Realistic $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ model with a type II Dirac neutrino seesaw mechanism

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Here we propose a realistic $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ electroweak gauge model with enlarged Higgs sector. The scheme allows for the natural implementation of a type II seesaw mechanism for Dirac neutrinos, while charged lepton and quark masses are reproduced in a natural way thanks to the presence of new scalars. The new $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ energy scale characterizing neutrino mass generation could be accessible to the current LHC experiments.

I. INTRODUCTION

Despite the fierce experimental effort over the last decades the long-standing challenge concerning the question of whether neutrinos are their own antiparticles remains [1–3]. Although neutrinos could be Dirac or Majorana fermions, the leading theoretical expectation is that they are Majorana, the general belief being that the smallness of neutrino masses relative to the other standard model fermion masses is due to their charge neutrality. This fits naturally to the idea that neutrinos acquire Majorana masses from Weinberg's dimension five operator. Realizations include various varieties of type I [4–8] or type II [7–10] seesaw mechanisms, irrespective of whether the seesaw is realized at high or at low mass scale [11]. Until the observation of neutrinoless double beta decay [12] becomes a reality [1] we must keep an eye open to the possibility that neutrinos can be Dirac particles as well.

There are two issues that Dirac neutrino models must face, namely predicting the Dirac nature of neutrinos, and understanding their small mass. Dirac neutrinos require extra symmetries beyond $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ gauge symmetry; otherwise massive neutrinos are generally expected to be Majorana fermions. Within the standard $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ electroweak gauge structure this can be ensured by imposing a conserved lepton numberlike symmetry, continuous [13] or discrete [14]. Likewise, one may consider schemes based on flavor symmetries, as suggested in [15] or appealing to the existence of extra dimensions [16,17]. Alternatively one may extend the gauge group, so as to include the lepton number symmetry [18], for example, by using the extended $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ gauge structure [19] although most formulations lead to Majorana neutrinos [20–31].

Because of the special features of the $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ based models with respect to other electroweak extensions based, for example, on left-right symmetry, in this paper we focus on the possibility of having naturally light Dirac neutrinos with seesaw-induced masses, within the $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ gauge structure. However, in contrast to Ref. [32] in order to predict realistic quark masses in a natural way, two extra antitriplet scalar multiplets are included. As before, we have a lepton numberlike symmetry $\mathcal{L}$, preserved in the leptonic and quark sector. This symmetry is softly broken in the scalar sector by the term $f\phi_0\phi_1\phi_2$ so that in the limit $f \to 0$ the Lagrangian symmetry gets enhanced. We show that the smallness of f is related to the smallness of neutrino mass. This way we recover the new variant of the type II seesaw mechanism for Dirac neutrinos recently considered in [32]. However, in contrast to the previous Ref. [32], here the naturally small induced vacuum expectation values (vevs) responsible for neutrino mass generation are decoupled from the quark sector. This eliminates the need of fine-tuning the Yukawa couplings so that all fermion masses are naturally reproduced in a realistic way, since exotic quarks and standard model quarks are now decoupled. The scales associated to neutrino mass generation and the new $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ gauge interactions could be accessible to the current LHC experiments.

II. MODEL

The model is a modified version of the one presented in Ref. [32]. In the new setup, two extra scalar triplets are included and an auxiliary $\mathbb{Z}_4$ symmetry is implemented in order to decouple heavy quarks. The matter content and the transformation properties of the fields are contained in Table I.

In terms of the gauge group generators, the electric charge is expressed as

$$Q = T_3 + \frac{1}{\sqrt{3}}T_8 + X, \qquad (1)$$

while lepton number is defined as

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TABLE I. Matter content of the model, where $\ell = (e, \mu, \tau)$, $\hat{u} \equiv (u, c, t)$ and $\hat{d} \equiv (d, s, b)$.

	ψ_L^ℓ	ℓ_R	S_R^ℓ	$\tilde{S}_R^\ell$	$Q_L^{1,2}$	Q_L^3	$\hat{u}_R$	U_R	$\hat{d}_R$	$D_R^{1,2}$	ϕ_0	ϕ_1	ϕ_2	ϕ_3	ϕ_4
$SU(3)_c$	1	1	1	1	3	3	3	3	3	3	1	1	1	1	1
$SU(3)_L$	3^*	1	1	1	3	3^*	1	1	1	1	3^*	3^*	3^*	3^*	3^*
$U(1)_X$	$-\frac{1}{3}$	-1	0	0	0	$+\frac{1}{3}$	$+\frac{2}{3}$	$+\frac{2}{3}$	$-\frac{1}{3}$	$-\frac{1}{3}$	$+\frac{2}{3}$	$-\frac{1}{3}$	$-\frac{1}{3}$	$-\frac{1}{3}$	$-\frac{1}{3}$
$\mathcal{L}$	$-\frac{1}{3}$	-1	$+1$	$+1$	$-\frac{2}{3}$	$+\frac{2}{3}$	0	1	0	-1	$+\frac{2}{3}$	$-\frac{4}{3}$	$-\frac{4}{3}$	$-\frac{1}{3}$	$+\frac{2}{3}$
$\mathbb{Z}_3$	ω	ω	ω	ω	ω^2	ω^2	ω^2	ω^2	ω^2	ω^2	1	1	1	1	1
$\mathbb{Z}_4$	1	1	$-i$	i	1	1	1	-1	1	-1	1	i	$-i$	-1	1

$$L = \frac{4}{\sqrt{3}} T_8 + \mathcal{L}. \tag{2}$$

The symmetry breaking pattern is assumed to be of the form

$$\langle \phi_0 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} k_0 \\ 0 \\ 0 \end{pmatrix},$$

$$\langle \phi_1 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \delta_1 \\ n_1 \end{pmatrix},$$

$$\langle \phi_2 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \delta_2 \\ n_2 \end{pmatrix}, \tag{3}$$

$$\langle \phi_3 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ 0 \\ n_3 \end{pmatrix}, \qquad \langle \phi_4 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ k_4 \\ 0 \end{pmatrix}, \tag{4}$$

with $n_{1,2,3} \gg k_{0,4}$, $\delta_{1,2}$ and $\delta_1 \sim \delta_2$.

The Yukawa interactions invariant under all the defining symmetries of the model are

$$\begin{aligned} -\mathcal{L}_f = {}& y^\ell \bar{\psi}_L \ell_R \phi_0 + y_1 \bar{\psi}_L S_R \phi_1 + \tilde{y}_2 \bar{\psi}_L \tilde{S}_R \phi_2 \\ & + y^u \bar{Q}_L^{1,2} \hat{u}_R \phi_0^* + \hat{y}^u \bar{Q}_L^3 \hat{u}_R \phi_4 + y^U \bar{Q}_L^3 U_R \phi_3 \\ & + y^d \bar{Q}_L^3 \hat{d}_R \phi_0 + \hat{y}^d \bar{Q}_L^{1,2} \hat{d}_R \phi_4^* + y^D \bar{Q}_L^{1,2} D_R^{1,2} \phi_3^* + \text{H.c.} \end{aligned} \tag{5}$$

In this setup, ϕ_1 and ϕ_2 are completely decoupled from the quark sector and are responsible for the neutrino mass generation, whereas ϕ_3 and ϕ_4 contribute exclusively to the quark masses. After spontaneous symmetry breaking, the above interactions lead to the following mass matrices for the fermion fields:

(i) charged leptons: $m_\ell = \frac{k_0}{\sqrt{2}} y^\ell$,

(ii) neutrinos:

$$\begin{aligned} m_\nu &= \frac{1}{\sqrt{2}} \begin{pmatrix} \delta_1 y_1 & \delta_2 \tilde{y}_2 \\ n_1 y_1 & n_2 \tilde{y}_2 \end{pmatrix} \\ &= \frac{1}{\sqrt{2}} \begin{pmatrix} \delta_1 \mathbb{1} & \delta_2 \mathbb{1} \\ n_1 \mathbb{1} & n_2 \mathbb{1} \end{pmatrix} \begin{pmatrix} y_1 & 0 \\ 0 & \tilde{y}_2 \end{pmatrix}, \end{aligned} \tag{6}$$

(iii) up-type quarks, basis (u, c, t, U):

$$m_u = \frac{1}{\sqrt{2}} \begin{pmatrix} k_0 y_1^u & k_0 y_1^c & k_0 y_1^t & 0 \\ k_0 y_2^u & k_0 y_2^c & k_0 y_2^t & 0 \\ k_4 \hat{y}^u & k_4 \hat{y}^c & k_4 \hat{y}^t & 0 \\ 0 & 0 & 0 & n_3 y^U \end{pmatrix}, \tag{7}$$

(iv) down-type quarks, basis (d, s, b, D^1, D^2):

$$m_d = \frac{1}{\sqrt{2}} \begin{pmatrix} k_4 \hat{y}_1^d & k_4 \hat{y}_1^s & k_4 \hat{y}_1^b & 0 & 0 \\ k_4 \hat{y}_2^d & k_4 \hat{y}_2^s & k_4 \hat{y}_2^b & 0 & 0 \\ k_0 y^d & k_0 y^s & k_0 y^b & 0 & 0 \\ 0 & 0 & 0 & n_3 y_1^{D^1} & n_3 y_1^{D^2} \\ 0 & 0 & 0 & n_3 y_2^{D^1} & n_3 y_2^{D^2} \end{pmatrix}. \tag{8}$$

Realistic quark masses can be readily obtained as the standard model and exotic sectors are independent by virtue of the $\mathbb{Z}_4$ symmetry. This also implies that the Cabibbo-Kobayashi-Maskawa matrix describing quark mixing is unitary.

Moreover, from the diagonalization of $m_\nu m_\nu^\dagger$, the mass of the light neutrino in the one-family approximation is given by

$$m_{\text{light}} \approx \frac{1}{\sqrt{2}} \frac{|y_1 \tilde{y}_2 (\delta_2 n_1 - \delta_1 n_2)|}{\sqrt{n_1^2 y_1^2 + n_2^2 \tilde{y}_2^2}}. \tag{9}$$

Notice that the Dirac nature of neutrinos is ensured by the discrete $\mathbb{Z}_3$ group. In what follows, we specialize to the

case $\delta_1 = \delta_2 \equiv \delta$ and show that the smallness of the neutrino mass can be understood as emerging from a type II seesaw mechanism for Dirac neutrinos. This links the emergence of the small induced scale δ to the breaking of the global symmetry $\mathcal{L}$ [32].

The scalar potential of the model is

$$V = \sum_{i=0}^{4}(\mu_i^2|\phi_i|^2 + \lambda_i|\phi_i|^4) + \sum_{i<j}[\lambda_{ij}|\phi_i|^2|\phi_j|^2$$
$$+ \tilde{\lambda}_{ij}(\phi_i^\dagger\phi_j)(\phi_j^\dagger\phi_i)] + [\lambda(\phi_1^\dagger\phi_2)(\phi_1^\dagger\phi_2)$$
$$+ f\phi_0\phi_1\phi_2 + \text{H.c.}], \tag{10}$$

where the term $f\phi_0\phi_1\phi_2$ breaks explicitly the $\mathcal{L}$ symmetry. Assuming real vevs and parameters in the potential, the tadpole equations can be solved for the parameters μ_i^2, $\tilde{\lambda}_{13}$, $\tilde{\lambda}_{14}$ and f as follows:

$$\mu_0^2 = -\frac{1}{2}[2\lambda_0 k_0^2 + \lambda_{01}(\delta^2 + n_1^2) + \lambda_{02}(\delta^2 + n_2^2)$$
$$+ n_3^2\lambda_{03} + k_4^2\lambda_{04} - \frac{\delta^2(n_1 - n_2)}{2k_0^2}\left[(n_1 - n_2)(\tilde{\lambda}_{12} + 2\lambda)\right.$$
$$\left.+ \frac{n_2(n_3^2\tilde{\lambda}_{23} - k_4^2\tilde{\lambda}_{24})}{(\delta^2 + n_1 n_2)}\right], \tag{11}$$

$$\mu_1^2 = -\frac{1}{2}[k_0^2\lambda_{01} + n_3^2\lambda_{13} + k_4^2\lambda_{14} + 2(\delta^2 + n_1^2)\lambda_1$$
$$+ (\delta^2 + n_2^2)(\lambda_{12} + \tilde{\lambda}_{12} + 2\lambda)]$$
$$+ \frac{n_2(\delta^2 k_4^2\tilde{\lambda}_{24} + n_1 n_2 n_3^2\tilde{\lambda}_{23})}{2n_1(\delta^2 + n_1 n_2)}, \tag{12}$$

$$\mu_2^2 = -\frac{1}{2}[k_0^2\lambda_{02} + k_4^2\lambda_{24} + 2(\delta^2 + n_2^2)\lambda_2 + \lambda_{23}n_3^2$$
$$+ (2\lambda + \lambda_{12} + \tilde{\lambda}_{12})(\delta^2 + n_1^2)]$$
$$- \frac{\delta^2 k_4^2\tilde{\lambda}_{24} + n_1 n_2 n_3^2\tilde{\lambda}_{23}}{2(\delta^2 + n_1 n_2)}, \tag{13}$$

$$\mu_3^2 = -\frac{1}{2}[k_0^2\lambda_{03} + k_4^2\lambda_{34} + \lambda_{13}(\delta^2 + n_1^2) + \lambda_{23}(\delta^2 + n_2^2)$$
$$+ 2\lambda_3 n_3^2 + n_2(n_2 - n_1)\tilde{\lambda}_{23}], \tag{14}$$

$$\mu_4^2 = -\frac{1}{2}\left[k_0^2\lambda_{0,4} + 2k_4^2\lambda_4 + \lambda_{14}(\delta^2 + n_1^2) + \lambda_{24}(\delta^2 + n_2^2)\right.$$
$$\left.+ \lambda_{34}n_3^2 + \frac{\delta^2(n_1 - n_2)\tilde{\lambda}_{24}}{n_1}\right], \tag{15}$$

$$\tilde{\lambda}_{13} = -\frac{n_2\tilde{\lambda}_{23}}{n_1}, \tag{16}$$

$$\tilde{\lambda}_{14} = -\frac{n_2\tilde{\lambda}_{24}}{n_1}, \tag{17}$$

$$f = \frac{\delta}{\sqrt{2}k_0}\left[(n_2 - n_1)(2\lambda + \tilde{\lambda}_{12}) + \frac{n_2(k_4^2\tilde{\lambda}_{24} - n_3^2\tilde{\lambda}_{23})}{\delta^2 + n_1 n_2}\right]. \tag{18}$$

Assuming that the bracketed factor in the rhs of Eq. (18) is nonvanishing, δ can be interpreted as an induced vev, whose smallness is related to the scale of $\mathcal{L}$ symmetry breaking, characterized by f, in the sense that in the limit $\delta \to 0$, the symmetry of the potential is enhanced by $f \to 0$.

We conclude this section with an estimate of the scales involved and the resulting light neutrino mass,

$$m_{\text{light}} \approx \frac{1}{\sqrt{2}}\frac{|y_1\tilde{y}_2\delta(n_1 - n_2)|}{\sqrt{n_1^2 y_1^2 + n_2^2\tilde{y}_2^2}}. \tag{19}$$

Taking $f \sim \mathcal{O}(1)$ keV, $k_0 \sim k_4 \sim \mathcal{O}(10^2)$ GeV, $n_1 \sim n_2 \sim n_3 \sim \mathcal{O}(1)$ TeV and quartic couplings of $\mathcal{O}(1)$ in Eq. (18), the resulting scale $\delta \sim \mathcal{O}(10)$ eV can accommodate easily a neutrino mass of $\mathcal{O}(10^{-1})$ eV without invoking tiny Yukawa couplings or large values for the vacuum expectation values n_1, n_2. Finally, note that the light neutrino mass in Eq. (19) is mostly insensitive to the values of n_1 and n_2, and therefore new physics in this model can lie within reach of the LHC experiments.

III. SUMMARY AND DISCUSSION

Here we have proposed a realistic SU(3)$_c$ ⊗ SU(3)$_L$ ⊗ U(1)$_X$ electroweak gauge model with an enlarged Higgs sector. The scheme leads to Dirac neutrinos, and allows for the natural implementation of a type II seesaw mechanism. The model substantially improves the one previously considered in Ref. [32], in that here the small vacuum expectation values associated to neutrino mass generation are decoupled from the quark sector. The charged lepton and quark masses are reproduced in a realistic way, avoiding the mixing between exotic and standard model quarks. The energy scales characterizing neutrino mass generation and the new gauge interactions arising from the enlarged SU(3)$_c$ ⊗ SU(3)$_L$ ⊗ U(1)$_X$ gauge symmetry could be accessible to the current LHC experiments.

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[1] A. S. Barabash, DAE. Symp. Nucl. Phys. **55,** I13 (2010).

[2] F. T. Avignone, III, S. R. Elliott, and J. Engel, Rev. Mod. Phys. **80,** 481 (2008).

[3] S. Blot (NEMO-3, SuperNEMO), J. Phys. Conf. Ser. **718,** 062006 (2016).

[4] M. Gell-Mann, P. Ramond, and R. Slansky, CERN Report No. 80-0576, 1979.

[5] T. Yanagida, *KEK Lectures, 1979*, edited by O. Sawada and A. Sugamoto (KEK, Ibaraki, Japan 1979).

[6] R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. **44,** 912 (1980).

[7] J. Schechter and J. W. F. Valle, Phys. Rev. D **22,** 2227 (1980).

[8] J. Schechter and J. W. F. Valle, Phys. Rev. D **25,** 774 (1982).

[9] G. Lazarides, Q. Shafi, and C. Wetterich, Nucl. Phys. **B181,** 287 (1981).

[10] W. Grimus, L. Lavoura, and B. Radovcic, Phys. Lett. B **674,** 117 (2009).

[11] S. M. Boucenna, S. Morisi, and J. W. F. Valle, Adv. High Energy Phys. **2014,** 831598 (2014).

[12] H. V. Klapdor-Kleingrothaus, I. V. Krivosheina, A. Dietz, and O. Chkvorets, Phys. Lett. B **586,** 198 (2004).

[13] J. T. Peltoniemi, D. Tommasini, and J. W. F. Valle, Phys. Lett. B **298,** 383 (1993).

[14] S. C. Chuliá, E. Ma, R. Srivastava, and J. W. F. Valle, arXiv:1606.04543.

[15] A. Aranda, C. Bonilla, S. Morisi, E. Peinado, and J. W. F. Valle, Phys. Rev. D **89,** 033001 (2014).

[16] N. Arkani-Hamed, S. Dimopoulos, G. R. Dvali, and J. March-Russell, Phys. Rev. D **65,** 024032 (2001).

[17] P. Chen, G.-J. Ding, A. D. Rojas, C. A. Vaquera-Araujo, and J. W. F. Valle, J. High Energy Phys. 01 (2016) 007.

[18] E. Ma, N. Pollard, R. Srivastava, and M. Zakeri, Phys. Lett. B **750,** 135 (2015).

[19] M. Singer, J. W. F. Valle, and J. Schechter, Phys. Rev. D **22,** 738 (1980).

[20] M. B. Tully and G. C. Joshi, Phys. Rev. D **64,** 011301 (2001).

[21] T. Kitabayashi and M. Yasue, Phys. Rev. D **63,** 095002 (2001).

[22] T. Kitabayashi and M. Yasue, Phys. Rev. D **63,** 095006 (2001).

[23] J. C. Montero, C. A. De S. Pires, and V. Pleitez, Phys. Rev. D **65,** 095001 (2002).

[24] D. Chang and H. N. Long, Phys. Rev. D **73,** 053006 (2006).

[25] P. V. Dong, H. N. Long, and D. V. Soa, Phys. Rev. D **75,** 073006 (2007).

[26] J. K. Mizukoshi, C. A. de S. Pires, F. S. Queiroz, and P. S. Rodrigues da Silva, Phys. Rev. D **83,** 065024 (2011).

[27] S. M. Boucenna, S. Morisi, and J. W. F. Valle, Phys. Rev. D **90,** 013005 (2014).

[28] S. M. Boucenna, R. M. Fonseca, F. Gonzalez-Canales, and J. W. F. Valle, Phys. Rev. D **91,** 031702 (2015).

[29] H. Okada, N. Okada, and Y. Orikasa, Phys. Rev. D **93,** 073006 (2016).

[30] S. M. Boucenna, J. W. F. Valle, and A. Vicente, Phys. Rev. D **92,** 053001 (2015).

[31] V. V. Vien, A. E. C. Hernández, and H. N. Long, arXiv:1601.03300.

[32] J. W. F. Valle and C. A. Vaquera-Araujo, Phys. Lett. B **755,** 363 (2016).