

Baryogenesis via Neutrino Oscillations

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We propose a new mechanism of leptogenesis in which the asymmetries in lepton numbers are produced through the CP -violating oscillations of “sterile” (electroweak singlet) neutrinos. The asymmetry is communicated from singlet neutrinos to ordinary leptons through their Yukawa couplings. The lepton asymmetry is then reprocessed into baryon asymmetry by electroweak sphalerons. We show that the observed value of baryon asymmetry can be generated in this way, and the masses of ordinary neutrinos induced by the seesaw mechanism are in the astrophysically and cosmologically interesting range. Except for singlet neutrinos, no physics beyond the standard model is required. [S0031-9007(98)06856-2]

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The origin of the excess of baryons over antibaryons in the Universe remains one of the fascinating problems of particle physics and cosmology. A number of mechanisms have been proposed to date to explain this asymmetry (for recent reviews see, e.g., Ref. [1]). One of the simplest possibilities, suggested by Fukugita and Yanagida [2], is that the baryon asymmetry has originated from physics in the leptonic sector. Namely, it was assumed that, at temperatures well above the electroweak scale, lepton asymmetry was produced, which was then reprocessed into the baryon asymmetry by nonperturbative electroweak effects [3]—sphalerons [4]. According to Ref. [2] the lepton asymmetry is generated in out-of-equilibrium, CP , and lepton number nonconserving decays of heavy Majorana neutrinos (for recent discussions see, e.g., Ref. [5] and references therein).

In this Letter we propose a new realization of baryogenesis through leptogenesis which also makes use of the electroweak reprocessing of the lepton number into the baryon number. Like the Fukugita–Yanagida mechanism, our proposal requires only a mild extension of the standard model by introducing “sterile” (i.e., electroweak singlet) heavy neutrinos. However, our mechanism of leptogenesis is entirely different from that of Ref. [2]: We suggest that asymmetries in lepton numbers were generated due to oscillations of these singlet neutrinos and their interactions with ordinary matter in the early Universe. Moreover, the novel feature of our scenario is that the total lepton number is not violated in these oscillations and/or interactions; an important ingredient is separation (rather than generation) of the lepton number, i.e., its redistribution between different species of singlet neutrinos.

For this reason we do not necessarily require that singlet neutrinos be Majorana particles; Dirac sterile neutrinos are equally suitable (and even better in some respect) for our mechanism to work. Furthermore, in our case the values

of the masses and couplings of singlet neutrinos are very different from those of Ref. [2].

Let us consider the standard model extended by adding three types of Majorana neutrinos N_a , $a = A, B, C$, which interact with other particles only through their Yukawa couplings [6]. The corresponding Lagrangian can be written in the “Yukawa basis” (where the matrix of Yukawa coupling constants has been diagonalized) as follows:

$$\mathcal{L} = \bar{N}_{Ra} i \not{\partial} / N_{Ra} + h_a \bar{l}_a N_{Ra} \Phi + \frac{M_{ab}}{2} N_{Ra}^T C N_{Rb} + \text{H.c.}$$

Here N_{Ra} are right-handed components of N_a , h_a are (real) Yukawa coupling constants, l_a are three independent linear combinations of the usual leptonic doublets, Φ is the Higgs doublet, and M_{ab} is the mass matrix.

We suggest that the baryogenesis proceeds in the following way:

(i) In the course of the evolution of the Universe, singlet neutrinos are produced through their Yukawa couplings. The production mechanism of singlet neutrinos conserves CP , i.e., for each type equal numbers of particles and antiparticles (particles of opposite helicities in the Majorana case) are produced.

(ii) Once created, singlet neutrinos oscillate, and also interact with ordinary matter. None of these processes violates the total lepton number $L^{\text{tot}} = L + L_A + L_B + L_C$, where L is the usual lepton number (we assume that Majorana masses are small enough, see below). However, CP is not conserved due to mixing in the singlet neutrino sector. Therefore the initially created state with individual lepton numbers $L_A = L_B = L_C = 0$ evolves through the oscillations into a state in which $L_A \neq 0$, $L_B \neq 0$, $L_C \neq 0$, but still $L^{\text{tot}} = 0$. That is, the total lepton number gets unevenly distributed between different species.

(iii) Singlet neutrinos communicate their lepton asymmetry to ordinary neutrinos and charged leptons through their Yukawa couplings. We assume that the Yukawa couplings of N_A , N_B , and N_C have certain hierarchy, so that neutrinos of at least one type, N_A , come into thermal equilibrium before the time t_{EW} at which sphalerons become inoperative (the corresponding temperature is $T_{EW} \sim 100$ GeV) and those of at least one other type, N_C , do not equilibrate by t_{EW} . The neutrinos of the third type, N_B , may or may not come into thermal equilibrium by t_{EW} . To be specific, we discuss mostly the case when the Yukawa coupling of N_B is relatively large so that N_B equilibrates at temperatures exceeding T_{EW} (the opposite case is treated in a similar way). In this case, the lepton numbers L_A and L_B are communicated to the ordinary leptons before t_{EW} , whereas L_C is not. Therefore (a fraction of) $L_A + L_B$ is reprocessed into baryon asymmetry by electroweak sphalerons; the lepton number L_C is transferred to the active leptons only after the sphalerons have already switched off, so it has no effect on the baryogenesis.

Let us stress that, if not for the sphaleron freezing effect, no baryon asymmetry would have been obtained: All three singlet neutrino species eventually transform into ordinary leptons, and, since $L^{\text{tot}} = 0$, no net lepton charge would have been generated in the sector of ordinary leptons. Hence, the requirement that N_C does not get into thermal equilibrium before t_{EW} is crucial for our mechanism.

The system of three types of singlet neutrino of a given momentum $k(t) \propto T(t)$, interacting with cosmic plasma, is described by a 3×3 density matrix $\rho(t)$ which obeys the evolution equation [7]

$$i \frac{d\rho}{dt} = [\hat{H}, \rho] - \frac{i}{2} \{\Gamma, \rho\} + \frac{i}{2} \{\Gamma^p, 1 - \rho\}. \quad (1)$$

Here \hat{H} is a Hermitian effective Hamiltonian, Γ and Γ^p are destruction and production rates, respectively, and we have neglected the relaxation due to elastic scattering; see below. In what follows we will use the approximation of Boltzmann statistics for order-of-magnitude analysis; in this approximation the last term in Eq. (1) becomes merely $i\Gamma^p$.

In the Yukawa basis one has, at $T \gg M_a$,

$$\hat{H} = V(t) + U \frac{\hat{M}^2}{2k(t)} U^\dagger. \quad (2)$$

Here U is the mixing matrix which relates the Yukawa states and the mass eigenstates N_i ($i = 1, 2, 3$): $N_a = U_{ai} N_i$ (we assume that mixing is small, and choose the standard parametrization for U [8]); $\hat{M}^2 = \text{diag}(M_1^2, M_2^2, M_3^2)$ is the matrix of mass eigenvalues; $k(t)$ is the neutrino momentum which depends on time due to the expansion of the Universe. The CP violation in the system is described by the CP -odd phase δ in the mixing matrix U .

The medium effects are described by real potential V and rates Γ and $\Gamma^p = \exp(-k/T)\Gamma$ [9] whose matrices

are diagonal in the Yukawa basis,

$$V = \text{diag}(V_A, V_B, V_C), \quad \Gamma = \text{diag}(\Gamma_A, \Gamma_B, \Gamma_C).$$

The potentials V_a are due to the coherent forward scattering processes, the main contribution coming from the 1-loop self-energy diagrams with ordinary lepton and Higgs doublets in the intermediate states. For our estimates we use their thermal averages which, at temperatures above T_{EW} , are [10]

$$V_a = \frac{1}{8} h_a^2 T. \quad (3)$$

Consider now the rates Γ_a . We will be interested in the temperatures far exceeding the masses of singlet neutrinos. Therefore the rates of the $1 \leftrightarrow 2$ reactions which correspond to the absorptive parts of the self-energy diagrams are suppressed by the factor M_a/T , and $2 \leftrightarrow 2$ reactions are more important. The main contributions to Γ_a come from the Higgs exchange reactions $Q_L N_{Ra} \leftrightarrow t_R l_a$, $t_R N_{Ra} \leftrightarrow Q_L l_a$, and $l_a^c N_{Ra} \leftrightarrow t_R Q_L^c$, where Q_L is the third-generation quark doublet. The result [11] for the average destruction rates at $T \gg M_a$ (corrected to include the color factor for quarks) is

$$\Gamma_a \simeq \frac{9h_t^2}{64\pi^3} h_a^2 T. \quad (4)$$

Here $h_t \simeq 1$ is the top quark Yukawa coupling. The rates of the elastic $2 \rightarrow 2$ scattering processes are proportional to $h_a^2 h_b^2$ instead of $h_a^2 h_t^2$. We will need very small Yukawa couplings of singlet neutrinos, so the elastic processes can be safely neglected.

At late times the lepton asymmetry is stored in the least interacting species of a singlet neutrino, which at $t \sim t_{EW}$ coincides with the mass eigenstate $N_3 \approx N_C$. The conditions that $N_1 (\approx N_A)$ and $N_2 (\approx N_B)$ come into thermal equilibrium before the time t_{EW} , while N_3 does not, are

$$\Gamma_{1,2}(T_{EW}) > H(T_{EW}), \quad \Gamma_3(T_{EW}) < H(T_{EW}), \quad (5)$$

where $H(T) = T^2/M_{\text{Pl}}^*$ is the Hubble parameter, $M_{\text{Pl}}^* \equiv M_{\text{Pl}}/1.66\sqrt{g_*} \simeq 10^{18}$ GeV, and $\Gamma_3 \approx \Gamma_C + s_{13}^2 \Gamma_A + s_{23}^2 \Gamma_B$. Here $s_{13} \equiv \sin \theta_{13}$ and $s_{23} \equiv \sin \theta_{23}$ determine the admixtures of N_A and N_B in N_3 .

The conditions (5) translate into bounds on the Yukawa couplings,

$$h_{A,B}^2 > 2 \times 10^{-14}, \quad h_C^2 < 2 \times 10^{-14}, \quad (6)$$

and on mixing angles,

$$s_{13}^2 < 2 \times 10^{-14} h_A^{-2}, \quad s_{23}^2 < 2 \times 10^{-14} h_B^{-2}. \quad (7)$$

Equation (6) implies a certain hierarchy between the Yukawa couplings, which, however, need not be very strong.

As described above, to find the baryon asymmetry we should calculate the asymmetry $\Delta_L(t_{EW})$ which was communicated to the usual leptons by the time t_{EW} at which sphalerons switch off. This asymmetry emerges because singlet neutrinos N_3 do not transfer their asymmetry to active neutrinos by the time t_{EW} due to the smallness of h_C , s_{13} , and s_{23} . Since the total lepton number is conserved in all processes of interest, we have $\Delta_L(t_{EW}) = -\Delta_3(t_{EW})$ (up to a factor of order 1 that accounts for the distribution of the asymmetry between B , L , and $L_A + L_B$; we will not write this factor in formulas below), where $\Delta_3(t_{EW})$ is the asymmetry stored in N_3 . The asymmetry $\Delta_3(t_{EW})$ can be found as follows. Let $S(t, t_0)$ be the evolution matrix corresponding to the operator $\tilde{H} = \hat{H} - (i/2)\Gamma$ (notice that $S(t, t_0)$ is not unitary since this operator is non-Hermitian). The density matrix $\rho(t)$ can be expressed through $S(t, t_0)$ assuming $\rho(t_i) = 0$, where t_i is the time at which the production of singlet neutrinos begins [12]. The ratio of the number density of N_3 , $n_3 = \rho_{33}$, to the equilibrium density of one spin degree of freedom at time t_{EW} is

$$\frac{n_3(t_{EW})}{n_{eq}(t_{EW})} = \sum_{a,b} \int_{t_i}^{t_{EW}} dt_0 \Gamma_a(t_0) |U_{3b}^\dagger S_{ba}(t_{EW}, t_0)|^2. \quad (8)$$

The asymmetry $\Delta_3(t_{EW})$ is the CP -odd part of the quantity (8).

The integration over the production time t_0 can be performed in a closed form. Indeed, the matrices S and S^\dagger obey $\partial_{t_0} S = i\tilde{H}(t_0)S$, $\partial_{t_0} S^\dagger = -i\tilde{H}^\dagger(t_0)S^\dagger$. From these equations and $\Gamma = i(\tilde{H} - \tilde{H}^\dagger)$, one finds that $\partial_{t_0}(SS^\dagger) = S\Gamma(t_0)S^\dagger$. Using this relation, one can readily perform the integration over the production time t_0 in Eq. (8),

$$\frac{n_3(t_{EW})}{n_{eq}(t_{EW})} = 1 - [S^M(t_{EW}, t_i)S^M(t_{EW}, t_i)^\dagger]_{33}.$$

Here $S^M \equiv U^\dagger S U$ is the evolution matrix in the mass eigenstate basis. The CP -odd part of this expression determines the asymmetry transferred to the usual leptons by t_{EW} (and hence the generated baryon asymmetry)

$$\begin{aligned} \Delta_L(t_{EW}) &\equiv (n_L - n_{\bar{L}})/n_\gamma = -\Delta_3(t_{EW}) \\ &= \frac{1}{2} \sum_i |S_{3i}^M(t_{EW}, t_i)|_{CP\text{-odd}}^2, \end{aligned} \quad (9)$$

where the factor 1/2 accounts for two helicity states of photon. The production of singlet neutrinos starts at very early times, so we set $t_i = 0$. Since T_{EW} is much smaller than all relevant energy parameters in the problem, one can formally let $t_{EW} \rightarrow \infty$ in actual calculations.

The lepton asymmetry is produced mainly at the epoch t_L when the differences of the eigenvalues of the Hamiltonian $\Omega_{ij} \equiv \Omega_i - \Omega_j$ become of the order of the Hubble parameter H : $\Omega_{ij}(t_L) \sim 1/t_L$. Indeed, at $t \ll t_L$ the elements of $\rho(t)$ essentially stay constant, whereas at $t \gg t_L$ they undergo fast oscillations and, due to averaging effects, the asymmetry is strongly suppressed.

In what follows we will present the results for the most interesting range of the parameter space, where the mass differences are relatively large, $\Delta M^2 \gg (h_{ab}^2/8)^3 M_{\text{Pl}}^{*2}$ ($h_{ab}^2 \equiv h_a^2 - h_b^2$), ΔM^2 being the typical value of ΔM_{ij}^2 . In this case, the mass terms dominate over potentials in \hat{H} at the epoch t_L , so that $\Omega_{ij} \sim \Delta M_{ij}^2/2T$ and the leptogenesis temperature is $T_L \equiv T(t_L) \sim (M_{\text{Pl}}^* \Delta M_{ij}^2)^{1/3}$. At $t = t_L$ we have

$$\frac{|V_{ab}|}{H} = \frac{h_{ab}^2}{8} \left(\frac{M_{\text{Pl}}^{*2}}{\Delta M^2} \right)^{1/3} \equiv \lambda \ll 1. \quad (10)$$

This means that the potentials V_a and rates Γ_a can be treated perturbatively, with λ being the expansion parameter. The lepton asymmetry (9) appears in the third order of perturbation theory, and therefore is suppressed by the cube of λ . This can be seen in the mass eigenstate basis, where the potential has the form $U^\dagger V U$. Indeed, Δ_L is a CP -violating observable, so it should be proportional to the invariant $J = s_{12}c_{12}s_{13}c_{13}^2s_{23}c_{23} \sin \delta$; this invariant can be collected from $[U^\dagger V U]^3$ only. It is clear from Eqs. (8) and (9) that Δ_L vanishes in the limit $\Gamma \rightarrow 0$, so one expects it to be proportional also to $\sin \phi \equiv \Gamma_a/2V_a \simeq 2 \times 10^{-2}$. Therefore, up to a numerical constant, we have an estimate

$$\Delta_L \sim J \lambda^3 \sin \phi. \quad (11)$$

The calculations are simplified if $|\Delta M_{13}^2| \ll |\Delta M_{12}^2|$, $|\Delta M_{23}^2|$. In this case, using Eq. (9), we obtain

$$\Delta_L \simeq \frac{[\Gamma(\frac{1}{3})]^3}{384} J \sin \phi \frac{h_{AC}^2 h_{AB}^2 h_{BC}^2 M_{\text{Pl}}^{*2}}{(\Delta M_{13}^2)^{1/3} (\Delta M_{12}^2)^{2/3}}, \quad (12)$$

where $\Gamma(1/3) \simeq 2.68$. The estimate (11) as well as the formula (12) are valid both in the case when N_2 equilibrates before t_{EW} and in the opposite case.

The asymmetry increases when the parameter λ approaches 1, i.e., the maximal effect for given h_a is expected when $\Delta M^2 \sim (h_{ab}^2/8)^3 M_{\text{Pl}}^{*2}$. Notice, however, that these and smaller values of ΔM^2 correspond to singlet neutrinos strongly degenerate in mass.

Let us present constraints on the parameters of singlet neutrinos and discuss the value of the asymmetry.

(i) To be consistent with the standard mechanism of nucleosynthesis, all singlet neutrinos, including the most weakly interacting one N_3 , should decay before the nucleosynthesis epoch. The decay of N_3 at $T \ll T_{EW}$ occurs due to its mixing with an ordinary neutrino. Requiring that the decay rate of N_3 exceeds the inverse lifetime of the Universe at temperatures of the order of a few MeV, and recalling Eq. (6), we obtain a lower bound on the mass, $M_3 \geq 1$ GeV. Alternatively, for $M_3 \gg 1$ GeV, we have a lower bound on the Yukawa constant, $h_3^2 \geq 10^{-16} (1 \text{ GeV}/M_3)^3$, where $h_3^2 \simeq h_C^2 + s_{13}^2 h_A^2 + s_{23}^2 h_B^2$.

(ii) If the mass of N_3 is close to 1 GeV and/or the Yukawa couplings of N_A and N_B are close to the bound (6), the decays of singlet neutrinos may lead to the reheating

of the Universe after the electroweak epoch (but before nucleosynthesis) and, hence, to the dilution of the baryon asymmetry. This reheating is rather modest, however: Given the constraints already imposed, the entropy density may increase at most by a factor of 10. In this case the baryon asymmetry produced before the reheating should be an order of magnitude larger than the observed one.

(iii) Baryon and lepton asymmetries should not be washed out before $T = T_{EW}$ by the Majorana mass itself. At $T \gg M_A$ the lepton number equilibration rate is suppressed with respect to the lepton charge conserving rate Γ_A , given by Eq. (4), by a factor M_A^2/T^2 , so we have to require that $\Gamma_A(T_{EW})(M_A^2/T_{EW}^2) \ll H(T_{EW})$, and similarly for N_B . The parameters of Majorana singlet neutrinos should therefore satisfy

$$1 \text{ GeV} \lesssim M_a \ll 100 \text{ GeV}, \quad h_A^2, h_B^2 \ll 10^{-10}. \quad (13)$$

The upper bounds here do not apply to Dirac singlet neutrinos.

The Lagrangian of the model leads, via the seesaw mechanism, to the generation of masses of the light (active) neutrinos: $m_{\nu_a} \equiv m_a = h_a^2 v^2/M_a$, where v is the Higgs vacuum expectation value and ν_a are mass eigenstates—combinations of ν_e , ν_μ , and ν_τ . The constraints (6) and (13) imply that the mass of the heaviest active neutrino is in the range $m_a = (10^{-2}-10^3)$ eV. From the cosmological bound $m_a \lesssim 10$ eV, we get a constraint which is somewhat stronger than Eq. (13), $h_A^2 \lesssim 10^{-11}$. For the lightest active neutrino, the constraints (6) and (13) lead to $m_C = (10^{-6}-10^{-1})$ eV. (Notice that the mixing parameters of active neutrinos are unconstrained by our scenario as they are not directly related to the mixing matrix of singlet neutrinos.)

The above constraints imply that the condition (10) is indeed satisfied in a large part of the allowed parameter space. In particular, it holds in the two cases which we now turn to.

As follows from Eq. (12), in the case when two of the singlet neutrinos are relatively strongly interacting the desired lepton (and baryon) asymmetry $\Delta_L \sim (\text{a few}) \times 10^{-9}$ is obtained for the generic values of the parameters subject to the above constraints. For example, for $h_A^2 \sim h_B^2 = 10^{-12}$, $\Delta M^2 \sim M_a^2$, and $M_a = 10$ GeV, the correct asymmetry is generated provided that $J \gtrsim 10^{-3}$ which is certainly consistent with Eq. (7). The temperature of leptogenesis is $T_L \sim 10^7$ GeV. In this case, *two* active neutrinos are relatively heavy, $m_A \sim m_B = (\text{a few})$ eV, so that they can constitute the hot dark matter of the Universe. Moreover, oscillations between them can solve either atmospheric or solar neutrino problems, provided that their mass splitting is small.

A variant of our scenario makes use of weakly interacting N_C and N_B . In this case, two of the usual neutrino species have masses in the range $(10^{-6}-10^{-1})$ eV,

and the remaining one is relatively heavy. As an example, let us take $h_A^2 = 5 \times 10^{-14}$, $h_B^2 = 10^{-15}$, $h_C^2 < h_B^2$, and $M_a = 20$ GeV, which corresponds to the masses of the usual neutrinos $m_A \sim 0.1$ eV, $m_B \sim 2 \times 10^{-3}$ eV, and $m_C < m_B$. This variant fits particularly well into the mass pattern suggested by the solar and atmospheric neutrino data [13]. Given that $J \lesssim 10^{-2}$ due to the constraints analogous to Eq. (7), the correct value of asymmetry is obtained for $\Delta M^2 \lesssim 10^{-2}$ GeV². Thus, in this case the singlet neutrinos should be degenerate in mass. The temperature of leptogenesis is lower, $T_L \sim 3 \times 10^5$ GeV. Let us note in passing that the degeneracy of masses M_a is helpful also for obtaining the desired baryon asymmetry for very small mixing angles θ_{ij} .

The crucial feature of the suggested mechanism is that it works only if Yukawa couplings of all singlet neutrinos are small, $h_a \sim (10^{-8}-10^{-6})$. This smallness can be explained, e.g., by the mixing of N_a with very heavy right-handed neutrinos having Yukawa couplings of the same order of magnitude as those of quarks (and charged leptons) h_q . In this case, $h_a \sim h_q \sqrt{M_a/M_R}$, and for our values of h_a and M_a the mass scale M_R may be close to the grand unification scale, $M_R \sim 10^{16}$ GeV.

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