Blocking active-sterile neutrino oscillations in the early universe with a Majoron field

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We propose a new mechanism to block the active-sterile neutrino oscillations in the early universe. We show that a typical consequence of theories where the lepton number is spontaneously broken is the existence of a coherent cosmological Majoron field with a strength proportional to the lepton and baryon numbers of the universe. This field interacts with leptons and changes the potentials relevant for neutrino oscillations. If the scale of lepton number symmetry breaking is of the order of 1 GeV then a Majoron field and lepton number asymmetry of the order of the baryon asymmetry are strong enough to block the active-sterile neutrino oscillations with the atmospheric neutrino mass gap which otherwise would bring the sterile neutrino into equilibrium at the big bang nucleosynthesis epoch.

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I. INTRODUCTION

The explanation of the present neutrino puzzles [1] may require the existence of one or several species of extra light sterile neutrinos [2]. In particular, the sterile neutrino could be relevant for the explanation of the atmospheric neutrino problem (ANP) [3] in the presence of a significant ν_μ-ν_s mixing.[1] The typical required values are $\delta m^2_{\rm atm} \sim 3 \times 10^{-3}$ eV^2 and a large mixing angle, $\sin^2 2\theta_{\rm atm} \simeq 1$.

On the other hand, for such a parameter range one can encounter a contradiction with the big bang nucleosynthesis (BBN) bounds [6] on the number of extra light particle species; namely, according to the analyses of Ref. [7], the sterile neutrino comes into equilibrium with the particle thermal bath via ν_μ-ν_s oscillation, unless the condition $\delta m^2 \sin^4 2\theta \lesssim 3 \times 10^{-6}$ eV^2 is satisfied (updated constraints are given in Ref. [8] for small mass differences $\delta m^2 \lesssim 10^{-7}$ eV^2), which is certainly out of the range of parameters needed to explain the ANP.

However, it was found [9] that the ν_μ-ν_s oscillations are suppressed at temperatures $T \lesssim 3$ MeV (the decoupling temperature of $\nu_{\mu,\tau}$) if the lepton number asymmetry at these temperatures is very high, namely, $L_a \gtrsim 10^{-5}$ (lepton number to photon number ratio). But this is 4–5 orders of magnitude larger than the observed baryon asymmetry of the universe ($B \lesssim 10^{-9}$) and in the most generic baryogenesis context one can expect that $L \sim B$ [e.g., in the context of grand unified theory (GUT) baryogenesis or leptogenesis [10,11] this is because the $B+L$ nonconserving sphaleron processes redistribute B and L among each other [12–14]]. The same is true in the context of the electroweak baryogenesis $(B-L=0)$. In the Affleck-Dine mechanism B and L can in principle be independent of each other, but still of the same order.

It has been shown [15,8] that at much lower temperatures ($T < 100$ GeV) neutrino oscillations can actually produce a rapid increase of the lepton asymmetries from the initial very small values up to the order of 0.1. This, however, only occurs for negative $\delta m^2 \cos 2\theta$ and very small active-sterile mixing angles, not directly relevant for the ANP.

In this paper we show that a lepton asymmetry as small as the present baryon asymmetry may be enough to block the sterile neutrino oscillations. The necessary new ingredient is the existence of a coherent Majoron field in the early universe and a low scale of spontaneous breaking of lepton number, $F_L \sim 1$ GeV. (The Majoron [16] is the massless Nambu-Goldstone boson in models where the total or any partial lepton number is spontaneously broken.)

While it has been common wisdom that due to their derivative coupling nature [17] Nambu-Goldstone bosons cannot mediate long range interactions, it has recently been demonstrated [18] that a coherent source of a Majoron field, to be specific, is formed whenever the corresponding broken lepton number suffers a net increase or decrease in a certain region of space. The processes that violate this lepton number can be the very neutrino oscillations as exemplified in previous papers [18], or any other reactions. In the present work we show that a Majoron field can be produced due to lepto- and baryogenesis processes in the early universe. The Majoron field interacts with neutrinos with a strength inversely proportional to the lepton breaking scale F_L. If F_L is around 1 GeV, a lepton asymmetry as small as $L \sim B \sim 10^{-9}$ can block the ν_μ oscillation into sterile neutrinos with $\delta m^2 \sim 3 \times 10^{-3}$ eV^2 no matter how large their mixing angle is. That is our thesis.

The origin of the Majoron field is elaborated in Sec. V and its role in neutrino oscillations into a sterile neutrino in Sec. VI. But first we build in Sec. III a specific model of neutrino

[1]The recent Super-Kamiokande data can be explained by ν_μ-ν_τ oscillations while the situation where the ANP is exclusively due to ν_μ-ν_s oscillations is disfavored [4]. However, the more general case where ν_μ oscillates into ν_τ and ν_s with comparable rates is completely consistent with the data [5].

masses with spontaneous breaking of lepton number, and in Sec. IV we derive the relations between the particle asymmetries in the early universe and the present baryon number. The model aims to fit the present known observations from solar, atmospheric, and terrestrial neutrino experiments, including the Liquid Scintillation Neutrino Detecter (LSND) result [19], with an extra sterile neutrino. However, it should be emphasized that the mechanism we propose of suppression of oscillations into a sterile neutrino at the BBN epoch, based on the existence of a Majoron field, does not depend on the particular model or set of neutrino mass parameters. The only fundamental assumptions are that the lepton number (or a partial lepton number) is spontaneously broken and the breaking scale is around the 1 GeV magnitude. In the next section we also make the point that in the absence of the LSND neutrino mass gap the oscillations of atmospheric and solar neutrinos into a sterile neutrino are no longer correlated with each other. In the last section we draw our conclusions.

II. NEUTRINO MASSES AND MIXING

The existence of a fourth, sterile, neutrino has been suggested as it is the only way of reconciling the atmospheric, solar, and LSND neutrino oscillation evidence and their very different δm^2 mass gap scales. The wide mass gap $[\mathcal{O}(1\text{ eV})]$ that is necessary to explain the LSND result in terms of ν_e-ν_μ mixing requires that the neutrino mass pattern should have a two-doublet structure [20]: one of the doublets consists of ν_e and ν_s or ν_τ, or a linear combination of both, and is responsible for the solar neutrino deficit, and the other one, responsible for the atmospheric neutrino anomaly, consists of ν_μ and ν_τ (or a linear combination of ν_s and ν_τ). These doublets are separated by the LSND mass gap.

It has been shown [21,22] that even after the recent SNO observations [23] both the sterile and active neutrino oscillations are viable solutions of the solar neutrino problem as well as a more general superposition of both. On the other hand, the atmospheric neutrino data seem [4] to favor the ν_τ solution against the sterile neutrino case but the analysis [5] of the most recent data still allows a quite large relative probability, more than 50%, of oscillation into ν_s (the larger the probability $\sin^2\xi$ the smaller the allowed δm^2 range).

A consequence of the LSND large mass gap and the limits from reactor disappearance experiments such as Chooz [24] is that the solar neutrinos ν_e and the atmospheric neutrinos ν_μ must oscillate into states that are essentially orthogonal to each other. In other words, if the solar electron neutrinos oscillate into the linear combination $\nu_{\bar{e}} \equiv \cos\xi\,\nu_s - \sin\xi\,\nu_\tau$, then the atmospheric muon neutrinos necessarily oscillate into the state $\nu_{\bar{\mu}} \equiv \sin\xi\,\nu_s + \cos\xi\,\nu_\tau$. However, the situation would be totally different if the LSND evidence was not present.

To be more specific, let ν_1, ν_2, and $\theta_\odot$ be the mass eigenstates and mixing angle responsible for the solar neutrino deficit and ν_3, ν_4, and $\theta_{\rm atm}$ the states and mixing angle relevant for atmospheric neutrinos. The two pairs are separated by the LSND mass gap and no other specific mass hierarchy has to be assumed. The reactor experiments constrain the mixing matrix elements U_{e3}, U_{e4}, $U_{\mu 1}$, and $U_{\mu 2}$ to be small [20], but not $U_{\tau i}$ or U_{si}. If one neglects all the mixing angles that are necessarily small and irrelevant to explain the present bulk of data, the mixing matrix is given as

$$\nu_1 = \cos\theta_\odot \nu_e - \sin\theta_\odot(\cos\xi\,\nu_s - \sin\xi\,\nu_\tau), \tag{1a}$$

$$\nu_2 = \sin\theta_\odot \nu_e + \cos\theta_\odot(\cos\xi\,\nu_s - \sin\xi\,\nu_\tau), \tag{1b}$$

$$\nu_3 = \cos\theta_{\rm atm}\nu_\mu - \sin\theta_{\rm atm}(\sin\xi\,\nu_s + \cos\xi\,\nu_\tau), \tag{1c}$$

$$\nu_4 = \sin\theta_{\rm atm}\nu_\mu + \cos\theta_{\rm atm}(\sin\xi\,\nu_s + \cos\xi\,\nu_\tau). \tag{1d}$$

The mass eigenstates ν_1 and ν_2 are separated by the gap $\delta m^2_\odot$ and ν_3, ν_4 by $\delta m^2_{\rm atm}$.

Clearly, the less the atmospheric neutrinos oscillate into the sterile neutrino the more the solar neutrinos have to oscillate into ν_s. There is a potential clash in the future if both solar and atmospheric neutrino experiments happen to constrain the respective sterile neutrino solutions to less than a 50% probability. In that case the conflict with the LSND data will be insoluble, which will call for new results from MiniBooNE [25], the next new independent accelerator experiment. Suppose for a moment that the LSND evidence does not exist or is going to be ruled out by the MiniBooNE experiment. We would like to stress that this does not rule out the sterile neutrino as a possible protagonist in the other solar and atmospheric neutrino problems, not even if both of them exclude dominant sterile neutrino solutions. On the contrary, the absence of the LSND mass gap increases the freedom in the neutrino mixing parameters. Then the two-doublet mass pattern is no longer inevitable and the role of ν_s in the solar neutrino deficit is completely decoupled from its role in the atmospheric neutrino oscillations.

As a matter of proof we make explicit an extreme case, namely, where the atmospheric neutrinos oscillate into ν_τ or ν_s with arbitrary relative probabilities, while the solar neutrinos oscillate exclusively to ν_τ and ν_μ but not to ν_s. The mixing matrix can be described as follows:

$$\nu_1 = \cos\theta_\odot \nu_e - \sin\theta_\odot(-\sin\alpha\,\nu_\mu + \cos\alpha\,\nu_\tau), \tag{2a}$$

$$\nu_2 = \sin\theta_\odot \nu_e + \cos\theta_\odot(-\sin\alpha\,\nu_\mu + \cos\alpha\,\nu_\tau), \tag{2b}$$

$$\nu_3 = \sin\beta\,\nu_s + \cos\beta(\cos\alpha\,\nu_\mu + \sin\alpha\,\nu_\tau), \tag{2c}$$

$$\nu_4 = \cos\beta\,\nu_s - \sin\beta(\cos\alpha\,\nu_\mu + \sin\alpha\,\nu_\tau). \tag{2d}$$

As far as the mass spectrum is concerned, the mass eigenstate ν_3 is separated from the other three by mass gaps that are in the atmospheric neutrino range $\sim\delta m^2_{\rm atm} \sim 3\times 10^{-3}\ \text{eV}^2$. ν_1 and ν_2 are almost degenerate and separated by the solar neutrino mass gap and, finally, ν_4 is only subject to the condition $m_3^2 - m_4^2 \sim \delta m^2_{\rm atm}$, as it, like ν_3, does not participate in the solar neutrino oscillations. The mixing angles relate to the atmospheric mixing angle as $\cos\theta_{\rm atm} = \cos\alpha\cos\beta$. The atmospheric neutrinos ν_μ oscillate into ν_τ with a probability proportional to $\sin^2\alpha\cos^2\beta$ whereas the probability of oscillation into ν_s is proportional to $\sin^2\beta$. The ratio between them is given by $\tan^2\xi = \tan^2\beta/\sin^2\alpha$. It is clear that the solar neutrinos do not oscillate into ν_s.

This just shows that the potential problem raised by the possibility that the atmospheric neutrinos oscillate significantly into a sterile neutrino with its consequences for BBN is not necessarily linked to the solar neutrino solutions and does not depend on the LSND evidence, although it has been motivated by the coexistence of all three kinds of observation. In the present work we want to present a solution and a mechanism to block the oscillations of muon neutrinos into sterile neutrinos in the early universe at the time of BBN. The idea does not depend crucially on the particular neutrino mixing pattern but the actual numbers vary, of course, from model to model. We worked out in detail a particular model that is suitable to encompass all three types of neutrino oscillation evidence, including LSND.

III. NEUTRINO MASS MODEL

The seesaw mechanism [26] can be incorporated within a model where the lepton number is spontaneously broken at a relatively low energy scale by adding to the standard lepton doublets l_i and charged singlets e_i two heavy sterile neutrinos per lepton generation, N_{Li} and N_{Ri}^C (left-handed), with lepton numbers $+1$ and -1, respectively. The additional light sterile neutrino ν_s (left-handed) has lepton number $L_s = -3$. The most general Yukawa interaction Lagrangian in the lepton sector is written in Majorana matrix form as

$$\mathcal{L}_Y = \frac{1}{2}\psi^T C \mathcal{M} \psi + \text{H.c.}, \tag{3}$$

where $\psi \equiv (e_i^C, l_i, \nu_s, N_{L\,i}, N_{R\,i}^C)$ and $\mathcal{M}$ is the symmetric matrix

$$\mathcal{M} = \begin{array}{c|ccccc} & e_j^C & l_j & \nu_s & N_{L\,j} & N_{R\,j}^C \\ \hline e_i^C & 0 & h_e^T H_1 & 0 & 0 & 0 \\ l_i & - & 0 & 0 & 0 & h_N H_2 \\ \nu_s & - & - & 0 & h_s^T \sigma^* & 0 \\ N_{L\,i} & - & - & - & h_L \sigma & M \\ N_{R\,i}^C & - & - & - & - & h_R \sigma^* \end{array}. \tag{4}$$

The omitted elements are obtained by symmetrization. H_1 and H_2 are two standard Higgs doublets under SU(2) and σ is the singlet scalar field with lepton number $L_\sigma = -2$. h_s is a 3×1 column and h_e, h_N, h_L, h_R, and M are 3×3 matrices.

Before lepton number spontaneous breaking the heavy sterile neutrinos form Dirac particles, namely, $N_i = N_{L\,i} + N_{R\,i}$, with lepton number equal to 1 and masses M_i in the basis where M is diagonal: $M = \text{diag}(M_i)$. After lepton and gauge symmetry breaking the light neutrinos acquire masses and mix with the sterile neutrino in a 4×4 Majorana mass matrix. Denoting the 3×3 active, 3×1 active-sterile, and 1×1 sterile neutrino blocks, respectively, as $m_{\nu\nu}$, $m_{\nu s}$, and m_{ss}, we obtain in leading order in any basis where M is a real matrix

$$m_{\nu\nu} = h_N M^{-1} h_L (h_N M^{-1})^T \langle\sigma\rangle v_2^2, \tag{5a}$$

$$m_{\nu s} = -h_N M^{-1} h_s \langle\sigma\rangle v_2, \tag{5b}$$

$$m_{ss} = (M^{-1} h_s)^T h_R M^{-1} h_s \langle\sigma\rangle^3, \tag{5c}$$

where $v_2 = \langle H_2^0\rangle$. We take as reference scales $m_{\nu\nu} \sim 0.05$ eV to account for the atmospheric neutrino anomaly and $m_{\nu s} \sim 1$ eV for the LSND $\nu_\mu(\overline{\nu_\mu}) \to \nu_e(\overline{\nu_e})$ evidence. Since we also assume $\langle\sigma\rangle \sim 1$ GeV the element $m_{ss} \sim 10^{-13}$ eV is completely negligible.

IV. ASYMMETRIES IN THE EARLY UNIVERSE

At temperatures below the heavy neutrino masses $M_i \gtrsim 10^6$ GeV, the Dirac masses M_i still mediate scattering processes capable of producing the light singlet particles ν_s and σ like $l H_2 \to \overline{\nu_s}\sigma$. They can be studied in terms of the effective operators

$$\mathcal{L}_{\text{eff}} = l_i H_2 \frac{m_{ij}}{2\langle\sigma\rangle v_2^2} l_j H_2 \sigma + l_i H_2 \frac{m_{is}}{\langle\sigma\rangle v_2} \nu_s \sigma^* + \nu_s \frac{m_{ss}}{2\langle\sigma\rangle^3 v_2} \nu_s \sigma^{*3} + \text{H.c.}, \tag{6}$$

which also give rise to the light neutrino masses after spontaneous breaking of lepton number. One obtains the c.m. cross sections of the scattering processes (1) $\overline{l_i l_j} \to H_2 H_2 \sigma$, (2) $\sigma \overline{H_2} \to l_i \nu_s$, and (3) $\sigma\sigma \to \overline{\sigma} \nu_s \nu_s$ as

$$\sigma_1 = \frac{6\,s}{(8\pi)^3}\frac{|m_{ij}|^2}{\langle\sigma\rangle^2 v_2^4} \approx \frac{T^2}{\langle\sigma\rangle^2 v_2^4} \times 10^{-5}\ \text{eV}^2, \tag{7}$$

$$\sigma_2 = \frac{1}{8\pi}\frac{|m_{is}|^2}{\langle\sigma\rangle^2 v_2^2} \approx \frac{4}{\langle\sigma\rangle^2 v_2^2} \times 10^{-2}\ \text{eV}^2, \tag{8}$$

$$\sigma_3 = \frac{6\,s}{(8\pi)^3}\frac{|m_{ss}|^2}{\langle\sigma\rangle^6} \approx \frac{T^2}{\langle\sigma\rangle^6} \times 10^{-28}\ \text{eV}^2, \tag{9}$$

respectively, where we have summed over initial and final weak isospin states ($\sqrt{s}$ is the c.m. energy).

In each case one compares the rate of collisions per particle $\Gamma = \sigma n \approx 0.1\sigma T^3$ (the boson number density is $n_b \simeq 0.122 T^3$ and the fermion number density $n_f \simeq 0.091 T^3$) with the Hubble rate $H \approx T^2/10^{18}$ GeV, assuming a total number of degrees of freedom around 100. The scalar singlet σ is produced through the processes (1) $\overline{l_i l_j} \to H_2 H_2 \sigma$, $\overline{l_i H_2} \to l_j H_2 \sigma$, and $\overline{H_2 H_2} \to l_i l_j \sigma$ with cross sections σ_1, $2\sigma_1/3$, and $\sigma_1/3$, respectively (for $i \neq j$) which gives a total rate per Hubble time $\Gamma_\sigma H^{-1} \approx 2(v/v_2)^4 T^3/10^{15}$ GeV^3 ($v \simeq 174$ GeV is the electroweak breaking scale). This shows that σ is in thermal equilibrium at temperatures larger than $T_\sigma \approx 10^5$ GeV, if one takes $v_2 = v$.

The sterile neutrino is produced in the processes (2) $\sigma \bar{H}_2 \to l_i \nu_s$, $\bar{l}_i \sigma \to H_2 \nu_s$, and $\bar{l}_i \bar{H}_2 \to \bar{\sigma} \nu_s$, with cross sections σ_2, $\sigma_2/2$, and $\sigma_2/2$, respectively, and a total rate $\Gamma_s H^{-1} \approx (v/v_2)^2 T/T_s$, which makes the decoupling temperature of the light sterile neutrino $T_s \approx 4\times 10^6$ GeV. If one or more heavy neutrinos N_i have masses under that value, ν_s may decouple when some of the N_i degrees of freedom are still present in the universe (note that $M_i \gtrsim 10^6$ GeV). Finally, processes like $\sigma\sigma \to \bar{\sigma} \nu_s \nu_s$ are too weak to be relevant.

Above T_s the sterile neutrino and scalar singlet σ are in chemical equilibrium with the lepton and Higgs doublets and their number asymmetries are constrained by the equations of detailed balance. The precise relations between the particle asymmetries depend on which particles and processes are in thermodynamical equilibrium at a given time. To be definite we assume that by the time the sterile neutrino decouples B and L are only violated by electroweak instanton processes while $B-L$ is conserved. On the other hand, the right-handed electrons e_R are not yet in chemical equilibrium and the quarks u_R, d_R may or may not be in equilibrium depending on the exact values of their Yukawa couplings and temperature T_s. In either case the equations of detailed balance yield the particle asymmetries as functions of the $B-L$ asymmetry.

At temperatures above T_s the operators of Eq. (6) yield the chemical potential constraints

$$\mu_\sigma + 2\mu_l + 2\mu_H = 0, \tag{10a}$$

$$\mu_s - \mu_\sigma + \mu_l + \mu_H = 0. \tag{10b}$$

The other constraints come from standard model reactions [13,14]. To be definite we assume that u_R and d_R are in equilibrium at T_s (the temperature at which they come into equilibrium increases with increasing Yukawa couplings and therefore with increasing number of Higgs doublets). The electroweak and QCD instantons and the Yukawa interactions imply that

$$3\mu_q + \mu_l = 0, \tag{11a}$$

$$2\mu_q - \mu_u - \mu_d = 0, \tag{11b}$$

$$\mu_q - \mu_d - \mu_H = 0, \tag{11c}$$

$$\mu_q - \mu_u + \mu_H = 0, \tag{11d}$$

$$\mu_l - \mu_\tau - \mu_H = 0, \tag{11e}$$

where μ_q, μ_l, and μ_H designate the flavor universal chemical potentials of the quark, lepton, and Higgs doublets, respectively, μ_u, μ_d those of the right-handed quark isosinglets, and μ_τ the common chemical potential of the lepton isosinglets μ_R and τ_R. Since the electron singlet e_R is not in chemical equilibrium its chemical potential μ_e is an independent variable. We may assume that a baryon asymmetry originally produced in a GUT baryogenesis scenario is later communicated through electroweak instantons to the lepton sector ($T \lesssim 10^{12}$ GeV) but not to e_R. In that case μ_e remains zero until the e_R Yukawa interactions come into equilibrium at temperatures lower than T_s.

A vanishing weak hypercharge implies

$$3(\mu_q + 2\mu_u - \mu_d - \mu_l) - 2\mu_\tau - \mu_e + 2n_H \mu_H = 0. \tag{12}$$

Here n_H is the total number of Higgs doublets; $n_H = 1$ if H_1 and H_2 are the same field and $n_H = 2$ otherwise. The above constraints and the condition $\mu_e = 0$ leave only one independent variable. It is convenient to choose this as μ_s because the ν_s abundance and number asymmetry are conserved after its decoupling. The other quantity that is conserved is $B-L$. Denoting the baryon number density as $dB/dV = \bar{B} T^3/6$ and likewise for L and $B-L$, one has

$$\bar{B} = 6\mu_q + 3\mu_u + \mu_d, \tag{13a}$$

$$\bar{L} = \bar{L}_{\not{s}} - 3\mu_s, \tag{13b}$$

$$\bar{L}_{\not{s}} = 6\mu_l + 2\mu_\tau + \mu_e - 4\mu_\sigma + 2n_N \mu_N, \tag{13c}$$

where $L_{\not{s}}$ stands for the lepton number of all particles except ν_s and n_N is the number of relativistic Dirac heavy neutrinos at a given moment. The decay processes $N_i \to l_j H_2$ set the equation $\mu_N = \mu_l + \mu_H$. Putting everything together, one obtains

$$B - L_{\not{s}} = \frac{1}{3}\left(8 + 2n_N - \frac{15 + 4n_H}{9 + n_H}\right)(N_s - N_{\bar{s}}), \tag{14}$$

$$\frac{B-L}{B-L_{\not{s}}} = \frac{(9+n_H)(17+2n_N) - 15 - 4n_H}{(9+n_H)(8+2n_N) - 15 - 4n_H}. \tag{15}$$

After ν_s decoupling, $B-L$, $B-L_{\not{s}}$, and the ν_s number asymmetry $N_s - N_{\bar{s}}$ are all conserved. Although the above relations are strictly valid only when ν_s is in thermal equilibrium, one expects that the decoupling process does not introduce very large perturbations and one may use these results as a first approximation. They give the ν_s number asymmetry and $B-L_{\not{s}}$ as functions of the primordial $B-L$ and number n_N of heavy neutrinos that are relativistic when ν_s decouples.

The next transition is the decoupling of the scalar singlet σ at a temperature T_σ around 10^5 GeV. This is close to the e_R coupling epoch which starts at $T_e \sim (v/v_1)^2 \times 10^4$ GeV. This temperature rises with increasing electron Yukawa coupling and if one assumes the existence of two Higgs doublets and $v_1 \lesssim v/3$ ($v \simeq 174$ GeV) then e_R is already in equilibrium when σ decouples. To be definite we assume so. One repeats the exercise with Eqs. (11) and (12), complemented with $\mu_\tau = \mu_e$ and Eq. (10a), to obtain all chemical potentials in terms of μ_σ. Then Eqs. (13a) and (13c) with $n_N = 0$ yield the relation between the σ number asymmetry and the baryon and lepton numbers. Denoting the total lepton number of the standard model particles as L_l ($L_l = L_e + L_\mu + L_\tau$) one derives, for $n_H = 2$,

$$N_\sigma - N_{\bar\sigma} = \frac{12}{23}(B - L_l), \tag{16}$$

$$N_s - N_{\bar s} = \frac{47}{23}\,\frac{33}{65+22n_N}(B - L_l), \tag{17}$$

$$\frac{B-L}{B-L_l} = \frac{47}{23}\,\frac{164+22n_N}{65+22n_N}, \tag{18}$$

where in the last two equations we used Eqs. (14) and (15), keeping in mind that n_N is the number of Dirac neutrinos N_i that are relativistic when ν_s decouples. Again, one expects that the above results remain a reasonable approximation when σ decouples. From then on, $B - L_l$, $B - L$, and the σ and ν_s abundances are conserved by all effective interactions.

When the temperature drops down to the electroweak phase transition the weak isospin and hypercharge are no longer conserved, contrary to the electric charge. The quarks and charged leptons form Dirac mass eigenstates whose well defined chemical potentials are subject to a new set of constraints [14], together with the neutrinos and charged Higgs and W bosons as follows:

$$\mu_{W^+} = \mu_u - \mu_d = \mu_\nu - \mu_e = \mu_{H^+}, \tag{19}$$

$$3\mu_u + 3\mu_d + \mu_\nu + \mu_e = 0. \tag{20}$$

The latter constraint is due to sphaleron processes. On the other hand, the net electric charge is zero:

$$3(4\mu_u - 2\mu_d - 2\mu_e) + 2(2 + n_H)\mu_{W^+} = 0. \tag{21}$$

Definite B and L numbers are predicted in terms of the preexisting $B - L_l$ number, namely [14],

$$B = \frac{32+4n_H}{98+13n_H}(B - L_l), \tag{22}$$

$$L_l = -\frac{66+9n_H}{98+13n_H}(B - L_l). \tag{23}$$

These relations are preserved during the phase transition if there is no intrinsic electroweak baryogenesis. After that the sphaleron processes stop being effective and the baryon number is separately conserved. This allows us to predict the scalar σ particle and sterile neutrino asymmetries in terms of the present baryon number. From Eqs. (16) and (17), valid for two Higgs doublets ($n_H = 2$), one derives the asymmetries and lepton numbers $L_{(\sigma)} = -2(N_\sigma - N_{\bar\sigma})$ and $L_{(s)} = -3(N_s - N_{\bar s})$ carried by σ and ν_s as

$$L_{(\sigma)} = -A_\sigma B = -\frac{24}{23}\,\frac{31}{10}B, \tag{24}$$

$$L_{(s)} = -\frac{47}{23}\,\frac{31}{10}\,\frac{99}{65+22n_N}B. \tag{25}$$

It is important to notice that, after spontaneous breaking of lepton number, the $B - L$ violating processes are too weak to be in equilibrium, in particular during the electroweak phase transition if it occurs after L breaking. If that was not the case the baryon and lepton numbers would be washed out. As soon as the sphalerons decouple the baryon and lepton numbers start to be separately conserved. This is also true after L spontaneous breaking because the L violating reactions are weak. The lepton number may only be significantly violated at much lower temperatures of the order of 1 to 10 MeV when neutrino oscillations from active to sterile neutrinos become possible. Another point is that, after L breaking, the lepton number $L_{(\sigma)}$ carried by the scalar singlet σ still exists but is then associated with a coherent Majoron field. This is the subject of the next section.

V. MAJORON FIELD

The Majoron equation of motion is determined by the equation of conservation of the lepton number Noether current. The lepton current of the scalar field σ with lepton number $L_\sigma = -2$ is

$$J^\mu_\sigma = L_\sigma i\langle \sigma^* \nabla^\mu \sigma - \sigma \nabla^\mu \sigma^* \rangle. \tag{26}$$

At the classical level the total lepton number of charged leptons, neutrinos, and scalar σ is conserved but electroweak instanton effects break L explicitly. $B - L$ remains conserved and its equation of conservation reads as

$$\nabla_\mu J^\mu_\sigma + \nabla_\mu J^\mu_f = 0, \tag{27}$$

where J^μ_f is the $L - B$ current of all the other particles, in our case leptons and quarks:

$$J^\mu_f = -\sum (B_f - L_f)(n_f - n_{\bar f})v^\mu, \tag{28}$$

where n_f and $n_{\bar f}$ are the particle and antiparticle densities and $v^\mu = (1, \mathbf{v})$ the macroscopic velocity vector.

Before spontaneous breaking of the lepton number the σ current is related to the σ particle asymmetry, $J^\mu_\sigma = L_\sigma (n_\sigma - n_{\bar\sigma})v^\mu$, but after lepton symmetry breaking, the mass eigenstates are no longer the complex field σ but rather the massive Higgs particle ρ and massless Majoron boson φ. They relate to each other as

$$\sigma = \frac{1}{\sqrt{2}}(v_\sigma + \rho)\exp(-i\,\varphi/v_\sigma) \tag{29}$$

and the lepton current is expressed as

$$J^\mu_\sigma = L_\sigma v_\sigma \left\langle \left(1 + \frac{\rho}{v_\sigma}\right)^2 \nabla^\mu \varphi \right\rangle. \tag{30}$$

As emphasized in Refs. [18,17], after symmetry breaking the global symmetry is realized as an invariance under translations of the Majoron field whose equation of motion is determined by the still valid Eq. (27) of lepton number conservation.

After lepton breaking the current J^μ_σ can be realized only through a coherent Majoron field φ; in other words, the expectation value of σ has a variable phase:

$$\langle\sigma\rangle=\frac{v_\sigma}{\sqrt{2}}\exp(-i\,\varphi/v_\sigma). \tag{31}$$

We obtain for the current

$$J^\mu_\sigma=L_\sigma v_\sigma\left(1+\frac{\langle\rho^2\rangle}{v_\sigma^2}\right)\nabla^\mu\varphi, \tag{32}$$

where the term $\langle\rho^2\rangle$ is an average over quantum fluctuations, i.e., the thermal bath of massive Higgs particles ρ. When the temperature of the ρ bosons, lower than the photon temperature when the number of relativistic degrees of freedom drops down one order of magnitude, is much smaller than the breaking scale v_σ, the $\langle\rho^2\rangle$ term can be neglected. Then, $J^\mu_\sigma=L_\sigma v_\sigma\nabla^\mu\varphi$. In a homogeneous and isotropic universe φ depends only on time and

$$J^0_\sigma=L_\sigma v_\sigma\dot{\varphi}=F_L\dot{\varphi}. \tag{33}$$

The value of $\dot{\varphi}$ is subject to the equation of conservation (27). Integrating over space, the σ lepton charge $L_{(\sigma)}=\int dV\,J^0_\sigma$ is determined at a given time by its initial value and the variation of the $B-L$ number carried by leptons and quarks:

$$L_{(\sigma)}(t)=L_{(\sigma)}(t_i)+(B-L)_f(t)-(B-L)_f(t_i). \tag{34}$$

The initial value of $L_{(\sigma)}$ is the lepton charge carried by the complex bosons σ before spontaneous lepton breaking. This value is proportional to the initial $B-L$ or to the present baryon number, as Eq. (24) shows for the particular model we worked out. On the other hand, $B-L$ is possibly violated by neutrino oscillations only at very low temperatures when B is conserved. As a result, $L_{(\sigma)}(t)=-A_\sigma B-\Delta L_f$ and the Majoron time derivative is obtained from Eq. (33) as

$$\dot{\varphi}=-\frac{n_\gamma}{F_L}(A_\sigma\hat{B}+\Delta\hat{L}) \tag{35}$$

in terms of the baryon number and lepton number variation per photon

$$\hat{B}=\frac{B}{N_\gamma},\quad \Delta\hat{L}=\frac{\Delta L_f}{N_\gamma}. \tag{36}$$

Notice that L_f and $\Delta\hat{L}$ count only the fermion particles, charged leptons and neutrinos, but not the σ field.

VI. NEUTRINO OSCILLATIONS

Neutrino oscillations [27] are governed by the neutrino masses and mixing angles and interactions with the background medium [28], which in turn depend on the temperature and particle number asymmetries. As far as standard model weak interactions are concerned the electron, proton, and neutron asymmetries, closely related to the baryon asymmetry, are too small to play a role in the oscillation of active neutrinos into sterile neutrinos. However, the neutrino asymmetries can in principle be much larger than the baryon asymmetry. Normalizing them to the photon density as

$$\hat{L}_a=\frac{N_{\nu_a}-N_{\bar{\nu}_a}}{N_\gamma}, \tag{37}$$

the potential of the flavor $\nu_a=\nu_e,\nu_\mu,\nu_\tau$ induced by electroweak interactions is given at low temperatures ($T<m_\mu\ll M_W$) by [29]

$$V_{\rm EW}=\pm\sqrt{2}G_F n_\gamma(\hat{L}_a+\hat{L}_e+\hat{L}_\mu+\hat{L}_\tau\mp A_a T^2 M_W^{-2}), \tag{38}$$

where $A_e=55$ and $A_{\mu,\tau}=15.3$ (the electron and nucleon asymmetries are neglected). The upper sign holds for neutrinos and the lower sign for antineutrinos. The sterile neutrino has no standard model potential by definition.

In the case of interest, $\delta m^2\sim 3\times 10^{-3}$ eV^2, the thermal contribution proportional to $n_\gamma T^2$ prevents the oscillation into a sterile neutrino at temperatures above ~ 10 MeV. At smaller temperatures that term becomes ineffective and the active neutrino oscillates into the sterile flavor as in vacuum, violating the bounds on the number of light degrees of freedom at BBN [6]. Foot and Volkas [9] pointed out that, if there is an initial asymmetry $\hat{L}^a=\hat{L}_a+\hat{L}_e+\hat{L}_\mu+\hat{L}_\tau$ larger than 7×10^{-5} and $\delta m^2\lesssim 10^{-2}$ eV^2, the active neutrino ν_a cannot significantly oscillate into the sterile ν_s and the initial neutrino asymmetries are preserved until the active neutrino ν_a decouples or, in the case of ν_e, until the protons and neutrons stop being in equilibrium. In these conditions the BBN bounds on the extra light degrees of freedom are satisfied. These straightforward considerations have one price, which is the assumption of an initial neutrino asymmetry five orders of magnitude larger than the baryon asymmetry.[2]

The situation changes if there is a Majoron field. Majorons, like any Nambu-Goldstone boson, have only derivative couplings. As a result, a coherent Majoron field produces neutrino potentials proportional to its gradient [18]. If Λ is the spontaneously broken lepton number, in general, any combination of partial lepton numbers, and Λ_a the quantum number of the flavor ν_a, a Majoron field φ produces the potential

$$V_\Lambda=-\frac{1}{F_\Lambda}\Lambda_a v^\mu\partial_\mu\varphi \tag{39}$$

[2]In Ref. [30] Foot and Volkas explored the case where $\nu_\tau\to\nu_s$ oscillations with $-\delta m^2\gtrsim 10$ eV^2 and $\sin^2 2\theta\lesssim 10^{-5}$ create a lepton asymmetry large enough to block $\nu_\mu\to\nu_s$ oscillations with ANP parameters. In any case, this cannot be a generic situation and requires some agreement in the parameter space — the masses and mixing of all neutrino species.

for the neutrino ν_a and the symmetric one for the antineutrino $\overline{\nu_a}$, where $v^{\mu}=(1,\mathbf{v})$ is the neutrino four-velocity ($|\mathbf{v}|=1$ in leading order). In the present case the Majoron is associated with the total lepton number L and the Majoron field is a uniform field in the early universe given by Eq. (35). Hence, it induces the potentials[3]

$$V_L=F_L^{-2}n_{\gamma}L_a(A_{\sigma}\hat{B}+\Delta\hat{L}), \tag{40}$$

where A_{σ} is a model dependent coefficient of order 1; $A_{\sigma}\simeq 3.2$ in the case we are considering. The quantum numbers are $L_a=+1$ (-1) for an active neutrino (antineutrino) and $L_a=-3$ $(+3)$ for the sterile neutrino ν_s $(\overline{\nu_s})$.

The variation of the lepton number in the neutrino sector can be caused only by oscillations into the sterile neutrino because this is the only one with lepton number different from ± 1 and the types of oscillations we are considering conserve chirality. The oscillations $\nu_a\leftrightarrow\nu_s$ $(\overline{\nu_a}\leftrightarrow\overline{\nu_s})$ produce a lepton number variation $\Delta L=\Delta N_a-3\Delta N_s=-4\Delta N_s$ $(\Delta L=4\Delta N_{\bar{s}})$; hence,

$$\Delta\hat{L}=-4\,\frac{\Delta N_s-\Delta N_{\bar{s}}}{N_{\gamma}}. \tag{41}$$

Combining Eqs. (38) and (40), the difference between active and sterile neutrino potentials is

$$\begin{aligned} V_a-V_s=&\sqrt{2}G_F n_{\gamma}(\pm\hat{L}^a-A_aT^2M_W^{-2})\\ &\pm 4F_L^{-2}n_{\gamma}(A_{\sigma}\hat{B}+\Delta\hat{L}), \end{aligned} \tag{42}$$

where $\hat{L}^a=\hat{L}_a+\hat{L}_e+\hat{L}_{\mu}+\hat{L}_{\tau}$ and the lower signs apply to antineutrinos. It is now clear that if, in the case of standard weak interactions, an asymmetry $\hat{L}^a>7\times 10^{-5}$ is enough to block the $\nu_a\leftrightarrow\nu_s$, $\overline{\nu_a}\leftrightarrow\overline{\nu_s}$ oscillations then, in the presence of a Majoron field, the known baryon asymmetry $\hat{B}=(4\text{–}7)\times 10^{-10}$ can do the same if the scale of lepton number breaking F_L obeys the condition

$$F_L^2<\frac{4A_{\sigma}\hat{B}}{7\sqrt{2}G_F}\times 10^5=5\text{–}9\ \ \mathrm{GeV}^2. \tag{43}$$

Such a low scale of lepton number breaking is perfectly consistent with the existing bounds for this kind of singlet Majoron model [16,32,18] including the astrophysical bounds [33]. The reason is that in scattering processes the Majorons couple primarily to neutrinos with strengths proportional to the neutrino masses $g\sim m_{\nu}/F_L$, which are therefore negligibly small for the assumed neutrino mass spectrum even if $F_L\sim 1$ GeV.

[3]They should not be confused with the potentials induced by a thermal bath of Majoron particles, proportional to the neutrino masses, which could be significant for rather large neutrino masses like 17 keV as considered in Ref. [31].

VII. CONCLUSIONS

It is well known that an explanation of the atmospheric neutrino problem [3] in terms of oscillations of the muon neutrino into a, at least in part, sterile neutrino is in contradiction with the BBN limits on the number of extra light degrees of freedom [6]. Indeed, for such a mass gap, $\delta m^2\sim 3\times 10^{-3}$ eV^2, in view of the large mixing angle, the sterile neutrino would come into equilibrium with the particle thermal bath via ν_{μ}-ν_s oscillations [7].

In this paper we formulate a mechanism capable of blocking the active-sterile neutrino oscillations that operates in the framework of theories where the lepton number is spontaneously broken. It has been shown [18] that a generic feature of these theories is the production of coherent, long-range Majoron fields. While they can appear in stars as a result of neutrino oscillations or any other lepton number violating large scale process, we found that in the early universe a cosmological Majoron field emerges also as a result of a primordial lepton asymmetry carried by the scalar particles and complex scalar field (σ) whose expectation value spontaneously breaks the lepton number. The Majoron field amplitude is thus naturally proportional to the lepton number of the universe, and baryon number as well, due to the L and B violating sphaleron processes.

The leptons, and neutrinos in particular, interact with the derivatives of the Majoron field, which gives rise to new neutrino potentials that are relevant for the oscillation phenomena. The potentials are inversely proportional to the second power of the scale of lepton number symmetry breaking ($\langle\sigma\rangle$) but this scale can be much smaller than the electroweak breaking scale. For a lepton number breaking scale of the order of 1 GeV, a Majoron field associated with a lepton number asymmetry of the same order of magnitude as the baryon asymmetry can block the active-sterile neutrino oscillations at temperatures in the MeV range for a neutrino mass gap $\delta m^2\sim 3\times 10^{-3}$ eV^2.

This provides an interesting way to block the sterile neutrino oscillations as it does not require the extraordinarily high lepton asymmetries, $\sim 10^{-5}$, that are necessary [9] in the framework of the standard model weak interactions. In fact, the standard model neutrino potentials that are proportional to the background lepton asymmetries are suppressed by the Fermi constant. In contrast, the potentials due to a Majoron field vary as the inverse square lepton number symmetry breaking scale. For a 1 GeV energy scale one immediately obtains the five orders of magnitude increase factor that brings 10^{-5} down to the baryon number asymmetry $B\sim 10^{-10}$.

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