Electroweak Baryogenesis from Late Neutrino Masses

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Electroweak baryogenesis, given a first-order phase transition, does not work in the standard model because the quark Yukawa matrices are too hierarchical. On the other hand, the neutrino mass matrix is apparently not hierarchical. In models with neutrino mass generation at low scales, the neutrino Yukawa couplings lead to large CP violation in the reflection probability of heavy leptons by the expanding Higgs bubble wall, and can generate the observed baryon asymmetry of the universe. The mechanism predicts new vectorlike leptons below the TeV scale and sizable $\mu \to e$ processes.

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The neutrino physics has been undergoing a revolutionary progress in the past several years. The observations of neutrino flavor conversion in various experiments provided evidence for non-standard model (SM) physics [see, e.g., [1,2]]. The data on the distance/energy dependence of atmospheric [3] and reactor neutrino experiments [4] indicate oscillatory behavior. This strongly favors the presence of tiny but nonzero neutrino masses.

The most popular explanation for the origin of neutrino masses is the seesaw mechanism [5]. In its minimal form, right-handed (RH) neutrinos are introduced with lepton-number violating Majorana mass terms, $\Lambda_L NN$, as well as Yukawa couplings to the SM lepton doublets l and the Higgs h, $Y_L hNl$. The oscillation data require $\Lambda_L \sim 10^{14}$ GeV if $Y_L \sim O(1)$. Moreover, the seesaw mechanism naturally provides a way to account for the observed baryon asymmetry of the universe (BAU) [1,6], $n_B/s \sim 8 \times 10^{-11}$, where n_B/s is the baryon to entropy ratio. The out-of-equilibrium decay of the RH neutrinos creates a lepton asymmetry, which is partially converted to a baryon asymmetry by the electroweak sphaleron process (leptogenesis) [7]. Thermal leptogenesis typically requires, however, $\Lambda_L \sim 10^9$ GeV [8] with somewhat small Yukawa couplings Y_L and large hierarchies in the RH masses.

While the seesaw mechanism is very appealing theoretically, it is unlikely that it will be subject to a direct experimental test in the near future. It is important to explore other possibilities for the origin of neutrino masses. One such example is the late neutrino mass framework that induces small neutrino masses due to a low scale of symmetry breaking [9]. When this symmetry is broken, say, by a set of symmetry breaking vacuum expectation values, $\langle \phi \rangle = f$, the neutrinos acquire masses from operators

$$\left(\frac{\phi}{M_F}\right)^n lNh(\text{Dirac}) \quad \text{or} \quad \left(\frac{\phi}{M_F}\right)^n \frac{(lh)^2}{\Lambda_L}(\text{Majorana}). \quad (1)$$

We stress that this does not depend on the details of the model or whether one uses global [9] or gauge [10–12] symmetries. The strongest limits on f arise from big bang nucleosynthesis [9,11–14] and from observation of supernova neutrinos [15]: for $n = 1$, $f \gtrsim 10$ keV, while more powerful limits apply to higher n. It is remarkable that new physics at such a low scale is not excluded by direct experimental data.

In this Letter we show that the late neutrino mass framework can naturally realize leptogenesis at the electroweak (EW) phase transition [16]. Heavy vectorlike leptons that give rise to the operators (1) bounce off the expanding bubble walls with $O(1)$ Yukawa couplings, and acquire a large asymmetry. They quickly decay to the standard-model leptons and the sphaleron process partially converts their asymmetry to the baryon asymmetry. Thus our mechanism leads to various predictions that can be tested by near future experiments. Here we focus on the qualitative features of our scenario while a more detailed analysis will be presented in a following publication. We simply assume that the phase transition is first order and focus on the size of the baryon asymmetry as well as phenomenological constraints on the model. We briefly comment on the origin of the first-order phase transition towards the end of this Letter.

The model and mechanism.—As well as introducing a low-scale f, theories of late neutrino masses introduce flavor scales that are much lower than the scale Λ_L of the seesaw mechanism. One economical possibility, that we explore in this Letter, is that these flavor scales are all of order the EW scale: $M_F, \Lambda_L \sim v \equiv \langle h \rangle$. In this case the non-SM states that generate the operators of (1) have masses of order v and are available to take part in EW baryogenesis. The flavor symmetry breaking scale becomes $f \sim v(m_\nu/v)^{1/n} \sim (m_\nu, 100\ \text{keV}, 30\ \text{MeV} \ldots)$ for $n = 1, 2, 3 \ldots$. The $n = 1$ case is excluded by big band nucleosynthesis (BBN) and supernova constraints. Large n theories may be preferred in the sense that the scale f grows and requires less protection. In this Letter, we describe a simple $n = 2$ theory that illustrates our baryogenesis mechanism. We focus on the Majorana case since, as shown below, the Dirac case is disfavored by the direct experimental data for the parameter range that produce enough baryon asymmetry.

The minimal model for late Majorana neutrino masses is

$$\mathcal{L}_\nu = YhNL + ML^cL + y\phi L^c l + \frac{M_N}{2}NN + \text{H.c.}, \quad (2)$$

where L, L^c are vectorlike lepton doublets, and the couplings Y, y, and the masses M, M_N are 3×3 flavor matrices. We consider all eigenvalues of M to be comparable, the same for M_N, and those of Y to be of $O(1)$. This is suggested by the lack of hierarchy [anarchy [17]] in the neutrino masses. We refer to their eigenvalues as $\bar{M}$, $\bar{M}_N$, and $\bar{Y}$. There may be a moderate hierarchy in the eigenvalues of y: y_i. Note that there is a leptonic analogue of the rephasing invariant Jarlskog determinant even for just two generations ($N_G = 2$) [18],

$$J_L = \text{Im}\,\text{Tr}\,(YY^\dagger M_N^* Y^* Y^T M_N M_N^* M_N). \quad (3)$$

Below the scales $\bar{M}$ and $\bar{M}_N$, neutrino masses are described by the operator $(y_i^2\bar{Y}^2/\bar{M}^2\bar{M}_N)l^2h^2\phi^2$, corresponding to the $n = 2$ case of (2), so that the flavor symmetry breaking scale is predicted to be

$$f^2 \sim \frac{m_{\nu i}v}{y_i^2\bar{Y}^2}\frac{\bar{M}^2}{v^2}\frac{\bar{M}_N}{v} \sim \frac{(100\ \text{keV})^2}{y_i^2\bar{Y}^2}\frac{\bar{M}^2}{v^2}\frac{\bar{M}_N}{v}. \quad (4)$$

A low $f \sim 100$ keV would imply that the ϕ states contribute to the energy density of the universe during BBN. However, our baryogenesis mechanism will require $\bar{M}, \bar{M}_N$ somewhat larger than v, and lepton flavor violation will constrain y_i to be somewhat small. Hence we expect that the ϕ states are heavier than 1 MeV and decay to neutrinos before BBN.

Let us now describe our main mechanism. With $\bar{M} \sim v$, during the electroweak phase transition L particles and their CP conjugates, $\bar{L}$, are reflected differently from the Higgs bubble wall. This is due to the presence of unsuppressed CP violating phases in Y, M, and M_N [for a related idea, see [18,19]]. Thus an asymmetry in L is induced in the region just in front of the wall. As shown below, the size of this asymmetry is expected to be of order $J_L/\bar{M}_N^4$. We assume the $L \to Nh$ decay process is kinematically forbidden (i.e., $\bar{M}_N > \bar{M}$) so that the asymmetry is transferred to SM leptons via $L \to l\phi$ decays. The presence of the asymmetry in the SM lepton doublets biases the sphaleron rate to induce a B production in the vicinity of the wall. When the expanding bubble passes over this region, the sphaleron processes decouple, freezing in a B asymmetry. Outside the bubble the sphaleron rate $\Gamma_{\text{Sp}} \sim \alpha_W^4 T_c$, T_c being the critical temperature, is much slower than other dynamical scale near the bubble wall [20]. Thus the baryon asymmetry could be as large as

$$\frac{n_B}{s} \sim \frac{1}{g_*}\alpha_W^4\frac{J_L}{\bar{M}_N^4} \sim 10^{-8}, \quad (5)$$

where $g_* \simeq 100$ is the number of relativistic degrees of freedom, and anarchical neutrino masses suggest that $J_L/\bar{M}_N^4$ is of order unity. Below we study what other factors might suppress the baryon asymmetry.

Estimating the baryon asymmetry.—To estimate BAU we apply the thin wall approximation [its validity depends on the details of the first order phase transition [20]]. The L and $\bar{L}$, n_L density difference is induced by the reflection asymmetry [21]

$$\frac{n_L}{s} \sim \frac{1}{45T_c^2}\int\frac{d\omega}{2\pi}n_0(\omega)[1-n_0(\omega)]\Delta(\omega)\Delta\vec{p}\cdot\vec{v}_w, \quad (6)$$

where $v_w \sim 0.1$ is the wall velocity and $n_0(\omega) = 1/(e^{\omega/T_c}+1)$ is the Fermi-Dirac distribution. The difference between N and L momenta for a given energy, $\Delta\vec{p} \equiv \vec{p}_L - \vec{p}_N$, is large, due to $O(1)$ mass differences among $\bar{M}$ and $\bar{M}_N$. This is welcome because, in the SM, Δp arises only due to the electroweak thermal corrections to the masses, and is suppressed by α_W [21]. The reflection asymmetry $\Delta(\omega)$ is given by [20] $\Delta(\omega) = \text{Tr}\,(R_{NL}^\dagger R_{NL} - \bar{R}_{NL}^\dagger\bar{R}_{NL})$, where R_{NL} is the reflection coefficient for $N \to L$ and the bars stand for the CP conjugated process.

In order to calculate the reflection asymmetry from the bubble wall, we have computed the Green function for our model in the simplifying limit $\bar{M}_N \gg \bar{M}$. Using a perturbative expansion in Yv/T_c, M/ω, and for $N_G = 2$, we find $\Delta(\omega) \sim (\ell_T\frac{v}{\omega})^{2N_G}J_L$, where $\ell_T \sim 10/T$ is the mean free path for the leptons [20] (see Fig. 1). Below we find that phenomenological constraints favor somewhat heavier RH Majorana masses, $\bar{M}_N \gtrsim \bar{Y}v$. In this case a suppression in the reflection asymmetry is expected since the heavy incoming particle is hardly affected by the potential barrier. This effect cannot be captured using an expansion in M_N/ω. We estimated the corresponding $1/M_N$ suppression by analyzing a single generation reflection problem which can be solved analytically. We indeed found that the reflection amplitude is further suppressed via $(\bar{Y}v/\bar{M}_N)^{2N_G}$. Consequently, in the relevant region of parameter space the reflection asymmetry is given by

$$\Delta(\omega) \sim \left(\frac{\ell_T\bar{Y}^2v^2}{\omega\bar{M}_N}\right)^{2N_G}J_L\theta(\omega-\bar{M}_N). \quad (7)$$

Below we use the estimate (7) even in the case of interest where Yv/T_c and M/ω are of order unity.

The rate for baryon production is approximately $dn_B/dt \sim 3T_c^2\Gamma_{\text{Sp}}\Delta F$ [20,22]. ΔF is the free energy difference between two neighboring zero field-strength con-

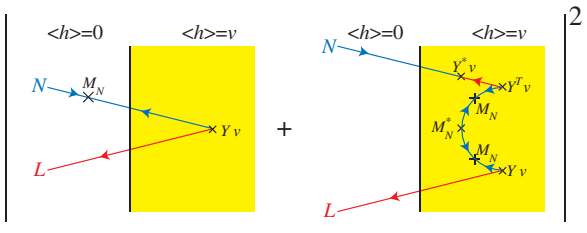

FIG. 1 (color online). Perturbative calculation of the reflection coefficients which picks up the Jarlskog invariant $J_L = \text{Im}\,\text{Tr}\,[(M_NY)^\dagger(Y^*Y^TM_NM_N^*M_NY)]$ from the one-particle cuts in the amplitudes.

figurations, for which $\Delta B = \Delta L = 3$ [provided that $\Gamma(L \to l\phi)$ is not much slower than other thermalization rates; see below]. ΔF is calculated in the presence of a hypercharge density $n_Y \sim -n_L/2$ [22]. [More correctly, one should define a global approximately conserved charge orthogonal to hypercharge $B' = B - xY$ where in our case $x = 1/7$ [22]. We verified that this will hardly modify our results.] In our case we find that $\Delta F \simeq n_L/T_c^2$. The BAU is obtained by integrating dn_B/dt, which we estimate via $n_B \sim dn_B/dt \times \ell_w/v_w$, where ℓ_w is the typical penetration length for the nonzero global charge which flows through the plasma in the unbroken phase. With fast massless leptons one expects, from estimation of energy loss in the plasma [23], $T_c\ell_w = \mathcal{O}(100)$ [20,22]. In our case, with semirelativistic leptons, the penetration distance is shorter even though the energy loss rate is very low, because the massive leptons rather quickly lose their directionality (away from the wall). For instance, elastic scattering of a lepton of mass $4T_c$ with a plasmon that carries a perpendicular momenta, T_c, roughly results in 25% momentum loss in the original direction of motion for the lepton. Thus in our case we estimate $\ell_w \sim N_{\rm coll}\ell_T$, where $N_{\rm coll}$ is the average number of collisions that a lepton undergoes before its directionality is lost. Below, we assume $N_{\rm coll} = 2$.

Using Eqs. (6) and (7) we obtain

$$\frac{n_B}{s} \sim 3\Gamma_{\rm Sp} N_{\rm coll} \frac{\ell_T}{v_w} \frac{n_L}{2s}. \tag{8}$$

Our next step is to identify the dependence of the resultant BAU on the model fundamental parameters. Assuming that CP violation is maximal and taking for simplicity $(\bar{Y}v)^2\ell_T/T_c \sim 1$ we find $\Delta(\omega) = (T_c/\omega)^4\theta(\omega - \bar{M}_N)$. In addition, fixing $N_{\rm coll} = 2$ and $\ell_T \sim 10/T_c$ we can numerically compute the resultant BAU as a function of $\bar{M}_N/T_c$. As $\bar{M}_N/T_c$ increases, the BAU rapidly decreases due to both a Boltzmann suppression factor [see Eq. (6)] and also a polynomial one (7). In Fig. 2 we plot n_B/s as a function of $\bar{M}/T_c$. We find that the observed asymmetry can be accounted for when

$$\bar{M}_N \lesssim 4T_c. \tag{9}$$

In the Dirac case $N_G = 3$ and the resultant asymmetry is further suppressed.

Direct constraints and tests.—In the following we briefly discuss the direct phenomenological constraints on our scenario and discuss several ways to directly test our model in the near future.

Electroweak precision measurements.—Our model requires additional heavy lepton doublets. Because they are vectorlike, the S parameter is hardly affected. The Yukawa couplings to the SM Higgs break the custodial symmetry and modify the T parameter. The corresponding shift in T, T_L, is similar to a single extra vectorlike top quark [24] $T_L \sim T_t^{\rm SM}\frac{\bar{Y}^4v^4}{\bar{m}_t^2M^2} \sim 1.2$, where $T_t^{\rm SM}$ is the SM-top contribution to T [1] and m_t is the top mass. Requiring the $T_L \lesssim 0.2$

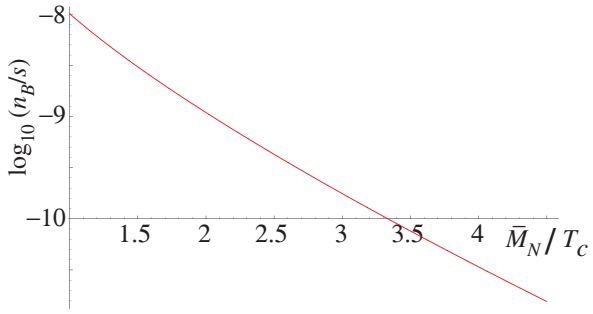

FIG. 2 (color online). n_B/s as a function of $\bar{M}_N/T_c$, for $\ell_w \sim 20/T_c$, $N_G = 2$, and $\Delta(\omega) = (T_c/w)^4\theta(\omega - \bar{M}_N)$.

we find

$$\frac{\bar{M}}{v} \gtrsim 2.5\bar{Y}\left(\frac{\bar{Y}v}{m_t}\right). \tag{10}$$

In the *Dirac* case, the light, RH neutrinos are partially active and the Z can decay invisibly to them. The 3σ allowed range [25] implies that $\bar{M} \gtrsim 4\bar{Y}v$.

Lepton flavor violation.—The Yukawa couplings to the SM leptons, y, are generically not aligned with M, M_N, and Y and hence induce lepton flavor violation. This will contribute to processes such as $\mu \to e$ conversion which are highly constrained by data [1,26]. In our model the SM leptons couple to the additional fields only through y. To avoid these constraints naturally a mild hierarchy in y (in the appropriate basis), $y \sim \mathrm{diag}\,(10^{-2}, 10^{-1}, 10^{-1})$, which is consistent with neutrino data is required. The smaller the eigenvalues of y, y_i, however, the smaller is the decay rate $L \to l\phi$. Comparing the decay rate into SM doublets to the thermalization time we find $\Gamma_L\ell_w/v_w \sim 10y_i^2\bar{M}/T_c \sim 0.25 - 1$ which implies further suppression in the resultant BAU (this is still consistent with the observed value; see Fig. 2). We find that there is a tension between producing sizable BAU and suppressing the contribution to lepton flavor violation. Therefore, our mechanism typically predicts that the rates for $\mu \to e$ processes are within the reach of near future experiments. Note that with the above ansatz for y, contributions to neutrinoless double beta decay are evaded [12].

Collider physics.—From Eqs. (9) and (10), our scenario predicts the presence of SU(2) doublets with masses below the $\mathcal{O}$ (1 TeV) range. These contain charged particles, with masses above 100 GeV from current bounds [1], that may be produced and detected at the LHC, or better at the international linear collider. A clear signal at the LHC is expected if the vectorlike leptons are sufficiently light, via production of $L^{\pm}L^0$ which decay to a single, energetic SM charged lepton. If the Higgs boson is heavier than the new particles it can also decay invisibly to N and L.

The presence of new lepton doublets is crucial to our mechanism. Implementing our idea with SM singlets would yield light, partially sterile, neutrinos. This leads to nonuniversality in weak processes as β decay and μ decay. As these are measured at the 0.1% level, the lower

bound on $\bar{M}$ would strengthen, significantly suppressing the BAU. Similarly difficulty would be encountered in the case when neutrino masses are induced by the vacuum expectation value of an SU(2) triplet.

Discussion.—In this Letter, we showed that the framework of late neutrino masses naturally leads to a viable model of electroweak baryogenesis that can be directly tested in near future experiments.

The additional vectorlike leptons with mass close to the EW scale play a crucial role in our scenario. This raises a coincidence problem since a vectorlike mass term is unprotected by the symmetries of our model and therefore is unrelated to the EW scale, similar to the μ problem in the minimal supersymmetric SM. There are, however, many proposed solutions to this problem: vectorlike mass may be induced from supersymmetry breaking, additional Higgs fields, or strong dynamics.

The first-order EW phase transition is mandatory. It may be induced, for instance, by higher order terms in the effective potential that arise from integrating out additional heavy states [28].

Our baryogenesis mechanism relies only on physics at the TeV scale, and hence is compatible with any scheme for new physics above the TeV scale. For example, there could be flat extra dimensions just above the TeV scale, or a warped extra dimension as in the Randall-Sundrum scheme; in supersymmetric theories our mechanism allows for a low reheating temperature after inflation, solving the various moduli/gravitino problems. Finally, our scenario may accommodate unification by completing the vectorlike leptons to full SU(5) representations.

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