Dirac neutrinos from Peccei-Quinn symmetry: two examples

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We aim to explain the nature of neutrinos using Peccei-Quinn symmetry. We discuss two simple scenarios, one based on a type-II Dirac seesaw and the other in a one-loop neutrino mass generation, which solve the strong CP problem and naturally lead to Dirac neutrinos. In the first setup latest neutrino mass limit gives rise to axion which is in the reach of conventional searches. Moreover, we have both axion as well as WIMP dark mater for our second set up.

1. INTRODUCTION

Over the past few decades there has been remarkable progress in the field of particle physics. The discovery of neutrino oscillations [1] provides a major milestone to understand some intriguing aspects of neutrinos, which clarify the fact that neutrinos possess non-zero mass and their different flavors are mixed. Apart from these, various observed phenomena provide some hints for the existence of a non-baryonic form of matter, known as dark matter [2], in the Universe. Both of these issues are the most serious drawbacks of the Standard Model (SM) of particle physics. Thus, they provide a clear evidence of new physics beyond the SM. Besides these, the SM also fails on shedding light on the strong CP problem of QCD, suggested by an experimental bound of the electric dipole moment of a neutron [3]. Peccei-Quinn (PQ) [4] symmetry has been the most appreciated approach to explain the strong CP problem. The PQ symmetry predicts the existence of the associated pseudo-Nambu-Goldstone (pNG) boson, the axion [5, 6] which can be a good cold dark matter candidate [7–10]. Another puzzling challenge in the neutrino sector is whether neutrinos are Dirac or Majorana particles. Despite the ongoing experimental effort on the search for the neutrinoless double beta decay [11], which if observed will indicate the Majorana nature of neutrinos [12], no signal of this process has been detected. Connecting these seemingly unrelated puzzles with the smallness of the neutrino masses is the scope of the present manuscript.

Axion models are mainly categorized into two classes, depending on whether quarks are charged under PQ symmetry or not, namely, DFSZ [13, 14] and KSVZ [15, 16]. In axion models, where quarks carry PQ charge, one needs two Higgs doublets $H_u$ and $H_d$ both charged under PQ symmetry in such a way that they couple to the PQ field $\sigma$ (singlet under the SM gauge group). There are two possibilities for such a coupling, namely it can be trilinear or quartic. When the coupling is quartic, the spontaneous breaking of the Peccei-Quinn symmetry can be connected to the breaking of lepton number by two units [17, 18] leading to Majorana neutrinos [19].

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Recently, it has been shown that in a specific DFSZ axion scenario that neutrinos can only be Dirac particles [20]. In order to explain the small (effective) neutrino Yukawa coupling, the tree-level coupling with the Higgs is forbidden by means of the PQ symmetry and allowed at the dimension-5 level. The PQ field plays a role in generating the Dirac neutrino mass which is proportional to the PQ breaking scale. In [20] the Dirac neutrino masses were generated through a type-I Dirac seesaw. Scenarios where the neutrino mass mechanism and the PQ symmetry breaking are related has attracted attention\(^1\), both in the Majorana \([23–29]\) and Dirac \([30–33]\) frameworks.

In this letter, we continue with the same approach by giving two different ways for naturally generating small Dirac neutrino masses. In the first case, we extend the idea by using a type-II seesaw with exactly the same PQ charges for the RH neutrino and by including an extra Higgs doublet to give the small Dirac neutrino masses. Another option we explore is to generate neutrino masses at one loop level \(^2\). In the second case, we kept exactly the same PQ charges for the right-handed (RH) neutrino as well as for the Higgs doublets \(H_u\) and \(H_d\) as discussed in [20]. Moreover, to serve our purpose, we also include SM singlets \(N_R, N_L, \zeta\) with different PQ charges and an extra SM doublet \(\eta_u\) (see Table IV for details). In both frameworks, we show how it leads to a novel class of minimal axion models that effectively imply Dirac neutrinos.

2. FRAMEWORK

It has been pointed out in [20] that one can explain Dirac nature of neutrinos in the context of the DFSZ axion model. If we consider the model in Table I, there is no way to generate a \(\Delta L = 2\) operator at any order at the perturbative level. If we include the RH neutrino transforming as a \(-1\) as the rest of the RH fields, a Dirac mass is generated through the Yukawa coupling with the Higgs, in such a case, the Yukawa Lagrangian is given by

\[
\mathcal{L}_Y = y_{ij}^{u} \bar{Q}_i H_u u_j + y_{ij}^d \bar{\nu}_i H_d d_j + y_{ij}^{l} \bar{L}_i H_d l_j + y_{ij}^\nu \bar{L}_i H_u \nu_R j + h.c.,
\]

in such a way that the Yukawa couplings \(y_{ij}^\nu\) must be \(O(10^{-12})\) in order to account for the recent KATRIN bound \([34]\). As was shown in [20], if the RH neutrino transforms as \(-5\) the direct Yukawa coupling is forbidden \(^3\). The effective dimension-5 operator can be generated as follows

\[
\mathcal{L}_{\text{dim} 5}^D = y_{ij}^\nu \bar{L}_i H_u \nu_R j \frac{\sigma}{\Lambda_{\text{UV}}} + h.c.,
\]

where \(\sigma\) can be expressed as

\[
\sigma(x) = \frac{1}{\sqrt{2}} (\rho(x) + f_a) e^{ia(x)/f_a}.
\]

\(^1\) For alternative solution to the strong CP problem connected with neutrino masses, see for instance \([21, 22]\).

\(^2\) Following the reasoning in [20] it can also be extended to more than one loop but this is not the scope of the present manuscript.

\(^3\) A charge of +3 is also possible but in that case a Dirac mass is generated by the \(\tilde{H}_d\) field.
Here, $a(x)$ is the QCD axion [5, 6], $f_a$ is the PQ breaking scale and $\rho(x)$ is the radial part that will gain a mass of order of the PQ symmetry breaking scale.

As was pointed out previously, it can be easily UV completed through a type-I Dirac seesaw. In the following sections we will give two other frameworks to UV complete such an operator.

### 2.1. Case-I: Type-II Dirac seesaw

Here we will consider a concrete example of an extension of the model in Table I by including the RH neutrino transforming as $-5$ under the PQ symmetry and an extra $SU(2)_L$ doublet, $\Phi_u$, with PQ charge 6, see Table II. In this case there are terms in the potential of the form

$$ V \sim \kappa H_u H_d \sigma^* + \lambda' H_d \Phi_u \sigma^*^{2}, \quad (4) $$

where the couplings $\kappa$ and $\lambda'$ are dimensionful and dimensionless, respectively. Now, by inspecting the PQ charge assignments of different fields content as given in Table II, we notice that there is no way to form the dimension-5 Weinberg operator for the light neutrino masses, nor any other operator with powers of $\sigma$ and or powers of $H_u$, $H_d$ and $\Phi_u$. We first write down the possible dimension-5 Weinberg operators in the presence of Higgs doublets

\[ \mathcal{L}_{\text{dim} 5} \sim \{ \frac{LLH_u H_d}{\Lambda_{UV}} \}_{1 + 1 + (0) = +2} \]

It is apparent from Eq. (5) that none of these operators are invariant under the PQ symmetry. Moreover, notice that all these operators transform as $m \mod(4) = 2$ under PQ. Hence, from

| Fields/Symmetry | $Q_i$ | $u_i$ | $d_i$ | $L_i$ | $l_i$ | $H_u$ | $H_d$ | $\sigma$ |
|----------------|------|------|------|------|------|------|------|--------|
| $SU(2)_L \times U(1)_Y$ | (2,1/6) | (1,2/3) | (1,-1/3) | (2,-1/2) | (1,-1) | (2,-1/2) | (2,1/2) | (1,-1) |
| $U(1)_{PQ}$ | 1 | -1 | -1 | 1 | -1 | 2 | 2 | 4 |

TABLE I: Quantum numbers in the DFSZ axion model.
Eq. (5), there is no way to construct an operator invariant under $U(1)_{PQ}$ and the SM symmetries simultaneously. We further realize that this argument also extends to all the higher order effective operators that could potentially generate Majorana neutrino masses. In the following, we give all possible gauge invariant contractions of the scalar fields and their PQ charges:

$$\sigma^n (4n); \quad (\sigma^*)^n (-4n);$$
$$\left(H_u H_d\right)^n (4n); \quad \left(H_u H_d\right)^* n (-4n);$$
$$\left(H_u^\dagger H_u\right)^n (0); \quad \left(H_d^\dagger H_d\right)^n (0);$$
$$\left(H_u^\dagger \Phi_u\right)^n (4n); \quad \left(H_u^\dagger \Phi_u\right)^* n (-4n);$$
$$\left(H_d^\dagger \Phi_u\right)^n (8n); \quad \left(H_d^\dagger \Phi_u\right)^* n (-8n);$$
$$\left(\Phi_u^\dagger \Phi_u\right)^n (0);$$

where as we can see, all these contractions (and their combinations) are 0 or multiples of 4 under PQ symmetry. Therefore, there is no way to make a combination of operators on Eqs. (5) and (6) invariant under PQ symmetry and hence, neutrinos must be Dirac particles.

The relevant part of Yukawa Lagrangian that generates Dirac neutrino masses is given by

$$L_Y \supset y_{ij}\nu_i H_u \nu_j H_d \sigma + \mu H_u^\dagger \sigma + h.c. \quad (7)$$

where the scalar particles of the model mix through a unitary matrix $K$ into a mass eigenstate basis as

$$\phi_i = K_{ij} S_j, \quad (8)$$

where $\phi$ are the real neutral components of the scalars $H_u$, $H_d$, $\Phi_u$ and $\sigma$. We consider the limit where one of the mass eigenstates is mostly composed of $\Phi_u$, with a mass $M_{\Phi_u}^2 >> v_u^2$. The largest
FIG. 2: Exclusion region plots (colored regions are excluded) in \((M_{\Phi_u} - f_a)\) plane for a type-II Dirac seesaw mechanism. Three benchmark values for \(\mu y = 1\) GeV, 1 MeV, 1 keV have been adopted, respectively. The plots are presented by using the limits on neutrino mass from KATRIN [34] which gives \(m_\nu < 1.1\) eV at 90% C.L. (left panel) and Planck [35] \(\sum m_\nu < 0.12\) eV (at 95% confidence level using TT, TE, EE + lowE + lensing + BAO) (right panel).

corresponding to the mixing between \(\Phi_u\) and other fields is the \(\mu\) term of Eq. (7), the large vev of \(\sigma\) can induce a large mixing between \(H_u\) and \(\Phi_u\). This mixing can raise the mass of one of the light eigenstates above the EW scale, excluding the possibility that the scalar which is predominantly \(H_u\) is the 125 GeV Higgs boson. At leading order the mixing between these fields goes as

\[
\sin \theta \sim \frac{\mu f_a}{M_{\Phi_u}^2} .
\]  

The neutrino masses resulting from the breaking of the \(SU(2)_L \times U(1)_Y\) and \(U(1)_{PQ}\) symmetries by the scalar vevs \(\langle H_u \rangle = v_u\) and \(\langle \sigma \rangle = f_a\) are

\[
(m_\nu)_{ij} = y^{\nu}_{ij} \frac{\mu v_u f_a}{M_{\Phi_u}^2} \sim y^{\nu}_{ij} v_u \sin \theta ,
\]  

where in the last term we have used Eq. (9).

To explain the tiny masses of neutrinos, a large scale for the scalar mediator must be introduced. Notice that in this case the neutrino masses depend on the inverse squared of the mediator mass, unlike the type-I case where the dependence is only on the inverse mass. This suggests that a lower mass scale for the mediator than in type-I seesaw [20] may be allowed. The measurement of the tritium beta decay spectrum at KATRIN [34] currently yields a direct limit for neutrino masses of \(m_\nu < 1.1\) eV at 90% C.L., while the indirect limit from the Planck telescope data [35] (at 95% confidence level using TT, TE, EE + lowE + lensing + BAO) constrains them further to \(\sum m_\nu < 0.12\) eV. The bounds of tritium beta decay and cosmology are translated into bands for the allowed scales for \(f_a\) and \(M_\Phi\) as shown in Fig. 2.
2.2. Case-II: One-loop Dirac seesaw

In this section, we discuss a one-loop mechanism to UV complete the effective coupling of Eq. (2). A detailed discussion of all the possible topologies to explain Dirac neutrino masses with four external lines was outlined in [36]. In what follows, we consider the most economical scenarios for the dimension-5 operator which can lead to Dirac neutrino masses. All the necessary fields carry $SM \otimes PQ$ charges are presented in Table (III). Under this assignment of PQ charges, and the subsequent spontaneous symmetry breaking, the residual symmetry is $Z_2$. Note that the leptons, quarks, $\eta_u$ and $\zeta$ are odd under the $Z_2$ residual symmetry. Dark matter stability is achieved for the lightest odd scalar or for $N$ by the interplay of $Z_2$ and Lorentz invariance [37].

![Feynman diagram for Dirac neutrino masses in alternative one-loop DFSZ scenario. Here left panel respects PQ charge assignment, whereas right panel respects the remnant $Z_2^{PQ}$ charge assignment that arises due to PQ symmetry breaking.](image)

**TABLE III:** Field content and transformation properties under PQ symmetry in the alternative one-loop mechanism.

| Symmetry/Fields | $L_1$ | $\nu_R$ | $H_u$ | $H_d$ | $N_R$ | $N_L$ | $\sigma$ | $\eta_u$ | $\zeta$ |
|-----------------|-------|---------|-------|-------|-------|-------|---------|---------|-------|
| $SU(2)_L \times U(1)_Y$ | (2, 1/2) | (2, -1/2) | (2, 1/2) | (2, -1/2) | (2, 1/2) | (2, -1/2) | (2, 1/2) | (2, -1/2) | (2, 1/2) |
| $U(1)_{PQ}$ | 1 | -5 | 2 | 2 | 2 | 4 | -1 | 7 |
| $Z_2^{Q}$ | -1 | -1 | +1 | +1 | +1 | +1 | -1 | -1 |

Now we write down the scalar potential that is allowed by the $SM \times U(1)_{PQ}$ symmetries as below,

$$V = \mu_u^2 H_u^\dagger H_u + \lambda_u (H_u^\dagger H_u)^2 + \mu_{\eta_u}^2 \eta_u^\dagger \eta_u + \lambda_{\eta_u} (\eta_u^\dagger \eta_u)^2$$

$$+ \mu_\sigma^2 \sigma^* \sigma + \lambda_\sigma (\sigma^* \sigma)^2 + \mu_\zeta^2 \zeta^* \zeta + \lambda_\zeta (\zeta^* \zeta)^2$$

$$+ \lambda_{\eta_u \eta_u} (H_u^\dagger H_u)(\eta_u^\dagger \eta_u) + \lambda_{\eta_u \sigma} (H_u^\dagger \eta_u)(\eta_u^\dagger H_u) + \lambda_{\eta_u \sigma} (H_u^\dagger H_u)(\sigma^* \sigma)$$

$$+ \lambda_{\eta_u \zeta} (H_u^\dagger H_u)(\zeta^* \zeta) + \lambda_{\eta_u \zeta} (\eta_u^\dagger \eta_u)(\zeta^* \zeta) + \lambda_{\eta_u \sigma} (\eta_u^\dagger \eta_u)(\sigma^* \sigma) + \lambda_{\sigma \zeta} (\sigma^* \sigma)(\zeta^* \zeta)$$

$$+ \kappa H_u H_d \sigma^* + \lambda_1 [H_u H_d \zeta^* \sigma + h.c.] .$$

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4 We also provide an alternate loop model considering half-integral PQ charges for the particles running inside the loop in appendix A. There dark matter stability is obtained under $Z_4$ residual symmetry.
Note that all the terms that are allowed for \( H_u \) are also allowed for \( H_d \) except for \( \lambda_{u_{th}}^2 \), \( \lambda_1 \) terms, which are not invariant under the SM gauge group. After the electroweak symmetry breaking, the SM Higgs doublet \( H_u \) acquires its vev \( v_u \). Then, considering the \( \eta_u - \zeta \) mixing matrix, the mass matrix for the neutral scalars in \((\eta_u^0, \zeta)\) basis can be written as

\[
\mathcal{M} = \begin{pmatrix}
\mu_{u}^{2} + \lambda_{u_{th}} v_{u}^{2} + \lambda_{\eta_{u}} f_{u}^{2} & \lambda_{1} v_{u} f_{u} \\
\mu_{1}^{2} + \lambda_{\eta_{u}} v_{u}^{2} + \lambda_{\sigma_{u}} f_{u}^{2} & \mu_{\sigma}^{2} + \lambda_{\eta_{u}} v_{u}^{2}
\end{pmatrix}.
\]

(12)

where \( f_u \) is the PQ symmetry breaking scale and we used \( \lambda_{u_{th}} = \lambda_{1}^{u_{th}} + \lambda_{2}^{u_{th}} \).

To find DM masses i.e., the mass eigenvalues of Eq. (12), we make a simple numerical estimation. We take \( v_u \sim 100 \) GeV, \( f_u \sim 10^{12} \) GeV and \( \lambda_1 = 10^{-3} \) in our calculation. We write down \( \mathcal{M} \) as

\[
\mathcal{M} = \begin{pmatrix}
10^{24} + 10^{-7} \times 10^{4} - 10^{-4} \times 10^{24} & 10^{-3} \times 10^{2} \times 10^{12} \\
* & 10^{20} + 10^{-7} \times 10^{-4} - 10^{-4} \times 10^{24}
\end{pmatrix} \text{[GeV]}.
\]

(13)

For simplicity, we assume \( \lambda_{u_{th}} = \lambda_{u_{z}} \) and \( \lambda_{\eta_{u}} = \lambda_{\sigma_{u}} \). We also consider \( \mu_{\eta_{u}} = 10^{12} \) GeV, \( \mu_{\sigma} = 10^{10} \) GeV together with \( \lambda_{u_{th}} = 10^{-7} \) and \( \lambda_{\eta_{u}} = -10^{-4} \). Diagonalization of Eq. (13) leads to the lightest DM of \( \mathcal{O}(1) \) TeV, whereas the heaviest DM came out as \( \sim 10^{12} \) GeV for the given set of numerical values.

We describe the relevant part of Yukawa Lagrangian for the leptons as

\[
\mathcal{L}_Y \supset y_{ij}^c \bar{L}_i \eta_R N_{Rj} + M_{jk} \bar{N}_{Rj} N_{Lk} + y_{ij}^l \bar{N}_{Lk} \nu_{Ri} \zeta + h.c.
\]

(14)

The neutrino mass obtained from the one-loop diagram in Fig. 4 is given by [32, 38]

\[
m_{ij}^{\nu} = \frac{1}{64\pi^2} \sum_{X=R,L} \frac{\lambda_{1} v_{u} f_{u}}{m_{X}^2 - m_{S}^2} \sum_{k} y_{ik}^{c} y_{kj}^{c} m_{N_{k}} \left[ F \left( m_{S}^{2} x \right) - F \left( m_{S}^{2} x \right) \right].
\]

(15)

where \( F(x) = x \log(x)/(x - 1) \) and \( S_{Ri/Li} \) (for \( i = 1, 2 \)) denote the CP even and odd mass scalar eigenstates, obtained from the \( \eta_u - \zeta \) mixing. In the \( m_{N_{light}} \ll m_{S} \) limit, and considering \( m_{S_1}^2 \sim m_{S_2}^2 = m_{S}^2 \gg 1 \), \( f_u v_u \), the neutrino mass matrix can be written as

\[
m_{ij}^{\nu} \sim \frac{\lambda_{1} v_{u} f_{u}}{32\pi^2} \sum_{k} y_{ik}^{c} y_{kj}^{c} m_{N_{k}} \left[ \log \left( m_{N_{k}}^{2} m_{S}^2 \right) - 1 \right].
\]

(16)

Alternatively, considering \( m_{S} \ll m_{N_{k}} \), we find

\[
m_{ij}^{\nu} \sim \frac{\lambda_{1} v_{u} f_{u}}{32\pi^2} \sum_{k} y_{ik}^{c} y_{kj}^{c} \left[ \log \left( m_{N_{k}}^{2} m_{S}^2 \right) - 1 \right].
\]

(17)

Because of the loop suppression factor we can accommodate one of the heavy mediators in the TeV range, not far from the EW scale, which can be tested at LHC [39] while having neutrino masses at the KATRIN limit. In this particular scenario, we will have two DM candidates, one being the axion and the other one the lightest particle running inside the loop. For the scalar DM
case, \( m_S << m_{N_e} \), the DM is a mixture of a electroweak doublet and a singlet, as in [40]. Two limiting cases are found in the small mixing regime, when DM is mostly composed of \( \eta_u \) or of \( \zeta \). The mixings between \( \eta_u \) and \( \zeta \) are given, at leading order by

\[
\sin 2\theta_X \sim \frac{\lambda_1 v_u f_\mu}{m^2_{S_{X_2}} - m^2_{S_{X_1}}}.
\]

When the lightest dark scalar is mostly a singlet, \( S_{\text{light}} \approx \zeta \), the annihilation channel for thermal production is through the scalar couplings in the potential [41, 42]. This means that the annihilation processes \( S_{\text{light}} S_{\text{light}} \rightarrow \text{SM SM} \) that determine the relic abundance of \( S_{\text{light}} \) are mediated by scalar channels. On the other hand, when the lightest dark scalar is mostly a doublet, \( S_{\text{light}} \approx \eta_u \), the annihilation channel for thermal production is determined by the scalar couplings in the potential and the gauge couplings. When gauge interactions dominate, a \( \mathcal{O}(1) \) TeV mass is needed to obtain the correct relic density [43]. The last possibility is that a fermion \( N \) is the DM candidate. Here t-channel annihilation to leptons determine the annihilation cross section, as in the original Scotogenic model [38], with the addition of an annihilation channel into the right handed component of Dirac neutrinos.

Given the possibility of having mixed dark matter in the model, in the form of a scalar or fermionic WIMP and the axion, we expect that the constraints on the model parameters for pure axion models, or pure WIMP models are relaxed. The combined WIMP and axion density, therefore has to obey [35]

\[
(\Omega_{\text{WIMP}} + \Omega_a) h^2 \leq 0.12,
\]

where the WIMP relic density is produced by the freeze-out mechanism, and the axion relic density can be produced by thermal or non-thermal mechanisms, such as the misalignment mechanism. For example, the DM direct detection experiments [44] constrain only the WIMP component of DM, while the axion detection experiments [45] constrain the axion component. On the other hand, the decreased abundance of the WIMP component of DM necessitates larger couplings for it to augment the annihilation cross section. This same couplings are involved in Lepton Flavor Violating (LFV) processes, such as the \( e \rightarrow \mu \gamma \) decay, which are strongly constrained. In this model, the additional interaction with \( \nu_R \) may be exploited to enhance the annihilation rate while keeping the LFV inducing couplings low. This can result in an increased production of right handed neutrinos in the early Universe, which may oversaturate the effective number of relativistic degrees of freedom, \( N_{\text{eff}} \) [46]. This may disfavour the fermionic DM case.

3. SUMMARY

We have discussed the DFSZ model where neutrinos are Dirac particles due to the PQ symmetry [20]. In this context, we propose two different scenarios to generate naturally small effective Yukawa coupling for neutrinos. In order to explain the smallness of the Yukawa coupling, the tree-level coupling is forbidden by the PQ symmetry while an effective dimension-5 operator with the
PQ field is allowed. This means that the Dirac neutrino mass is proportional to the PQ breaking scale. The first scenario is based on the Type-II Dirac seesaw where an extra heavy $SU(2)_L$ doublet allows Dirac neutrino mass when acquires a small vev once the PQ and the EW symmetry are broken. We summarize our results for this scenario in Fig. 2, considering the latest KATRIN [34] and the Planck data [35]. These constraints set limits on the PQ breaking scale $f_a$ and the mass of the heavy scalar.

We also discuss the UV completion of the dimension-5 operator at one-loop level. In this context once the PQ is broken, a residual $Z_2$ symmetry remains. The SM fermions are odd while the scalars are even, making the lightest field inside the loop stable by means of the residual $Z_2$ and Lorentz invariance giving a potential DM candidate. This residual symmetry is crucial, otherwise the the scalar particles inside the loop (odd under $Z_2$) acquire a vev, and the loop would be a correction to the type-I Dirac seesaw. Therefore, in this scenario we have a potentially rich phenomenology with two dark matter components, a stable WIMP running inside the neutrino mass loop and the axion. We have also discussed what are the parameters of such a dark sector in order to avoid an overclosed Universe. It is worth mentioning that in both cases, the UV scale is more relaxed that in the type-I case [20], where the KATRIN neutrino bound set it to be $\mathcal{O}(M_{GUT}) - \mathcal{O}(M_{PLANCK})$.

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Appendix A: An Alternate Loop Model

In this appendix we present an alternate loop model to generate Dirac neutrino mass. The $SM \otimes PQ$ charges of all the necessary fields are presented in Table IV.

| Symmetry/Fields | $L_i$ | $\nu_{Ri}$ | $H_u$ | $H_d$ | $N_{Ri}$ | $N_{Li}$ | $\sigma$ | $\eta_u$ | $\zeta$ |
|-----------------|-------|------------|-------|-------|----------|----------|--------|---------|-------|
| $SU(2)_L \times U(1)_Y$ | (2, -1/2) | (1, 1) | (2, -1/2) | (2, 1/2) | (1, 0) | (1, 0) | (1, 0) | (2, -1/2) | (1, 0) |
| $U(1)_{PQ}$ | 1 | -5 | 2 | 2 | 1/2 | 1/2 | 4 | 1/2 | 11/2 |

TABLE IV: Fields content and transformation properties under PQ symmetry in the one-loop mechanism.

In this case, the vevs of $\sigma$, $H_u$ and $H_d$ breaks the PQ symmetry into a $Z_4$ symmetry. The fields transforms are given in Table V. Here, the particles inside the loop are automatically stable [37].

We notice from the right panel of Fig. (4) that all the particles inside the loop carry $Z_4$ odd charges, whereas the SM particles are even under $Z_4$. Therefore, one can see that any combination
FIG. 4: Feynman diagram for Dirac neutrino masses in one-loop DFSZ scenario. Here left panel respects PQ charge assignment, whereas right panel respects remnant $Z_4^{PQ}$ charge assignment, arises due to PQ symmetry breaking.

| Symmetry/Fields | $L_i$ | $\nu_{Ri}$ | $H_u$ | $H_d$ | $N_R$ | $N_L$ | $\sigma$ | $\eta_u$ | $\zeta$ |
|-----------------|------|-----------|------|------|------|------|--------|--------|------|
| $SU(2)_L \times U(1)_Y$ | (2, -1/2) | (1, 1) | (2, -1/2) | (2, 1/2) | (1, 0) | (1, 0) | (1, 0) | (2, -1/2) | (1, 0) |

TABLE V: Fields content and transformation properties under PQ symmetry in the one-loop mechanism.

of SM fields will be even under the $Z_4$ charges. Further forbidding all effective operators to dark matter decay.

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