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Electromagnetic leptogenesis

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We present a new leptogenesis scenario, where the lepton asymmetry is generated by *CP*-violating decays of heavy electroweak singlet neutrinos via electromagnetic dipole moment couplings to the ordinary light neutrinos. Akin to the usual scenario where the decays are mediated through Yukawa interactions, we have shown, by explicit calculations, that the desired asymmetry can be produced through the interference of the corresponding tree-level and one-loop decay amplitudes involving the effective dipole moment operators. We also find that the relationship of the leptogenesis scale to the light neutrino masses is similar to that for the standard Yukawa-mediated mechanism.

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I. INTRODUCTION

Baryogenesis via thermal leptogenesis [1] provides an elegant explanation of the cosmic baryon asymmetry [2]. The conventional setup involves minimally extending the standard model (SM) by adding three heavy right-handed (RH) Majorana neutrinos, which are electroweak singlets, and allowing them to interact with ordinary left-handed (LH) lepton doublets via complex Yukawa couplings. As a result, when the heavy neutrinos decay out-of-equilibrium in a *CP*-violating way, a lepton asymmetry is generated in the early Universe. This asymmetry is then partially converted by sphaleron processes to the baryon asymmetry we detect today.

An attractive aspect of this scenario is that it naturally allows the implementation of the see-saw mechanism [3] which gives the light neutrinos a tiny but nonzero mass.[1] Consequently, a strong link between some neutrino properties and successful asymmetry generation can be established. For instance, to simultaneously obtain a sufficiently large lepton asymmetry and the correct light neutrino masses, the heavy neutrino masses must be larger than 10^9 GeV in most viable leptogenesis scenarios [6].

In this paper, we will consider a scenario where leptogenesis is mediated not by the standard Yukawa couplings, but instead by electromagnetic dipole moment couplings. One motivation for introducing such couplings is to explore whether new sources of *CP* violation can lead to a significant lepton asymmetry. A natural question we may ask is whether the introduction of *CP*-violating dipole moment couplings will allow leptogenesis to occur at a lower scale, closer to experimentally accessible energies. However, we find this not to be possible due to the connection between the dipole moment couplings and the neutrino mass.

[1]Note that similar results can be instead achieved with the addition of heavy Higgs triplets [4], or with other mechanisms [5].

The general form of a dipole moment coupling of the light neutrinos, ν, to the heavy neutrinos, N, is given by $\bar{\nu}(\mu + id\gamma_5)\sigma^{\alpha\beta}NF_{\alpha\beta}$, where μ and d are the magnetic and electric transition moments, respectively. These dimension-five effective operators may be assumed to be generated by some new physics beyond the electroweak scale. While we do not speculate on the nature of this new physics, the inclusion of these operators permits a new leptogenesis mechanism. Radiative decays of the heavy neutrinos, $N \to \nu + \gamma$, can now produce the required lepton asymmetry in the early Universe, provided that the complex electromagnetic dipole moment couplings violate *CP*.

Below, we outline the relevant properties of the electromagnetic dipole moment (EMDM) couplings, and discuss the necessary requirements for a decay process to manifestly violate *CP*. We shall then explicitly calculate the decay rates and *CP* asymmetry for electromagnetic leptogenesis, and compare with the standard Yukawa-mediated leptogenesis scenario.

II. ELECTROMAGNETIC COUPLING BETWEEN LIGHT AND HEAVY NEUTRINOS

We extend the particle content of the minimal SM by adding heavy RH neutrinos, N_R, which are assumed to have large Majorana masses. Since we are interested in leptogenesis energy scales above the electroweak phase transition, we will take the usual light neutrinos to be massless LH states. The most general electromagnetic dipole moment coupling of the heavy RH neutrinos to the light LH neutrinos is then given by

$$\mathcal{H}_{\rm EM} = g_{ij}\bar{\nu}_{Li}\sigma^{\alpha\beta}N_{Rj}F_{\alpha\beta} + \text{H.c.}, \tag{1}$$

where g is a complex (dimensionful) matrix, and i, j, are flavor indices. Note that there is only one distinct electromagnetic dipole moment coupling when expressed in term of LH and RH chiral fields (rather than distinct magnetic and electric dipole moment terms) since $\gamma_5 P_{L,R} \propto P_{L,R}$.

III. *CP* VIOLATION IN DECAYS

If the dipole coupling of Eq. (1) had less-constrained chiral structure, then for a given ν_i and N_j there could be independent magnetic and electric transition moments, μ_{ij} and d_{ij}. It might be thought that in this situation, a tree-level interference between the amplitude induced by μ_{ij} and that induced by d_{ij} could lead to a *CP*-violating difference between the rates for $N_j \to \nu_i + \gamma$ and its *CP* conjugate. However, there can never be a difference between the rates for *CP*-conjugate decay modes until one goes beyond first order in the underlying Hamiltonian. This fact is well known, but it is interesting to see that it can be proved very simply by using *CPT* invariance. Consider, in the rest frame of the parent particle Q, the decay $Q \to a_1 + a_2 + \cdots$. If *CPT* invariance holds, the amplitude for this decay obeys the constraint

$$|\langle a_1(\vec{p}_1, \lambda_1) a_2(\vec{p}_2, \lambda_2) \cdots |\mathcal{T}| Q(\hat{m})\rangle|^2 = |\langle \bar{a}_1(\vec{p}_1, -\lambda_1) \bar{a}_2(\vec{p}_2, -\lambda_2) \cdots |\mathcal{T}^\dagger| \bar{Q}(-\hat{m})\rangle|^2. \quad (2)$$

Here, $\vec{p}_i$ and λ_i are, respectively, the momentum and helicity of daughter particle a_i, $\hat{m}$ is the z-axis projection of the spin of Q, and $\mathcal{T}$ is the transition operator for the decay. If S is the S-matrix operator, $\mathcal{T} = i(S - I)$. To first order in the Hamiltonian $\mathcal{H}$ for the system, $\mathcal{T} = \mathcal{H}$, so that, to this order, $\mathcal{T}^\dagger = \mathcal{T}$. From the latter relation and Eq. (2), it follows that, after summing over the final helicities and integrating over the outgoing momenta,

$$\Gamma[\bar{Q} \to \bar{a}_1 + \bar{a}_2 + \cdots] = \Gamma[Q \to a_1 + a_2 + \cdots]. \quad (3)$$

This equality must hold to first order in $\mathcal{H}$ regardless of whether $\mathcal{H}$ contains numerous terms and *CP*-violating coupling constants.

In the special case of a two-body decay, $Q \to a_1 + a_2$, we have $\vec{p}_1 = -\vec{p}_2 \equiv \vec{p}$. For this case, let us rotate the system of particles on the right-hand side of Eq. (2) by 180° about the axis perpendicular to the z-axis and to $\vec{p}$. Equation (2) then states that, to first order in $\mathcal{H}$ (so that $\mathcal{T} = \mathcal{T}^\dagger$),

$$|\langle a_1(\vec{p}, \lambda_1) a_2(-\vec{p}, \lambda_2) |\mathcal{T}| Q(\hat{m})\rangle|^2 = |\langle \bar{a}_1(-\vec{p}, -\lambda_1) \bar{a}_2(\vec{p}, -\lambda_2) |\mathcal{T}| \bar{Q}(\hat{m})\rangle|^2. \quad (4)$$

The processes whose amplitudes appear on the two sides of this constraint are the *CP*-mirror images of each other. Thus, in two-body decays, to first order in $\mathcal{H}$, the rates for *CP*-mirror-image decay processes must be equal even before one sums over final helicities and integrates over outgoing momenta.

We conclude that *CP*-violating rate differences between *CP*-conjugate electromagnetic decays of heavy neutrinos can only appear once amplitudes involving loops are included.

IV. A TOY MODEL

In this section we illustrate, by means of a toy model, the viability of generating a lepton asymmetry through EMDM interactions between ordinary light neutrinos and the postulated heavy Majorana neutrinos in the early Universe. In order to illustrate the physics as transparently as possible, we begin by considering a simplistic model which is not invariant under the SM gauge symmetry, $SU(2)_L \times U(1)_Y$, but is instead invariant only under the electromagnetic symmetry $U(1)_Q$. We will generalize to a realistic model in which invariance under the SM gauge group is enforced in Sec. V below.

We assume the EMDM couplings are generated by new physics at an energy scale $\Lambda > M$, where M denotes a heavy Majorana neutrino mass. We work with an effective theory that is valid below the scale Λ, obtained after integrating out all new heavy degrees of freedom. The lowest-dimension EMDM operator of interest in such a scenario is given by

$$\mathcal{L}_{\rm EM}^{\rm 5D} = -\frac{1}{\Lambda} e^{-i\varphi_k/2} \lambda_{jk} \bar{\nu}_{Lj} \sigma^{\alpha\beta} P_R N_k F_{\alpha\beta} + \text{H.c.}, \quad (5)$$

where $j = e, \mu, \tau$ and $k = 1, 2, 3$. The electromagnetic field strength tensor is $F_{\alpha\beta} = \partial_\alpha A_\beta - \partial_\beta A_\alpha$, with A_α being the photon field. The light neutrinos are denoted by ν_j, while N_k is the heavy neutrino field in the mass eigenbasis which satisfies the Majorana condition: $N_k = e^{i\varphi_k} N_k^c$, for some arbitrary phase φ_k. We have defined λ as a dimensionless 3×3 matrix of complex coupling constants, while Λ is the cutoff scale of our effective theory. For convenience, we have factored out $e^{-i\varphi_k/2}$ for each N_k in Eq. (5) so that the Majorana phases will not appear explicitly in any of our final expressions.

To ascertain whether leptogenesis is possible, the key quantity of interest is the *CP* asymmetry in the decays of N_k:

$$\varepsilon_{k,j}^{(5)} = \frac{\Gamma_{(N_k \to \nu_j \gamma)} - \Gamma_{(N_k \to \bar{\nu}_j \gamma)}}{\Gamma_{(N_k \to \nu \gamma)} + \Gamma_{(N_k \to \bar{\nu} \gamma)}}, \quad (6)$$

where $\Gamma_{(N_k \to \nu\gamma)} \equiv \sum_j \Gamma_{(N_k \to \nu_j \gamma)}$. The lowest-order contribution to the decay rate, shown in Fig. 1, is given by

$$\Gamma_{(N_k \to \nu\gamma)} = \Gamma_{(N_k \to \bar{\nu}\gamma)} = \frac{(\lambda^\dagger \lambda)_{kk}}{4\pi} \frac{M_k^3}{\Lambda^2}. \quad (7)$$

The leading contribution to the *CP* asymmetry, $\varepsilon_{k,j}^{(5)}$, comes from the interference of the tree-level process of Fig. 1

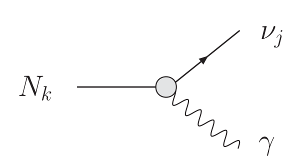

FIG. 1. The tree-level diagram for the decay of N_k via the EMDM interaction of Eq. (5).

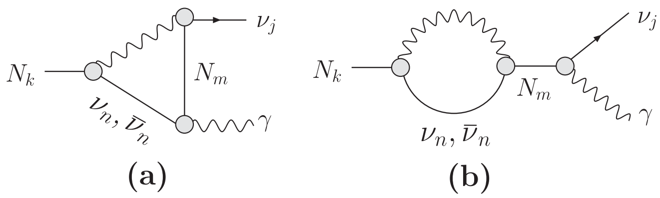

FIG. 2. (a) Vertex and (b) self-energy diagrams which contribute to the CP asymmetry of N_k decay via the interaction of Eq. (5). Note that since weak isospin is violated in this model, both ν_n and $\bar{\nu}_n$ are allowed in the loop of (a).

with the 1-loop diagrams with on-shell intermediate states depicted in Fig. 2. As with standard (Yukawa interaction) leptogenesis, the CP asymmetry receives two contributions: the self-energy and vertex correction. We have calculated these explicitly and obtain

$$\varepsilon^{(5)}_{\text{self-}k,j} = \frac{-(M_k/\Lambda)^2}{2\pi(\lambda^\dagger\lambda)_{kk}} \sum_{m\neq k} \text{Im}\bigg[\lambda^*_{jk}\lambda_{jm} \times \bigg\{(\lambda^\dagger\lambda)_{km}\frac{\sqrt{z}}{1-z} + (\lambda^\dagger\lambda)_{mk}\frac{1}{1-z}\bigg\}\bigg], \quad (8)$$

for the self-energy contribution and

$$\varepsilon^{(5)}_{\text{vert-}k,j} = \frac{-(M_k/\Lambda)^2}{2\pi(\lambda^\dagger\lambda)_{kk}} \sum_{m\neq k} \text{Im}[\lambda^*_{jk}\lambda_{jm}(\lambda^\dagger\lambda)_{km} f(z)] \quad \text{with}$$

$$f(z) = \sqrt{z}\bigg[1 + 2z\bigg(1 - (z+1)\ln\bigg[\frac{z+1}{z}\bigg]\bigg)\bigg], \quad (9)$$

for the vertex piece, where $z \equiv M_m^2/M_k^2$. Note that in Eq. (8) and (9), we have not yet summed over the final lepton flavor j.

The expressions given in Eq. (8) and (9) are very akin to those in standard leptogenesis [7]. The first and second terms of Eq. (8) correspond to the interference terms involving $\bar{\nu}_n$ and ν_n as the intermediate state in Fig. 2(b) respectively. It should be emphasized that upon summing over j (i.e., ignoring flavor effects [8]), the second term vanishes. Furthermore, explicit calculations have shown that in Fig. 2(a), the contribution from the ν_n intermediate state actually evaluates to zero, and hence there is no term proportional to $(\lambda^\dagger\lambda)_{mk}$ in Eq. (9).

The total CP asymmetry from the decays of N_k's into light neutrinos and photons in this model is simply given by $\varepsilon^{(5)}_{k,j} = \varepsilon^{(5)}_{\text{self-}k,j} + \varepsilon^{(5)}_{\text{vert-}k,j}$. This asymmetry would be nonzero as long as there are phases in the coupling matrix λ which cannot be removed by redefinitions of the neutrino fields. As λ is an arbitrary complex matrix, it is not hard to see that one cannot eliminate all the relevant phases to render both Eqs. (8) and (9) zero. Hence, this type of EMDM interaction between light and heavy neutrinos will in general generate a lepton asymmetry in the early Universe. Before discussing the magnitude of this asymmetry, we will first generalize this scenario to one in which the EMDM couplings respect the SM gauge symmetries.

V. A MORE REALISTIC EXTENSION

While the simplistic model in Sec. IV can demonstrate the viability of lepton generation through EMDM operators, it is nonetheless unrealistic as it is incompatible with the SM. We now overcome this by considering only EMDM type operators that respect the SM gauge group. Again, we construct an effective theory by taking the usual minimally extended SM Lagrangian with three generations of heavy Majorana neutrinos, and augmenting it with EMDM operators. The most economical of such operators involving only (the minimally extended) SM fields are of dimension six [9], and the interaction Lagrangian of interest is

$$\mathcal{L}_{\text{EM}} = -\frac{1}{\Lambda^2}e^{-i\varphi_k/2}\bar{\ell}_j[\lambda'_{jk}\phi\sigma^{\alpha\beta}B_{\alpha\beta} + \tilde{\lambda}'_{jk}\tau_i\phi\sigma^{\alpha\beta}W^i_{\alpha\beta}]P_R N_k + \text{H.c.}, \quad (10)$$

where the τ_i are the $SU(2)_L$ generators, $\ell_j = (\nu_j, e_j^-)^T$ is the lepton doublet, and $\phi = (\phi^{0*}, -\phi^-)^T$ is the SM Higgs doublet. The field strength tensors of $U(1)_Y$ and $SU(2)_L$ are given by $B_{\alpha\beta} = \partial_\alpha B_\beta - \partial_\beta B_\alpha$ and $W^i_{\alpha\beta} = \partial_\alpha W^i_\beta - \partial_\beta W^i_\alpha - g\epsilon_{imn}W^m_\alpha W^n_\beta$, respectively, where g' and g are the corresponding coupling constants. Again, Λ is the high-energy cutoff of our effective theory, while the matrices of dimensionless coupling constants, λ' and $\tilde{\lambda}'$, are in general complex.

The higher dimension (nonrenormalizable) operators of Eq. (10) are assumed to be generated at the energy scale Λ, beyond the electroweak scale. The presence of these operators would imply the existence of some new physics at a high energy. After $SU(2)_L \otimes U(1)_Y$ breaking, these operators will give rise to electromagnetic dipole transition moments of N and ν. However, for the purposes of leptogenesis, we are interested here in the regime above the electroweak symmetry breaking scale.

The decay of N will now produce 3-body final states, namely $N_k \to \ell_j\phi W^i_\alpha$ and $N_k \to \ell_j\phi B_\alpha$, as shown in Fig. 3(a). Likewise, the self-energy and vertex corrections now become two-loop processes, an example of which is shown in Fig. 3(b). As before, a lepton asymmetry will be

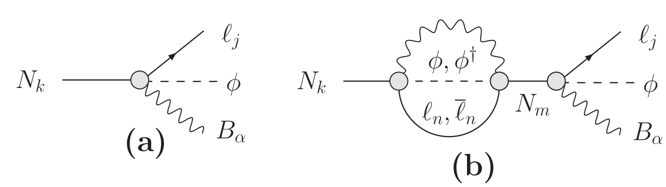

FIG. 3. (a) The tree-level diagram for the the 3-body decay: $N_k \to \ell_j\phi B_\alpha$ induced by the first term in Eq. (10). (b) The corresponding self-energy diagram.

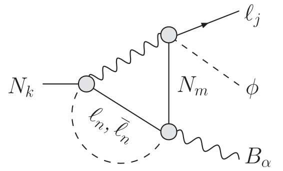

FIG. 4. An example of a vertex correction to the 3-body decay process.

generated through the interference of the tree-level amplitude with the on-shell part of the vertex and self-energy amplitudes.

To demonstrate that the EMDM interactions in this model can indeed generate a lepton asymmetry, we have explicitly calculated the tree-level decay rate, $\Gamma_{(N_k \to \ell_j \phi B_\alpha)}$, and self-energy contribution to the CP asymmetry, $\varepsilon^{(6)}_{\text{self-}k,j}$, for the couplings which involve the hypercharge boson B (see Fig. 3). The $SU(2)$ gauge boson will make a contribution to the decay rate and CP asymmetry of similar magnitude. For simplicity, we have not explicitly calculated the vertex corrections, an example of which is shown in Fig. 4. The vertex diagrams can again have on-shell intermediate states and, barring accidental cancellations, would contribute to the lepton asymmetry generated.

For the diagrams shown in Fig. 3, our explicit calculations of Γ and ε lead to similar forms to those found earlier[2]:

$$\Gamma_{(N_k \to \ell_j \phi B_\alpha)} = \left(\frac{M_k}{8\pi\Lambda}\right)^2 \Gamma_{(N_k \to \ell_j \gamma)}, \tag{11}$$

and

$$|\varepsilon^{(6)}_{\text{self}-k,j}| = \left(\frac{M_k}{8\pi\Lambda}\right)^2 |\varepsilon^{(5)}_{\text{self}-k,j}|, \tag{12}$$

where we must replace $\lambda \to \lambda'$ in the RHS of Eqs. (11) and (12). It is thus possible to generate a nonzero CP asymmetry via the EMDM type interactions of Eq. (10), provided λ' contains complex phases.

VI. NEUTRINO MASS

We now address the connection between the dipole moment operators and the neutrino mass. The new physics that gives rise to the effective operators in Eq. (10) might also be expected to give rise to neutrino mass terms. It is well known that via a careful choice for the new physics, one can obtain large neutrino dipole moments without correspondingly large mass terms [11–15]. However, even if neutrino masses are absent or suppressed at lowest order, radiative corrections involving the dipole moment operators generically induce neutrino mass terms [9,16,17].

[2]It should be noted that these 3-body decay processes are inherently different from those discussed in [10] for these are 2-loop rather than 1-loop diagrams.

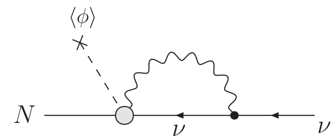

FIG. 5. Contribution to the neutrino Dirac mass, m_D, induced by the electromagnetic dipole moment operator.

As in Ref. [9], the operators in Eq. (10) will lead to a Dirac mass term for the neutrinos via the diagram shown in Fig. 5. There is no model-independent way of calculating the exact size of this neutrino mass contribution, since the exact relationship between the dipole moments and the mass requires a UV completion of the theory (i.e., it depends on the nature of the physics at scale Λ). However, we may estimate the size of the contribution to the neutrino mass using naïve dimensional analysis. The neutrino Dirac mass arising from Fig. 5 is thus estimated to be

$$m_D \sim \frac{\lambda'}{\Lambda^2} \frac{g'}{16\pi^2} \langle\phi\rangle \Lambda^2, \tag{13}$$

where the Λ^2 in the numerator arises from the cutoff of the loop integral, and g' is the gauge coupling constant. This Dirac mass term will lead to a contribution to the light neutrino mass via the see-saw mechanism of

$$m_\nu^A = m_D^T M^{-1} m_D \sim \lambda'^T M^{-1} \lambda' \langle\phi\rangle^2 \left(\frac{g'}{16\pi^2}\right)^2. \tag{14}$$

However, there will also be a direct contribution to Majorana mass of the light neutrinos via the diagram in Fig. 6, which we estimate as

$$m_\nu^B \sim \lambda'^T M \lambda' \frac{\langle\phi\rangle^2}{\Lambda^2} \frac{1}{16\pi^2}. \tag{15}$$

In the standard Yukawa-mediated leptogenesis scenario, the light neutrino masses are linked with the leptogenesis parameters, since the Yukawa coupling constants that control the decay rate of the N also appear in the Dirac mass terms and thus (via the see-saw mechanism) in the ν masses. In order for these Yukawa coupling constants to give rise to the correct values for both ε and m_ν, the heavy

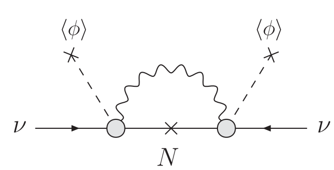

FIG. 6. Contribution to the light neutrino mass, m_ν^B, induced by the electromagnetic dipole moment operator.

neutrinos are required to have masses $M \gtrsim 10^9$ GeV [6].[3] The presence of the neutrino mass term arising from the diagrams of Figs. 5 and 6 signifies that a similar connection will hold for the couplings λ' that control electromagnetic leptogenesis. Hence, we will also need a relatively high N mass scale in order to produce a sufficiently large asymmetry together with sufficiently small ν masses.

VII. DISCUSSION

If we assume there is at least a mild hierarchy in the masses of the heavy neutrinos, the asymmetry will be predominantly generated from the decay of the lightest state, N_1. In Table I, we compare the decay rate, CP asymmetry, and light neutrino masses for electromagnetic leptogenesis with the corresponding expressions for the standard (Yukawa) scenario. The coupling constants and the N_1 mass enter into the expressions for Γ_1, ε, and m_ν in essentially the same way for the two scenarios. Modulo the additional suppression factors in the RHS of the table, we see that the region of viable parameter space for electromagnetic leptogenesis must be similar to that for the standard Yukawa mechanism. It is also clear that although we require $\Lambda > M$ in order for our effective operator approach to be valid, we do not want $\Lambda \gg M$, as it would suppress both Γ and ε. We therefore require a moderate hierarchy between M and Λ in order to obtain an appropriately sized CP asymmetry.

In general, both the Yukawa and EMDM interactions will contribute to the decay rate, CP asymmetry, and neutrino mass. Depending on the relative size of h and λ', either one mechanism will dominate or there will be an interplay between the two. In what follows, we suppose that the Yukawa couplings are negligible. For simplicity, we also ignore the flavor structure of the matrices λ', m_ν, and M, and make the crude assumption that all elements are of similar magnitude. If we then take

$$\Lambda \sim 10M_2 \sim 20M_1 \quad \text{and} \quad \lambda' \sim 35, \tag{16}$$

we obtain an asymmetry of $\varepsilon \sim 10^{-6}$, while smaller values of λ' would lead to a correspondingly smaller asymmetry according to $\varepsilon \propto (\lambda')^2$. We define the decay parameter as

$$K \equiv \Gamma_1/H|_{T=M_1}, \tag{17}$$

where H is the Hubble expansion rate. Note that K controls whether the N_1 decays are in equilibrium, and is also a measure of washout effects via inverse decays. If we now take

[3]This condition may be alleviated in degenerate leptogenesis scenarios, in which the CP asymmetry can be enhanced and thus the N mass lowered through a resonance in the self-energy contribution when $M_k \simeq M_m$ [18].

TABLE I. Comparison of key quantities in standard and electromagnetic leptogenesis, where h denotes Yukawa coupling constants. We have assumed there is at least a mild hierarchy in the masses of the heavy neutrinos, such that the asymmetry is predominantly generated from the decay of the lightest state, N_1.

Yukawa	Electromagnetic
$\Gamma_1 = \frac{1}{8\pi}(h^\dagger h)_{11} M_1$	$\Gamma_1 = \frac{1}{2\pi}(\lambda'^\dagger \lambda')_{11} M_1 \left(\frac{M_1^2}{8\pi\Lambda^2}\right)^2$
$\varepsilon \sim \frac{1}{8\pi}\frac{\mathrm{Im}(h^\dagger h)^2_{1m}}{(h^\dagger h)_{11}}\frac{M_1}{M_m}$	$\varepsilon \sim \frac{1}{2\pi}\frac{\mathrm{Im}(\lambda'^\dagger \lambda')^2_{1m}}{(\lambda'^\dagger \lambda')_{11}}\frac{M_1}{M_m}\left(\frac{M_1^2}{8\pi\Lambda^2}\right)^2$
$m_\nu \sim h^T M^{-1} h \langle\phi\rangle^2$	$m_\nu^A \sim \lambda'^T M^{-1} \lambda' \langle\phi\rangle^2 \left(\frac{g'}{16\pi^2}\right)^2$
	$m_\nu^B \sim \frac{\lambda'^T M \lambda'}{\Lambda^2}\langle\phi\rangle^2 \frac{1}{16\pi^2}$

$$M_1 \sim 5 \times 10^{12}\ \text{GeV} \tag{18}$$

together with Eq. (16), we obtain $K \sim 0.3$. Since $K \ll 1$ ($K \gg 1$) corresponds to the weak (strong) washout regime, this parameter choice should lead to moderate washout. Moreover, for these parameters the light neutrino mass terms become $m_\nu^A \sim 0.04$ eV and $m_\nu^B \sim 0.1$ eV, such that the neutrino mass is dominated by the contribution from the diagram in Fig. 6. For a larger M_1/Λ hierarchy, m_ν^B would be suppressed with respect to m_ν^A, but this is undesirable as the asymmetry ε would also be suppressed by 4 powers of M_1/Λ.

Finally, we note that effective dipole moment interactions of two light neutrino states are induced by our Lagrangian. The largest contributions arise from two-loop diagrams for which we estimate

$$\mu_{\nu_{\rm eff}} \sim \left(\frac{\lambda'}{16\pi^2}\right)^2 \frac{g'}{\Lambda} \simeq 5 \times 10^{-19}\mu_B, \tag{19}$$

where the second approximate equality assumes the parameter values specified above. These induced dipole moments are thus well below current experimental upper limits [19] which are of order $10^{-11}\mu_B$.

VIII. CONCLUSIONS

In summary, we have presented a new leptogenesis mechanism in which electromagnetic dipole moment couplings induce CP-violating decays of heavy RH neutrinos. Via explicit calculation of the decay rates, we have demonstrated that a sufficient asymmetry can be produced through such decays to make this scenario viable. However, since the electromagnetic dipole moment operators induce neutrino mass terms, the leptogenesis parameters are constrained by the masses of the usual light neutrinos. For this reason, we find that leptogenesis must take place at a high mass scale, comparable to that for the standard scenario.

Washout effects, which reduce the asymmetry, will differ from the standard scenario. For example, the inverse decay is now a $3 \to 1$ process, and there will be a different set of L-violating scattering processes. We leave the detailed study of these washout effects to future work. Finally, it is possible that the interference of the tree-level and vertex amplitudes, which we have not explicitly calculated, may enhance the asymmetry. As with the standard mechanism, a near degeneracy in the masses of the heavy neutrinos would also enhance the self-energy contribution to the asymmetry [7,18].

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