Abelian extension of standard model with four generations

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An Abelian gauge extension of the standard model is proposed with a fourth generation. The fourth-generation fermions obtain their masses from a heavier Higgs doublet which makes no tree- level contributions to the first three generations' masses. Light first-three-generations' neutrino masses continue to have a type-I seesaw explanation, whereas the fourth-generation neutrino turns out to be a heavy Dirac neutrino. In the minimal version of such a model with no off-diagonal Yukawa couplings between the fourth and the first three generations, such a heavy Dirac neutrino can be long-lived on cosmological time scales. In this model, the stated LHC exclusion range 120 GeV $< m_H <$ 600 GeV on the lighter Higgs placed in the context of a generic fourth-generation standard model is evaded. Also, the Dirac fourth-generation neutrino in this model, if stable, would constitute up to 1% of the cold dark matter in the Universe.

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I. INTRODUCTION

The standard model of particle physics has been phenomenologically the most successful low-energy effective theory for the last few decades. The predictions of the standard model have been verified experimentally with a very high accuracy, except the missing Higgs boson. Despite its phenomenological success, we all now know that this model neither addresses many theoretical issues like the gauge hierarchy problem nor provides a complete understanding of various observed phenomena like nonzero neutrino masses, dark matter, etc. A great deal of works has been done so far on various possible extensions of the standard model, although none of them can be called a complete phenomenological model. Such extensions generally involve incorporating some extra symmetries into the standard model. These symmetries may be an extra gauge symmetry like in the left-right symmetric model [1–5], grand unified theories [6], etc. Another highly motivating symmetry is supersymmetry, the symmetry between bosons and fermions. The inclusion of supersymmetry into the standard model has many advantages, among which stabilizing the Higgs mass against the radiative corrections, providing a cold dark matter candidate (which can be made stable by incorporating R-parity), and making the gauge coupling constants unify at high energy are significantly important.

One very nonconventional extension of the standard model is to go beyond three generations of quarks and leptons [7–9]. Although the number of light neutrinos are constrained to three from big bang nucleosynthesis as well as precision measurement of Z boson decay width, there is absolutely nothing which prevents us from adding a heavy fourth generation with the corresponding fourth neutrino heavier than $M_Z/2 = 45$ GeV. Similar lower bounds will exists for charged fermions also [10]. We know that the smallness of three standard model neutrino masses [11–14] can be naturally explained via the see-saw mechanism [15–18]. After incorporating a fourth-generation neutrino into the standard model (SM4), the see-saw mechanism should be such that it gives one very heavy and three light neutrinos. Such analysis within the context of SM4 was done in Ref. [19]. Motivated by the idea of introducing a separate Higgs with larger vacuum expectation value (VEV) to account for the heavier fermion masses such as top-quark mass [20] or fourth-generation fermion masses [21], here, also, we propose a model with an extended Higgs sector and a fourth chiral family of quarks and leptons. However, our model differs from the earlier models in the sense that we incorporate an additional Abelian gauge sector which couples to the first three generations differently than it does to the fourth-generation fermions. Such nonuniversal gauge couplings automatically force one to have at least two different Higgs doublets to give masses to the fermions. In our model, Majorana neutrino masses of the first three generations arise after spontaneous gauge symmetry breaking, whereas the fourth-generation neutrino turns out to be a Dirac neutrino. We also point out that our model reproduces the Higgs-fermion structure considered by the authors in Ref. [22]. Because of the existence of a heavy Higgs doublet, which couples only to the fourth-generation fermions, and a lighter Higgs, which couples only to the first three generations, our model can evade the LHC exclusion range 120 GeV $< m_H <$ 600 GeV placed within the context of generic SM4 [23]. The fourth - generation Dirac neutrino can be long-lived in this minimal model if we set the off-diagonal Yukawa couplings between first three generations and the fourth generation to zero. We further show that the long-lived heavy Dirac fourth-generation neutrino can contribute up to 1% of the total dark matter in the Universe.

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Fourth-generation chiral fermions can have many other interesting phenomenological consequences, for example, in rare B and K decays [24]. It can also account for like-sign dimuon charge asymmetry observed by the D0 Collaboration recently [25], as was discussed in Ref. [26]. However, here, we restrict ourselves to the issue of neutrino mass and dark matter only.

This paper is organized as follows. In Sec. II, we discuss the $U(1)_X$ extended standard model with four generations, the spontaneous gauge symmetry breaking, and neutrino mass. We briefly comment on the fourth-generation and LHC Higgs search in Sec. III. In Sec. IV, we calculate the relic abundance of a stable fourth-generation Dirac neutrino within the $U(1)_X$ model framework and then conclude in Sec. V.

II. $U(1)_X$ EXTENDED MODEL WITH A PURELY DIRAC FOURTH-GENERATION NEUTRINO

The Abelian gauge extension of standard model is one of the best motivating examples of beyond-standard-model physics. For a review, see Ref. [27]. Such a model is also motivated within the framework of grand unified theory models, for example, E_6. The supersymmetric versions of such models have an additional advantage in the sense that they provide a solution to the minimal supersymmetric standard model μ problem. Here, we consider an extension of the standard model gauge group with one Abelian $U(1)_X$ gauge symmetry. Thus, the model we are going to work on is an $SU(3)_c \times SU(2)_L \times U(1)_Y \times U(1)_X$ gauge theory with four chiral generations. We will consider family non-universal $U(1)_X$ couplings such that the first three generations and the fourth generation have different charges under $U(1)_X$. Since the coupling is universal in the first three generations and we set the off-diagonal Yukawa couplings between first three generation and the fourth generation to zero in this minimal model, there will not be any severe constraints from flavor-changing-neutral-current limits. As it will be clear later, our purpose in choosing such nonuniversal couplings is to allow Majorana mass terms for the first-three-generation right-handed neutrinos only and not for the fourth-generation right-handed neutrino.

A. The matter content

The fermion content of our model is

$$Q_i = \begin{pmatrix} u \\ d \end{pmatrix} \sim \left(3, 2, \frac{1}{6}, n_1\right), \quad L_i = \begin{pmatrix} \nu \\ e \end{pmatrix} \sim \left(1, 2, -\frac{1}{2}, n_4\right),$$
$$u_i^c \sim \left(3^*, 1, \tfrac{2}{3}, n_2\right), \quad d_i^c \sim \left(3^*, 1, -\tfrac{1}{3}, n_3\right),$$
$$e_i^c \sim (1, 1, -1, n_5), \quad \nu_i^c \sim (1, 1, 0, n_6),$$

where $i = 1, 2, 3$ goes over the three generations of the standard model and the numbers in the parentheses correspond to the quantum number under the gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y \times U(1)_X$. Similarly, the fourth-generation fermions are

$$Q_4 = \begin{pmatrix} u \\ d \end{pmatrix} \sim \left(3, 2, \frac{1}{6}, n_8\right), \quad L_4 = \begin{pmatrix} \nu \\ e \end{pmatrix} \sim \left(1, 2, -\frac{1}{2}, n_{11}\right),$$
$$u_4^c \sim \left(3^*, 1, \tfrac{2}{3}, n_{10}\right), \quad d_4^c \sim \left(3^*, 1, -\tfrac{1}{3}, n_9\right),$$
$$e_4^c \sim (1, 1, -1, n_{12}), \quad \nu_4^c \sim (1, 1, 0, n_{13}).$$

The quantum numbers of the fermions under the new gauge symmetry $U(1)_X$ should satisfy the following anomaly cancellation conditions:

$$[SU(3)_c]^2 U(1)_X\colon\ 3(2n_1 - n_2 - n_3) + (2n_8 - n_9 - n_{10}) = 0, \qquad [SU(2)_L]^2 U(1)_X\colon\ 3\left(\frac{3n_1}{2} + \frac{n_4}{2}\right) + \left(\frac{3n_8}{2} + \frac{n_{11}}{2}\right) = 0,$$
$$[U(1)_Y]^2 U(1)_X\colon\ 3\left(\frac{n_1}{6} - \frac{4n_2}{3} - \frac{n_3}{3} + \frac{n_4}{2} - n_5\right) + \left(\frac{n_8}{6} - \frac{4n_{10}}{3} - \frac{n_9}{3} + \frac{n_{11}}{2} - n_{12}\right),$$
$$U(1)_Y[U(1)_X]^2\colon\ 3(n_1^2 - 2n_2^2 + n_3^2 - n_4^2 + n_5^2) + (n_8^2 - 2n_{10}^2 + n_9^2 - n_{11}^2 + n_{12}^2) = 0,$$
$$[U(1)_X]^3\colon\ 3(6n_1^3 - 3n_2^3 - 3n_3^3 + 2n_4^3 - n_5^3) + (6n_8^3 - 3n_{10}^3 - 3n_9^3 + 2n_{11}^3 - n_{12}^3) - 3n_6^3 - n_{13}^3 = 0,$$
$$U(1)_X\colon\ 3(6n_1 - 3n_2 - 3n_3 + 2n_4 - n_5) + (6n_8 - 3n_9 - 3n_{10} + 2n_{11} - n_{12}) - 3n_6 - n_{13} = 0.$$

We consider one possible solution to the above anomaly matching conditions, which allows us to choose a minimal scalar sector for our purposes:

$$n_1 = n_4 = n_8 = n_{11} = 0, \qquad n_2 = -n_3 = -n_5 = n_6,$$
$$n_9 = -n_{10} = n_{12} = -n_{13}.$$

Here, we are interested in a solution where $n_2 \neq n_9$.

B. Fermion and gauge boson masses

For the particular solution of the anomaly matching conditions mentioned in the previous subsection, the Higgs fields required to give rise to the fermion Dirac masses are $H_1(1, 2, -\frac{1}{2}, n_2)$ and $H_2(1, 2, -\frac{1}{2}, n_{10})$, where H_1 gives rise to the first-three-generation Dirac masses, and the latter gives rise to the fourth-generation Dirac masses. Since the Higgs scalars do not contribute to the anomalies, we can choose their representation under the gauge group independent of the above anomaly matching conditions. Since $n_{13} \neq 0$, it is clear that the Majorana mass term for the fourth-generation singlet neutrino is not allowed in the Lagrangian, although the same is allowed for the first three generations if we include one more Higgs $S(1, 1, 0, -2n_2)$, whose VEV will give rise to the Majorana mass term of the first-three-generation neutrinos. This singlet Higgs field also plays a crucial role in the gauge symmetry breaking, resulting in a heavy $U(1)_X$ boson. It also keeps the scale of $U(1)_X$ symmetry breaking higher than the electroweak symmetry breaking. Because of the absence of Majorana mass term of the fourth-generation neutrino, we arrive at three light ($\sim$ eV) standard-model Majorana neutrinos and one Dirac neutrino. This is exactly what we want: a fourth-generation Dirac neutrino whose mass can be easily adjusted to lie above the experimental lower bounds. The Yukawa Lagrangian is

$$\begin{aligned}\mathcal{L}_Y = {} & Y_u \bar{Q}_L H_1 u_R + Y_d \bar{Q}_L H_1^\dagger d_R + Y_\nu \bar{L} H_1 N_R \\ & + Y_e \bar{L} H_1^\dagger e_R + f S N_R N_R + Y_u^{(4)} \overline{Q^{(4)}}_L H_2 u_R^{(4)} \\ & + Y_d^{(4)} \overline{Q^{(4)}}_L H_2^\dagger d_R^{(4)} + Y_\nu^{(4)} \overline{L^{(4)}} H_2 N_R^{(4)} \\ & + Y_e^{(4)} \overline{L^{(4)}} H_2^\dagger e_R^{(4)}.\end{aligned}$$

Let $\langle H_{1,2}\rangle = v_{1,2}$, $\langle S\rangle = s$. For the minimal version of our model, we set the off-diagonal Yukawa couplings between the first three generations and the fourth generation to zero; that is, $Y_{4i} = 0$ where $i = 1, 2, 3$. Thus, the first-three-generation right-handed neutrinos acquire Majorana neutrino mass proportional to $f\langle S\rangle = fs$, whereas the fourth-generation right-handed neutrino does not acquire any Majorana mass. Hence, the first-three-generation neutrino mass arises from the see-saw mechanism, whereas the fourth-generation neutrino is a Dirac neutrino due to the absence of a corresponding Majorana mass term. Under the assumption that the off-diagonal Yukawa couplings between first three and the fourth generations are zero, the fourth generation neutrino can be a long-lived particle and hence can play a nontrivial role in cosmology. We pursue this study in the next section in the context of dark matter. It may be noted that, choosing a nonminimal solution to the anomaly matching conditions and hence an extended Higgs sector, the off-diagonal Yukawa couplings can naturally be set to zero.

The gauge boson masses will come from the kinetic terms of the Higgs fields. Denoting the $SU(2)_L$, $U(1)_Y$, $U(1)_X$ gauge fields as $(W_1^\mu, W_2^\mu, W_3^\mu)$, Y^μ, X^μ respectively, we write the neutral gauge boson mass matrix in the (W_3^μ, Y^μ, X^μ) basis as

$$M = \frac{1}{2}\begin{pmatrix} g_2^2(v_1^2 + v_2^2) & g_1 g_2(v_1^2 + v_2^2) & -g_2 g_x(n_2 v_1^2 + n_{10} v_2^2) \\ g_1 g_2(v_1^2 + v_2^2) & g_1^2(v_1^2 + v_2^2) & -g_1 g_x(n_2 v_1^2 + n_{10} v_2^2) \\ -g_2 g_x(n_2 v_1^2 + n_{10} v_2^2) & -g_1 g_x(n_2 v_1^2 + n_{10} v_2^2) & 4g_x^2(n_2^2 v_1^2 + n_{10}^2 v_2^2 + 4n_2^2 s^2) \end{pmatrix}. \tag{1}$$

The off-diagonal elements of the above mass matrix indicate nonzero mixings between the standard model gauge bosons and the $U(1)_X$ boson. The nonzero mixings arise since the doublet Higgs fields H_1, H_2 are charged under both the standard model gauge group as well as the extra $U(1)_X$. However, these mixings have to be very small so as not to be in conflict with the electroweak precision measurements. The simplest way to evade all these restrictions is to consider zero mixings, which can be achieved simply by the following constraint:

$$n_2 v_1^2 + n_{10} v_2^2 = 0.$$

The charged W boson mass is $M_W^2 = \frac{1}{2} g_2^2(v_1^2 + v_2^2)$. Using this and the above constraint, we get

$$v_1^2 = \frac{2M_W^2}{g_2^2} \frac{n_{10}}{n_{10} - n_2}, \tag{2}$$

$$v_2^2 = \frac{2M_W^2}{g_2^2} \frac{-n_2}{n_{10} - n_2}. \tag{3}$$

Here, we restrict our discussion to this simple situation of no mixing with the extra $U(1)_X$ boson. Thus, the first-three-generation neutrino mass comes from the usual type-I see-saw formula:

$$m_\nu = -\frac{v_1^2}{M_R} Y_\nu Y_\nu^T,$$

and the fourth-generation neutrino has a Dirac mass $m_{\nu 4} = Y_\nu^{(4)} v_2$. The lower bounds on the fourth-generation charged fermion masses [10] are

$$m_{t'} \geq 256\,\text{GeV}, \quad m_{b'} \geq 199\,\text{GeV}, \quad m_{\tau'} \geq 100\,\text{GeV},$$

which can be satisfied by a suitable choice of Yukawa couplings. The requirement of perturbativity of Yukawa couplings ($Y^2 < 4\pi$) at the electroweak scale restricts the VEV of H_2 to be, at most, $256/\sqrt{4\pi}$, and, hence,

$$\frac{-n_2}{n_{10} - n_2} > \frac{g_2^2}{2M_W^2} \frac{256}{\sqrt{4\pi}}. \tag{4}$$

In this particular model, we assume that the fourth-generation fermions have no mixing with the first three generations at tree level, and, hence, the fourth-generation neutrino (if lighter than the corresponding charged fermion) can be stable or long-lived and may play a role as dark matter. We do the analysis of relic density of such a stable heavy neutrino in the next section.

III. FOURTH GENERATION AND HIGGS SEARCHES AT LHC

For the last year, the LHC has been producing a large number of data sets to confirm many of the already-established facts as well as to rule out a lot of parameter space for many beyond-standard-model frameworks. For the standard model with four generations, the CMS Collaboration has ruled out the Higgs boson mass in the range of ~120–600 GeV at 95% C.L [23].

Some considerations on the implications of a fourth generation that may evade this bound have appeared recently in Ref. [22]. As pointed out in Ref. [22], it is possible to reduce the tension between Higgs searches at the LHC and a heavy fourth generation by extending the scalar sector of the standard model. The LHC exclusion range 120 GeV $< m_H <$ 600 GeV [23] corresponds to the process $gg \rightarrow H \rightarrow VV$, where V can be either $W^\pm$ or Z bosons. The authors of Ref. [22] considered two Higgs doublets H_1, H_2, with the first one coupling to the first three generations and the second one coupling to the fourth generation only. Choosing the Higgs mixing angle α to be same as $\beta = \tan^{-1} v_2/v_1$ makes the coupling of one of the Higgs mass eigenstates to the vector boson zero and, hence, can be light without conflicting with the LHC exclusion range coming from the $gg \rightarrow H \rightarrow VV$ process. The other Higgs, which couples to vector bosons, can be as heavy as the generic unitarity bound for two Higgs doublet models ~700 GeV [28–31]; such a Higgs can still be outside the LHC search range.

Our model has a different motivation and differs from Ref. [22] in that the extra $U(1)_X$ gauge sector prevents bilinear mixings of the Higgs doublets introduced. The dimensionless coupling of the quartic mixing terms can be chosen to be small enough so that the lightest neutral Higgs mass eigenstate has negligible coupling to the fourth-generation fermions. Thus, although both of them couple to vector bosons, the lighter Higgs have no coupling to the fourth-generation fermions and, hence, can evade the LHC exclusion range coming from the $gg \rightarrow H \rightarrow VV$ process. The heavier Higgs couple to the fourth generation only and can be as heavy as the unitarity bound to stay outside the current LHC search range. It may be noted that in Ref. [22], there is no physical mechanism to prevent the extra Higgs doublet VEV from feeding into the lighter fermion masses, whereas our model provides an explanation for this by the nonuniversal couplings of the extra $U(1)_X$ gauge boson to the fermions, which allows only H_1-first-three-generation and H_2-fourth-generation couplings.

IV. FOURTH GENERATION DIRAC NEUTRINO AS DARK MATTER

The relic abundance of a dark matter particle χ is given by the Boltzmann equation

$$\frac{dn_\chi}{dt} + 3Hn_\chi = -\langle\sigma v\rangle[n_\chi^2 - (n_\chi^{eqb})^2], \tag{5}$$

where n_χ is the number density of the dark matter particle χ and n_χ^{eqb} is the number density when χ was in thermal equilibrium. H is the Hubble rate, and $\langle\sigma v\rangle$ is the thermally averaged annihilation cross section of the dark matter particle χ. In terms of partial wave expansion, $\langle\sigma v\rangle = a + bv^2$. The numerical solution of the Boltzmann equation above gives [32]

$$\Omega_\chi h^2 \approx \frac{1.04 \times 10^9 x_F}{M_{\rm Pl}\sqrt{g_*}(a + 3b/x_F)}, \tag{6}$$

where $x_F = m_\chi/T_F$, T_F is the freeze-out temperature and g_* is the number of relativistic degrees of freedom at the time of freeze-out. Dark matter particles with electroweak scale mass and couplings freeze out at temperatures approximately in the range $x_F \approx 20$–30. This again simplifies to [33]

$$\Omega_\chi h^2 \approx \frac{3 \times 10^{-27}\ \mathrm{cm}^3\ \mathrm{s}^{-1}}{\langle\sigma v\rangle}. \tag{7}$$

The thermal averaged annihilation cross section $\langle\sigma v\rangle$ is given by [34]

$$\langle\sigma v\rangle = \frac{1}{8m^4 T K_2^2(m/T)} \int_{4m^2}^{\infty} \sigma(s - 4m^2)\sqrt{s} K_1(\sqrt{s}/T) ds, \tag{8}$$

where K_i's are modified Bessel functions of order i, m is the mass of dark matter particle, and T is the temperature. We consider two annihilation cross sections of cosmological importance, namely $\sigma(\nu_4\bar{\nu}_4 \rightarrow f\bar{f})$ and $\sigma(\nu_4\bar{\nu}_4 \rightarrow W^+W^-)$, where f is any standard model fermion. The first process can take place via s-channel exchange of Z and X bosons. The second process can take place via s-channel Z or Higgs boson H exchange or t-channel exchange of charged leptons. Since there is no mixing of the fourth generation with the first three generations, this charged lepton is the fourth charged lepton τ'. In our model, we always consider $m_{\tau'} > m_{\nu 4}$ so as to make the fourth-generation neutrino perfectly stable. Thus, only the s-channel annihilation processes are of interest. The cross section for the $f\bar{f}$ final state through Z-boson exchange and the W^+W^- final has been calculated in Ref. [35], and we use their standard result. The s-channel X boson exchange cross section is

$$\begin{aligned}\sigma_X(\nu_4\bar{\nu}_4 \rightarrow f\bar{f}) &= \frac{N_c g_x^4}{32\pi s}\frac{\beta_f}{\beta_{\nu 4}}|D_x|^2\bigg[(v_f^2 + a_f^2)(v_{\nu 4}^2 + a_{\nu 4}^2)\frac{s^2}{4}\left(1 + \frac{\beta^2}{3}\right) \\ &\quad + (v_f^2 - a_f^2)(v_{\nu 4}^2 + a_{\nu 4}^2)m_f^2(s - 2m_{\nu 4}^2) \\ &\quad + (v_f^2 + a_f^2)(v_{\nu 4}^2 - a_{\nu 4}^2)m_{\nu 4}^2(s - 2m_f^2) \\ &\quad + 4(v_f^2 - a_f^2)(v_{\nu 4}^2 - a_{\nu 4}^2)m_f^2 m_{\nu 4}^2\bigg],\end{aligned} \tag{9}$$

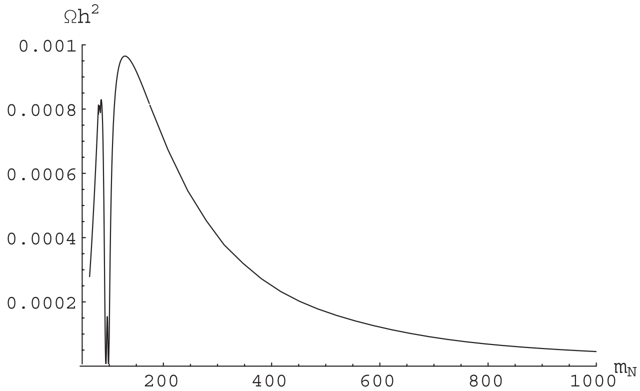

FIG. 1. Relic abundance of a fourth-generation Dirac neutrino as a function of its mass.

where v,a are the vector and axial couplings of fermions to the X boson, respectively; N_c is the color factor which is 3 for quarks and 1 for leptons, $|D_x|^2 = 1/[(s - M_x^2)^2 + \Gamma_x^2 M_x^2]$; and β's are defined as

$$\beta_f = \left(1 - \frac{4m_f}{s}\right)^{1/2}, \quad \beta_{\nu 4} = \left(1 - \frac{4m_{\nu 4}}{s}\right)^{1/2}, \quad \beta = \beta_f \beta_{\nu 4}.$$

The relic density as a function of the fourth-generation neutrino mass m_N is shown in Fig. 1. In this mass range $M_Z/2 < m_N < M_W$, the only possible annihilation channels are the $N\bar{N} \to f\bar{f}$ through s-channel exchange of Z boson or X boson (depending on the mass of M_X). In this particular example, the $U(1)_X$ gauge charges are chosen as $n_2 = -1$, $n_{10} = \frac{1}{2}$. Also, we have taken the extra $U(1)_X$ coupling to be 10^{-2} and the $U(1)_X$ symmetry breaking scale to be 5 TeV, which results in $M_X = 142$ GeV. Thus, the X-boson mediating channel opens only when $m_N \geq 71$ GeV. As we go beyond this mass range, more and more annihilation channels become important and, hence, reduce the relic abundance further. It is seen from Fig. 1 that the fourth-generation Dirac neutrino can at most give rise to 1% of the total dark matter abundance estimated by Seven-Year Wilkinson Microwave Anisotropy Probe (WMAP) observations [36]:

$$\Omega_{\rm DM} h^2 = 0.1123 \pm 0.0035. \quad (10)$$

Thus, the fourth-generation Dirac neutrino if stable or long-lived, like in the minimal version of our model, can give rise to a very small fraction of total dark matter in the Universe and, hence, be of little cosmological significance. However, it can have distinct collider signatures. Being heavy and stable, it can give rise to a large missing transverse energy in the colliders.

V. CONCLUSION

We have studied one possible framework which gives rise to the experimentally allowed neutrino mass spectra in four-generation models. The standard model gauge group is enhanced to include one extra Abelian gauge symmetry $U(1)_X$, under which the fourth-generation fermions transform differently from the first three generations. We have assumed zero mixing between the fourth and first three generations and showed that there can be a common seesaw mechanism, which can generate three light standard model neutrinos and one heavy stable fourth-generation Dirac neutrino. We also point out that our model reproduces the Higgs-fermion structure considered by the authors of Ref. [22] with certain differences. Gauge structure of our model naturally prevents bilinear mixing between the two Higgs doublets and provides a physical mechanism whereby one doublet couples only to the first three generations and the other couples only to the fourth generation. Thus, the Higgs which couple to the fourth generation only can be as heavy as the unitarity bound so as to evade the current LHC exclusion range for the generic standard model with four generations. The lighter Higgs do not couple to fourth generation at tree level and, hence, can also evade the LHC exclusion range.

We also consider the possibility of such a heavy neutrino as dark matter and found out the relic density. However, we find that for the mass range $M_Z/2 < m_{\nu 4} < 1$ TeV, a fourth-generation Dirac neutrino can at most give rise to 1% of the total dark matter in the Universe and, hence, can play a role in multicomponent dark matter formalisms. It can also have important collider signatures in terms of missing transverse energies, the details of which we have skipped in our present work.

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