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CP violation from flavor symmetry in a lepton quarticity dark matter model

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ABSTRACT

We propose a simple $\Delta(27) \otimes Z_4$ model where neutrinos are predicted to be Dirac fermions. The smallness of their masses follows from a type-I seesaw mechanism and the leptonic CP violating phase correlates with the pattern of $\Delta(27)$ flavor symmetry breaking. The scheme naturally harbors a WIMP dark matter candidate associated to the Dirac nature of neutrinos, in that the same Z_4 lepton number symmetry also ensures dark matter stability.

1. Introduction

Currently the information on neutrino properties comes mainly from neutrino oscillation experiments [1]. These are insensitive to whether neutrinos are Dirac or Majorana fermions [2–4]. The fact that the weak interaction is V-A turns the quest for lepton number violation and the Majorana nature of neutrinos into a major experimental challenge [5–7]. The detection of neutrinoless double beta decay would signify a major step forward in this endeavor. According to the black-box theorem [8,9] its observation would demonstrate that neutrinos are Majorana fermions and thus lepton number is violated in nature.

Concerning the mechanism responsible for generating small neutrino masses, little is known regarding the nature of its associated messenger particles, their characteristic mass scale or other detailed features of the effective operator [10]. The smallness of neutrino masses almost always assumes them to be Majorana fermions. For example, this is the case in the conventional high-scale (type-I) [2,11–13] seesaw mechanism. Likewise in low-scale variants of the seesaw mechanism as well as in radiative schemes, neutrinos turn out to be Majorana fermions, as reviewed in [14].

Having naturally light Dirac neutrinos requires extra assumptions beyond the standard $\mathrm{SU(3)_c \otimes SU(2)_L \otimes U(1)_Y}$ electroweak gauge invariance. One possibility is to extend the electroweak gauge group, for example, by using the $\mathrm{SU(3)_c \otimes SU(3)_L \otimes U(1)_X}$ group to exploit its peculiar features [16]. In this framework it has recently been shown that one can obtain a type-II seesaw mechanism for Dirac neutrinos [17,18]. One can also use a $B-L$ gauge extension with unconventional charges for right handed neutrinos, which leads to Dirac neutrinos obtained from type-I seesaw mechanism [19–21]. Alternatively one may stick to the simplest $\mathrm{SU(3)_c \otimes SU(2)_L \otimes U(1)_Y}$ gauge structure but use extra symmetries implying a conserved lepton number, so as to obtain Dirac neutrinos, as suggested in [22]. One can also adopt extra-dimensional theory frameworks [15].

Here we focus on the possibility of requiring that neutrino masses arise from a simple type-I seesaw mechanism within a flavor-symmetric scenario. Moreover we will also require that the existence of a viable dark matter particle arises from the same symmetry which ensures that neutrinos do not acquire Majorana masses and remain Dirac fermions. In Sect. 2 we sketch in some detail the extended particle content required to realize the non-Abelian flavor symmetry of the model and demonstrate how the Dirac nature of neutrinos and the smallness of their seesaw-induced masses follow from our non-Abelian discrete flavor symmetry. In Sect. 3 we present numerical predictions for CP violation in terms of the scalar boson alignment patterns. Towards the end of the paper, in Sect. 4, we discuss the appearance of viable dark matter in this model and give a brief discussion of its direct detection potential. Finally we summarize our results in Sect. 5.

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Table 1
The $\Delta(27)$ and Z_4 charge assignments for leptons, the Higgs scalars (Φ_i, χ_i) and the dark matter sector scalars (ζ and η). Here $\mathbf{z}$ is the fourth root of unity, i.e. $\mathbf{z}^4 = 1$.

Fields	$\Delta(27)$	Z_4	Fields	$\Delta(27)$	Z_4
$\bar{L}_e$	$\mathbf{1}$	$\mathbf{z}^3$	$\nu_{e,R}$	$\mathbf{1}$	$\mathbf{z}$
$\bar{L}_\mu$	$\mathbf{1}''$	$\mathbf{z}^3$	$\nu_{\mu,R}$	$\mathbf{1}'$	$\mathbf{z}$
$\bar{L}_\tau$	$\mathbf{1}'$	$\mathbf{z}^3$	$\nu_{\tau,R}$	$\mathbf{1}''$	$\mathbf{z}$
$l_{i,R}$	$\mathbf{3}$	$\mathbf{z}$	$\bar{N}_{i,L}$	$\mathbf{3}$	$\mathbf{z}^3$
$N_{i,R}$	$\mathbf{3}'$	$\mathbf{z}$			
Φ_i	$\mathbf{3}'$	$\mathbf{1}$	χ_i	$\mathbf{3}'$	$\mathbf{1}$
ζ	$\mathbf{1}$	$\mathbf{z}$	η	$\mathbf{1}$	$\mathbf{z}^2$

2. The model

Our model is based on the discrete flavor symmetry $\Delta(27) \otimes Z_4$ where Z_4 is the cyclic group of order four and $\Delta(27)$ is a discrete non-Abelian symmetry group isomorphic to $(Z_3 \otimes Z_3) \ltimes Z_3$. Before presenting the details of the model, we briefly discuss the most relevant features of the $\Delta(27)$ group. The $\Delta(27)$ group belongs to the general class of discrete groups denoted by $\Delta(3N^2)$, with N being a positive integer. The smallest member of $\Delta(3N^2)$ is $\Delta(3)$ which is nothing but the Abelian group Z_3. The next member is $\Delta(12)$ which is isomorphic to the well-known group A_4. The third smallest member of the group is $\Delta(27)$ and has 27 elements divided into 11 conjugacy classes [22–25]. It has nine singlet irreducible representations $\mathbf{1}_i$; $i = 1, \ldots 9$ and two triplet irreducible representations $\mathbf{3}$ and $\mathbf{3}'$.[1] The multiplication rules for $\Delta(27)$ are given by

$$\mathbf{3} \otimes \mathbf{3} = \mathbf{3}' \oplus \mathbf{3}' \oplus \mathbf{3}'\,; \qquad \mathbf{3} \otimes \mathbf{3}' = \sum_{i=1}^{9} \mathbf{1}_i\,. \tag{1}$$

The particle content of our model along with the $\Delta(27) \otimes Z_4$ charge assignments of the particles are as shown in Table 1.

In Table 1 $L_i = (\nu_i, l_i)^T$; $i = e, \mu, \tau$ are the lepton doublets which also transform as singlets under $\Delta(27)$ and have charge $\mathbf{z}$ under Z_4. The $l_{i,R}$; $i = e, \mu, \tau$ are the charged lepton singlets which transform as a $\mathbf{3}$ under $\Delta(27)$ and have Z_4-charge $\mathbf{z}$. Apart from the Standard Model fermions, the model also includes three right-handed neutrinos $\nu_{i,R}$ transforming as singlets under the $\mathrm{SU(3)_c \otimes SU(2)_L \otimes U(1)_Y}$ gauge group and as singlets under $\Delta(27)$, with charge $\mathbf{z}$ under Z_4. We also add three gauge singlet Dirac fermions $N_{i,L}$, $N_{i,R}$; $i = 1, 2, 3$ transforming as triplets of $\Delta(27)$ and with charge $\mathbf{z}$ under Z_4, as shown in Table 1.

In the scalar sector the $\Phi_i = (\phi_i^+, \phi_i^0)^T$; $i = 1, 2, 3$ transform as $\mathrm{SU(2)_L}$ doublets, as triplet under $\Delta(27)$ and trivially under Z_4. On the other hand the scalars χ_i; $i = 1, 2, 3$ are gauge singlets transforming as a triplet under $\Delta(27)$ and trivially under Z_4. We also add two other gauge singlet scalars ζ and η both of which transform trivially under $\Delta(27)$ but carry Z_4 charges $\mathbf{z}$ and $\mathbf{z}^2$ respectively. Since $\mathbf{z}^2 = -1$, the field η can be real and is taken to be real. We comment on the important role of the scalars ζ and η towards the end of the paper. The $\mathrm{SU(3)_c \otimes SU(2)_L \otimes U(1)_Y} \otimes \Delta(27) \otimes Z_4$ invariant Yukawa term for the charged leptons is given by

$$\begin{aligned}
\mathcal{L}_{\mathrm{Yuk,l}} = {} & y_1\, (\bar{L}_e)_1 \otimes \left[\begin{pmatrix} l_{e,R} \\ l_{\mu,R} \\ l_{\tau,R} \end{pmatrix}_3 \otimes \begin{pmatrix} \Phi_1 \\ \Phi_2 \\ \Phi_3 \end{pmatrix}_{3'} \right]_1 \\
& + y_2\, (\bar{L}_\mu)_{1''} \otimes \left[\begin{pmatrix} l_{e,R} \\ l_{\mu,R} \\ l_{\tau,R} \end{pmatrix}_3 \otimes \begin{pmatrix} \Phi_1 \\ \Phi_2 \\ \Phi_3 \end{pmatrix}_{3'} \right]_{1'} \\
& + y_3\, (\bar{L}_\tau)_{1'} \otimes \left[\begin{pmatrix} l_{e,R} \\ l_{\mu,R} \\ l_{\tau,R} \end{pmatrix}_3 \otimes \begin{pmatrix} \Phi_1 \\ \Phi_2 \\ \Phi_3 \end{pmatrix}_{3'} \right]_{1''}
\end{aligned} \tag{2}$$

where y_i, $i = 1, 2, 3$, are the Yukawa couplings which, for simplicity, we take them to be real. After symmetry breaking the scalars acquire vacuum expectation values (vevs) $\langle \Phi_i \rangle = v_i$; $i = 1, 2, 3$ so the charged lepton mass matrix is given by

$$M_l = \begin{pmatrix} y_1 v_1 & y_1 v_2 & y_1 v_3 \\ y_2 v_1 & \omega y_2 v_2 & \omega^2 y_2 v_3 \\ y_3 v_1 & \omega^2 y_3 v_2 & \omega y_3 v_3 \end{pmatrix}. \tag{3}$$

The corresponding Yukawa term, relevant for generating masses for the neutrinos and the heavy neutral fermions N_L, N_R is given by

$$\begin{aligned}
\mathcal{L}_{\mathrm{Yuk},\nu} = {} & a_1\, (\bar{L}_e)_1 \otimes \left[\begin{pmatrix} \tilde{\Phi}_1 \\ \tilde{\Phi}_2 \\ \tilde{\Phi}_3 \end{pmatrix}_3 \otimes \begin{pmatrix} N_{1,R} \\ N_{2,R} \\ N_{3,R} \end{pmatrix}_{3'} \right]_1 \\
& + a_2\, (\bar{L}_\mu)_{1''} \otimes \left[\begin{pmatrix} \tilde{\Phi}_1 \\ \tilde{\Phi}_2 \\ \tilde{\Phi}_3 \end{pmatrix}_3 \otimes \begin{pmatrix} N_{1,R} \\ N_{2,R} \\ N_{3,R} \end{pmatrix}_{3'} \right]_{1'} \\
& + a_3\, (\bar{L}_\tau)_{1'} \otimes \left[\begin{pmatrix} \tilde{\Phi}_1 \\ \tilde{\Phi}_2 \\ \tilde{\Phi}_3 \end{pmatrix}_3 \otimes \begin{pmatrix} N_{1,R} \\ N_{2,R} \\ N_{3,R} \end{pmatrix}_{3'} \right]_{1''} \\
& + b_1 \left[\begin{pmatrix} \bar{N}_{1,L} \\ \bar{N}_{2,L} \\ \bar{N}_{3,L} \end{pmatrix}_3 \otimes \begin{pmatrix} \chi_1 \\ \chi_2 \\ \chi_3 \end{pmatrix}_{3'} \right]_1 \otimes (\nu_{e,R})_1 \\
& + b_2 \left[\begin{pmatrix} \bar{N}_{1,L} \\ \bar{N}_{2,L} \\ \bar{N}_{3,L} \end{pmatrix}_3 \otimes \begin{pmatrix} \chi_1 \\ \chi_2 \\ \chi_3 \end{pmatrix}_{3'} \right]_{1''} \otimes (\nu_{\mu,R})_{1'} \\
& + b_3 \left[\begin{pmatrix} \bar{N}_{1,L} \\ \bar{N}_{2,L} \\ \bar{N}_{3,L} \end{pmatrix}_3 \otimes \begin{pmatrix} \chi_1 \\ \chi_2 \\ \chi_3 \end{pmatrix}_{3'} \right]_{1'} \otimes (\nu_{\tau,R})_{1''} \\
& + M \begin{pmatrix} \bar{N}_{1,L} \\ \bar{N}_{2,L} \\ \bar{N}_{3,L} \end{pmatrix}_3 \otimes \begin{pmatrix} N_{1,R} \\ N_{2,R} \\ N_{3,R} \end{pmatrix}_{3'}
\end{aligned} \tag{4}$$

where a_i, b_i; $i = 1, 2, 3$ are the Yukawa couplings which are taken to be real. The parameter M is the gauge and flavor-invariant mass term for the heavy leptons. After symmetry breaking the scalars χ_i also acquire vevs $\langle \chi_i \rangle = u_i$; $i = 1, 2, 3$. The 6×6 mass matrix for the neutrinos and the heavy fermions in basis $(\bar{\nu}_{e,L}, \bar{\nu}_{\mu,L}, \bar{\nu}_{\tau,L}, \bar{N}_{1,L}, \bar{N}_{2,L}, \bar{N}_{3,L})$ and $(\nu_{e,R}, \nu_{\mu,R}, \nu_{\tau,R}, N_{1,R}, N_{2,R}, N_{3,R})^T$ is given by

$$M_{\nu,N} = \begin{pmatrix} 0 & 0 & 0 & a_1 v_1 & a_1 v_2 & a_1 v_3 \\ 0 & 0 & 0 & a_2 v_1 & \omega a_2 v_2 & \omega^2 a_2 v_3 \\ 0 & 0 & 0 & a_3 v_1 & \omega^2 a_3 v_2 & \omega a_3 v_3 \\ b_1 u_1 & b_2 u_1 & b_3 u_1 & M & 0 & 0 \\ b_1 u_2 & \omega b_2 u_2 & \omega^2 b_3 u_2 & 0 & M & 0 \\ b_1 u_3 & \omega^2 b_2 u_3 & \omega b_3 u_3 & 0 & 0 & M \end{pmatrix} \tag{5}$$

The invariant mass term M for the heavy leptons N_L, N_R can be naturally much larger than the symmetry breaking scales appearing in the off-diagonal blocks, i.e. $M \gg v_i, u_i$. In this limit the mass matrix in (5) can be easily block diagonalized by the perturbative

[1] Here we denote the irreducible representations of $\Delta(27)$ as in [22], instead of the alternative "two index" notation used in [25]. The two are related by: $\mathbf{1} \equiv \mathbf{1}_{(0,0)}$, $\mathbf{1}' \equiv \mathbf{1}_{(2,0)}$, $\mathbf{1}'' \equiv \mathbf{1}_{(1,0)}$, $\mathbf{3} \equiv \mathbf{3}_{(0,1)}$, $\mathbf{3}' \equiv \mathbf{3}_{(0,2)}$.

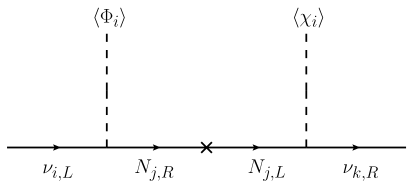

Fig. 1. The Dirac type-I seesaw mechanism. Φ_i and χ_i are triplets under $\Delta(27)$.

seesaw diagonalization method given in [26]. The resulting 3×3 mass matrix for light neutrinos can be viewed as the Dirac version of the well known type-I seesaw mechanism. The above mass generation mechanism can also be represented diagramatically as shown in Fig. 1.

The 3×3 light neutrino mass matrix along with the charged lepton mass matrix (3) have enough free parameters to account for all the observed mass and mixing parameters in the lepton sector. As has been discussed in several previous works [22–24,27–29], for $\Delta(27)$ we focus on the vev alignment $v_1=v_2=v_3=v$ and $u_1=u_2=u_3=u$ as a reference case. Taking this "double alignment" limit for the vevs of the scalars the charged lepton mass matrix M_l is given by

$$M_l = v\begin{pmatrix} y_1 & y_1 & y_1 \\ y_2 & \omega y_2 & \omega^2 y_2 \\ y_3 & \omega^2 y_3 & \omega y_3 \end{pmatrix}, \tag{6}$$

and can be easily diagonalized from right by the familiar "magic matrix" U_ω given by

$$U_\omega = \frac{1}{\sqrt{3}}\begin{pmatrix} 1 & 1 & 1 \\ 1 & \omega & \omega^2 \\ 1 & \omega^2 & \omega \end{pmatrix}. \tag{7}$$

This leads to

$$M_l \,.\, U_\omega^\dagger = \begin{pmatrix} \sqrt{3}v y_1 & 0 & 0 \\ 0 & \sqrt{3}v y_2 & 0 \\ 0 & 0 & \sqrt{3}v y_3 \end{pmatrix}. \tag{8}$$

Likewise, the neutral fermion mass matrix $M_{\nu,N}$ in the above alignment limit is given by

$$M_{\nu,N} = \begin{pmatrix} 0 & 0 & 0 & a_1 v & a_1 v & a_1 v \\ 0 & 0 & 0 & a_2 v & \omega a_2 v & \omega^2 a_2 v \\ 0 & 0 & 0 & a_3 v & \omega^2 a_3 v & \omega a_3 v \\ b_1 u & b_2 u & b_3 u & M & 0 & 0 \\ b_1 u & \omega b_2 u & \omega^2 b_3 u & 0 & M & 0 \\ b_1 u & \omega^2 b_2 u & \omega b_3 u & 0 & 0 & M \end{pmatrix} \tag{9}$$

As mentioned before the invariant mass term M for the heavy fermions N_L, N_R is naturally expected to be much larger than the symmetry breaking scale i.e. $v, u \ll M$. In such limit the mass matrix in (10) can be easily block-diagonalized. The resulting 3×3 mass matrix for the light neutrinos assuming such simplest alignment is given by

$$M_\nu = \frac{u\, v}{M}\begin{pmatrix} 3a_1b_1 & 0 & 0 \\ 0 & 0 & 3a_2b_3 \\ 0 & 3a_3b_2 & 0 \end{pmatrix} \tag{10}$$

Clearly the light neutrino mass matrix in Eq. (10) is inconsistent with the current neutrino oscillation data [30] and needs to be modified.

In order to obtain a realistic light neutrino mass spectrum one must generalize the above vev-alignment pattern i.e. $v_1=v_2=v_3=v$ and $u_i=u_j=u$, $u_k\neq u$ where $i, j, k = 1, 2, 3$. Thus in our generalized ansatz we keep the alignment for the isodoublet scalar vevs unchanged, but modify the isosinglet scalars vev alignment. Such a generalization is not unfounded since the scalar sector of our model is much richer than that characterizing the simpler case of only one type of scalars transforming as $\Delta(27)$ triplets, discussed in [22–24,27]. In contrast to previous models we have two different types of scalars namely Φ_i and χ_i both transforming as triplets under $\Delta(27)$. The resulting scalar potential is rich enough to allow for other possible vev alignments to be realized.

We find that any of the three possible choices namely $u_1=u_2=u$, $u_3\neq u$; $u_2=u_3=u$, $u_1\neq u$; $u_1=u_3=u$, $u_2\neq u$ can give realistic neutrino mass matrices. However, for definiteness and to avoid unnecessary repetition henceforth we focus on the choice $u_1=u_3=u$, $u_2\neq u$. Towards the end of the discussion we will comment on the similarities and differences in results for other possibilities.

Since we have kept the vev alignment for the Φ_i fields unchanged it follows that the charged lepton mass matrix Eq. (6) also remains unchanged. As a result it can still be diagonalized by a "magic" rotation from the right as shown in (8). The 6×6 neutral fermion mass matrix now becomes

$$M_{\nu,N} = \begin{pmatrix} 0 & 0 & 0 & a_1 v & a_1 v & a_1 v \\ 0 & 0 & 0 & a_2 v & \omega a_2 v & \omega^2 a_2 v \\ 0 & 0 & 0 & a_3 v & \omega^2 a_3 v & \omega a_3 v \\ b_1 u & b_2 u & b_3 u & M & 0 & 0 \\ b_1 u_2 & \omega b_2 u_2 & \omega^2 b_3 u_2 & 0 & M & 0 \\ b_1 u & \omega^2 b_2 u & \omega b_3 u & 0 & 0 & M \end{pmatrix} \tag{11}$$

As before this mass matrix can be block-diagonalized in the approximation $v, u, u_3 \ll M$. The resulting light three-neutrino mass matrix is

$$M_\nu = \frac{v}{M} \times \begin{pmatrix} a_1b_1(2u+u_2) & a_1b_2(u+\omega^2 u+\omega u_2) & a_1b_3(u+\omega u+\omega^2 u_2) \\ a_2b_1(u+\omega^2 u+\omega u_2) & a_2b_2(u+\omega u+\omega^2 u_2) & a_2b_3(2u+u_2) \\ a_3b_1(u+\omega u+\omega^2 u_2) & a_3b_2(2u+u_2) & a_3b_3(u+\omega^2 u+\omega u_2) \end{pmatrix} \tag{12}$$

3. CP violation

The neutrino mass matrix in Eq. (12) can be diagonalized numerically and leads to neutrino masses and mixing angles consistent with neutrino oscillation experiments [30] as well as cosmological limits [31]. Here we present our numerical results for CP violation in this model. Notice that from the beginning, we have assumed real Yukawa couplings. If we also take a real scalar potential, leptonic CP violation must arise solely by the complex nature of the $\Delta(27)$ flavor symmetry. Indeed, one finds that, with our generalized alignment the resulting neutrino mass matrix (12) leads to no CP violation and in terms of standard parametrization of neutrino mixing matrix [30], one has $\delta_{CP}=0,\pm\pi$ for the CP phase. The latter implies that the Jarlskog invariant J_{CP}, which in the standard PDG parametrization[2] is given by

$$J_{CP} = \frac{1}{8}\sin 2\theta_{12}\sin 2\theta_{23}\sin 2\theta_{13}\cos\theta_{13}\sin\delta_{CP},$$

vanishes.

Recent experimental results have predicted a slight preference for $\delta_{CP}\neq 0,\pm\pi$ implying CP violation in lepton sector [30]. If this indeed is the case then one must consider deviations from the

[2] For a recent discussion of fermion mixing parametrizations see [32].

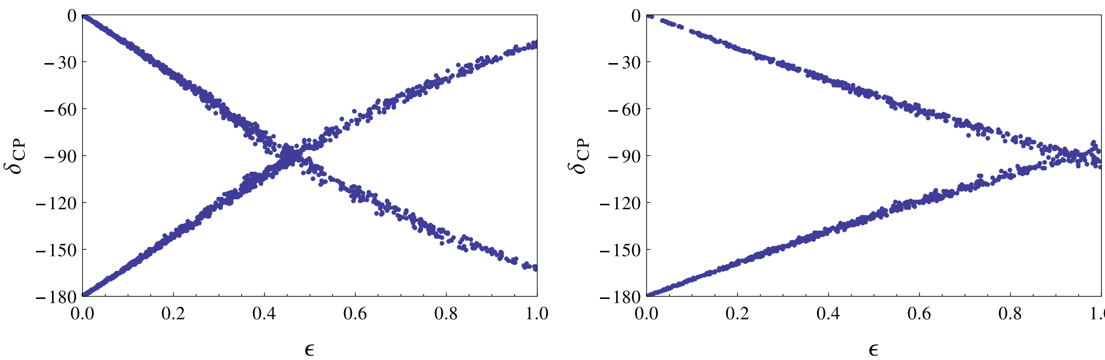

Fig. 2. Leptonic CP violation phase δ_{CP} versus ϵ, the deviation from the reference alignment. For the left panel we have taken $\alpha = 1.2$ whereas in the right panel the $\alpha = 2.5$ is taken. See text.

generalized alignment limit. For example, if we consider small deviation of the type $u_1 = u$, $u_3 = u(1+\epsilon)$, $u_2 = u(1+\alpha)$ then finite CP violation can indeed be generated even for real ϵ and α, as shown in Fig. 2 and 3. The source of CP violation can be traced to the complex parameter ω, where ω is cube root of unity with $\omega^3 = 1$.

As can be seen from the figures, when $\epsilon = 0$, $\delta_{CP} = 0, -\pi$ and $J_{CP} = 0$ implying no CP violation. As we deviate from our reference alignment limit CP violation is generated with $J_{CP} \neq 0$. The magnitude of the CP violation parameter is directly proportional to the deviation ϵ from the alignment limit as well as the parameter α which measures the deviation of u_2 from u_1 i.e. $u_2 = u_1(1+\alpha)$. In plotting Fig. 2 we have randomly varied all other free parameters, namely the vevs and Yukawa couplings. All the Yukawa couplings are varied between -0.5 to 0.5, the u_1 vev is varied between 700 to 800 GeV and the u_2 vev is taken to be $u_1(1+\alpha)$.

For a given value of α the magnitude of CP violation is directly correlated to ϵ as is clear from Fig. 2 and Fig. 3. In Fig. 2 we show the deviation of δ_{CP} with respect to ϵ for fixed values of α. The dependence of the Jarlskog invariant J_{CP} with respect to ϵ, for fixed values of α, is shown in Fig. 3. For the left panel of both figures, we have fixed $\alpha = 1.2$ while for right panel we took $\alpha = 2.5$. As is clear from a comparison of the two panels, the magnitude of CP violation not only depends on ϵ but also on the value of α. For smaller values of α the deviation is sharper than for larger values. In the left panels of the two figures, where a relatively smaller value of α is taken, the δ_{CP} as well as J_{CP} changes rapidly with ϵ and maximal CP violation corresponding to $\delta_{CP} = -\pi/2$ is obtained for $\epsilon \approx 0.45$. Further increase in ϵ values results in decrease in CP violation as can be inferred from the decreasing value of J_{CP} in Fig. 3. The J_{CP} eventually falls back to zero with $\delta_{CP} = 0, -\pi/2$, when $\epsilon = \alpha$ which again corresponds to the reference alignment with u_3 now being equal to u_2. In the right panels of Fig. 2 and Fig. 3, the δ_{CP} and J_{CP} are plotted with respect to ϵ for a fixed values of $\alpha = 2.5$. The nature of the departures of both δ_{CP} and J_{CP} is similar to what is seen in the left panels, but now the slope of the deviation is smaller. For $\alpha = 2.5$ maximal CP violation is achieved for higher value of $\epsilon \approx 0.95$. Just like for the left panels, further increase in ϵ beyond 0.95 leads to decrease in CP violation with the case of no CP violation i.e. $J_{CP} = 0$ with $\delta_{CP} = 0, -\pi/2$ again achieved for $\epsilon = \alpha$ corresponding to the alignment $u_3 = u_2$. Notice that, although here we are presenting results only for positive values of ϵ we mention that negative values of ϵ are equally viable. If we take $\epsilon < 0$ then the essential features of Fig. 2 and 3 are reproduced but for positive values of δ_{CP} and J_{CP}. This means that as ϵ deviates more and more from zero on the negative side, both δ_{CP} and J_{CP} start deviating more and more from the CP conserving case but along the positive direction. Again the departure depends also on the value of α with smaller values of α leading to sharper deviation with respect to ϵ.

Finally before closing this section let us briefly remark on other possible alignment choices, e.g. $u_1 = u$, $u_2 = u(1+\epsilon')$, $u_3 = u(1+\alpha')$ where ϵ' and α' parametrize the deviations of u_2, u_3 from u, respectively. As in the previous case, here also for the case of perfect alignment i.e. for $\epsilon' = 0$ we have no CP violation with $\delta_{CP} = 0, \pm\pi$ and $J_{CP} = 0$. Also as before when ϵ' deviates from zero in either direction we generate CP violation. However, unlike the previous case, the nature of the correlation in this case is different, since for $\epsilon' > 0$ both δ_{CP} and J_{CP} acquire positive values, whereas for $\epsilon' < 0$ both $\delta_{CP}, J_{CP} < 0$. This behaviour is opposite to that found in previous, case where for $\epsilon > 0$ we had $\delta_{CP}, J_{CP} < 0$ and for $\epsilon < 0$ we had $\delta_{CP}, J_{CP} > 0$. Apart from this, other features of the previous case like the dependence on ϵ' and α' are qualitatively realized in this case also. Finally, for the third alignment choice i.e. $u_2 = u$, $u_3 = u(1+\epsilon')$, $u_1 = u(1+\alpha')$ the qualitative nature of CP violation with respect to alignment deviation is essentially the same as shown in Fig. 2 and Fig. 3. To avoid unnecessary repetition we refrain from discussing these two alignment choices in more detail.

4. WIMP scalar dark matter candidate

Here we recall the dark matter features of the model, which employs similar ingredients as the simplest prototype model considered in [33]. In this section we briefly consider the role of the scalars ζ and η, which are singlets under the $\mathrm{SU(3)_c} \otimes \mathrm{SU(2)_L} \otimes \mathrm{U(1)_Y}$ gauge group, transform trivially under $\Delta(27)$, but carry Z_4 lepton quarticity charges $\mathbf{z}$ and $\mathbf{z}^2$ respectively. If ζ and η are removed, the Lagrangian of the model presents a larger symmetry associated to $\mathrm{SU(3)_c} \otimes \mathrm{SU(2)_L} \otimes \mathrm{U(1)_Y} \otimes \Delta(27) \otimes U(1)$ where $U(1)$ is a continuous global symmetry which may be interpreted as a generalized global lepton number. However, in the presence of the scalars ζ and η one can write following Z_4 invariant terms in the scalar potential

$$\eta^2,\ \eta\zeta^2,\ \eta^4,\ \zeta^4,\ \eta^2\zeta^*\zeta \quad + \quad \text{h.c.} \tag{13}$$

Notice that all these terms are Z_4 invariant but break the global $U(1)$ invariance so the remaining family symmetry group is just $\Delta(27) \otimes Z_4$.[3] On the other hand note that the field η also couples to the right handed neutrinos through a Z_4 invariant term

$$\bar{\nu}^c_{i,R}\,\nu_{j,R}\,\eta \ + \ \text{h.c.} \tag{14}$$

[3] We do not bother writing the other scalar potential terms which are also invariant under the global U(1).

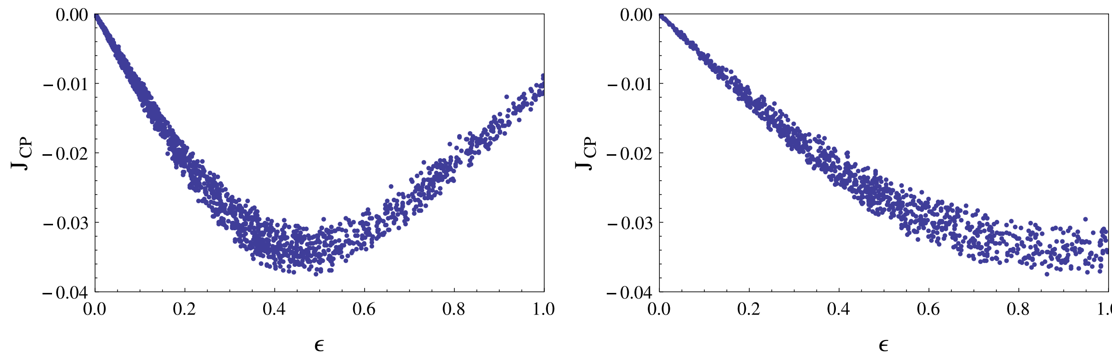

Fig. 3. The Jarlskog invariant J_{CP} versus the deviation from alignment ϵ. The range of variation of other free parameters is the same as in 2. For the left panel we have taken $\alpha = 1.2$ whereas in the right panel we take $\alpha = 2.5$.

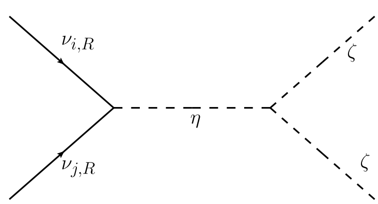

Fig. 4. Interaction between the dark matter candidate ζ and the right handed neutrinos, mediated by the exchange of the scalar η.

Since this Yukawa coupling is only Z_4 invariant, it breaks the continuous $U(1)$ symmetry. Due to the couplings of η to the scalar ζ in Eq. (13) and to right handed neutrinos as in Eq. (14), the latter also couple to ζ as shown in Fig. 4.

Note that the flavor symmetry $\Delta(27)$ breaks spontaneously when the $\Delta(27)$ triplet scalars Φ_i and χ_i acquire nonzero vevs. However, since neither Φ_i nor χ_i carries the Z_4 charge, and ζ and η are assumed not to acquire any vev, one finds that the Z_4 remains unbroken. This implies that the neutrinos retain their Dirac nature, since Majorana mass terms are forbidden by the unbroken Z_4.

As a result one finds that the field ζ can act as a stable particle and hence a potential candidate for the cosmological dark matter. This implies that there is no term of the form $\zeta\rho_i\rho_j$ or of the form $\zeta\psi_i\psi_j$, where ρ_i, ρ_j stand for other scalars and ψ_j, ψ_i denote generic fermions. Thus, the residual Z_4 symmetry responsible for the Dirac nature of neutrinos also ensures the stability of the ζ making it a potentially viable dark matter candidate.

Although ζ is stable, and without direct tree level coupling to fermions, due to the model symmetry, it still interacts with other scalars through quartic potential terms of the type $\zeta^*\zeta\rho_i^\dagger\rho_j$ and also couples to right handed neutrinos through exchange of η as shown in Fig. 4. These terms imply that two dark matter particles can annihilate into two other scalars, potentially leading to the correct relic density for dark matter [21,37]. Also, the dark matter interaction with the Higgs (h)[4] can be used to detect it by experiments searching for nuclear recoil [21] induced by Higgs boson exchange. Moreover, if the dark matter mass obeys $m_\zeta < m_h/2$ then it can lead to invisible decay of Higgs. Both nuclear recoil experiments such as LUX [34] and PandaX [35] as well as LHC searches for invisibly decaying Higgs boson [36,38] lead to stringent constraints on the Higgs dark matter coupling as shown in Fig. 5. In plotting Fig. 5, we have taken the constraints from the latest ATLAS searches for invisible Higgs decays [36], since the ATLAS constraint is more stringent than that of CMS [38]. Concerning constraints from nuclear recoil experiments, the LUX [34] and PandaX [35] experimental constraints are taken, assuming that the nucleon Higgs coupling and the nucleon mass parameters are the same as in [33]. Our treatment for dark matter constraints follows closely Ref. [33] which should be consulted for further details. Thus ζ realizes a "Higgs portal" dark matter scenario. This type of dark matter, charged under a given discrete symmetry, has been previously studied in several papers and shown to provide a viable dark matter scenario [21,37,39,40]. Another implication of our model is the conservation of the Z_4 charge in the presence of lepton number violation [41,42]. The fact that η is a real scalar field which couples to right handed neutrinos, means that its decay to two neutrinos or two antineutrinos would potentially generate a lepton asymmetry in the Universe. The possibility of leptogenesis with a conserved Z_4 lepton number has indeed been pointed out in [43]. Clearly this scenario deserves more work.

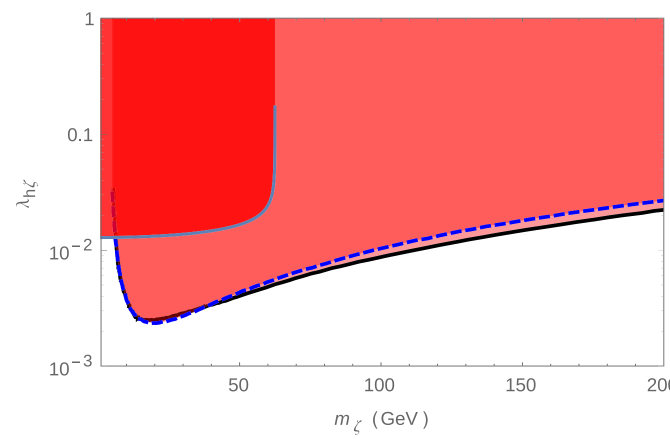

Fig. 5. The experimental sensitivity of our WIMP scalar dark matter candidate to invisible Higgs decay and direct detection. The light shaded region is ruled out by LUX (black continuous line) [34] and PandaX (blue dashed line) [35] data whereas the dark shaded region is ruled out by the bound on the Higgs invisible decay width from the LHC [36].

5. Discussion and summary

We have suggested a simple flavor model based on the $\Delta(27)$ group, in which the light neutrinos are Dirac fermions and the smallness of their masses results from a type-I seesaw mechanism. Leptonic CP violation is related to the pattern of flavor symmetry breaking, described through the Higgs vacuum expectation values alignment, as shown in Figs. 2 and 3 above. The scheme naturally leads to a WIMP dark matter candidate which is made stable by

[4] We denote the 125 GeV scalar discovered at LHC in 2012 as the Higgs. In our model it will be an admixture of the scalars Φ_i and χ_i.

the same discrete lepton number Z_4 symmetry which makes neutrinos to be Dirac particles. In short, dark matter stability emerges from the lepton quarticity which also ensures the Dirac nature of neutrinos. A detailed study of its discovery potential in direct and indirect detection experiments will be presented elsewhere. Before closing let us also mention that our model can easily be generalized by including vector-like quarks, so as to accommodate the recent diphoton hint seen by the ATLAS and CMS collaborations. It would be identified with one of the scalars in the χ multiplet, very much along the lines of Refs. [44,45]. In this paper we have discussed leptons only. Quarks can be introduced in a trivial way as flavor singlets, along with a new Higgs scalar multiplet. This Higgs scalar can be forbidden to couple with leptons by an additional Z_2 symmetry in a way akin to the lepton specific two Higgs doublet model [46]. This way the quark and lepton sectors would be clearly independent, without any predictions for the CKM matrix. In contrast, obtaining successful CKM predictions by assigning non-trivial charges in the quark sector constitutes a challenge beyond the scope of this paper.

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