Significant neutrinoless double beta decay with quasi-Dirac neutrinos

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A significant signal of neutrinoless double beta decay can be consistent with the existence of light quasi-Dirac neutrinos. To demonstrate this possibility, we consider a realistic model where the neutrino masses and the neutrinoless double beta decay can be simultaneously generated after a Peccei-Quinn symmetry breaking.

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I. INTRODUCTION

The process of neutrinoless double beta decay [1] is motivated by the conjecture of Majorana neutrinos [2] because they both need a lepton number violation of two units. Specifically, the amplitude of the neutrinoless double beta decay is proportional to the Majorana mass of the electron neutrino. For a Majorana neutrino, its small mass [3] can be naturally understood by the seesaw mechanism [4]. There have been a lot of works (e.g. [14–19]) studying the standard scenario of the neutrinoless double beta decay. The electron neutrino can mediate a neutrinoless double beta decay even if it is a quasi-Dirac neutrino [20, 21]. However, the quasi-Dirac neutrino cannot induce a significant neutrinoless double beta decay since its Majorana mass is much smaller than the tiny neutrino mass itself. Alternatively, the neutrinoless double beta decay may not be directly governed by the Majorana neutrino masses [13, 22–28]. In this case, the lepton number violation for the neutrinoless double beta decay eventually will result in a Majorana neutrino mass term at loop level according to the Schechter-Valle theorem [22, 23, 30].

In this paper, we shall consider an interesting scenario where the quasi-Dirac neutrinos and the neutrinoless double beta decay are simultaneously generated after a Peccei-Quinn (PQ) symmetry breaking. In our model, the PQ symmetry breaking will lead to a lepton number violating interaction between two TeV-scale colored scalars. This lepton number violation can induce a neutrinoless double beta decay through the Yukawa couplings of the colored scalars to the right-handed up-type quarks, down-type quarks and charged leptons. The Majorana masses of the left-handed neutrinos can be produced at four-loop order. However, their magnitude is highly suppressed by the loop and chirality suppression factors. The loop-induced Majorana masses hence are too small to explain the neutrino oscillations [3]. Fortunately, the Dirac masses between the left- and right-handed neutrinos can arrive at the desired values as our model accommodates a variant seesaw mechanism [32].

II. THE MODEL

We extend the standard model (SM) gauge symmetry by a PQ global symmetry. The field content is summarized in Table 1 where the gauge-singlet fermions $\nu_R$, (right-handed neutrinos), the isodoublet scalar $\phi_v$, the isosinglet scalars $\delta, \omega$ and the gauge-singlet scalar $\chi$, respectively, carry the baryon and lepton numbers $0, (0, 2), (0, 1), \left(\frac{1}{2}, 0\right)$ and $(0, -2)$, the gauge-singlet scalar $\sigma$ and the isodoublet scalars $\phi_\mu, \phi_\tau$ have no baryon number or lepton number, while the SM quarks $q_L, u_R, d_R$, and leptons $l_L, e_R$ carry the usual baryon number $\frac{1}{3}$ or lepton number $1$. The global symmetry of baryon number is exactly conserved while the global symmetry of lepton number is softly broken. Under the imposed gauge and global symmetries, the allowed Yukawa interactions only include

$$L_Y = -y_{\nu_{ij}} \bar{q}_{L_i} \tilde{\nu}_{R_j} \phi_d + y_{d_{ij}} \bar{q}_{L_i} \phi_d \nu_{R_j} - y_{\nu_{ij}} \bar{L}_i \phi_d e_R,$$

As for the other gauge invariant Yukawa and mass terms,

$$L' = -h'^{ij}_d \delta \bar{q}_{L_i} t_2 \nu_{R_j} - h'^{\nu}_{ij} \delta \tilde{\nu}_{R_i} \nu_{R_j} - h'^{\nu}_{ij} \delta \tilde{q}_{L_i} \tau_2 \nu_{R_j} - \frac{1}{2} \delta y_{ij} \chi \bar{\nu}_{R_i} \nu_{R_j} + \frac{1}{2} \delta y_{ij} \chi \bar{\nu}_{R_i} \nu_{R_j} + \frac{1}{2} M_{ij} \bar{\nu}_{R_i} \nu_{R_j} + H.c.,$$

they are forbidden by the conservation of the baryon number, the lepton number or the PQ charge. We will not write down the full scalar potential. Instead, we only give the following terms,

$$V \supset \rho \sigma \phi^\dagger_0 \phi_d + \xi \chi \sigma^2 + \eta \chi \phi^\dagger_0 \phi_d + \kappa \chi \omega^2 + H.c.,$$

which are essential to our demonstration. Clearly, the $\xi$-term and $\eta$-term softly break the lepton number. Note that there are no other lepton number violating terms. For convenience and without loss of generality, we will take the parameters $\rho, \xi, \eta$ and $\kappa$ to be real by a proper phase rotation.

The scalars $\sigma$ and $\phi_d$ are responsible for the PQ symmetry breaking and the electroweak symmetry breaking,
The vacuum expectation values (VEVs) \( \langle \sigma \rangle \) and \( \langle \phi_u^0 \rangle \) can induce the VEVs of the scalars \( \phi_u \), \( \phi_d \) and \( \chi \), respectively. The vacuum expectation values (VEVs) \( \langle \sigma \rangle \) and \( \langle \phi_u^0 \rangle \) can induce the VEVs of the scalars \( \phi_u \), \( \phi_d \) and \( \chi \),

\[
\langle \phi_u^0 \rangle \simeq -\frac{\rho(\sigma)}{m_{\phi_u^0}},
\]

\[
\langle \chi \rangle \simeq -\frac{\xi(\sigma)^2}{m_{\chi}^2},
\]

\[
\langle \phi_u^0 \rangle \simeq -\frac{\eta(\chi)}{m_{\phi_u^0}^2},
\]

like the type-II seesaw \([2, 3] \). Clearly, our model can result in an invisible \([33, 36] \) axion \([31, 37, 38] \) to solve the strong CP problem. We should keep in mind the constraints on the VEVs \([3] \),

\[
10^9 \text{GeV} \lesssim \langle \sigma \rangle \lesssim 10^{12} \text{GeV},
\]

\[
\sqrt{(\langle \phi_u^0 \rangle)^2 + (\langle \phi_d^0 \rangle)^2 + (\langle \phi_{\nu}^0 \rangle)^2} \simeq 174 \text{GeV}.
\]

### III. QUASI-DIRAC NEUTRINOS

When the isodoublet scalars \( \phi_u \), \( \phi_d \) and \( \phi_{\nu} \) acquire their VEVs, the up-type quarks, the down-type quarks, the charged leptons and the neutrinos can obtain their Dirac masses,

\[
\mathcal{L} \supset -m_{u_{ij}} \bar{u}_i L_j - m_{d_{ij}} \bar{d}_i R_j - m_{e_{ij}} \bar{e}_i \nu_j - m_{\nu_{ij}} \bar{\nu}_i \nu_j + \text{H.c.}
\]

with

\[
m_{u_{ij}} = y_{u_{ij}} \langle \phi_u^0 \rangle, \quad m_{d_{ij}} = y_{d_{ij}} \langle \phi_d^0 \rangle, \quad m_{e_{ij}} = y_{e_{ij}} \langle \phi_{\nu}^0 \rangle, \quad m_{\nu_{ij}} = y_{\nu_{ij}} \langle \phi_{\nu}^0 \rangle.
\]

To generate the known charged fermion masses \([3] \), we should perform

\[
\langle \phi_u^0 \rangle = \mathcal{O}(100 \text{ GeV}),
\]

\[
\langle \phi_d^0 \rangle = \mathcal{O}(1-10 \text{ GeV}) \left[ \frac{\langle \phi_u^0 \rangle}{\mathcal{O}(100 \text{ GeV})} \right] \frac{\rho}{\mathcal{O}(0.1 m_{\phi_u^0})} \left[ \frac{m_{\phi_u^0}}{\mathcal{O}(\langle \sigma \rangle)} \right]^{-2} \frac{\langle \sigma \rangle}{\mathcal{O}(10^{12} \text{ GeV})}.
\]

In the presence of the \( \kappa \)-term in the potential \([3] \), the VEV \( \langle \chi \rangle \) given by Eq. \([3] \) will result in a lepton number violating interaction between the colored scalars \( \delta \) and \( \omega \),

\[
\mathcal{L} \supset -\mu \omega \delta^2 + \text{H.c.} \text{ with } \mu = \kappa(\chi).
\]

Remarkably, the VEV \( \langle \chi \rangle \) can be close to the TeV scale for a nice parameter choice,

\[
\langle \chi \rangle \simeq \mathcal{O}(10^5 \text{ GeV}) \left[ \frac{\langle \sigma \rangle}{\mathcal{O}(10^{12} \text{ GeV})} \right]^2 \frac{\xi}{\mathcal{O}(0.1 m_{\chi})} \frac{m_{\chi}}{\mathcal{O}(m_{\text{Pl}})} \left[ \frac{m_{\chi}}{\mathcal{O}(\langle \sigma \rangle)} \right]^{-2} \frac{\langle \sigma \rangle}{\mathcal{O}(10^{12} \text{ GeV})}.
\]

Here \( m_{\text{Pl}} \simeq 2.43 \times 10^{18} \text{ GeV} \) is the reduced Planck mass. We then can immediately read

\[
\mu = \mathcal{O}(10 \text{ TeV}) \quad \text{for} \quad \kappa = \mathcal{O}(0.1).
\]

Due to the lepton number violation \([19] \), the left-handed neutrinos can obtain their Majorana masses through the
four-loop diagrams as shown in Fig. 1. We can rapidly estimate the loop-induced Majorana masses as below,

\[
\mathcal{L} \supset -\frac{1}{2} \delta m_{\nu_{ij}} \bar{\nu}_L \nu_R + \text{H.c.},
\]

(19)

where

\[
\delta m_{\nu_{ij}} \simeq \frac{g^4}{(16\pi^2)^4} \left( \frac{g}{\sqrt{2}} \right)^4 \mu \sum_{abcdefgh} f_{gh} h_{ca} h_{db} \times (m_{\nu_{ia} m_{\nu_{jib}} + m_{\nu_{ijb} m_{\nu_{ia}}} m_{u_{iab}} m_{d_{ie}} m_{d_{ie}}) \delta_{h,f},
\]

(20)

with \( g \) being the \( SU(2)_L \) gauge coupling. It is easy to find that the loop-induced Majorana masses should have an upper bound,

\[
\delta m_{\nu_{ij}} \lesssim \frac{g^4}{2^{15/2}} \frac{\mu^2 m_R^2 m_{\nu_{Rb}}}{m_{m_{\delta}}^3}
\]

(21)

which yields

\[
\delta m_{\nu_{ij}} \lesssim \mathcal{O}(10^{-8} \text{eV}) \left( \frac{\mu}{\mathcal{O}(10 \text{ TeV})} \right) \left( \frac{m_{\nu_{\delta}}}{\mathcal{O}(\text{TeV})} \right)^2 \times \left( \frac{m_{\nu_{\delta}}}{\mathcal{O}(\text{TeV})} \right)^{-2}.
\]

(22)

In the above numerical estimation we have taken into account \( g = 0.653, m_\tau = 1.78 \text{ GeV}, m_\chi = 171.2 \text{ GeV} \) and \( m_\nu = 4.20 \text{ GeV} \). Therefore, the Majorana masses of the left-handed neutrinos are far below the values to explain the phenomenon of neutrino oscillations.

We then check the Dirac masses between the left- and right-handed neutrinos. It is striking that the Dirac neutrino masses can arrive at the desired level without fine tuning the Yukawa couplings \( y_{\nu_{ij}} \) because the VEV \( \langle \phi^0 \rangle \) given by Eq. (16) is just

\[
\langle \phi^0 \rangle \simeq \mathcal{O}(0.01 - 0.1 \text{eV}) \left( \frac{\langle \phi^0 \rangle}{\mathcal{O}(100 \text{ GeV})} \right) \left( \frac{\langle \chi \rangle}{\mathcal{O}(10^5 \text{ GeV})} \right) \times \left( \frac{\eta}{\mathcal{O}(m_\phi^0)} \right) \left( \frac{m_\phi^0}{\mathcal{O}(m_{\delta}^i)} \right)^{-2}.
\]

(23)

Obviously, the generation of Dirac neutrino masses retains the essence of the conventional type-II seesaw, as shown in Fig. 2.

We now can conclude that the neutrinos in our model are quasi-Dirac particles.

IV. NEUTRINOLESS DOUBLE BETA DECAY

We can integrate out the heavy colored scalars \( \delta \) and \( \omega \) from Eqs. (11) and (16). The leading dimension-9 op-
generator should be
\[
\mathcal{O}_9 = \frac{\mu}{m_3^2 m_\psi^3} f_{ij} h_{kl} h_{mn} \bar{R}_i^* R_j^* d_R^* e_R^* \epsilon_{R_l^* R_m^* R_n} + \text{H.c.},
\]
which can result in a neutrinoless double beta decay as shown in Fig. 3. With the cubic coupling \( \mu = \mathcal{O}(10\,\text{TeV}) \) and the Yukawa couplings \( h_{11, ij} \ll \mathcal{O}(1) \), the induced neutrinoless double beta decay can be verified by the ongoing and planned experiments [30] if the colored scalars are at the TeV scale, i.e. \( m_\delta, m_\omega = \mathcal{O}(\text{TeV}) \). Note that we have the flexibility to avoid the stringent constraints from other rare processes [3] by choosing the values of the Yukawa couplings \( f_{ij}, h_{ij} \neq (1, 1) \). With the TeV-scale colored scalars, we can also expect to check our scenario at colliders such as the LHC.

V. CONCLUSION

In this work, we have demonstrated a scenario where the neutrinoless double beta decay can be observed by the forthcoming experiments within the next few years even if the light neutrinos are quasi-Dirac particles. In our model, a spontaneous PQ symmetry breaking will induce a lepton number violating interaction between two colored scalars. Because the colored scalars do not have any Yukawa interactions involving the left-handed fermions or the right-handed neutrinos, the lepton number violation can significantly result in the neutrinoless double beta decay at tree level while can negligibly contribute to the left-handed Majorana neutrino mass at four-loop order. Through a variant seesaw mechanism, the quasi-Dirac neutrinos can naturally obtain the desired masses to explain the neutrino oscillations.

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