Electromagnetic leptogenesis at the TeV scale

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Abstract

We construct an explicit model implementing electromagnetic leptogenesis. In a simple extension of the Standard Model, a discrete symmetry forbids the usual decays of the right-handed neutrinos, while allowing for an effective coupling between the left-handed and right-handed neutrinos through the electromagnetic dipole moment. This generates correct leptogenesis with resonant enhancement and also the required neutrino mass via a TeV scale seesaw mechanism. The model is consistent with low energy phenomenology and would have distinct signals in the next generation colliders, and, perhaps even the LHC.

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I. INTRODUCTION

Several recent experiments have cited convincing evidence in favor of non-zero neutrino masses and mixing. While both could be admitted in the Standard Model (SM) by the simple expedient of adding right-handed neutrino fields (omitted, at the inception of the SM, only on account of the then apparent masslessness of the neutrinos), many theoretical challenges persist. Indeed, some authors have claimed neutrino masses to be the evidence of physics beyond the SM. The very smallness of the masses accompanied by the largeness of one of the mixing angles, as also several “anomalies” that appear periodically are indicative of the same. Furthermore, it is not inconceivable that these properties are related to other unexplained puzzles such as the presence of dark matter and/or dark energy, as also the matter-antimatter asymmetry in the universe. It is the last aspect that we shall concentrate on.

The seesaw mechanism\[1\] and the associated mechanism of leptogenesis \[2\] are very attractive means to explain the origin of the small neutrino masses and the baryon asymmetry of the universe. Leptogenesis provides an elegant mechanism to consistently address the observed Baryon Asymmetry in the Universe (BAU) \[3\] in minimal extensions of the SM \[4\]. In standard leptogenesis, at least two of the right handed neutrinos should be heavy with masses close to the GUT scale ($\sim 10^{15}$ GeV) and their out-of-equilibrium lepton number violating decay would create a net lepton asymmetry which, subsequently, would get converted into the observed baryon asymmetry via the $(B + L)$-violating sphaleron interactions \[5, 6\]. At the same time, the inclusion of the right handed (Majorana) fields with lepton number violating Majorana masses can explain the observed smallness of light neutrinos through the seesaw mechanism.

Although the aforementioned scheme is theoretically very attractive, it suffers from the lack of direct detectability, e.g. at high-energy colliders such as the LHC or ILC, or in any other foreseeable experiment. This has, naturally, led to efforts towards alternative routes to leptogenesis. A phenomenologically interesting solution to this problem may be obtained within the framework of resonant leptogenesis \[7–12\]. Characterized by the presence of two (or more) nearly degenerate (moderately) heavy Majorana neutrinos, in such scenarios the corrections to the self-energies play a pivotal role in determining the lepton asymmetry \[5, 13\]. Indeed, if the mass difference be comparable to their decay widths, the resonant
enhancement could render asymmetries to be as large as $O(1)$ [7, 12].

Recently, a very interesting possibility of electromagnetic leptogenesis [14] has been proposed, wherein the source of CP violation has been identified with the electromagnetic dipole moment(s) of the neutrino(s). As is well known, the electric neutrality of the neutrino does not preclude its having non-zero dipole moments. And while, naively, the presence of a magnetic dipole moment would seem to call for the presence of a nonzero mass, even this is not strictly necessary[15]. Originally mooted to account for the then apparent correlation of the solar neutrino flux with the sunspot activity, various schemes have been proposed to generate large magnetic moments for neutrinos [16, 17]. It should be noted at this stage that while Dirac neutrinos can have both direct and transition magnetic moments, only the latter are allowed for Majorana neutrinos. For a collection of neutrino fields of the same chirality, the most general form of such couplings is given by $\bar{\nu}^c_j(\mu_{jk} + i\gamma_5\mathcal{D}_{jk})\sigma_{\alpha\beta}\nu_kB^{\alpha\beta}$, where $B^{\alpha\beta}$ denotes the $U(1)$ field strength tensor. The magnetic and electric transition moment matrices, $\mu_{jk}$ and $\mathcal{D}_{jk}$, each need to be antisymmetric. For two Majorana neutrinos combining to give a Dirac particle, the resultant matrices, clearly, do not suffer from such restrictions.

The aforementioned dimension-five operators are, presumably, generated by some new physics operative beyond the electroweak scale. With $CP$-violation being encoded in the structure of the dipole moments, the decays of heavier neutrinos to lighter ones and a photon, can, in principle, lead to a lepton asymmetry in the universe. Although the proposal is a very interesting one, thus far it has not been incorporated in any realistic model. Indeed, the plethora of constraints suggests that some amount of fine tuning would be unavoidable in any realistic model. In this paper, we discuss the generic problems of any models for electromagnetic leptogenesis and suggest possible means to evade them. Considering all these issues, we point out that on allowing some fine tuning and imposing the resonant condition it may be possible to construct models of resonant electromagnetic leptogenesis, but that direct detection would need at least few more years.

II. THE MODEL

Retaining the gauge symmetry of the SM, we augment the fermion content by including three right-handed singlet fields $N_{iR}$ and, in addition, a singly charged vector-like fermion
Also added are a singly charged scalar \((H^+)\) and a pair of Higgs doublets \((\Sigma, D)\). In keeping with our stated paradigm of only one new scale, all the new masses are assumed to be around a few TeVs. While it could be arranged that all these masses arise from the vacuum expectation value of a single scalar field, for simplicity, we incorporate explicit mass terms. The entire particle content, along with the quantum number assignments, is displayed in Table I.

**Table I: Particle content of the proposed Model**

| Field | \(SU(3)_C \times SU(2)_L \times U(1)_Y\) | \(Z_2\) |
|-------|----------------------------------------|-------|
| Fermions | \(Q_L \equiv (u,d)^T_L\) | \((3, 2, 1/6)\) | + |
|        | \(u_R\) | \((3, 1, 2/3)\) | + |
|        | \(d_R\) | \((3, 1, -1/3)\) | + |
|        | \(\ell_L \equiv (\nu, e)^T_L\) | \((1, 2, -1/2)\) | + |
|        | \(e_R\) | \((1, 1, -1)\) | + |
|        | \(E_L\) | \((1, 1, -1)\) | - |
|        | \(E_R\) | \((1, 1, -1)\) | - |
|        | \(N_R\) | \((1, 1, 0)\) | - |
| Scalars | \(\Phi\) | \((1, 2, +1/2)\) | + |
|        | \(\Sigma\) | \((1, 2, +1/2)\) | - |
|        | \(D\) | \((1, 2, +1/2)\) | - |
|        | \(H^+\) | \((1, 1, +1)\) | + |

At this stage, we are faced with a problem generic to electromagnetic leptogenesis. While the effective \(\bar{N} \nu \gamma\) coupling has to be present (so as to allow the mandatory \(N \rightarrow \nu + \gamma\)), the coupling of this fermion pair to the SM Higgs, viz. \(\bar{N} \ell \Phi\), needs to be highly suppressed on two counts: (i) to ensure that the light neutrino mass, accruing from the seesaw mechanism, is not too large, and (ii) to prevent the \(N\) from decaying dominantly to \(\ell + \Phi\). While this could, nominally, be ensured by invoking some symmetry wherein the photon and the \(\Phi\) transform differently, such an assignment would adversely impact the phenomenology of the charged particles. We, rather, choose to introduce a discrete \(Z_2\) symmetry. All of the SM particles as well as the charged singlet scalar \(H^+\) are even under this \(Z_2\) symmetry, while
the rest are odd (see Table I).

The $Z_2$ symmetry allows both the (effective) Majorana mass terms $\bar{\nu}^c \nu$ and $\bar{N}^c N$ but the former is precluded if we limit ourselves to a renormalizable Lagrangian. On the other hand, the coupling of the neutrinos with the SM Higgs $\Phi$, namely a term of the form $\bar{N} \ell \Phi$ is disallowed, thereby preventing an effective Dirac mass term of the form $\bar{N} \nu$. More importantly, it also forbids the magnetic moment term $\bar{N} \nu \gamma$. Each of these can be generated only when the $Z_2$ is broken. Rather than break it spontaneously, and thereby risk domain walls, we choose to break it explicitly, but only through a soft term. While preserving the essential features of the model, this, then, allows the generation of both Dirac neutrino mass terms as well as magnetic moments and, thereby, driving resonant leptogenesis successfully.

While the Yukawa Lagrangian for the quarks remains unchanged from the SM, that for the leptonic sector can be written as

$$\mathcal{L}_{\text{Yuk}} \ni \left[ y_H N_R E_L H^+ + y_{\Sigma} \bar{\ell}_L \Sigma E_R + y_D \bar{\ell}_L D E_R ight. + h_{\Sigma} \bar{\ell}_L \Sigma N_R + h_D \bar{\ell}_L D N_R + y_e \bar{\ell}_L \Phi e_R + h.c. \right]$$

$$+ \left[ \frac{1}{2} (N_R)^C M_N N_R - M_E \bar{E}_R E_L + h.c. \right]$$

where the last two terms $(M_N, M_E)$ represent gauge- and $Z_2$--invariant bare mass matrices.

In the above, $\Phi = i \sigma_2 \Phi^*$ (similarly for $\bar{\Phi}$ and $\Sigma$) with $y_H$, $y_{\Sigma}$, $y_D$, $h_{\Sigma}$ and $h_D$ being the Yukawa coupling matrices.

The scalar potential can be parametrized as

$$V(\Phi, \Sigma, D, H^+) = -\mu^2_\Phi |\Phi|^2 + m^2_D |\Sigma|^2 + m^2_D |D|^2 + m^2_H |H|^2 + \lambda_1 |\Phi|^4 + \lambda_2 |\Sigma|^4$$

$$+ \lambda_3 |D|^4 + \lambda_6 |H|^4 + \lambda_{\Phi H} (\Phi^\dagger \Phi) |H|^2 + \lambda_{DH} (D^\dagger D) |H|^2$$

$$+ \lambda_{\Sigma H} (\Sigma^\dagger \Sigma) |H|^2 + \lambda_{D \Sigma H} (D^\dagger \Sigma) |H|^2 + \frac{\lambda_{\Phi \Sigma}}{2} [(\Phi^\dagger \Sigma)^2 + h.c.]$$

$$+ \lambda_{D E} (D^\dagger \Sigma) (\Phi^\dagger \Phi) + f_1 (\Phi^\dagger \Phi) (D^\dagger D) + f_2 (\Phi^\dagger \Phi) (\Sigma^\dagger \Sigma)$$

$$+ f_3 |\Phi^\dagger D|^2 + f_4 |\Phi^\dagger \Sigma|^2 + f_5 (D^\dagger D) (\Sigma^\dagger \Sigma) + f_6 |D^\dagger \Sigma|^2$$

$$+ \left[ \mu_s \Sigma \cdot D (H^+)^* + h.c. \right] .$$

Note that the two fields $D$ and $\Sigma$ are being ascribed a positive mass-squared each so that $Z_2$ is left unbroken at this stage. Furthermore, we assume that $m_{2,3}$ are large enough ($\sim \mathcal{O}(10 \text{ TeV})$) so that the decays $N \to \nu + D/\Sigma$ are kinematically disallowed.
As argued earlier, the $Z_2$ symmetry needs to be broken, and we achieve this through an explicit soft term. This has the advantage of obviating any domain wall problem without introducing any qualitative changes to the rest of the phenomenology. To this end, we posit terms of the form

$$V_{soft} = \mu^2_{soft}\Phi^\dagger D + \cdots$$

without going into their origin. It should be noted that although this solves the problem, in a realistic model one must explain the origin of such terms, which is somewhat nontrivial and may plague the model. The ellipses above denote other possible terms, such as $\Phi^\dagger \Sigma$ etc. that do not concern us directly. The scale of the soft symmetry breaking $\mu_{soft}$ needs to be significantly lower than the electroweak symmetry breaking scale. This naturally leads to a large gradation in the vacuum expectation values, namely $\langle D \rangle, \langle \Sigma \rangle \ll \langle \Phi \rangle$. The breaking of the $Z_2$ symmetry allows for a non-zero value of the effective magnetic moment term $\bar{N}\ell\gamma$, which is necessary for leptogenesis to go through. Also introduced is a Dirac mass term $\bar{N}\nu \langle D \rangle$. On the other hand, this breaking now permits the decay $N \rightarrow \nu + \Phi_0$ which proceeds through the mixing of $\Phi$ with $D$ and/or $\Sigma$. The twin facts of $N$ being heavy and $\Phi_0$ being light (unitarity of the SM as well as consonance with LEP data) implies that this cannot be wished away on kinematic grounds. Note, however, that this interaction is suppressed by a factor of $\langle D \rangle / \langle \Phi \rangle$ and, as we show in the next section, a value commensurate with light neutrino masses provides adequate suppression.

### III. NEUTRINO MASS

An exact $Z_2$ symmetry in the Lagrangian prevents the Yukawa term $\bar{\ell}\Phi N$. On the other hand, the fact of $m^2_{2,3} > 0$ prevents a vacuum expectation value for both the fields that do couple to the $\bar{\ell}N$ current, namely $D$ and $\Sigma$. Consequently, there is no Dirac neutrino mass at this level. However, once the soft-breaking term of eqn.(3) is included, the field $D$ may receive a non-zero vev, despite positive $m^2_{2,3}$. This, in turn, gives a Dirac mass to the neutrinos viz.

$$M_{\text{Dirac}} = h_D \langle D \rangle = h_D v_D.$$  

(4)
This, together with the Majorana mass term $M_N$ for the heavy right-chirality fields, gives rise to a light neutrino Majorana mass via type-I seesaw mechanism, viz.

$$m_\nu = M_{\text{Dirac}} M_N^{-1} M_{\text{Dirac}}.$$  \hspace{1cm} (5)

For the choice of parameters we are interested in, $M_D \sim 10^{-3} h_D v \sim 10^{-4}$ GeV ($v = \langle \Phi \rangle \sim 100$ GeV and $h_D \sim 0.001$). The right-handed neutrinos are heavier than the SM Higgs scalar. For $M_N \sim$ few TeV, this gives the correct magnitude of the light neutrino masses, namely $m_\nu \sim 0.1$ eV. The hierarchy of masses could be obtained because of the different values of the elements of the matrices $M_N$ and $h_D$.

IV. DIPOLE COUPLING BETWEEN LIGHT AND HEAVY NEUTRINOS

The effective Lagrangian describing the interaction between photon and the light–heavy ($\bar{\nu} N$) neutrino-current can, in general, be parametrized as

$$\mathcal{L}_{EM} = \overline{\nu}_{Lj} \lambda_{jk} \sigma_{\alpha \beta} P_R N_k F^{\alpha \beta} + \text{h.c.}$$  \hspace{1cm} (6)

The effective coupling matrix $\lambda_{jk}$ is, in general, a complex one, and needs to be calculated in terms of the parameters of the model. The Feynman diagrams which will quantify the EMDM coupling strength are shown in Fig. 1. In Ref.\cite{14}, no concrete model was suggested wherein the numbers required for successful leptogenesis could arise naturally. The main motivation of this paper is to show that it is possible to construct a simple extension of the SM, where it will be possible to calculate this effective coupling, which will lead to resonant electromagnetic leptogenesis. It should, however, be noted that, without the resonance condition, it is not possible to have the correct amount of leptogenesis in these models in view of the smallness of the effective couplings.

The effective dimension-5 coupling constant matrix $\lambda$ can, thus, be expressed in a simple form under the assumption of almost equal mass for the particles in the loop ($M_E \sim M_H \sim M_\Sigma \sim M_{eq}$) as

$$\lambda = - \frac{y_{\Sigma}^* y_H \mu_s v_D}{64 \pi^2 M_{eq}^3}.$$  \hspace{1cm} (7)

For a representative set of parameters, namely $M_N \sim$ few TeV, $M_{eq} \sim$ TeV, $y_\Sigma = y_H \sim \mathcal{O}(1)$, $\mu_s \sim 10$ GeV and $v_D = 0.1$ GeV, we have

$$\lambda \sim 10^{-12} \text{GeV}^{-1}.$$
FIG. 1: Feynman diagrams leading to the effective EMDM coupling strength between light neutrino $\nu_j$ and heavy Majorana neutrino $N_k$.

Note that such values are typical for each of the terms $\lambda_{jk}$, while the exact values would depend on the exact flavour structure. However, large hierarchies and/or texture zeroes are unexpected.

Now we shall investigate the viability of electromagnetic leptogenesis. We must first check that the out-of-equilibrium decay of the RH neutrinos can give rise to a nonzero $CP$ asymmetry under the most general situations. In addition, it is also necessary to examine whether the parameters considered in our model can produce an asymmetry of the correct magnitude via the dimension-five dipole moment operator through the self-energy enhancement.

V. RESONANT ELECTROMAGNETIC LEPTOGENESIS

As has been described above, leptogenesis, in this scenario, is driven by the electromagnetic dipole moment terms appearing in the effective Lagrangian. Specifically, the lepton asymmetry is generated by the CP-violating decays of heavy singlet neutrinos to the SM-like light neutrinos and a photon. As should be apparent from the discussion in the last section,
the size of the EMDM that is generated and the extent of CP-violation in them is inadequate for thermal leptogenesis. Indeed, this is a generic problem for all models of electromagnetic leptogenesis that seek to be consistent with observed physics and yet be natural.

Given this, we investigate the possibility of a resonant enhancement. As is well-known, this mechanism is contingent upon the existence of at least two neutrino species that are very closely degenerate, and this is what we shall assume. Aesthetically, the extent of degeneracy needed may seem uncomfortable. While it can, in principle, be motivated on the imposition of additional global symmetries, it should be noted that, in all models of resonant leptogenesis, the subsequent breaking of the same would, naturally, lead to a lifting of the degeneracy by a degree that negates the conditions for resonant enhancement. Hence, rather than introduce additional symmetries and a host of fields for additional mechanisms of compressing the spectrum adequately, we just assume that the said heavy neutrinos are highly degenerate. We will return to this point later in this section.

The key quantity of interest is the CP-asymmetry for the decay of $N_k$ to a photon and a light neutrino given by

$$\varepsilon_k \equiv \frac{\Gamma(N_k \to \nu \gamma) - \Gamma(N_k \to \nu \gamma)}{\Gamma(N_k \to \nu \gamma) + \Gamma(N_k \to \nu \gamma)}.$$  (8)

We begin by calculating the lowest order contribution to the decay rate $\Gamma(N_k \to \nu_j \gamma)$. Since we are interested in energy scales above the electroweak phase transition, we shall identify the light neutrino $\nu$ to be a massless left-handed (SM-like) state while $N'$s are assumed to have Majorana mass of around 1 TeV. Driven by the effective Lagrangian of eqn.(6), the lowest order decay rate is, thus, given by

$$\Gamma(N_k \to \nu \gamma) = \frac{\lambda^\dagger \lambda}{4\pi} M_k^3,$$  (9)

where all species of (massless) neutrinos $\nu_i$ have been summed over. For effectively creating a lepton asymmetry of the universe, the decay of, say $N_1$, should be out of equilibrium, the necessary condition for which is described by $\Gamma(N_1) \lesssim H(T) |_{T=M_1}$ where $H(T) = 1.67 \, g^\star_1 T^2 / M_{Pl}$ is the Hubble parameter at that particular epoch with the Planck mass $M_{Pl} \simeq 1.2 \times 10^{19}$ GeV and the number of relativistic degrees of freedom $g^\star \simeq 100$. With the operative temperature $T \sim M_1$, we then have

$$\frac{\lambda^\dagger \lambda}{4\pi} M_1^3 \lesssim 1.67 g^\star_1^{1/2} \frac{M_1^2}{M_{Pl}}.$$  (10)

This is satisfied by the effective EMDM coupling $\lambda$ with $M_1 \sim$ few TeV, for the choice of parameters we have considered here.
The next task is to calculate the interference terms between the tree level process and the one-loop diagrams with on shell intermediate states as shown in Fig. 2. The usual contributions to lepton asymmetry coming from vertex diagram is found to be very small, i.e., \( \varepsilon_1 = (\lambda^2/4\pi)M_1^3 \sim (10^{-23} \text{ GeV}^{-2}) M_1^3 \sim 10^{-14} \) when \( M_1 \) is at the TeV scale and, hence, can be neglected. So, the self energy contribution will only be considered during the rest of the discussion.

\[
\varepsilon_k = -\frac{M_k^3}{2\pi (\lambda^\dagger \lambda)_{kk}} \sum_{m \neq k} \text{Im} \left[ (\lambda^\dagger \lambda)_{km}^2 \right] \frac{(M_k^2 - M_m^2)M_m}{(M_m^2 - M_k^2)^2 + M_k^2 \Gamma_m^2} \\
= -2 \frac{M_k^3}{(\lambda^\dagger \lambda)_{kk}} \sum_{m \neq k} \frac{\text{Im} \left[ (\lambda^\dagger \lambda)_{km}^2 \right]}{M_m^2 (\lambda^\dagger \lambda)_{mm}} \frac{(M_m^2 - M_k^2)\Gamma_m}{(M_m^2 - M_k^2)^2 + M_k^2 \Gamma_m^2},
\]

(11)

where the expression for the total width \( \Gamma_m \) has been used to get to the second line. Consider the case where \( M_1 \sim M_2 \ll M_3 \). From eqn.(9), it is clear that \( \Gamma_1 \sim \Gamma_2 \) for nearly degenerate right handed neutrinos with masses \( M_{1,2} \). Hence, \( \Gamma_2 \approx \Gamma_1 = (\lambda^\dagger \lambda)_{22} M_2^2/(4 \pi) \) and, for the \( N_1 \)-dominated case, the CP-asymmetry is

\[
\varepsilon_1 = -\frac{M_1^3}{2\pi} \frac{\sum_{m \neq 1} \text{Im} \left[ (\lambda^\dagger \lambda)_{1m}^2 \right]}{(\lambda^\dagger \lambda)_{11}^2} \frac{(M_2^2 - M_1^2)M_1M_2}{(M_2^2 - M_1^2)^2 + M_1^2 \Gamma_1^2}.
\]

(12)

As \( \Gamma_2 \ll \left| M_1 - M_2 \right| \), even when the heavy neutrinos are quite degenerate, this further
simplifies to

$$
\varepsilon_1 \approx -\frac{M_1^2}{2\pi} \sum_{m \neq 1} \text{Im} \left[ \frac{[(\lambda^\dagger \lambda)_1^2]_m}{(\lambda^\dagger \lambda)^2_{11}} \right] \frac{M_1 M_2}{M_2^2 - M_1^2}. \tag{13}
$$

Clearly, in the almost degenerate case, \( \varepsilon_1 \) is resonantly enhanced. Indeed, writing \( M_2^2 - M_1^2 \sim 2M_2(M_2 - M_1) \), we have

$$
\varepsilon_1 \approx -\frac{M_1^2}{4\pi} \sum_{m \neq 1} \text{Im} \left[ \frac{[(\lambda^\dagger \lambda)_1^2]_m}{(\lambda^\dagger \lambda)^2_{11}} \right] \mathcal{R} \tag{14}
$$

where \( \mathcal{R} \equiv M_1/|M_1 - M_2| \).

As described above, a non-zero \( \varepsilon_1 \) can give rise to a net lepton number asymmetry in the Universe, provided its expansion rate is larger than the decay rate of \( N_1 \). The nonperturbative sphaleron interaction may partially convert this lepton number asymmetry into a net baryon number asymmetry [6],

$$
\eta_B \simeq -2.96 \times 10^{-2} \varepsilon_1 k
$$

where \( k \) is the efficiency factor measuring the washout effects associated with the out-of-equilibrium decays of \( N_1 \). In our model, \( k \sim \mathcal{O}(10^{-3}) \). We, thus, need \( |\varepsilon_1| \sim 10^{-5} \) to generate the requisite baryon asymmetry in the Universe. This is achieved if \( |M_2 - M_1| \lesssim 10^{-7} \text{ GeV} \) where the mass of the right handed Majorana neutrinos is around TeV scale.

While such a small mass difference may seem unnatural, it need not be so. To start with, let us assume that some symmetry forces them to be exactly degenerate at the tree level. The question of interest, then, is the extent to which this degeneracy is lifted by quantum corrections. To this end, consider a diagram with a vertex \( \lambda_{HD}(D^\dagger D)(H^\dagger H) \) attached to the singly charged scalar \( H \) which runs in the loop contributing to the neutrino mass. This engenders a finite contribution to the mass and the consequent splitting is

$$
\Delta M_R \sim \frac{\lambda_{HD} y_H^* y_H}{(4\pi)^2} \frac{\langle D \rangle^2}{4M_E} \tag{15}
$$

Since \( M_E \sim \mathcal{O}(1 \text{ TeV}) \), \( v_D = \langle D \rangle \sim \mathcal{O}(0.1 \text{GeV}) \) and \( y_H \sim \mathcal{O}(1) \), a moderate value of \( \lambda_{HD} \) will generate the requisite mass splitting.

Before closing, it may be instructive to make a comparison with the standard leptogenesis scenario where the CP asymmetry is generated via the decay \( N_R \rightarrow \ell \varphi \), where \( \varphi \) denotes a generic scalar. The decay rate is given by

$$
\Gamma_{\text{standard}} = \frac{y^2 M_N}{4\pi} \left( 1 - \frac{m_{\varphi}^2}{m_N^2} \right)^2
$$

11
with $y$ being the relevant Yukawa coupling. To have leptogenesis proceed dominantly via the electromagnetic decay, the above decay rate should be smaller than the corresponding rate into the electromagnetic channel. This requires $y \lesssim 10^{-8}$ for $M_1 \sim \mathcal{O}$(few TeV). In the present case, $y$ would refer to the effective $\bar{N} \ell \Phi$ coupling. Since this is generated only through $V_{soft}$, we have (for $\mu_{soft} \sim 10$ GeV and $m_D \sim 10$ TeV)

$$y_{eff} \approx h_D \left( \frac{\mu_{soft}^2}{m_D^2} \right) \sim 10^{-9},$$

which is consistent with the value of $y$ estimated above. However, this does introduce some amount of fine tuning in the model and the parameters have to be marginally adjusted to allow the electromagnetic leptogenesis, a fact that we believe is generic to any realistic model of electromagnetic leptogenesis.

VI. SUMMARY

The idea of electromagnetic leptogenesis is a very interesting and appealing alternative to the standard scenario of leptogenesis. We have shown that it is indeed possible to have a viable model for leptogenesis proceeding through such a channel. However, there are several generic problems associated with the construction of any model for electromagnetic leptogenesis. Highlighting these problems, we showed that the choice of parameters has to be a fine tuned one in the sense that deviations from the values chosen may not lead to successful predictions. Also, to have leptogenesis proceed via the electromagnetic decay channel rather than the standard channel involving the $N\ell\phi$ Yukawa coupling, it is necessary to have the Yukawa coupling highly suppressed. This, in a way, leads to some additional fine tuning the stability of which under radiative corrections is slightly suspect. Moreover, the model works only if there is resonant enhancement of the asymmetry, which requires almost degenerate heavy neutrinos. While this might seem an additional fine tuning, it is not quite so, as it is essentially the same as that responsible for the heavy neutrino decaying electromagnetically rather than to a scalar. Nonetheless, such a fine tuning, perhaps, is a generic feature of any realistic model of electromagnetic leptogenesis and warrants a study in its own right. In spite of all the fine tuning of parameters, it is difficult to have any reasonable choice of parameters that allows for immediate detection of new physics at the LHC. While detection of the heavy lepton ($E$) and some of the scalars is, in principle, possible once the
LHC starts operating at its design energy and luminosity, it would, nonetheless, need a few years of accumulation. The CLIC, on the other hand, would stand a very good chance of directly observing these states. Deciphering the structure of the theory, unfortunately, is likely to prove nearly impossible.

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