PHYSICAL REVIEW D **80,** 053006 (2009)

On the origin of neutrino masses

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(Received 24 June 2009; published 9 September 2009)

We discuss the simplest mechanisms for generating neutrino masses at tree level and one loop level. We find a significant number of new possibilities where one can generate neutrino masses at the one-loop level by adding only two new types of representations. These models have renormalizable interactions that automatically conserve baryon number. Adding to the minimal standard model a scalar color octet with $SU(3) \otimes SU(2) \otimes U(1)$ quantum numbers, $(8, 2, 1/2)$, and a fermionic color octet in the fundamental or adjoint representation of $SU(2)$, one can generate neutrino masses in agreement with experiment. Signals at the LHC and constraints from flavor violation are briefly discussed.

DOI: 10.1103/PhysRevD.80.053006 PACS numbers: 14.60.Pq, 12.15.Ff, 12.60.Fr, 13.15.+g

I. INTRODUCTION

The existence of massive neutrinos is one of the main motivations for physics beyond the standard model (SM). As is well known the neutrinos can be Dirac or Majorana fermions. In the case of Majorana neutrinos there is a great variety of scenarios for the origin of neutrino masses. At tree level we can generate neutrino masses using the well-known type I, type II, or type III seesaw scenarios. In the type I seesaw mechanism one adds at least two SM singlets, $\nu^C \sim (1, 1, 0)$ [1], and once those singlets are integrated out the neutrino mass matrix is given by $\mathcal{M}^I = Y_\nu M_R^{-1} Y_\nu^T v^2$, where Y_ν is the Yukawa coupling between the SM leptonic doublet and the right-handed neutrinos, v is the vacuum expectation value of the SM Higgs, and M_R is the Majorana mass matrix for the right-handed neutrinos. In the type II seesaw mechanism [2] an $SU(2)$ scalar triplet is introduced, $\Delta \sim (1, 3, 1)$, and the neutrino mass matrix reads as $\mathcal{M}^{II} = h_\nu v_\Delta$. Here, h_ν is the Yukawa coupling between the leptons and the triplet, and v_Δ is the vacuum expectation value of the neutral component of the triplet. It is also possible to generate neutrino masses at tree level if one introduces at least two extra fermions in the adjoint representation of $SU(2)$, $\rho \sim (1, 3, 0)$ [3], and the mass matrix for neutrinos is similar to the type I case, where one replaces M_R by M_ρ, the Majorana mass matrix for the fermionic triplets. This is the type III seesaw scenario. These are the simplest mechanisms for generating neutrino masses at tree level since they add just one new type of representation to the minimal standard model. If one realizes the type III seesaw mechanism in the context of grand unified theories (GUT's), it is always a hybrid scenario using type I plus type III seesaw mechanisms [3]. Of course one can add more than two fields as in the case of R-parity violation in supersymmetric theories. See Ref. [4] for a review of different seesaw mechanisms.

In this paper we investigate the simplest possible scenarios that generate the neutrino masses at one-loop level. We stick to the cases where there are at most two new types of fields with different gauge quantum numbers and restrict our attention, for the most part, to singlet, fundamental, and adjoint representations of the non-Abelian gauge groups. We also focus on the case where the new particles have masses of order the TeV scale (or less) since then it may be possible to test the origin of neutrino masses experimentally.

This paper is organized as follows: In the second section we discuss and classify the simplest scenarios for the generation of neutrino masses at one-loop level. In Sec. III phenomenological predictions and signals at the LHC are briefly discussed. We summarize our findings in the last section.

II. NEUTRINO MASSES: ONE-LOOP MECHANISMS

In this section we outline the simplest mechanisms where neutrino masses are generated at the one-loop level. It is well known that by introducing one scalar SM singlet field, $h \sim (1, 1, 1)$, and an extra Higgs doublet, $(1, 2, 1/2)$, one can generate neutrino masses at the one-loop level. This is the so-called Zee model [5]. In previous studies it has been shown that it is not possible to generate neutrino masses in agreement with the experiment [6] in the version of this model where only one Higgs couples to the leptons. (This naturally suppresses flavor changing neutral Higgs couplings.) This scenario is called the Zee-Wolfenstein model [5,7]. Introducing two scalar leptoquarks, it is possible to generate neutrino masses at the one-loop level. One example introduces the scalar fields $LQ_1 \sim (3, 2, 1/6)$ and $LQ_2 \sim (3, 1, -1/3)$. However, these fields have renormalizable baryon number violating couplings, and proton decay occurs at tree level. A mechanism is needed to

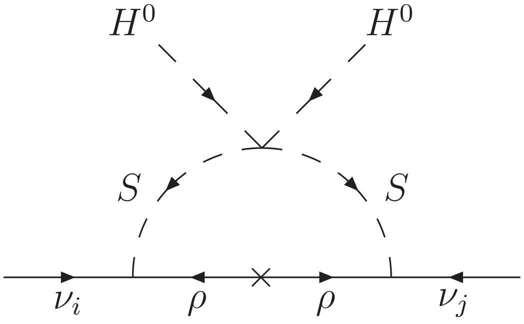

FIG. 1. Mechanism at one-loop level.

suppress baryon number violation if both fields are light. One can impose by hand some symmetries to forbid or suppress the proton decay rate or have the renormalizable coupling constants that violate baryon number be very small. In this paper we restrict our attention to models where baryon number conservation is an automatic symmetry of the renormalizable couplings. See Ref. [8] for recent studies of models where one generates neutrino masses at one-loop level using the leptoquark fields $LQ_{1,2}$.

If one introduces just two new types of fields, a scalar S and the other a fermion ρ, neutrino masses can be generated at one-loop level as shown in Fig. 1. Since one can generate neutrino masses at tree level using fermionic singlets, $\nu^C \sim (1, 1, 0)$, fermionic triplets, $\rho \sim (1, 3, 0)$, or scalar triplets, $\Delta \sim (1, 3, 1)$, we do not allow these representations. We are mainly interested in cases where colored fields play a role in the generation of neutrino masses since they can be produced at the LHC with large cross sections when their masses are below a TeV.

Let us now analyze the different scenarios where neutrino masses are generated through Fig. 1.

Case 1 In this case two fermionic, $\chi \sim (1, 2, 0)$, fields and two extra scalars, $S \sim (1, 3, 1/2)$, are added to the minimal standard model. The new fields, S and χ, occur inside the loop of Fig. 1. The simultaneous presence of the Yukawa interactions and the quartic interaction between S and H tells us that the lepton number is broken by two units generating the usual dimension-five lepton-number violating operator for neutrino masses. Notice that in this case the extra fields do not have direct couplings to the SM quarks. The interactions needed to realize this mechanism are

$$\begin{aligned} -\mathcal{L}_1 = {}& Y_1 l^T C i\sigma_2 S \chi + M_\chi \chi^T C i\sigma_2 \chi \\ & + \lambda_1 H^T i\sigma_2 S^\dagger S^\dagger H + \text{H.c.} \end{aligned} \tag{1}$$

Unfortunately, since the extra fields S and χ give rise to fractionally charged (color singlet) particles, there is always a stable charged particle in this scenario. Therefore this case is ruled out by cosmological constraints and searches of exotic nuclei.

Case 2 One can have alternative mechanisms where the extra fields live in nontrivial representations of $SU(3)$. In order to avoid new anomalies we stick to real representations of $SU(3)$, the one with the lowest dimension being the adjoint. Adding one (two) extra scalar $S_1 \sim (8, 2, 1/2)$ and two (one) fermionic fields $\rho_1 \sim (8, 1, 0)$, it is possible to generate neutrino masses at one-loop level via Fig. 1. Since the scalar octet has hypercharge $1/2$, one can use the quartic interactions between this field and the SM Higgs in order to generate the dimension-five operator for neutrino masses. Notice that ρ_1 has the same quantum numbers as the gluino in supersymmetric models and S_1 will have extra couplings to the standard model quark fields [9]. As in the previous cases we show explicitly the relevant interactions (here we write just one possible quartic interaction for simplicity) needed to generate neutrino masses,

$$\begin{aligned} -\mathcal{L}_2 = {}& Y_2 l^T C i\sigma_2 S_1 \rho_1 + M_{\rho_1} \mathrm{Tr}\rho_1^T C \rho_1 + \lambda_2 \mathrm{Tr}(S_1^\dagger H)^2 \\ & + \text{H.c.} \end{aligned} \tag{2}$$

In Eq. (2) the trace is over color matrices. In this case one can have a consistent scenario for cosmology since the scalar octet has couplings to the SM matter fields and we can satisfy all cosmological constraints. As far as we know, this mechanism has not been discussed in previous studies.

Case 3 One can generalize the previous mechanism using the extra scalar octet $S_1 \sim (8, 2, 1/2)$ and taking the extra fermion field in the adjoint representation of $SU(2)$, $\rho_2 \sim (8, 3, 0)$. One needs two (one) extra scalars and one (two) fermion in order to generate neutrino masses and mixings in agreement with experiment. In this case, in order to generate neutrino masses, one uses the following interactions:

$$\begin{aligned} -\mathcal{L}_3 = {}& Y_3 l^T C i\sigma_2 \rho_2 S_1 + M_{\rho_2} \mathrm{Tr}\rho_2^T C \rho_2 + \lambda_3 \mathrm{Tr}(S_1^\dagger H)^2 \\ & + \text{H.c.} \end{aligned} \tag{3}$$

As in the previous case one can satisfy all cosmological constraints since there are no stable charged particles and the extra scalar octet has renormalizable Yukawa couplings to the standard model quarks.

Case 4 It is also possible to introduce two copies of extra fermions which are in the adjoint representation of $SU(3)$ and in the fundamental of $SU(2)$, $\eta \sim (8, 2, 0)$, and two extra scalars in the adjoint of both gauge groups, $\Sigma \sim (8, 3, 1/2)$. Notice that in this case the extra scalar octets do not have Yukawa couplings to the standard model quark fields since they are in the adjoint of $SU(2)$. Using the following interactions, neutrino masses are generated at one loop via Fig. 1,

$$\begin{aligned} -\mathcal{L}_4 = {}& Y_4 l^T C i\sigma_2 \Sigma \eta + M_\eta \mathrm{Tr}\eta^T C i\sigma_2 \eta \\ & + \lambda_4 \mathrm{Tr} H^T i\sigma_2 \Sigma^\dagger \Sigma^\dagger H + \text{H.c.} \end{aligned} \tag{4}$$

However, this scenario is also ruled out since it has fractionally stable charged (color singlet) particles.

TABLE I. Different seesaw scenarios.

Seesaw scenario	Extra scalar representations	Extra fermion representations	Status
Tree level			
I		(1, 1, 0)	OK
II	(1, 3, 1)		OK
III		(1, 3, 0)	OK
One-loop level			
Zee model	(1, 1, 1)		OK[a]
Case 1	(1, 3, 1/2)	(1, 2, 0)	Ruled out
Case 2	(8, 2, 1/2)	(8, 1, 0)	OK
Case 3	(8, 2, 1/2)	(8, 3, 0)	OK
Case 4	(8, 3, 1/2)	(8, 2, 0)	Ruled out

[a] The Zee-Wolfenstein model, where only one of the Higgs doublets couples to the leptons, is ruled out [6].

In Table I we summarize the different scenarios showing the $SU(3) \otimes SU(2) \otimes U(1)$ gauge quantum numbers of the different representations needed. (The color representation 8 can be replaced by any real representation R. Higher dimension $SU(2)$ representations are also possible.) We have shown that to generate neutrino masses at one-loop level, adding the minimal number of new representations and imposing no extra symmetries, the most economical ways are the Zee model and models that introduce a scalar octet, $(8, 2, 1/2)$, and a fermionic octet which can be in the fundamental or adjoint representation of $SU(2)$.

Consider case 2 with two copies of the new fermions. Working in the mass eigenstate basis for the two new fermions ρ_1^α, the neutrino mass matrix reads as

$$\mathcal{M}_\nu^{ij} = Y_2^{i\alpha} Y_2^{j\alpha} \frac{\lambda_2}{16\pi^2} v^2 I(M_{\rho_1^\alpha}, M_{S_1}), \tag{5}$$

with $\alpha = 1, 2$. The loop integration factor, $I(M_{\rho_1^\alpha}, M_{S_1})$, is given by

$$I(M_{\rho_1^\alpha}, M_{S_1}) = 4M_{\rho_1^\alpha}\left(\frac{M_{S_1}^2 - M_{\rho_1^\alpha}^2 + M_{\rho_1^\alpha}^2 \ln(M_{\rho_1^\alpha}^2/M_{S_1}^2)}{(M_{S_1}^2 - M_{\rho_1^\alpha}^2)^2}\right). \tag{6}$$

With just this minimal number of copies of the new fields there is a massless neutrino. Therefore, there are two types of spectra: normal hierarchy with $m_1 = 0$, $m_2 = \sqrt{\Delta m_{\rm sol}^2}$, and $m_3 = \sqrt{\Delta m_{\rm sol}^2 + \Delta m_{\rm atm}^2}$, and inverted hierarchy with $m_3 = 0$, $m_2 = \sqrt{\Delta m_{\rm atm}^2}$, and $m_1 = \sqrt{\Delta m_{\rm atm}^2 - \Delta m_{\rm sol}^2}$. Here $\Delta m_{\rm sol}^2 \approx 8 \times 10^{-5}$ eV^2 and $\Delta m_{\rm atm}^2 \approx 2.5 \times 10^{-3}$ eV^2 are the solar and atmospheric mass squared differences. We can also have a minimal scenario with two extra octet scalars and one fermionic octet.

In the limit $M_{S_1} \gg M_{\rho_1}$ the neutrino mass matrix becomes

$$\mathcal{M}_\nu^{ij} = Y_2^{i\alpha} Y_2^{j\alpha} \frac{\lambda_2}{4\pi^2} v^2 \frac{M_{\rho_1^\alpha}}{M_{S_1}^2}. \tag{7}$$

Using as input parameters $M_{\rho_1} = 200$ GeV, $v = 246$ GeV, and $M_{S_1} = 2$ TeV, we find that in order to get the neutrino "scale," $\sim$1 eV, the combination of the couplings, $Y_2^2 \lambda_2 \sim 10^{-8}$. If $\lambda_2 \sim 1$ the elements of the Yukawa coupling matrix $Y_2 \sim 10^{-4}$. The Yukawa couplings can be larger if λ_2 is smaller. In the scenarios proposed by us it is possible to reproduce the measured neutrino masses and mixing using a mechanism that can be tested at the LHC.

As we have discussed before, the simultaneous presence of the Yukawa term proportional to Y_2 and the quartic interaction proportional to λ_2 in Eq. (2) violate lepton number. In this case, lepton flavor conservation is violated, even when $\lambda_2 = 0$, by the Yukawa couplings Y_2. Hence, even when λ_2 is very small, there are constraints on the size of these Yukawa coupling constants from limits on the rates for lepton flavor violating processes like $\mu \to e + \gamma$. We hope to investigate these constraints in a future publication.

III. PHENOMENOLOGICAL ASPECTS

Possible signals.—We now discuss a few of the phenomenological aspects of the scenarios discussed above assuming that the extra fields have masses of order a TeV (or less) so that they can be produced at the LHC. In case 2 neutrino masses are generated at one loop using the octet scalar $S_1 \sim (8, 2, 1/2)$ and the fermionic octet $\rho_1 \sim (8, 1, 0)$. The phenomenological aspects of the scalar octet have been studied in great detail by several groups [10]. If a neutral scalar is the lightest new particle it will decay directly to quark-antiquark pairs at tree level or at the one-loop level to gluons using the interaction term in the scalar potential, $\lambda_5 \,{\rm Tr} H^\dagger S S S^\dagger$ + H.c. [9]. In case 2 the extra fermion has the same quantum numbers as the gluino in supersymmetric theories. It can be produced in pairs through the strong interactions, and its decay width will be dominated by two-body decays, $\rho_1 \to l S_1$ if $M_{\rho_1} > M_{S_1} + M_l$, or three-body decays when S_1 is virtual. Since the ρ_1 fields are Majorana, one can have very exotic final state channels with two same-sign charged leptons and four jets. In particular, the channels $pp \to \rho_1 \rho_1 \to S_1^+ S_1^+ e_i^- e_j^- \to e_i^- e_j^- t t \bar{b} \bar{b}$, where one has leptons with the same electric charge, two tops, and two antibottom quarks, is the cleanest channel to test the mechanism for neutrino masses in case 2.

From the production and decay properties of the scalar octets, one may be able to determine their masses, and this information can be used to understand the three- or two-body decays of the fermionic octet. Once we impose the constraints coming from neutrino masses, one has some information on the Yukawa coupling between the octets and the leptonic doublets. Unfortunately, they are not completely determined since the coupling λ_2 multiplies the neutrino mass matrix [see Eq. (5)].

The phenomenology of case 3 is similar to case 2; however, one important difference is that there are both electrically charged and neutral color octet fermions. The splitting between the charged and neutral color octet fermions is small since it is generated at the one-loop level. One striking channel associated with the production of charged "gluinos" at the LHC is $pp \rightarrow \rho_2^+ \rho_2^- \rightarrow e_i^+ e_j^- S_1^0 S_1^0 \rightarrow e_i^+ e_j^- t\bar{t} t\bar{t}$.

We can generalize cases 2 and 3 above by, for example, changing the color octet representation to any other real representation R of the SU(3) gauge group. Interactions that break the $\rho \rightarrow -\rho$, $S \rightarrow -S$ symmetry are needed to allow for the new strongly interacting particles to decay. A term in the scalar potential of the form $H^\dagger SSS^\dagger$ breaks this discrete symmetry. It allows the neutral members of the S representation to decay to gluons at the one-loop level. The octet representation is the smallest real representation. It is also the only representation that allows for renormalizable Yukawa couplings of the scalars S to the quarks. In the case $R = 8$ the S scalars can decay at tree level to quark-antiquark pairs.

Other aspects.—If one does not impose minimal flavor [11] violation, the Yukawa couplings of the S_1 to quarks are constrained by the smallness of observed flavor changing neutral currents. For example, the measured value of the $K_L - K_S$ mass difference implies that the $d \rightarrow s$ Yukawa coupling of neutral S to quarks is less than about $10^{-5}(M_{S_1}/\text{TeV})$. This might seem like a very strong constraint; however, it is important to remember that the electron Yukawa coupling of the standard model Higgs doublet is about 10^{-5}.

Before finishing, we would like to comment on possible constraints coming from neutrinoless double beta decay. Cases 2 and 3 have the usual contribution to this rare process due to the existence of light Majorana neutrino masses. In addition, the Yukawa coupling of the scalar octet to the quarks gives rise to new contributions. However, they are highly suppressed by the masses of the scalar and fermionic octets.

IV. SUMMARY

We have discussed the simplest mechanisms that generate neutrino masses either at tree level or at one loop, where one introduces at most two types of representations beyond those that are in the minimal standard model. We found new possibilities where neutrino masses are generated at the one-loop level and the renormalizable interactions automatically conserve baryon number. The simplest cases have a scalar octet, $(8, 2, 1/2)$, and a fermionic octet in the fundamental or adjoint representation of $SU(2)$. Possible signals at the LHC and the constraints from flavor violation were briefly discussed. We hope to elaborate on some of the phenomenological implications of these models in a future publication.

ACKNOWLEDGMENTS

One of the authors (P. F. P.) would like to thank Caltech for hospitality. The work of P. F. P. was supported in part by the U.S. Department of Energy Contract No. DE-FG02-08ER41531 and in part by the Wisconsin Alumni Research Foundation. The work of M. B. W. was supported in part by the U.S. Department of Energy Contract No. DE-FG02-92ER40171.

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