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Probing the Majorana nature of TeV-scale radiative seesaw models at collider experiments

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ABSTRACT

A general feature of TeV-scale radiative seesaw models, in which tiny neutrino masses are generated via loop corrections, is an extended scalar (Higgs) sector. Another feature is the Majorana nature; e.g., introducing right-handed neutrinos with TeV-scale Majorana masses under the discrete symmetry, or otherwise introducing some lepton number violating interactions in the scalar sector. We study phenomenological aspects of these models at collider experiments. We find that, while properties of the extended Higgs sector of these models can be explored to some extent, the Majorana nature of the models can also be tested directly at the International Linear Collider via the electron–positron and electron–electron collision experiments.

1. Introduction

The neutrino data show that neutrinos have tiny masses as compared to the electroweak scale. This is clear evidence for physics beyond the standard model (SM). The data also indicate that the structure of flavor mixing for neutrinos is largely different from that for charged leptons. These facts would suggest that, while charged leptons have Dirac type masses, the neutrino masses are of the Majorana type. The tiny Majorana masses of left-handed neutrinos are generated from the dimension five effective operators

$$\mathcal{L} = \frac{c_{ij}}{2\Lambda}\overline{\nu^c}^i_L \nu^j_L \phi^0\phi^0, \tag{1}$$

where Λ represents a mass scale, c_{ij} are dimensionless coefficients, and ϕ^0 is the Higgs boson. After electroweak symmetry breaking, the mass matrix M_ν^{ij} for left-handed neutrinos appears as $M_\nu^{ij} = c_{ij}\langle\phi^0\rangle^2/\Lambda$. As the vacuum expectation value (VEV) $\langle\phi^0\rangle$ of the Higgs boson is $\mathcal{O}(100)$ GeV, the observed tiny neutrino masses ($M_\nu^{ij} \lesssim 0.1$ eV) are realized when $(c_{ij}/\Lambda) \sim \mathcal{O}(10^{-14})$ GeV^{-1}. It has been an interesting problem how we can naturally explain such a small number with less fine tuning.

If the operators in Eq. (1) appear at the tree level in the low energy effective theory, Λ has to be as large as $\mathcal{O}(10^8)$–$\mathcal{O}(10^{14})$ GeV for c_{ij} being $\mathcal{O}(10^{-6})$–$\mathcal{O}(1)$ to describe the data. For example, in the tree-level seesaw scenario where right-handed neutrinos are introduced, their Majorana masses have to be set much higher than the electroweak scale [1], corresponding to the scale Λ in Eq. (1). Although the scenario is simple, it requires another hierarchy between the mass of right-handed neutrinos and the electroweak scale, and in addition, physics at such a large mass scale is difficult to be tested at collider experiments.

Quantum generation of neutrino masses is an alternative way to obtain $(c_{ij}/\Lambda) \sim \mathcal{O}(10^{-14})$ GeV^{-1}. Due to the loop suppression factor, Λ in these models can be lower as compared to that in the tree-level seesaw models. Consequently, the tiny neutrino masses would be explained in a natural way by the TeV-scale dynamics without introducing very high mass scales. The original model of this line was proposed by Zee [2], where neutrino masses were generated at the one-loop level. Some variations were considered [3–7], for example, by Zee and Babu [3], Krauss, Nasri and Trodden [4], Ma [5], and the model in Ref. [6]. The last three models contain dark matter (DM) candidates with the odd quantum number under the discrete Z_2 symmetry. It must be a charming point in these TeV-scale radiative seesaw models that they are directly testable at the collider experiments such as Large Hadron Collider (LHC) and the International Linear Collider (ILC).

A general feature in radiative seesaw models is an extended Higgs sector, whose detail is strongly model dependent. The discovery of these extra Higgs bosons and detailed measurements of their properties at current and future collider experiments can give partial evidence for the radiative seesaw models. In the literature [8–14], phenomenology of these radiative seesaw models

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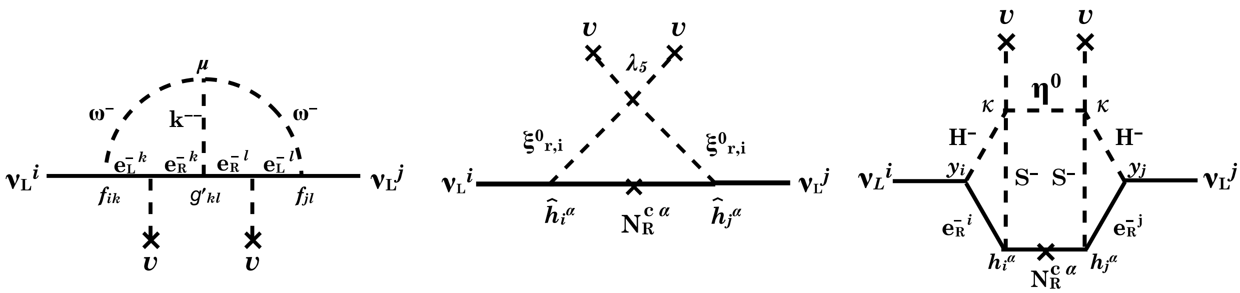

Fig. 1. Feynman diagrams for neutrino masses in the model by Zee and Babu [3] (left), that by Ma [5] (center) and that in Ref. [6] (right).

has already been studied extensively. Such previous works mainly discuss constraints on the flavor structure from the current data for such as neutrino physics and DM, and also study collider phenomenology of the Higgs sectors [15–17,19–24].

Another common feature in radiative seesaw models is the Majorana nature. In order to induce tiny Majorana masses for left-handed neutrinos, we need to introduce its origin such as lepton number violating interactions in the scalar sector [2,3] or right-handed neutrinos with TeV-scale Majorana masses [4–6]. When the future data would indicate an extended Higgs sector predicted by a specific radiative seesaw model, the direct detection of the Majorana property at collider experiments should be a fatal probe to identify the model.

In this Letter, we study the phenomenology in TeV-scale radiative seesaw models, in particular, a possibility of detecting the Majorana nature at collider experiments. We mainly discuss three typical radiative seesaw models as reference models; the model by Zee and Babu where neutrino masses are generated at the two-loop level [3], that by Ma with one-loop neutrino mass generation [5], and that in Ref. [6] where neutrino masses are generated at the three-loop level. Typical parameter regions where the data can be satisfied have been already studied in each model in the literature. We here study collider phenomenology in such typical parameter regions in each model, and discuss the discrimination of these models by measuring the details of the Higgs sector and the Majorana nature at the LHC and the ILC.

2. Radiative seesaw models

2.1. The Zee–Babu model

In the model proposed in Ref. [3] (we refer to as the Zee–Babu model), in addition to singly-charged singlet scalar bosons $\omega^\pm$, doubly-charged singlet scalar fields $k^{\pm\pm}$ are introduced, both of which carry the lepton number of two-unit, and their interactions are given by

$$\mathcal{L}_{\rm int} = f_{ab}\big(\overline{L^{ci}_{aL}}L^j_{bL}\big)\epsilon_{ij}\omega^+ + g'_{ab}\big(\overline{\ell^c_{aR}}\ell_{bR}\big)k^{++} - \mu k^{++}\omega^-\omega^- + \text{H.c.}, \qquad (2)$$

where L_L is the left-handed lepton doublet and ℓ_R is the right-handed lepton singlet. The matrices f_{ij} and g'_{ab} are respectively an anti-symmetric and a symmetric couplings and the lepton number is violated by the interaction with the parameter μ.

The neutrino mass matrix is generated at the two-loop level via the diagram in Fig. 1 (left);

$$M^\nu_{ij} = \sum_{k,\ell=1}^{3}\left(\frac{1}{16\pi^2}\right)^2\frac{4\mu}{m^2_\omega} f_{ik}(y_{\ell_k} g_{k\ell} y_{\ell_\ell}) f_{\ell j} v^2 I_1\big(m^2_k/m^2_\omega\big), \qquad (3)$$

where y_i $[=\sqrt{2}m_i/v\ (i = e,\ \mu,\ \tau)]$ are the SM Yukawa coupling constants of charged leptons with the masses m_i and the VEV v ($\simeq 246$ GeV), g_{ij} are defined as $g_{ii} = g'_{ii}$ and $g_{ij} = 2g'_{ij}$ $(i \neq j)$, m_ω and m_k are masses of $\omega^\pm$ and $k^{\pm\pm}$, and

$$I_1(r) = -\int_0^1 dx \int_0^{1-x} dy \frac{1}{x + (r-1)y + y^2} \ln\frac{y(1-y)}{x + ry}, \qquad (4)$$

where $I_1(r)$ takes the value of around 3–0.2 for $10^{-2} \lesssim r \lesssim 10^2$. The universal scale of neutrino masses is determined by the two-loop suppression factor $1/(16\pi^2)^2$ and the lepton number violating parameter μ. The charged lepton Yukawa coupling constants y_{ℓ_i} ($y_e \ll y_\mu \ll y_\tau \lesssim 10^{-2}$) give an additional suppression factor. Thus, any of f_{ij} or g_{ij} can be of $\mathcal{O}(1)$ when m_ω and m_k are at the TeV scale. The flavor structure of the mass matrix is determined by the combination of the coupling constants f_{ij} and $y_i g_{ij} y_j$.

The flavor off-diagonal coupling constants f_{ij} and g_{ij} induce lepton flavor violation (LFV). From the results of $\mu \to e\gamma$, $\tau \to e\gamma$ and $\tau \to \mu\gamma$, $|f_{\mu\tau}f_{\tau e}|$, $|f_{\tau\mu}f_{\mu e}|$ and $|f_{\tau e}f_{e\mu}|$ are respectively constrained as a function of m_ω. The data of rare decays of $\mu \to eee$, $\tau \to \mu\mu e$ and $\tau \to \mu ee$ are also used to constrain the combinations $|g_{\mu e}g_{ee}|$, $|g_{\tau e}g_{\mu\mu}| + |g_{\tau\mu}g_{\mu e}|$ and $|g_{\tau e}g_{\mu e}| + |g_{\tau\mu}g_{ee}|$, respectively, depending on m_k. The $g-2$ data can also be used to constrain a combination of these coupling constants with m_ω and m_k.[1]

In the scenario with hierarchical neutrino masses, f_{ij} satisfy $f_{e\mu} \simeq f_{e\tau} \simeq f_{\mu\tau}/2$. The typical relative magnitudes among the coupling constants g_{ij} can be $g_{\mu\mu} : g_{\mu\tau} : g_{\tau\tau} \simeq 1 : m_\mu/m_\tau : (m_\mu/m_\tau)^2$. For $g_{\mu\mu} \simeq 1$, the neutrino data and the LFV data give the constraints such as $m_k \gtrsim 770$ GeV and $m_\omega \gtrsim 160$ GeV [10]. On the other hand, the constraints on the couplings and masses are more stringent for the inverted neutrino mass hierarchy. The current data then gives $m_\omega \simeq 825$ GeV for $g_{\mu\mu} \simeq 1$ [10]. One of the notable things in this case is the lower bound on $\sin^2 2\theta_{13}$, which is predicted as around 0.002 [9].

2.2. Models with TeV-scale right-handed neutrinos with a discrete Z_2 symmetry

Similar to the tree-level seesaw model, tiny masses of left-handed neutrinos would also come from Majorana masses $M_{N^\alpha_R}$ of gauge-singlet right-handed neutrinos N^α_R in the radiative seesaw scenario [4–6]. One simple way to realize the absence of the tree-level Yukawa interaction $\bar{\nu}^i_L\tilde{\Phi}N^\alpha_R$ is introduction of a discrete Z_2 symmetry, with the assignment of the odd quantum number to N^α_R and the even to the SM particles. To obtain the dimension five operator in Eq. (1) at the loop level, we need to introduce additional Z_2-odd scalar fields. The lightest of all the Z_2-odd particles can be a candidate of DM if it is electrically neutral. The

[1] If we take $g_{ee} = 0$, then m_k is unbounded from the $\mu \to eee$ and $\tau \to \ell ee$ results ($\ell = e$ or μ), so that relatively light $k^{\pm\pm}$ ($m_k \sim 100$–200 GeV) are possible.