-loop Dirac neutrino mass is thus generated, as shown in Fig. 1. Note that χ^0 is essential here for a scotogenic Dirac neutrino mass, whereas the scalar singlet considered in Ref. [8] is not needed for a scotogenic Majorana neutrino mass. Whereas a scalar singlet was discussed as dark matter by itself many years ago [16–18], our proposal may be considered a natural justification of its existence.

Let the Yukawa interactions be given by $f_{\alpha k}\nu_\alpha N_k^c \eta^0$ and $h_{k\beta}N_k\nu_\beta^c\chi^0$. Without loss of generality, the A parameter of the trilinear $A\chi\bar{\phi}^0\eta^0$ term may always be chosen real, as well as the vacuum expectation value $\langle\phi^0\rangle = v$. Let $\eta^0 = (\eta_R + i\eta_I)/\sqrt{2}$, then there is a mixing between η_R and χ^0, but not between η_I and χ^0. Assuming in addition that η_I is a mass eigenstate and denoting the mass eigenstates of the (χ^0, η_R) sector as $\zeta_{1,2}$ with mixing angle θ, the one-loop Dirac neutrino mass matrix is then given by

$$(\mathcal{M}_\nu)_{\alpha\beta} = \frac{\sin\theta\cos\theta}{16\pi^2\sqrt{2}}\sum_k f_{\alpha k}h_{k\beta}m_{N_k} \times\left[\frac{m_{\zeta_1}^2}{m_{\zeta_1}^2 - m_{N_k}^2}\ln\frac{m_{\zeta_1}^2}{m_{N_k}^2} - \frac{m_{\zeta_2}^2}{m_{\zeta_2}^2 - m_{N_k}^2}\ln\frac{m_{\zeta_2}^2}{m_{N_k}^2}\right]. \quad (1)$$

This is in complete analogy to that of the radiative Majorana seesaw [2], with suppression of the neutrino mass from the usual assumption of very large m_N (now Dirac) as well as the loop factor. In addition, this diagram is only nonzero because of the soft breaking of $Z_2^{(A)}$. Thus, it is natural for the parameter A to be small. In the limit $A = 0$, the mixing angle θ in the above equation would be zero.

We assume that there are three copies of (N, N^c) so that all three neutrinos obtain scotogenic Dirac masses. If there is only one copy, then two neutrinos will be massless, which is clearly unrealistic. If there are two copies, one will be massless, which is acceptable as far as present neutrino phenomenology is concerned. From Table I, it can be easily confirmed that with three copies of ν^c, $U(1)_{B-L}$ will be anomaly-free.

In this model, Φ is the SM Higgs doublet with the usual Higgs boson H as its only physical degree of freedom. It has the usual SM decay modes, except for corrections due to its interactions with η and χ^0. For example, H may decay into $\zeta_1\zeta_1$ if kinematically allowed. If ζ_1 is dark matter, this decay would then be invisible. It would affect the search for the SM Higgs boson, as studied already in Ref. [19]. Another possible effect is that the coupling of H to $\eta^+\eta^-$ would change the one-loop decay of H to $\gamma\gamma$, thus affecting also the search for the SM Higgs boson via this channel. A third effect is the existence of the quartic $\chi\chi\Phi^\dagger\Phi$ coupling, which may contribute significantly to the effective potential of H and modify its stability condition as a function of mass. It may also induce a one-loop contribution to the H^3T term at finite temperature to cause a first-order phase transition needed for the electroweak baryogenesis.

The couplings $f_{\alpha k}L_\alpha N_k^c\eta$ contribute to radiative lepton flavor violating rare decays:

$$\Gamma(l_\alpha^- \to l_\beta^-\gamma) = \frac{m_\alpha^3}{16\pi}\sigma_R^2, \tag{2}$$

where

$$\sigma_R = \sum_k e f_{\alpha k} f_{\beta k}^* m_\alpha \frac{i}{16\pi^2 m_{\eta^+}^2}\left[\frac{t\ln t}{2(t-1)^4} + \frac{t^2-5t-2}{12(t-1)^3}\right], \tag{3}$$

with $t = (m_{N_k}^2/m_{\eta^+}^2)$. For $t\to 0$, $t\to\infty$ and $t\to 1$, the combination in the last parenthesis of Eq. (3) converges respectively to $1/6$, $1/(12t)$ and $1/24$. For $m_{N_k} \gg m_{\eta^+}$, which is the seesaw limit, we find

$$\left(\sum_k \frac{f_{\alpha k} f_{\beta k}^*}{m_{N_k}^2}\right)^{1/2} \sim 8\times 10^{-5}\left(\frac{\mathrm{B}(l_\alpha \to l_\beta\gamma)}{10^{-12}}\right)^{1/4}\ \mathrm{GeV}^{-1}. \tag{4}$$

We will consider first this scenario, so that the dark-matter candidate of our model is the lightest of the three exotic neutral scalars: $\zeta_{1,2}$ or η_I.

The most general scalar potential consisting of Φ, η, and χ is given by

$$\begin{aligned} V = {}& \mu_1^2\Phi^\dagger\Phi + \mu_2^2\eta^\dagger\eta + \frac{1}{2}\mu_3^2\chi^2 + \frac{1}{2}\lambda_1(\Phi^\dagger\Phi)^2 \\ & + \frac{1}{2}\lambda_2(\eta^\dagger\eta)^2 + \lambda_3(\Phi^\dagger\Phi)(\eta^\dagger\eta) + \lambda_4(\Phi^\dagger\eta)(\eta^\dagger\Phi) \\ & + \frac{1}{2}\lambda_5(\Phi^\dagger\eta)^2 + \mathrm{H.c.} + \frac{1}{4}\lambda_6\chi^4 + \frac{1}{2}\lambda_7(\Phi^\dagger\Phi)\chi^2 \\ & + \frac{1}{2}\lambda_8(\eta^\dagger\eta)\chi^2 + A\chi\Phi^\dagger\eta + \mathrm{H.c.} \end{aligned} \tag{5}$$

This potential preserves $Z_2^{(B)}$ and breaks $Z_2^{(A)}$ softly by the last term. The parameter A may be chosen real by a phase rotation of η relative to Φ, but then λ_5 is in general complex. For simplicity, we choose it to be real so that η_I is a mass eigenstate and decouples from the (χ^0, η_R) sector. The resulting mass spectrum is given by

$$m_H^2 = 2\lambda_1 v^2, \tag{6}$$

$$m_{\eta^+}^2 = \mu_2^2 + \lambda_3 v^2, \tag{7}$$

$$m_{\eta_I}^2 = \mu_2^2 + (\lambda_3+\lambda_4-\lambda_5)v^2, \tag{8}$$

$$m_{(\chi,\eta_R)}^2 = \begin{pmatrix} \mu_3^2 + \lambda_7 v^2 & \sqrt{2}Av \\ \sqrt{2}Av & \mu_2^2 + (\lambda_3+\lambda_4+\lambda_5)v^2 \end{pmatrix}. \tag{9}$$

This is very similar to previous studies such as Ref. [8] with an important conceptual difference. Since the parameter A breaks $Z_2^{(A)}$, it may be argued that it is small. This suppresses the radiative neutrino Dirac mass as well as the mixing between η_R and χ. Hence the dark-matter candidate of this model can be dominantly a singlet and as a result, it can naturally evade the constraints from the electroweak interactions of a doublet. If there is no $Z_2^{(A)}$ symmetry, the mixing between η_R and χ is then arbitrary, as in previous studies. Another difference is that η_I is not involved in the one-loop neutrino mass, contrary to the original Majorana case of Ref. [2]. The above possibility has also been discussed in Ref. [15]. In the following, we introduce a new possibility for dark-matter candidate within the present scenario.

Since m_{N_k} are assumed to be very large in this scenario, the annihilation of the dark-matter scalars in this model do not proceed via their Yukawa interactions, but rather through their gauge or scalar interactions. Examples of the latter have been discussed extensively in the literature [16–18,20–24].

We consider next the lightest N_k (call it N_1) as dark matter. As shown previously [25], this is subject to severe phenomenological constraints in the original model [2] of scotogenic Majorana neutrino mass. The reason is as follows. In order for N_1N_1 annihilation to account for the correct relic abundance, the η masses cannot be too heavy and the Yukawa couplings $f_{\alpha k}$ cannot be too small. However, these values are severely constrained by experimental upper limits on the $\mu\to e\gamma$ rate, as already discussed. It is thus not a viable option, without some detailed fine tuning of parameters. To retain N_1 as a natural dark-matter candidate, new interactions involving N_1 need to be postulated, such as a singlet scalar [26]. In our present model, the $h_{kj}N_k\nu_j^c\chi^0$ Yukawa couplings are exactly what are required. They are not constrained by flavor-changing charged-lepton radiative decays, so they can be large enough for a realistic $N_1\bar{N}_1$ annihilation cross section to account for the relic abundance of dark matter in the Universe today. In this scenario, the $f_{\alpha k}$ Yukawa couplings as well as the A parameter are small and the mass of ζ_1 (which is mostly composed of χ) is not much greater than m_{N_1}.

Combining (ν, ν^c) and (N, N^c) to form the four-component Dirac fermions ν and N, their Yukawa interactions are given by

$$\mathcal{L}_Y = f_{\alpha k}\bar{N}_k\left(\frac{1-\gamma_5}{2}\right)(\nu_\alpha\eta^0 - l_\alpha\eta^+) + h_{k\beta}\bar{N}_k\left(\frac{1+\gamma_5}{2}\right)\nu_\beta\chi^0 + \text{H.c.}, \qquad (10)$$

where in this four-component notation, $[(1+\gamma_5)/2]\nu$ represents ν^c going backwards. For the dark-matter candidate N_1, we assume $h_{1\beta}$ to be dominant, then

$$\sigma(N_1+\bar{N}_1\to\nu_\alpha+\bar{\nu}_\beta) = \sum_{\alpha,\beta}\frac{|h^*_{1\alpha}h_{1\beta}|^2}{32\pi v_{\rm rel}}\frac{m^2_{N_1}}{(m^2_{N_1}+m^2_{\zeta_1})^2} < \sum_{\alpha,\beta}\frac{|h^*_{1\alpha}h_{1\beta}|^2}{128\pi v_{\rm rel}m^2_{N_1}}, \qquad (11)$$

where to reach the last inequality we have used $m_{\zeta_1} > m_{N_1}$. Similarly,

$$\sigma(N_1+N_1\to\nu_\alpha+\nu_\beta) = \sum_{\alpha,\beta}\frac{|h^*_{1\alpha}h^*_{1\beta}|^2}{32\pi v_{\rm rel}}\frac{m^2_{N_1}}{(m^2_{N_1}+m^2_{\zeta_1})^2} < \sum_{\alpha,\beta}\frac{|h^*_{1\alpha}h^*_{1\beta}|^2}{128\pi v_{\rm rel}m^2_{N_1}}. \qquad (12)$$

Setting the sum of the two annihilation cross sections times the relative velocity equal to one picobarn, we find

$$m_{N_1} < \left(\sum_{\alpha,\beta}|h^*_{1\alpha}h_{1\beta}|^2\right)^{1/2}(1.4\ \text{TeV}). \qquad (13)$$

For $|h_{1\alpha}| < 1$, we then obtain $m_{N_1} < 4.2$ TeV. With such light N_1, the seesaw mechanism is not very effective. The smallness of the neutrino masses can be justified by the smallness of the trilinear A term which softly breaks $Z_2^{(A)}$ and the smallness of the f Yukawa couplings. If the h couplings were not available, the cross section must have then come from the f couplings, which are restricted by $\mu\to e\gamma$, so the annihilation cross section would in general be too small for N_1 to be a viable dark-matter candidate. If the $B-L$ symmetry is gauged, there should be another annihilation mode $N+\bar{N}\to Z'\to\nu+\bar{\nu},\ l+\bar{l},\ q+\bar{q}$. This cross section is given by Khalil *et al.* [27]

$$\sigma = \frac{g^4_{Z'}m^2_{N_1}}{\pi v_{\rm rel}(4m^2_{N_1}-m^2_{Z'})^2}. \qquad (14)$$

The present lower bound on $m_{Z'}$ from the Large Hadron Collider (LHC) [28] is estimated to be about 2 TeV. For $g_{Z'} = \sqrt{(5/8)}g_Y = 0.28$ (i.e., the $SO(10)$ limit), $m_{Z'} = 2$ TeV, and $\sigma v_{\rm rel} = 1$ pb, we find $m_{N_1} = 900$ GeV. In this case, $N_1\bar{N}_1$ production from Z' decay at the LHC is possible, as studied previously [29], except that N_1 is now dark matter. It may however be inferred from the increase of the Z' invisible width on top of the expected $Z'\to\nu\bar{\nu}$ mode. As N_1 is otherwise very difficult to produce, the existence of Z' seems to be the only realistic chance for it to be observed at the LHC, but still only indirectly. If η^+ is light enough, it can be produced at the LHC. The subsequent decay of η^+ into N_1 and a charged lepton is a possible signature, as discussed in Ref. [30].

As for direct detection of dark matter in underground experiments, if $B-L$ is not gauged, then N_1 has no interaction with nuclei. If $B-L$ is gauged, then the elastic scattering of N_1 with nuclei may proceed through Z' exchange. The cross section per nucleon is given by Khalil *et al.* [27]

$$\sigma_0 = \frac{4m_P^2}{\pi}\frac{g^4_{Z'}}{m^4_{Z'}}. \qquad (15)$$

For $g_{Z'} = 0.28$ and $m_{Z'} = 2$ TeV, this implies $\sigma_0 = 1.7\times 10^{-7}$ pb, which exceeds the XENON100 bound [31] of about 7×10^{-8} pb for $m_{N_1} = 900$ GeV. This means that in this case, $g_{Z'}/m_{Z'}$ should be reduced by a factor of 1.25 or more.

This minimal model is also very suitable for the implementation of the non-Abelian discrete A_4 symmetry [32] to the neutrino mass matrix [33]. In the charged-lepton sector, let $(\nu_i, l_i)\sim\underline{3}$ under A_4, and either $l^c_i\sim\underline{1},\underline{1}',\underline{1}''$ as in Ref. [32] or $l^c_i\sim\underline{3}$ as in Ref. [34], then with $\Phi\sim\underline{3}$ or $\underline{3}+\underline{1}$, and A_4 breaking to the residual symmetry Z_3, the charged-lepton mass matrix is diagonalized by the well-known unitary matrix

$$U_L = \frac{1}{\sqrt{3}}\begin{pmatrix}1 & 1 & 1\\ 1 & \omega & \omega^2\\ 1 & \omega^2 & \omega\end{pmatrix}, \qquad (16)$$

where $\omega = \exp(2\pi i/3)$. In the neutrino sector, let $\nu^c_i\sim\underline{3}$, $\eta\sim\underline{1}$, and $\chi\sim\underline{1}+\underline{3}$, with the soft scalar trilinear $\chi\Phi^\dagger\eta$ terms to break A_4, the neutrino mass matrix becomes [33]

$$\mathcal{M}_\nu = \begin{pmatrix}a & f & e\\ f & a & d\\ e & d & a\end{pmatrix}. \qquad (17)$$

If $e=f=0$, then neutrino mixing is tribimaximal, i.e., $\sin^2\theta_{12} = 1/3$, $\sin^2\theta_{23} = 1/2$, $\theta_{13} = 0$. This was known to be a good approximation of the measured neutrino mixing angles. However, two recent experiments have measured θ_{13} to be definitely nonzero, i.e.,

$$\sin^2 2\theta_{13} = 0.092 \pm 0.016(\text{stat}) \pm 0.005(\text{syst}) \qquad (18)$$

from the Daya Bay Collaboration [35], and

$$\sin^2 2\theta_{13} = 0.113 \pm 0.013(\text{stat}) \pm 0.019(\text{syst}) \qquad (19)$$

from the RENO Collaboration [36]. In that case, e and f should be nonzero. Let

$$\epsilon = \frac{e-f}{d\sqrt{2}}, \qquad \delta = \frac{e+f}{d\sqrt{2}}. \qquad (20)$$

The parameters a, d, e, f are complex, and for small e, f, the eigenvalues of $\mathcal{M}_\nu$ are $a+d$, a, and $a-d$. We can always choose a to be real, the phase of d is then determined by the absolute values of the three masses [37]. For the small values of e and f, we find

$$\theta_{13} = -\frac{\epsilon}{\sqrt{3}}, \qquad \tan^2\theta_{12} = \frac{1}{2}\left[\frac{(1-\sqrt{2}\mathrm{Re}\delta)^2 + 2(\mathrm{Im}\delta)^2}{(1+\mathrm{Re}\delta/\sqrt{2})^2 + (\mathrm{Im}\delta)^2/2}\right]. \tag{21}$$

Thus, a nonzero θ_{13} and a value of $\tan^2\theta_{12}$ smaller than 0.5 can be obtained. More precisely, the neutrino mass matrix in the tribimaximal basis is now of the form

$$\mathcal{M}_\nu^{(1,2,3)} = \begin{pmatrix} m_1 & m_6 & 0 \\ m_6 & m_2 & m_5 \\ 0 & m_5 & m_3 \end{pmatrix} = \begin{pmatrix} a+d & \delta d & 0 \\ \delta d & a & \epsilon d \\ 0 & \epsilon d & a-d \end{pmatrix}. \tag{22}$$

If $\delta = \epsilon = 0$, the tribimaximal mixing is then recovered. This differs from the originally proposed deviation [33] for A_4, which was updated recently [38], i.e.,

$$\mathcal{M}_\nu^{(1,2,3)} = \begin{pmatrix} m_1 & 0 & m_4 \\ 0 & m_2 & m_5 \\ m_4 & m_5 & m_3 \end{pmatrix} = \begin{pmatrix} a+d-(b+c)/2 & 0 & i\sqrt{3}/2(c-b) \\ 0 & a+b+c & \sqrt{2}e \\ i\sqrt{3}/2(c-b) & \sqrt{2}e & a-d-(b+c)/2 \end{pmatrix}. \tag{23}$$

Given that $m_4 = 0$ in Eq. (22), we obtain the approximate relationship

$$\sin^2 2\theta_{23} \simeq 1 - 8[\mathrm{Re}(U_{e3})]^2. \tag{24}$$

Using the experimental bound $\sin^2 2\theta_{23} > 0.92$, we find $|\mathrm{Re}(U_{e3})| < 0.1$. If we take the central value of $|U_{e3}|$ to be 0.16 (corresponding to $\sin^2 2\theta_{13} = 0.1$), we then obtain $|\tan\delta_{CP}| > 1.3$ in this model. Details are given elsewhere [39].

Below we also mention briefly how a supersymmetric model of scotogenic neutrino mass may be constructed. Consider the superfield content listed in Table II.

There are two one-loop diagrams contributing to the Dirac neutrino mass as shown in Fig. 2. Note that supersymmetry is broken by the soft scalar trilinear $\chi_2^0\phi_1^0\phi_2^0$ and bilinear $\tilde{N}\tilde{N}^c$ terms. There are now many particles of odd $Z_2^{(B)}$ as well as superpartners of odd R parity. There are thus at least two dark-matter candidates [40]. Obviously the details of the dark sector are much more complicated. We will not study them further in this paper.

TABLE II. Assignments of the particles of the supersymmetric model under $B-L$, $Z_2^{(A)}$, and $Z_2^{(B)}$.

superfields	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_{B-L}$	$Z_2^{(A)}$	$Z_2^{(B)}$
(u, d)	3	2	1/6	1/3	+	+
u^c	3*	1	−2/3	−1/3	+	+
d^c	3*	1	1/3	−1/3	+	+
(ν, e)	1	2	−1/2	−1	+	+
e^c	1	1	1	1	−	+
ν^c	1	1	0	1	−	+
(ϕ_1^0, ϕ_1^-)	1	2	−1/2	0	−	+
(ϕ_2^+, ϕ_2^0)	1	2	1/2	0	+	−
(ϕ_3^0, ϕ_3^-)	1	2	−1/2	0	+	+
(ϕ_4^+, ϕ_4^0)	1	2	1/2	0	+	+
χ_1^+	1	1	1	0	−	+
χ_1^0	1	1	0	0	−	+
χ_2^0	1	1	0	0	−	−
χ_2^-	1	1	−1	0	+	−
N	1	1	0	−1	+	−
N^c	1	1	0	1	+	−

In conclusion, we have studied a minimal model of radiative Dirac neutrino mass induced by dark matter. In order for the scotogenic Dirac neutrino mass to occur in one loop, we need to introduce a scalar singlet χ^0 which mixes with the neutral component of a new electroweak scalar doublet (η^+, η^0). It is thus a good theoretical justification for the existence of χ^0. In addition to the possibility of direct production at the LHC, the presence of η^+ can modify the Higgs decay mode to $\gamma\gamma$. As shown in Ref. [41], if the λ_3 coupling in Eq. (5) is negative, it can lead to the enhancement of $\mathrm{Br}(H \to \gamma\gamma)$ in conformity of the recent observation at the LHC [42]. Moreover, the quartic coupling of χ^0 with Higgs can stabilize its potential against radiative corrections.

This minimal model also requires three heavy neutral Dirac fermions N_i. Depending on the mass spectrum, the dark matter might be either the lightest Dirac fermion N_1 or one of the neutral scalars; i.e., the imaginary component η_I of η^0 or a linear combination of the real component η_R of η^0 and χ^0. In the latter case, depending on the mixing between η_R and χ^0, which should be small because of the soft breaking of $Z_2^{(A)}$, the annihilation rate due to the

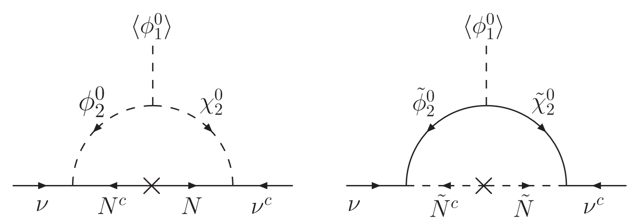

FIG. 2. One-loop generation of Dirac neutrino mass in the supersymmetric case.

electroweak interactions can be made equal to about 1 pb which is a value dictated by the dark-matter abundance in the thermal dark-matter scenario.

If N_1 is the dark-matter candidate, its annihilation can proceed via its Yukawa coupling with the right-handed neutrinos and χ^0. This is a possibility that does not exist within the scotogenic Majorana neutrino mass model because in that case the bounds from the $\mu \to e\gamma$ constraints restrict the annihilation cross section of the N_1 pair below the required value. At the LHC, N_1 can then be produced via the decay of η^+ and η^- along with a charged lepton [30].

The $B - L$ symmetry used to maintain the conservation of lepton number can be gauged. In that case, the present LHC lower bound on $m_{Z'}$ is about 2 TeV. The interaction with the Z' boson provides another route for the annihilation of the N_1 pair as well as a portal for the interaction with quarks and hence direct detection. The bound from the XENON100 experiment already constrains the parameter space.

This minimal model is also suitable for implementing an A_4 symmetry in such a way that nonzero θ_{13} and large δ_{CP} may be obtained. We have also briefly mentioned how a supersymmetric extension can be constructed.

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