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Hybrid seesaw leptogenesis and TeV singlets

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A B S T R A C T

The appealing feature of inverse seesaw models is that the Standard Model (SM) neutrino mass emerges from the exchange of TeV scale singlets with sizable Yukawa couplings, which can be tested at colliders. However, the tiny Majorana mass splitting between TeV singlets, introduced to accommodate small neutrino masses, is left unexplained. Moreover, we argue that these models suffer from a structural limitation that prevents a successful leptogenesis if one insists on having unsuppressed Yukawa couplings and TeV scale singlets. In this work we propose a *hybrid seesaw* model, where we replace the mass splitting with a coupling to a high scale seesaw module including a TeV scalar. We show that this structure achieves the goal of filling *both* the above gaps with couplings of order unity. The necessary structure automatically arises embedding the seesaw mechanism in composite Higgs models, but may also be enforced by new gauge symmetries in a weakly-coupled theory. Our hybrid seesaw models have distinguishing features compared to the standard high scale type-I seesaw and inverse seesaw. Firstly, they have much richer phenomenology. Indeed, they generally predict new TeV scale physics (including scalars) potentially accessible at present and future colliders, whereas weakly-coupled versions may also have cosmological signature due to the presence of a light Nambu–Goldstone boson coupled to neutrinos. Secondly, our scenario features an interesting interplay between high scale and TeV scale physics in leptogenesis and enlarges the range of allowed high scale singlet masses beyond the usual $\sim 10^9$–10^{15} GeV, without large hierarchies in the Yukawa couplings nor small mass splitting among the singlets.

1. Introduction

The seesaw mechanism [1] elegantly explains the extreme smallness of the SM neutrino masses. At the same time, the Majorana nature of SM neutrino, i.e., the presence of lepton-number violation, raises the highly attractive possibility of baryogenesis via leptogenesis [2,3]. In this work, we consider a class of seesaw models called inverse seesaw [4]. We first emphasize two inadequacies of the standard inverse seesaw scenario and then build an extended framework, which we will term *hybrid seesaw*, to overcome both issues.

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2. Inverse seesaw: μ-problem and leptogenesis

In the inverse seesaw one introduces a Dirac SM singlet (made up of two Weyl spinors: Ψ and Ψ^c) supplemented with an additional *tiny* Majorana mass term for one of the chiralities and Yukawa coupling of the *other* chirality to the SM Higgs and lepton doublet (denoted by H and ℓ respectively):

$$-\mathcal{L} \supset y\Psi^c H\ell + m_\Psi \Psi\Psi^c + \frac{\mu}{2}\Psi\Psi + \text{h.c.} \tag{1}$$

The generation indices have been suppressed for brevity (y, m_Ψ, μ are in general matrices). Here m_Ψ is assumed to be in the TeV range, while $\mu \ll m_\Psi$. Integrating out these *pseudo*-Dirac singlets generates a small neutrino mass:

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$$m_\nu \sim \frac{(y\,v)^2}{m_\Psi^2}\mu, \tag{2}$$

where $v = 174$ GeV is the vacuum expectation value (VEV) of H. The crucial point here is that we can obtain the observed size of the SM neutrino mass $m_\nu \sim 0.05$ eV with unsuppressed Yukawa couplings $y = O(0.01 - 1)$ and $m_\Psi \sim 1$ TeV, provided $\mu = O(10\text{ keV} - \text{eV})$.

An attractive feature of this scenario is that the singlets are potentially accessible at colliders because of their unsuppressed Yukawa coupling [5].[1] However, this set-up also has two drawbacks. Firstly, if the new physics resides at the TeV scale, there is a priori no reason to expect μ in the keV range or below. Although a small Majorana mass term μ is technically natural (since a symmetry, namely lepton number, is restored by its vanishing), the required value appears as unexpected within this picture: additional ingredients are needed. Secondly, as we will argue next, it appears difficult to achieve successful leptogenesis in this framework.

To study leptogenesis, we first calculate the CP asymmetry from decays of heavy singlet,

$$\epsilon_\Psi \equiv \frac{|\Gamma_\Psi - \overline{\Gamma}_\Psi|}{\Gamma_\Psi + \overline{\Gamma}_\Psi}, \tag{3}$$

where $\Gamma_\Psi(\overline{\Gamma}_\Psi)$ is the decay width of Ψ into $\ell H(\ell^* H^*)$. Assuming anarchic μ-terms and singlet masses, no hierarchies in Yukawa couplings and $O(1)$ CP-violating phases, we have:

$$\epsilon_\Psi \sim \frac{\mu}{m_\Psi}\frac{\mu}{\Gamma_\Psi}, \tag{4}$$

where the first factor may be interpreted as arising from the CP phase in Yukawa couplings, whereas the second comes from the on-shell propagator due to the near-degeneracy of the pseudo-Dirac pair Ψ, Ψ^c when calculating one-loop self-energy corrected decay width. The two powers of μ in eq. (4) can be understood in generality using the Nanopoulos–Weinberg theorem [6], that states that we need to go to at least second order in the lepton-number breaking parameter (namely, the μ-term in this model) in order to generate an asymmetry. This result was first obtained in the first reference in [7] and was backed up by a detailed analysis [7]. Crucially, it depends on the regulator used for the almost on-shell Ψ propagator in the self-energy diagram.

To determine the present-day asymmetry we should combine the above result with the effective washout factor from the *inverse* decay of SM leptons and Higgs into the singlets. This latter quantity was first estimated in [8]:

$$K_\Psi^{\text{eff}} \sim \frac{\Gamma_\Psi}{H_\Psi}\frac{\mu^2}{\Gamma_\Psi^2}, \tag{5}$$

where $K_\Psi \sim \Gamma_\Psi/H_\Psi$ is the "usual" washout factor [3] and H_Ψ is the Hubble parameter at $T = m_\Psi$, i.e., $H_\Psi \sim \sqrt{g_*}m_\Psi^2/M_{\text{Pl}}$, with g_* being the number of relativistic degrees of freedom at that temperature and M_{Pl} the Planck mass. The quadratic suppression in μ comes from the fact that the rate for lepton-number-violating processes, e.g. $\ell H \leftrightarrow (\ell H)^*$, should vanish in the lepton-number conserving limit.[2] If for definiteness we focus on the strong washout region, in which $K_\Psi^{\text{eff}} \gg 1$, the net lepton asymmetry can be obtained as[3]

$$Y_{\Delta\ell}^\Psi \sim 10^{-3}\frac{\epsilon_\Psi}{K_\Psi^{\text{eff}}}, \tag{6}$$

where $Y_X \equiv n_X/s$ $(Y_{\Delta X} \equiv (n_X - n_{X^*})/s)$ with n_X being number density of the corresponding species and s being total entropy density of the Universe. The numerical factor $\sim 10^{-3}$ in eq. (6) comes from relativistic number density of Ψ normalized to s.

Putting everything together, and assuming strong washout for simplicity, we find

$$Y_{\Delta\ell}^\Psi \sim 10^{-3}\frac{\sqrt{g_*}m_\Psi}{M_{\text{Pl}}} \sim 10^{-18}\left(\frac{g_*}{100}\right)^{\frac{1}{2}}\left(\frac{m_\Psi}{\text{TeV}}\right). \tag{7}$$

The final lepton asymmetry in eq. (7) is *in*dependent of the size of the Yukawa couplings. Furthermore, given that $Y_{\Delta B} \sim Y_{\Delta\ell}^\Psi$ after electroweak sphaleron processes are taken into account, we see that eq. (7) predicts too small baryon asymmetry to account for the observed one ($Y_{\Delta B}^{\text{obs}} \sim 10^{-10}$) for singlet masses in the TeV ballpark. In order to reach this conclusion, it is important to include the effect of washout: considerations based solely on ϵ_Ψ could suggest that larger μ than benchmark value shown below eq. (2) might suffice. For example, taking $\mu \sim 10$ MeV (and compensating this increase by reducing the size of y to $y \sim$ a few 10^{-3} to keep m_ν fixed) gives rise to $\epsilon_\Psi \sim 10^{-7}$. While the difficulty in getting required size of CP violation was pointed out in [7], to our knowledge, the parametric form of eq. (7) including washout effect has never been presented before.

A small baryon asymmetry is a very generic implication of TeV scale inverse seesaw.[4] We will show in a companion paper [11] that even allowing a departure from the above generic conditions, for example allowing a degeneracy among different generations of singlets, as well as considering the weak washout regime, the inverse seesaw scenario can at most barely reach the required asymmetry. Introducing other small sources of lepton-number violation as in the linear seesaw model [12] does not change this conclusion [11]. Similar conclusions are obtained in the numerical analysis of ref. [13].

3. A hybrid seesaw model

We now construct an extension of the original inverse seesaw model that features a high-scale module. We will see that, if the interactions between the low and high scale modules are properly chosen, the resulting scenario can simultaneously address *both* the smallness of neutrino masses and leptogenesis.

Our model is the following:

$$-\mathcal{L} \supset y\Psi^c H\ell + \kappa\Phi_\kappa\Psi\Psi^c + \lambda\Phi_\lambda\Psi N + \frac{M_N}{2}NN + \text{h.c.} \tag{8}$$

Here N is a *super*-heavy singlet with mass $M_N \gg$ TeV, whereas Ψ, Ψ^c, $\Phi_{\lambda,\kappa}$ acquire masses (and VEV's) of the order of TeV. Following the philosophy of inverse seesaw, we work with unsuppressed Yukawa couplings y, λ, κ. Furthermore, we will assume anarchical Yukawa couplings such that different generations are comparable, complex phases are of order unity, and the masses of N are not hi-

[1] The singlets may also be charged under new gauge symmetries broken at the TeV scale, giving additional production channels.

[2] Throughout the paper we take $\mu \ll \Gamma_\Psi$, as is expected given that Yukawa couplings are unsuppressed.

[3] The superscript Ψ is to remind the reader that the asymmetry originates from decays of Ψ. To be precise one should refer to the $B - L$ charge. However for simplicity we will work with a lepton asymmetry.

[4] Refs. [9] consider a scenario with GeV scale inverse seesaw where leptogenesis proceeds through oscillations [10].