Neutrino Masses from Loop-induced $d \geq 7$ Operators

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We propose a new scenario where neutrino masses are generated via operators with the mass dimension higher than five, which are induced at the loop level. The scenario is demonstrated with concrete models where neutrino masses are generated via a one-loop dimension-seven operator which is induced through TeV scale dynamics under the exact $Z_2$ symmetry. Tiny neutrino masses are naturally induced from the TeV scale dynamics without introducing any artificial assumption on magnitudes of coupling constants. The combination of one-loop factor $1/(4\pi)^2$ and the factor of the ratio $(v/\Lambda)^2$ between the electroweak scale $v$ and new physics scale $\Lambda$ provides sufficient suppression as compared to the model based on the dimension-five operator induced at the tree level. The reproduction of the data for neutrino masses and mixings are discussed under the constraint from experiments for lepton flavour violation. We also mention phenomenological implications at collider experiments and dark matter candidates.

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I. INTRODUCTION

Mystery is the origin of tiny neutrino masses that are indicated from the neutrino oscillation data. How can we understand the smallness of neutrino masses as compared to the electroweak scale? A simple way of the explanation may be based on the seesaw mechanism [1–4], introducing right-handed neutrinos with large Majorana masses at the scale such as that of grand unification. Although this is an attractive scenario, introduction of such large masses causes another hierarchy among mass scales. In addition, such a large mass scale is beyond the experimental reach and the theory would be untestable directly.

In the Standard Model (SM), the Higgs sector, on the other hand, is the last uncharted part. Although the SM Higgs sector is the simplest scenario with a scalar isospin doublet, the true Higgs sector may take a non-minimal form. Such an extended Higgs sector may be closely related to the mechanism to induce tiny neutrino masses at the TeV scale. Such a possibility is interesting because the model is in principle testable directly at on-going and future collider experiments, such as the Fermilab Tevatron, the CERN Large Hadron Collider (LHC) and the International Linear Collider (ILC).

If neutrino masses are of the Majorana type, they are generated through the lepton number violating effective operators. In the usual seesaw scenarios, the neutrino masses are derived from the dimension-five operator $\nu\nu\phi\phi/\Lambda$, where $\nu$ represents left-handed neutrinos, $\phi$ does the Higgs boson, and $\Lambda$ is a scale of the new physics. In a class of models where neutrino masses are radiatively generated, such a dimension-five operator is induced at the loop level by the TeV scale dynamics. For example, in the model proposed by A. Zee [5, 6], the dimension-five operator is generated at the one-loop level via the lepton number violating interaction and dynamics of the extended Higgs sector. In the model proposed by E. Ma [7], the dimension-five operator is also generated at the one-loop level via the physics of the extra scalar doublet and the TeV scale right-handed neutrinos, where the both of new fields are assigned odd quantum number under the discrete $Z_2$ symmetry. Such a one-loop generation of neutrino masses from the TeV scale dynamics, however, still requires unnaturally small coupling constants for reproducing the tiny neutrino masses. There are several models in which neutrino masses are generated at the two-loop level [8–11] and also the three-loop level [12, 13], where such fine tuning is not necessary because of the sufficient suppression by additional loop factors. In all these models, dimension-five operators are induced at the loop level.

Recently, a new idea has been proposed where tiny neutrino masses are generated via the operators whose dimension is higher than five [16–20]. In Ref. [18], some concrete examples are examined, in which neutrino masses are generated via the dimension-seven operator $\nu\nu\phi\phi\phi\phi/\Lambda^3$ which are induced at the tree level with the extend scalar dynamics. In this case, there is an additional suppression factor of $(v/\Lambda)^2$ as compared to neutrino masses via the dimension-five operators, where $v$ ($\sim 246$ GeV) is the vacuum expectation value (vev) of the Higgs boson. Although these models are interesting, a sort of fine tuning is still required especially to reproduce the scale of neutrino masses, when $\Lambda$ is assumed to be of TeV scale.

In this paper, we propose a scenario in which neutrino masses are generated via higher-dimensional operators $\nu\nu(\phi\phi)^{(d-3)/2}/\Lambda^{d-4}$ ($d = 7, 9, 11 \cdots$) which are induced by quantum effect. In general, the size of neutrino masses from the operator with the mass dimension
Before we come on to descriptions of the concrete models, let us look briefly at the essentials for the tree-level dimension-seven neutrino mass generation \[18\]. There are two key components to produce the effective dimension-seven operator for neutrino masses at the electroweak scale:

- An additional symmetry to forbid the dimension-five $\nu\nu\phi\phi/\Lambda$ operator. The simplest choice for non-supersymmetric models is $Z_5$.
- The extended Higgs sector with two Higgs doublets so that the combination $(H_1H_2)$ can carry a charge under the additional symmetry. Here the hypercharge of $H_1$ is given to be $-1/2$ and that of $H_2$ is $+1/2$.

Taking the setups and assigning appropriate charges to the standard model particles, we can forbid the dimension-five operator and make

$$\mathcal{L}_{\text{eff}} = \frac{C}{\Lambda^3}LLH_2H_2H_1$$

(2)

to be the leading contribution to neutrino masses, where $C$ is a mass dimensionless coefficient\(^1\). The possible models for this tree-level dimension-seven neutrino mass generation mechanism are listed in Ref. \[18\]. In the models including the SM singlet fermions at the high energy scale, one can see that the $Z_5$ symmetry forbids the fermions (=right-handed neutrinos) to have the Majorana mass term. Because of the absence of the Majorana mass term, the dimension-five operator cannot arise at the electroweak scale and the dimension-seven operator dominates the contribution to neutrino masses.

Extending these models, we consider the models, in which neutrino masses are generated via the dimension-seven operator but the effective operator is induced

\(^1\) The choice of dimension-seven operators which contribute to neutrino masses is not unique \[18\]. In this paper, we concentrate on the operator shown in Eq. (2).
through a one-loop diagram. To construct such loop-induced models, we follow the method with the exact $Z_2$ symmetry, which was developed in Refs. [2, 12]. The essential is introduction of the inert doublet with the odd parity under the $Z_2$ symmetry. Due to the exact $Z_2$ parity, the inert doublet cannot take a vacuum expectation value (vev). Let us describe the maneuver, taking the ordinary type-I seesaw model as an example. The procedure is schematically illustrated in Fig. 1. Assigning the odd parity to the right-handed neutrinos $N_R$ and substituting the inert doublet $\eta$ for the Higgs doublet $H_2$ in the neutrino Yukawa interaction, one can forbid the tree-level contribution to neutrino masses. The inert doublets in the diagram Fig. 1 are converted to the Higgs doublets through a quartic interaction,

$$\mathcal{L} = \frac{\lambda}{2} (\eta^\dagger H_2)(\eta H_2) + \text{H.c.}$$

(3)

In other words, the inert Higgs legs are closed by the quartic interaction and make a loop. This leads to the one-loop diagram for neutrino masses, which was proposed in Ref. [2]. In the following sections, we will apply this procedure to the models in which neutrino masses are generated through the dimension-seven operator which is induced via tree diagrams, and build the loop-induced dimension-seven models.

III. MODELS

We here consider two examples to illustrate the method to build the models in which neutrino masses are generated through the effective dimension-seven operator induced from a one-loop diagram.

A. Model A

The renormalizable models to induce the effective interaction Eq. (2) from tree diagrams at the electroweak scale are listed in Ref. [18]. In this subsection, we employ the model described as Decomposition #1 among them, in which the SM gauge singlet Dirac fermion $\psi$ and the singlet scalar $\varphi$ are introduced. The particle contents and the charge assignments are summarized in Table I. Here, the charges for the quarks and leptons are assigned so as to reproduce the Yukawa interactions of type-II two-Higgs-doublet model (THDM)\(^2\). For detailed arguments for this model and the (softly broken) $Z_5$ symmetry, see Sec. 3.1 in Ref. [18]. In this letter, we are interested in the neutrino masses induced from the dimension-seven operator Eq. (2) but the effective interaction is realized by a loop diagram. In order to forbid arising the dimension-seven operator from a tree diagram, we introduce the exact $Z_2$ parity and an inert doublet $\eta$, and assign the odd charge for the inert doublet and the singlet Dirac fermion $\psi$.

The Lagrangian of the fundamental interactions for the neutrino mass generation is given as

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \left[ (Y_\nu)_a \alpha \bar{\psi}_a P_L (\eta \tau^2 L_a) + (\kappa_L)^{ab} \bar{\varphi}_a P_L \psi_b H_2 \right] + \left[ (\kappa_R)^{ab} \varphi \psi_a P_R \psi_b H_2 + \mu \varphi^2 (H_1 \tau^2 H_2) + \text{H.c.} \right] + M_a \bar{\psi}_a \psi_b + m^2 \varphi^2 \varphi + m_\eta \eta^\dagger \eta + \left[ \frac{\lambda}{2} (\eta^\dagger H_2)(\eta H_2) + \text{H.c.} \right] - \mathcal{V}_{\text{scalar}},$$

(4)

where $a, b$ and $\alpha$ represent the flavour indices. Let us first focus on neutrino masses which are our main concern, and we will take up some phenomenological consequences of this model and the part $\mathcal{V}_{\text{scalar}}$ of the scalar potential later. With the interactions shown in Eq. (4), the dimension-seven operator for neutrino masses is induced by the one-loop diagram described in Fig. 2 which is evaluated as

$$\mathcal{L}_{\text{eff}} = \frac{1}{(4\pi)^2} \frac{\lambda \mu}{m_\eta^2 m^2_0} (Y_\nu)^{\alpha}_{\nu}\bar{\alpha}_a (Y_\nu)^{\beta}_{\nu} \times \left[ (\kappa_L)^{ab} \frac{M_a M_b}{m^2_0} \mathcal{I}(x_a, x_b) + (\kappa_R)^{ab} \mathcal{J}(x_a, x_b) \right] \times (\tau^2 H_2)(H_1 \tau^2 H_2),$$

(5)

where the functions $\mathcal{I}$ and $\mathcal{J}$ are defined as

$$\mathcal{I}(x_a, x_b) = \frac{1}{(1 - x_a)(1 - x_b)} \times \left[ 1 + \frac{(1 - x_b)x_a \ln x_a}{x_a - x_b} - \frac{(1 - x_b)x_b \ln x_b}{x_a - x_b} \right],$$

(6)

$$\mathcal{J}(x_a, x_b) = \frac{1}{(1 - x_a)(1 - x_b)} \times \left[ 1 + \frac{(1 - x_b)x_a^2 \ln x_a}{x_a - x_b} - \frac{(1 - x_b)x_b^2 \ln x_b}{x_a - x_b} \right],$$

(7)

with $x_a \equiv M_a^2/m^2_\eta$. We obtain neutrino masses

| $L$ e$^+ Q u^c d^c H_2 H_1 \psi(1^0)\eta(2^{1/2})\varphi(1^{3/2})$ | $Z_5$ | $1$ | $1$ | $0$ | $2$ | $0$ | $3$ | $1$ | $0$ | $3$ |
|---|---|---|---|---|---|---|---|---|---|---|
| exact $Z_2$ | $+$ | $+$ | $+$ | $+$ | $+$ | $+$ | $-$ | $-$ | $-$ | $+$ |

TABLE I: Particle contents and charge assignments for Model A. The symbol $X^5$ indicates the representations of the fields; $X$ for $SU(2)_L$, $Y$ for $U(1)_Y$, and $L$ for Lorenz group; i.e., Dirac spinor ($D$) and scalar ($s$).
broken discrete symmetry in this model. With the exact accessible to those fields.

One can expect that the collider experiments are fields without assuming extremely tiny couplings in the of one eV is compatible with TeV scale masses for new

dimensional type-I seesaw model. From the expressions Eqs. (8) produce all the features of the neutrino flavour in the canon-

izing contribution to the lepton flavour violating processes arises from a diagram with a loop between two Yukawa interactions. The contribution is exactly the same as that in the original dark doublet model \[ \eta \], which was calculated in Ref. \[ 32 \]:

\[ \text{Br}(\mu \rightarrow e\gamma) = \frac{3\alpha_{\text{em}}}{64\pi(G_Fm_\eta^2)^2} |(C_\lambda)_{e\mu}|^2, \]

where the mass-dimensionless coefficient \( C_\lambda \) for Model A is given as

\[ (C_\lambda)_{e\mu} = (Y_e^\dagger)_a a_F(x_a)(Y_\nu)_a^{\mu}, \]

and the function \( F \) is

\[ F(x_a) = \frac{1 - 6x_a + 3x_a^2 + 2x_a^3 - 6x_a^2 \ln x_a}{6(1 - x_a)^4}. \]

We can see that it might be essential to assume a large value for \( m_\eta \) enough to avoid a sizable \( \mu \rightarrow e\gamma \) effect. An alternative way to circumvent the large LFV process is discussed in Ref. \[ 33 \].

**B. Model B**

Let us show the second example with the different type of Decomposition (\# 13). We introduce two insert doublets. This allows to have two types of Yukawa interactions for neutrinos: one is the ordinary one with right-handed neutrinos \( \psi_R \), and the other appears with left-handed component \( \psi_L \) of the SM singlet fermion and violates the lepton number. The particle contents and

| \( L \) | \( e^+ \) | \( Q \) | \( u^c \) | \( d^c \) | \( H_2 \) | \( H_1 \) | \( \psi(1) \) | \( \eta(2_{1/2}) \) | \( \eta(2_{-1/2}) \) |
|-----|-----|-----|-----|-----|-----|-----|-----|-----|-----|
| soft br. | \( Z_5 \) | 1 | 1 | 0 | 0 | 2 | 0 | 3 | 1 | 0 | 2 |
| exact | \( Z_2 \) | + | + | + | + | + | + | + | + | + | + |

| TABLE II: Particle contents and charge assignments for the softly broken \( Z_5 \) and the exact \( Z_2 \) in Model B. |
their charge assignments are summarized in Tab. II. The interaction is given by

$$
\mathcal{L} = \mathcal{L}_{\text{SM}} + \left[ (Y_\nu)_{\alpha} \overline{\psi} \gamma^\mu P_L \eta \gamma^2 L_\alpha + (Y_{\nu}^\prime)_{\alpha} \overline{\psi} \gamma^\mu P_L \eta \gamma^1 L_\alpha \\
+ \zeta (H_1 \gamma^2 H_2)(\eta \gamma^2 \eta) + \frac{\lambda}{2} (\eta \gamma^2 H_2)(\eta \gamma^1 H_2) \\
+ \text{H.c.} \right] \\
+ M_{\alpha} \overline{\psi} \gamma^\mu \psi_b + m_\eta^2 \eta \gamma^1 \eta + m_\eta^2 \eta \eta \eta - \mathcal{V}_{\text{scalar}}. 
$$

The scalar potential is obviously different from that of Model A. However, we assume also that it includes the soft violation term of the $Z_5$ symmetry, which was shown in Eq. (11), to avoid the problem of the Nambu-Goldstone boson.

With the Lagrangian in Eq. (15), neutrino masses are constructed as shown in Fig. 3, and they are calculated to be

$$
(m_\nu)^{\alpha\beta} = - \frac{\lambda \zeta y^4}{8} \sin^2 \beta \sin 2\beta \\
\times \left[ (Y_\nu^T)^{\alpha a} M_{\alpha} \mathcal{I}(x_\alpha, y)(Y_\nu)^{a \beta} \\
+ (Y_{\nu}^T)^{\alpha a} M_{\alpha} \mathcal{I}(x_\alpha, y)(Y_{\nu}^\prime)^{a \beta} \right],
$$

where $y \equiv m_\eta^2/m_\eta^2$. The flavour structure of this model is rather involved, and it cannot be understood with the ordinary seesaw formula because of two independent Yukawa matrices $Y_\nu$ and $Y_{\nu}^\prime$. With the assumption that $Y_{\nu}^\prime$ takes the same flavour structure as $Y_\nu$, Eq. (16) is reduced to the ordinary type-I seesaw formula. Therefore, this model can obviously reproduce the mass matrices which are consistent with the observed mass squared differences and the mixings.

The lepton number violating Yukawa interaction gives an additional contribution to the LFV process $\ell_\alpha \rightarrow \ell_\beta \gamma$. The decay branching ratio in Model B can be obtained by substituting

$$(C_{\beta})_{e\mu} = (Y_{\nu}^\prime)^{\alpha e} \mathcal{F}(x_\alpha, y)(Y_\nu)^{\alpha \mu} + (Y_{\nu}^\prime)^{\alpha e} \mathcal{F}(x_\alpha, y)(Y_{\nu}^\prime)^{\alpha \mu},$$

for $C_A$ in Eq. (12).

### IV. SUMMARY AND DISCUSSION

We have proposed the new scenario in which tiny neutrino masses are generated via loop-induced $d > 5$ operators. In such a scenario, the scale of tiny neutrino masses can be reproduced from the TeV scale physics in a natural way without extreme fine tuning because the combination of the loop factor $1/(16\pi^2)^n$ and the additional coefficient of $(v/\Lambda)^d-5$ provides the sufficient suppression factor. We have in particular discussed as examples two concrete models where neutrino masses are generated via one-loop induced dimension-seven operators due to the dynamics of extended Higgs sector and a vector-like Dirac neutrino whose mass is assumed to be at the TeV scale under the imposed exact discrete $Z_2$ symmetry. We have shown that in these models neutrino masses can be reproduced and that the neutrino mixing data are also satisfied without contradicting the constraint from the LFV data $^{22,23}$.

We here give a comment on phenomenological implications in these models. However, the detailed discussion is beyond the scope of this paper, and it is given elsewhere $^{21}$. First of all, a common feature of these models is the extended Higgs sectors, in which there are two $Z_2$-even Higgs doublets and one or two $Z_2$-odd doublets. Phenomenology of the THDM has been discussed in literature. The Higgs potential is constrained by the perturbative unitarity $^{35,38}$, the vacuum stability $^{39,41}$, and also electroweak precision data $^{42,44}$. When the type-II THDM is assumed, the bounds from $b \rightarrow s\nu \nu$ $^{43}$, $B \rightarrow \tau\nu$ $^{46,48}$ and the leptonic tau decay $^{49}$ have also to be taken care. The discovery of extra Higgs bosons in addition to the lightest (SM-like) Higgs boson and the measurement of their properties are important to test these models. In these models, the induced neutrino masses are multiplied by the factor of $\sin^2 \beta \sin 2\beta$, so that a large value of $\tan \beta$ gives a further suppression factor. This may bring an interesting correlation between neutrino masses and the physics of the Higgs sector.

The experimental confirmation of the $Z_2$ odd sector is essentially important too. Especially, the lightest $Z_2$ odd particle can be a candidate of dark matter if it is electrically (and colour) neutral. In these models, there are two possibilities for the DM candidate; i.e., 1) the lightest $\eta^0$ boson is the DM or 2) the Dirac neutrino $\nu$ is the DM.

In Case 1), the phenomenology of such $Z_2$ odd sector has been studied with the physics of the DM candidate in the context of the dark doublet model $^{54}$ and the radiative seesaw model $^{6,12,51}$. An interesting signature of DM may be the invisible decay of the (SM-like) Higgs boson when DM is lighter than a half of the Higgs boson mass. It is expected that the branching ratio of the
Higgs boson invisible decay of greater than 50% (1%) can be detected at the LHC (at the ILC). The direct DM search is also important for the case of 1). The multi Higgs portal dark matter has been discussed in Ref. [52]. The detailed comprehensive study for models of the Higgs portal dark matter has been done in Ref. [52] in a specific scenario where only the Higgs boson and the DM candidate are electroweak scale and the other new particles are supposed to be decoupled. The collider phenomenology of the Higgs sectors with dark doublet fields has been supposed to be decoupled. The collider phenomenology has been discussed in Ref. [52]. For the test of our model, many parts of these previous studies can be applied.

If $\psi$ is dark matter corresponding to the case of 2), the situation may be similar to the case in the model by Ma where the right-handed neutrino is the DM candidate which has been studied in detail in Ref. [23]. However, $\psi$ is a Dirac neutrino, not a Majorana neutrino, so that the DM number can be assigned. The DM number may be dynamically generated in the context of asymmetric DM. Details of these issues are discussed in Ref. [54].

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