

# Neutrino masses, leptogenesis, and dark matter in a hybrid seesaw model

Pei-Hong Gu,<sup>1,\*</sup> M. Hirsch,<sup>2,†</sup> Utpal Sarkar,<sup>3,‡</sup> and J. W. F. Valle<sup>2,§</sup>

<sup>1</sup>*The Abdus Salam International Centre for Theoretical Physics, Strada Costiera 11, 34014 Trieste, Italy*

<sup>2</sup>*AHEP Group, Institut de Física Corpuscular–C.S.I.C./Universitat de València, Edificio Institutos de Paterna, Apt 22085, E-46071 Valencia, Spain*

<sup>3</sup>*Physical Research Laboratory, Ahmedabad 380009, India*

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We suggest a hybrid seesaw model where relatively light right-handed neutrinos give no contribution to neutrino mass matrix due to a special symmetry. This allows their Yukawa couplings to the standard model particles to be relatively strong, so that the standard model Higgs boson can decay dominantly to a left- and a right-handed neutrino, leaving another stable right-handed neutrino as cold dark matter. In our model neutrino masses arise via the type-II seesaw mechanism, the Higgs triplet scalars being also responsible for the generation of the matter-antimatter asymmetry via the leptogenesis mechanism.

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## I. INTRODUCTION

Laboratory experiments with reactors and accelerator neutrinos have confirmed the observations of solar and atmospheric neutrinos [1], establishing the phenomenon of neutrino oscillations and hence the need for small but nonzero neutrino masses [2]. Since neutrinos are massless in the  $SU(3)_c \times SU(2)_L \times U(1)_Y$  standard model (SM) this implies the need for new physics, whose detailed nature constitutes one of our deepest current challenges in particle physics [3]. The simplest extension of the SM to explain the neutrino masses is to include singlet right-handed neutrinos and/or triplet scalars [4]. The inclusion of the former is justified by the fact that there are right-handed partners for all other fermions of the SM. The singlet right-handed neutrinos will in general have Majorana masses as well as a Dirac mass term, of the order of the charged fermion masses. The latter arise from the usual Yukawa interaction between left- and right-handed leptons, once the SM Higgs doublet acquires a vacuum expectation value (VEV). The terms involving the right-handed neutrinos are

$$\mathcal{L} \supset -\frac{1}{2}(M_N)_{ij}\bar{N}_R^c N_{R_j} - h_{ij}\bar{\ell}_L \phi N_{R_j} + \text{H.c.} \quad (1)$$

Here  $N_R$ ,  $\ell_L$ , and  $\phi$  denote the right-handed neutrinos, the left-handed leptons, and Higgs doublet, respectively. In the basis  $(\nu_L, N_R^c)$ , the neutrino mass matrix reads

$$M_\nu = \begin{pmatrix} 0 & m_D \\ m_D^T & M_N \end{pmatrix}, \quad (2)$$

where  $(m_D)_{ij} = h_{ij}\langle\phi\rangle$  with  $\langle\phi\rangle \simeq 174$  GeV. After block-

diagonalization, one gets

$$m_\nu \simeq -m_D \frac{1}{M_N} m_D^T. \quad (3)$$

If the Majorana masses of the right-handed neutrinos are much larger than the Dirac masses, the left-handed neutrinos will naturally acquire a tiny Majorana through the type-I seesaw mechanism [4–6]. In order to account for neutrino masses in the eV range one needs

$$h_{ij}h_{kj}(M_N)_{jk}^{-1} \sim 10^{-23} \text{ GeV}^{-1} \quad (4)$$

so that the Majorana masses  $M_N$  are required to be orders of magnitude larger than the electroweak symmetry-breaking scale  $\sim\langle\phi\rangle$ , unless the coefficients  $h_{ij}$  are very small,

$$h_{ij} \lesssim 10^{-11}. \quad (5)$$

There is an attractive way to avoid this conclusion in the framework of the inverse seesaw model [7–9]. In this case the right-handed neutrinos pair-off with extra singlet leptons to form Dirac-type neutral heavy leptons in such a way that their mixing with the doublet neutrinos does not lead to light neutrino masses in the limit of conserved lepton number. As a result of this “symmetry-protection,” right-neutrinos can lie at the TeV scale and produce signals at accelerators [10,11], without conflict with the observed smallness of neutrino masses.

Here we focus on a seesaw mechanism without additional fermion degrees of freedom. If one assumes “generic” violation of lepton number through right-handed neutrino Majorana masses, acceptable neutrino masses require very tiny effective Yukawa couplings connecting the right-handed neutrinos with the left-handed neutrinos, Eqs. (4) and (5), hence the right-handed neutrinos cannot be produced at the LHC, nor can the Higgs boson decay into neutrinos. We explore the possibility that the neutrino mass  $m_\nu$  can be exactly zero even if the Majorana mass  $M_N$

\*pgu@ictp.it

†mahirsch@ific.uv.es

‡utpal@prl.res.in;

http://www.prl.res.in/~utpal

§valle@ific.uv.es

http://ahep.uv.es

and the Dirac mass  $m_D$  are both nonvanishing [12,13], as a result of a cancellation. We present a model where the type-I seesaw contribution to the neutrino masses vanishes identically due to a suitable symmetry, avoiding the main constraints on  $M_N$  and  $m_D$ . Neutrino masses arise from the type-II triplet seesaw mechanism [4]. As a result, relatively light right-handed neutrinos could have sizable Yukawa couplings. If the Majorana masses of the right-handed neutrinos are of the order of TeV or less, they may be produced at the LHC. We also note that, if it is heavier than the right-handed neutrino, the Higgs boson in this scenario will also have a distinct decay channel into a left- plus a right-handed neutrino. Our explicit symmetry will also protect one of the right-handed neutrinos from decaying, as it has no SM Yukawa interactions. This naturally accounts for a stable cold dark matter candidate [14,15]. Finally, our scenario for the origin of neutrino masses also provides successful leptogenesis [16,17] induced by the out-of-equilibrium decays of the heavy scalar triplets [18].

## II. NEUTRINO MASS MATRICES

We shall first discuss the structure of the Dirac and Majorana masses of the neutrinos and then present the model in detail. To demonstrate the basic idea, consider a two-generation scenario in which the type-I seesaw matrix takes a particular form,

$$M_\nu = \begin{pmatrix} 0 & 0 & a & a \\ 0 & 0 & b & b \\ a & b & M & 0 \\ a & b & 0 & -M \end{pmatrix}, \quad (6)$$

with  $M \gg a, b$ . Indeed, in this example, the two left-handed neutrinos remain massless despite the coexistence of Dirac and Majorana mass terms, as a result of the cancellation between the contributions from the two right-handed neutrinos. The latter may be due to some underlying symmetry. In this case the Yukawa couplings thus can be large even if the right-handed neutrinos have ‘‘low’’ Majorana masses.

Consider the mass matrix

$$M_\nu = \begin{pmatrix} 0 & 0 & a & 0 \\ 0 & 0 & b & 0 \\ a & b & 0 & M \\ 0 & 0 & M & 0 \end{pmatrix}. \quad (7)$$

It is easy to see that (i) this matrix reduces to the previous one after diagonalizing out the right-handed states and (ii) it emerges from a  $Z_3$  symmetry

$$N_{R_1} \rightarrow \omega N_{R_1}, \quad N_{R_2} \rightarrow \omega^2 N_{R_2}, \quad \nu_{L_i} \rightarrow \omega \nu_{L_i}, \quad (8)$$

where  $\omega$  is the cube root of 1,  $\omega^3 = 1$ , and  $1 + \omega + \omega^2 = 0$ . All other charged fermions also transform under this  $Z_3$  symmetry as  $f \rightarrow \omega f$  and the Higgs doublet  $\phi$  is invariant,

so that all the usual couplings allowed in the SM remain the same. The Lagrangian contributing to the neutrino masses is then given by

$$\mathcal{L} \supset -M \bar{N}_{R_1}^c N_{R_2} - h_{i1} \bar{\ell}_{L_i} \phi N_{R_1} + \text{H.c.}, \quad (9)$$

which clearly leads to the form in Eq. (7) once the VEV  $\langle \phi \rangle$  is generated. This symmetry makes the light neutrinos remain exactly massless. In order to generate the required neutrino masses one may either break this  $Z_3$  softly or, alternatively, introduce triplet scalars for implementing the type-II seesaw [4], without affecting the symmetry in the right-handed neutrino sector.

This mass matrix can be generalized to three generations. Consider the  $3 \times 3$  mass matrix

$$\mathcal{M}_\nu = \begin{pmatrix} 0 & 0 & 0 & a & a & 0 \\ 0 & 0 & 0 & b & b & 0 \\ 0 & 0 & 0 & c & c & 0 \\ a & b & c & M & 0 & 0 \\ a & b & c & 0 & -M & 0 \\ 0 & 0 & 0 & 0 & 0 & \tilde{M} \end{pmatrix}, \quad (10)$$

with  $M, \tilde{M} \gg a, b, c$ . Clearly, the three left-handed neutrinos remain exactly massless, while the masses of the right-handed neutrinos  $+M, -M, \tilde{M}$  can be low enough to be accessible to the LHC. Note, however, that the third-generation right-handed neutrino decouples. In another basis of the right-handed neutrinos one has

$$\mathcal{M}_\nu = \begin{pmatrix} 0 & 0 & 0 & a & 0 & 0 \\ 0 & 0 & 0 & b & 0 & 0 \\ 0 & 0 & 0 & c & 0 & 0 \\ a & b & c & 0 & M & 0 \\ 0 & 0 & 0 & M & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & \tilde{M} \end{pmatrix}. \quad (11)$$

This mass matrix could emerge from a  $Z_3$  symmetry

$$\begin{aligned} N_{R_1} &\rightarrow \omega N_{R_1}, & N_{R_2} &\rightarrow \omega^2 N_{R_2}, \\ N_{R_3} &\rightarrow N_{R_3}, & \nu_{L_i} &\rightarrow \omega \nu_{L_i}, \end{aligned} \quad (12)$$

or alternatively from a  $U(1)$  global symmetry. The Lagrangian containing the right-handed neutrinos

$$\mathcal{L} \supset -M \bar{N}_{R_1}^c N_{R_2} - \frac{1}{2} \tilde{M} \bar{N}_{R_3}^c N_{R_3} - h_{i1} \bar{\ell}_{L_i} \phi N_{R_1} + \text{H.c.} \quad (13)$$

implies, in the limit  $M \gg h_{i1} \langle \phi \rangle$ , two degenerate right-handed neutrinos

$$\frac{1}{\sqrt{2}} (N_{R_1} \pm N_{R_2}) \rightarrow N_{R_{1,2}} \quad (14)$$

with masses  $\pm M$  (opposite  $CP$  signs). For  $M \lesssim 1$  TeV these two states  $N_{R_{1,2}}$  would be accessible to LHC searches. On the other hand, a more interesting possibility may open up when the Higgs boson is heavier than  $N_{R_1}$  and  $N_{R_2}$ .

Given the mild constraint on the Dirac mass term, the Yukawa couplings could be large enough that the Higgs boson would dominantly decay into a left-handed neutrino and a right-handed one, posing a new challenge to the Higgs search program at the LHC. The third right-handed neutrino  $N_{R_3}$  has mass  $\tilde{M}$  and no decay modes. Hence it could serve as the dark matter if its relic density is consistent with the cosmological observations.

### III. LEPTOGENESIS, NEUTRINO MASS, AND DARK MATTER

We now propose a realistic model, which contains the previous phenomenology of the right-handed neutrinos and explains the matter-antimatter asymmetry, the neutrino masses, and the dark matter. We extend the SM by including three right-handed neutrinos, two triplet Higgs scalars  $\xi_{1,2} \equiv (\mathbf{1}, \mathbf{3}, 2)$ , and two singlet scalar fields  $\sigma, \chi \equiv (\mathbf{1}, \mathbf{1}, 0)$ . In addition to the SM gauge symmetry, we also impose a global  $U(1)_{\text{lep}}$  symmetry of lepton number, under which the different fields transform as shown in Table I.

The relevant part of the Lagrangian is given as

$$\begin{aligned} \mathcal{L} \supset & -\alpha_1 \sigma \bar{N}_{R_1}^c N_{R_2} - \frac{1}{2} \alpha_2 \sigma \bar{N}_{R_3}^c N_{R_3} - h_{i1} \bar{\ell}_{L_i} \phi N_{R_1} \\ & - \frac{1}{2} f_{ij} \bar{\ell}_{L_i}^c i \tau_2 \xi_1 \ell_{L_j} - \alpha_3 \chi \phi^T i \tau_2 \xi_2 \phi - \mu \chi \xi_2^\dagger \xi_1 \\ & + \text{H.c.} \end{aligned} \quad (15)$$

Here we have tacitly constrained ourselves to a renormalizable context. After the singlet scalar  $\sigma$  develops its VEV the first line will induce the Lagrangian (13), so that the right-handed neutrinos obtain their Majorana masses. The second line will generate the type-II seesaw in the presence of  $\langle \chi \rangle$ .

In our model, the global  $U(1)_{\text{lep}}$  is assumed to break at a very large scale by  $\langle \chi \rangle \sim 10^{13}$  GeV. The triplet scalars  $\xi_1$  and  $\xi_2$ , whose masses  $\sim M_\xi$  are of the order of  $\langle \chi \rangle$ , mix with each other and pick up tiny VEVs after the electro-weak symmetry breaking

$$\langle \xi_2 \rangle \sim -\alpha_3 \frac{\langle \chi \rangle \langle \phi \rangle^2}{M_\xi^2}, \quad \langle \xi_1 \rangle \sim -\frac{\mu \langle \chi \rangle \langle \xi_2 \rangle}{M_\xi^2}. \quad (16)$$

These triplet VEVs will give rise to the left-handed neutrino's Majorana mass matrix

$$m_{\nu_{ij}} = f_{ij} \langle \xi_1 \rangle. \quad (17)$$

TABLE I. Lepton number assignments. For simplicity, we do not show the right-handed charged leptons, which carry the same lepton number as their left-handed partners.

$N_{R_1}$	$N_{R_2}$	$N_{R_3}$	$\ell_{L_i}$	$\xi_1$	$\xi_2$	$\sigma$	$\chi$	$\phi$
1	3	2	1	-2	-1	-4	1	0

Clearly,  $\xi_2$  will have a trilinear coupling with  $\phi$  and will mix with  $\xi_1$  after  $\chi$  develops its VEV. Furthermore,  $\xi_1$  has the Yukawa couplings to  $\ell_L$ . So, we can diagonalize the mass matrix of  $\xi_1$  and  $\xi_2$  and then obtain two mass-eigenstates  $\hat{\xi}_1$  and  $\hat{\xi}_2$ , both of which simultaneously have the Yukawa couplings to  $\ell_L$  and the trilinear couplings with  $\phi$ . The  $CP$ -violating and out-of-equilibrium decays of the triplet scalars  $\hat{\xi}_1$  and  $\hat{\xi}_2$  into the SM lepton and Higgs doublets can generate a lepton asymmetry [18]. The sphaleron [17] processes active in the range  $100 \text{ GeV} \leq T \leq 10^{12} \text{ GeV}$  will partially convert this lepton asymmetry to a baryon asymmetry for explaining the matter-antimatter asymmetry of the Universe. In order to successfully induce leptogenesis and suppress washout processes, we require that the singlet scalar  $\sigma$  develop its VEV after the sphaleron epoch is over; for example, we take  $\langle \sigma \rangle \sim 100 \text{ GeV}$ . Through their Yukawa couplings to  $\sigma$  the right-handed neutrinos acquire their Majorana masses  $M = \alpha_1 \langle \sigma \rangle$  and  $\tilde{M} = \alpha_2 \langle \sigma \rangle$ , expected to lie below 100 GeV or so.

The spontaneous breakdown of the global lepton number global symmetry through  $\langle \chi \rangle$ , leads to a Goldstone boson, whose profile can be determined by the symmetry [6], leading to

$$\mathcal{G} = \frac{1}{N} [\langle \chi \rangle \text{Im}(\chi) + \langle \xi_1 \rangle \text{Im}(\xi_1^0) + \langle \xi_2 \rangle \text{Im}(\xi_2^0)], \quad (18)$$

where  $N$  is a suitable normalization, which is of the order of the lepton number breaking scale  $\sim \langle \chi \rangle$ . Clearly, the triplet component of the Goldstone boson is highly suppressed by the ratio of the triplet VEVs over the singlet VEV, suppressing its coupling to the Z bosons [6].

Note that since lepton number conservation forbids the couplings of  $\chi$  to the right-handed neutrinos,  $\langle \chi \rangle$  will not have any effect on them. The states  $N_{R_1}$  and  $N_{R_2}$  mix maximally and in the basis where their Majorana mass matrix is diagonal, the degenerate states  $N_{R_1}$  and  $N_{R_2}$  in Eq. (14) with mass  $\pm M$  both couple to the left-handed neutrinos through equal Yukawa couplings. In addition, the state  $N_{R_3}$  has no couplings to the left-handed neutrinos. Therefore, the resulting neutrino mass matrix is a null matrix. This implies that the strongest constraints on the couplings, for example, those set by the smallness of neutrino masses, are absent, very much like the case of the inverse seesaw model [7,9].

Another most severe constraint on the couplings of the right-handed neutrinos with the charged leptons comes from their contribution to  $\mu \rightarrow e\gamma$ , which is trivially removed now. Since here only the  $N_{R_1}$  and  $N_{R_2}$  would mediate the process and they are degenerate, the  $\mu \rightarrow e\gamma$  amplitude would depend on the effective light neutrino masses, which now vanish, hence avoiding the constraint on the couplings.

Note also that the physical Higgs boson can significantly decay into  $h \rightarrow \nu_L + N_{R_1}$   $h \rightarrow \nu_L + N_{R_2}$ , leading to a

monojet-like signal. In contrast, the decay  $h \rightarrow \nu_L + N_{R_3}$  does not take place.

The third right-handed neutrino  $N_{R_3}$  cannot decay at all because its only Yukawa coupling is with the singlet scalar  $\sigma$ . This means that  $N_{R_3}$  will contribute a sizable relic density to the Universe. One can indeed check that  $N_{R_3}$  with the mass of a few GeV can have a desired cross section to serve as the dark matter for  $\langle \sigma \rangle \sim 100$  GeV.

As a last comment we note that there may be a sizable quartic interaction between the singlet scalar  $\sigma$  and the SM Higgs doublet  $\phi$ , i.e.  $\lambda(\sigma^\dagger\sigma)(\phi^\dagger\phi)$ . This term cannot be forbidden by imposing extra symmetries, and there is no *a priori* reason for  $\lambda$  to be small. In the presence of such coupling this dark matter  $N_{R_3}$  can be searched in the decays of the Higgs boson produced at the LHC,  $h \rightarrow N_{R_3} + N_{R_3}$ , resulting in a missing momentum signal. As for other quartic interactions, i.e.  $(\chi^\dagger\chi)(\phi^\dagger\phi)$  and  $(\chi^\dagger\chi)(\sigma^\dagger\sigma)$ , they will have no significant effects since we have assumed the hierarchy of scalar VEVs  $\langle \chi \rangle \gg \langle \phi \rangle \sim \langle \sigma \rangle$  and hence the leading mixing of physical bosons should be very small.

#### IV. SUMMARY

Because of some special symmetry it may happen that the type-I seesaw mechanism does not generate the ob-

served neutrino masses, despite the coexistence of sizable Dirac mass terms and relatively low right-handed neutrino Majorana masses. Such *null seesaw mechanism* can be understood as a cancellation between the contributions from the right-handed neutrinos. In this case there is only a very mild constraint on the right-handed neutrinos, which can have sizable Yukawa couplings to the SM particles even if they are light. The leading Higgs boson decay mode into a left-handed neutrino and a right-handed neutrino could be probed at the LHC. In the model we have presented, one of the right-handed neutrinos has no Yukawa couplings to the SM states, a fact that follows from our assumed symmetry. Hence it can provide the relic density required to solve the puzzle of the dark matter. Finally, in our model the source of neutrino masses is the type-II seesaw contribution arising from the induced VEVs of scalar Higgs triplets. The *CP*-violating and out-of-equilibrium decays of these scalar triplets may also account for the matter-antimatter asymmetry of the Universe through the leptogenesis mechanism.

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