Long Range Majoron Fields and Implications for Neutrino Oscillations

Luís Bento

Centro de Física Nuclear da Universidade de Lisboa
Av. Prof. Gama Pinto 2, 1649-003 Lisboa, Portugal

Abstract

A long-range Nambu-Goldstone (NG) field can be produced when a large number of reactions violate simultaneously the quantum number associated with it. The source of the NG field is the time rate of variation of that quantum number. If the lepton numbers are spontaneously broken at a scale up to 1 TeV, long range majoron fields can have significant effects on neutrino oscillation processes. Some examples are discussed in the context of supernovae and Early Universe.

1 Introduction

When a global symmetry is spontaneously broken, one or more massless scalar bosons – so-called Nambu-Goldstone bosons – should exist, one per broken quantum number. Good candidates for spontaneously broken quantum numbers are the partial lepton numbers L_e, L_μ, L_τ and the total lepton number L. They are conserved by the standard model (SM) interactions, but several experiments going now from solar to atmospheric and laboratory neutrino experiments indicate that the lepton flavors are not conserved (for recent reviews see [1]).

The NG bosons do not have mass or any either scalar potential terms like ϕ^4. They only have derivative couplings [2]. For that reason, although they interact with fermions they do not mediate long range scalar forces. The very global symmetry that is broken by the vacuum is still a good symmetry of the Lagrangian and is realized as a translation of the NG field, $\phi \to \phi + \alpha$, $\alpha =$ constant. The ϕ equation of motion identifies with the conservation law of that symmetry,

$$\partial_\mu (\partial^\mu \phi + J^\mu_\Lambda / V_\Lambda) = 0 \, . \tag{1}$$

V_Λ is essentially the scale of symmetry breaking and the current J^μ_Λ is determined in leading order by the quantum numbers of the other particles. For instance, if Λ is one of the lepton numbers L_e or L_μ, the current and interaction Lagrangian are given by

$$\begin{aligned} \mathcal{L}_{\rm int} &= J^\mu_\Lambda \, \partial_\mu \phi / V_\Lambda + \text{quadratic terms in } \partial_\mu \phi \ , & (2) \\ J^\mu_\Lambda &= \Lambda_e (\bar{e} \gamma^\mu e + \bar{\nu}_e \gamma^\mu \nu_e) + \Lambda_\mu (\bar{\mu} \gamma^\mu \mu + \bar{\nu}_\mu \gamma^\mu \nu_\mu) + \cdots \ , & (3) \end{aligned}$$

where the dots stand for model dependent scalar boson contributions and higher order effective couplings.

This can be easily understood as follows [2]-[4]. Let Λ_a and Λ_i be the quantum numbers of the fermion fields χ^a and scalars σ_i under an abelian global symmetry $\mathrm{U}(1)_\Lambda$ that is spontaneously broken by the vacuum expectation values $\langle \sigma_i \rangle$. The fields can be expressed as

$$\begin{aligned} \chi^a &= e^{-i \, \Lambda_a \phi / V_\Lambda} \psi^a \ , & (4) \\ \sigma_i &= e^{-i \, \Lambda_i \phi / V_\Lambda} \left(\langle \sigma_i \rangle + \rho_i \right) \ , & (5) \end{aligned}$$

in terms of the *physical weak* eigenstates, $\psi^a = e, \, \nu_e, \, ...$, massive bosons ρ_i and gauge bosons (singlets of the global symmetry) all of them invariant under $\mathrm{U}(1)_\Lambda$. The only non-invariant field is the NG boson ϕ. Its couplings, or current J^μ_Λ, are derived from the kinetic Lagrangians of χ^a and σ_i by applying the transformations (4-5). The source term of the NG equation of motion, Eq. (1), is clearly not an ordinary scalar density. The divergency of a vector or axial-vector fermion current gives after using Dirac equation pseudo-scalar and *off-diagonal* scalar densities,

$$\partial_\mu \, \overline{f}_i \gamma^\mu \gamma_5 f_j = i(m_i + m_j) \overline{f}_i \gamma_5 f_j \ , \qquad \partial_\mu \, \overline{f}_i \gamma^\mu f_j = i(m_i - m_j) \overline{f}_i f_j \ . \tag{6}$$

but not *flavor-diagonal* scalar densities [2, 5]. The pseudo-scalars that are diagonal in flavor vanish for free particle states and the flavor violating densities depend on the relative phase of distinct flavors. The natural conclusion has been that these relative phases cancel each other when summed over a large number of particles and therefore long range '$1/r$' NG fields are not possible, only *spin-dependent* '$1/r^3$' potentials can exist [2, 5]. It has been shown that this is not true [3, 6, 7].

2 Long range Nambu-Goldstone fields

The source of a NG field does not vanish if the matter current J^μ_Λ is not conserved. The volume integral

$$\int d^3x \;\; \partial_\mu J^\mu_\Lambda = \frac{d\Lambda}{dt} \tag{7}$$

represents the rate of *creation* per unity of time of the Λ number carried by all particles. It constitutes the source of the NG boson. A solution of the equation of motion (1) is given by

$$\phi(t,\vec{r}) = \frac{-1}{V_\Lambda}\int d^3x \, \frac{\partial_\mu J^\mu_\Lambda(t-|\vec{r}-\vec{x}|\,,\vec{x})}{4\pi\,|\vec{r}-\vec{x}|}\;. \tag{8}$$

This is a long-range field. Its magnitude depends on the number of Λ violating processes ocurring in a time scale of ${V_\Lambda}^{-1}$. To produce a significant macroscopic field one needs thus a large rate of reactions. That calls for astrophysical sources. There is evidence that solar neutrinos change flavor as they come out of the Sun. If ν_e oscillates into ν_μ, for instance, then L_e and L_μ are not conserved in a large scale in stars and presumably in supernovae as well. If any of these quantum numbers is associated with a spontaneously broken global symmetry then, the respective NG bosons (so-called majorons [5]) acquire classical long-range field configurations.

For definiteness suppose that $\Lambda = L_e - L_\mu$ is spontaneously broken and that a fraction of ν_es emitted from a star are resonantly converted into ν_μ through the Mikheyev-Smirnov-Wolfenstein (MSW) mechanism [8] (also $\nu_\mu \to \nu_e$ in a supernova) in a certain shell inside the star. For simplicity assume spherical symmetry and stationary neutrino fluxes in which case the $L_e - L_\mu$ number carried by electrons and neutrinos decreases at a constant rate in time. The rate $\dot{L}_e - \dot{L}_\mu$ generates a long-range majoron field of Coulombian form – charge divided by $4\pi r$ – above the resonance shell,

$$\phi = \frac{1}{4\pi r}\frac{2}{V_\Lambda}\left[\dot{N}(\nu_e \to \nu_\mu) - \dot{N}(\nu_\mu \to \nu_e)\right]\;, \tag{9}$$

while ϕ is constant below the resonance shell.

This long-range '$1/r$' field configurations are in conflict with the conclusion traditionally taken by looking at the fermion bilinears (6) that appear as the source of a NG boson namely, that only '$1/r^3$' spin-dependent NG potentials could exist. In [7] we showed how the above result is actually derived by careful examining the fermion pseudo-scalar densities in a flavor oscillating system. In the interest of clearness we focus here [4] on a specific system with a pair of neutrino flavors, $\nu_e - \nu_\mu$. They are assumed to have weak interactions with a background medium, interact with a majoron associated with breaking of $\Lambda = L_e - L_\mu$ and mix through a 2×2 *real* Majorana mass matrix m, as follows:

$$\begin{aligned}\mathcal{L}_{\rm int} &= -\sqrt{2}G_{\rm F}n_e(\bar\nu_e\gamma^0\nu_e) - \frac{g\langle Z^0\rangle}{2\cos\theta_{\rm W}}(\bar\nu_e\gamma^0\nu_e + \bar\nu_\mu\gamma^0\nu_\mu)\\ &+ \frac{1}{V_\Lambda}\partial_\mu\phi(\bar\nu_e\gamma^\mu\nu_e - \bar\nu_\mu\gamma^\mu\nu_\mu) - (\nu_L^T C m \nu_L + \text{H. C.})\;.\end{aligned} \tag{10}$$

n_e is the electron number density. Notice that only m_{ee} and $m_{\mu\mu}$ break $L_e - L_\mu$. This Lagrangian determines the equations of motion of the (left-handed) neutrino operators and also of the single particle wave functions, related to them by a relation of the type $\nu_L + \nu_L^C = \int a\,\psi + a^\dagger\,\psi^C$. The wave functions obey the equations of motion

$$i\,\partial\!\!\!/\psi_a = m_{ab}\,\psi_b - V_a^\mu\,\gamma_\mu\gamma_5\psi_a\;, \tag{11}$$

where a stands for ν_e and ν_μ and V_a^μ represent the *flavor conserving* weak and majoron vector potentials ($V_a^\mu = -\Lambda_a \partial^\mu \phi / V_\Lambda$ in the latter case). When $\partial_\mu J_\Lambda^\mu$ is evaluated over a system of neutrino particles the equation of motion (1) becomes [7]:

$$\partial_\mu \partial^\mu \phi = \frac{2}{V_\Lambda} \sum_\nu i \left(m_{ee}\, \bar{\psi}_{\nu_e} \gamma_5 \psi_{\nu_e} - m_{\mu\mu}\, \bar{\psi}_{\nu_\mu} \gamma_5 \psi_{\nu_\mu} \right) . \tag{12}$$

As expected, the parameters m_{ee} and $m_{\mu\mu}$ that break $L_e - L_\mu$ in the Lagrangian (10) appear explicitly in the source terms. These are precisely the kind of pseudo-scalar densities identified in Eq. (6). They would vanish identically if the neutrinos were free particles. But they are not. When studying neutrino oscillations one usually separates the spin from the flavor degrees of freedom by expressing the wave function as a product of a left-handed *free massless* spinor ψ_0^α ($\gamma_5 \psi_0 = -\psi_0$, $i \not{\partial} \psi_0 = 0$) and a flavor-valued wave function φ_a. It obeys a well known evolution equation [9] as a function of the distance s travelled by each particle,

$$i \frac{d}{ds} \begin{pmatrix} \varphi_e \\ \varphi_\mu \end{pmatrix} = \frac{1}{2E} \begin{pmatrix} m_{ee}^2 + 2EV_e & \frac{1}{2}\Delta m^2 \sin 2\theta \\ \frac{1}{2}\Delta m^2 \sin 2\theta & m_{\mu\mu}^2 + 2EV_\mu \end{pmatrix} \begin{pmatrix} \varphi_e \\ \varphi_\mu \end{pmatrix} , \tag{13}$$

see Eqs. (17-19). However, pure chiral spinors give vanishing scalar densities in Eq. (12) so one has to go beyond the approximation $\psi_a = \psi_0 \varphi_a$.

For a spherical symmetric majoron field as in Eq. (9) the velocity of a radially moving neutrino is parallel to the vector potentials $\mathbf{V}_a \propto \nabla\phi$ and in that case an approximate solution of Eq. (11) is [7]

$$\psi_a^\alpha(x) \cong \psi_0^\alpha\, \varphi_a + \gamma_{\alpha\beta}^0 \frac{m_{ab}}{2E}\, \psi_0^\beta\, \varphi_b . \tag{14}$$

Applying to the majoron equation of motion (12) it becomes

$$\begin{aligned} \partial_\mu \partial^\mu \phi &= \frac{1}{V_\Lambda} \sum_\nu \frac{m_{e\mu}}{E} (m_{ee} + m_{\mu\mu})\, i\, (\varphi_e^\dagger \varphi_\mu - \varphi_\mu^\dagger \varphi_e)\, \psi_0^\dagger \psi_0 \\ &= \frac{1}{V_\Lambda} \sum_\nu \frac{\Delta m^2}{2E} \sin 2\theta\, i\, (\varphi_1^\dagger \varphi_2 - \varphi_2^\dagger \varphi_1)\, \psi_0^\dagger \psi_0 . \end{aligned} \tag{15}$$

where $\varphi_{1,2}$ refer to the mass eigenstate basis. It is now clear that the generation of the majoron field requires the interference between different flavor components of the neutrino wave functions. But that simply means the existence of neutrino oscillations. A crucial fact is that the source term does not depend on arbitrary initial phases of the neutrino wave functions. Hence, there is no automatic phase cancellation when the sum is taken over a large number of particles. This of course makes possible to have sources of macroscopic dimensions. Making use of Eq. (13), the last result can be written in terms of the probabilities of observing the ν_e and ν_μ flavors respectively, $P_{\nu_e} = \varphi_e^\dagger \varphi_e / \varphi^\dagger \varphi$ and $P_{\nu_\mu} = 1 - P_{\nu_e}$:

$$\partial_\mu \partial^\mu \phi_\Lambda = \frac{-1}{V_\Lambda} \sum_\nu \psi^\dagger \psi\, \frac{d(P_{\nu_e} - P_{\nu_\mu})}{ds} = \frac{-1}{V_\Lambda} \frac{d(L_e - L_\mu)}{dt\, d^3x} . \tag{16}$$

Again, the second member does not depend on the initial phases of the individual particles. It also tells under which conditions a large number of particles makes a charge with definite sign: it is when the $\Lambda = L_e - L_\mu$ number carried by neutrinos suffers a net increase or decrease in Λ violating processes.

We arrive at the starting point of this section: the source term of a long-range coherent NG field is the time rate of *macroscopic* variation of the respective quantum number carried by matter particles [3, 6]. In a star or reactor where only ν_e are produced out of electrons, the oscillations $\nu_e \to \nu_\mu$ lead necessarily to a total decrease of the lepton number $L_e - L_\mu$: the electrons captured in nuclear reactions are, after all, totally or partially transformed into ν_μ (the probability of observing the ν_μ flavor is not zero). The global variation of a lepton number does not necessarily require an oscillation process. The simple production of neutrino particles implies a violation of lepton number, but is ordinarily suppressed by the ratio Majorana mass over neutrino energy. Such suppression does not take place in processes of leptogenesis in the Early Universe where a heavy Majorana neutrino decays into light neutrinos. The implications of a coherent majoron field in the Early Universe will be discussed in section 3.3.

3 Majoron field effects on neutrino oscillations

Understood as they are the conditions under which a long-range NG field is produced the next question is what are the implications? A NG field only interacts through its gradient which implies a suppression factor of V_Λ^{-1} over the time or length scale of field variation. This is typically an extremely small number. However neutrino oscillations are sensitive to very tiny potentials as long as they are not flavor independent, which is in general the case of the majoron couplings, see Eqs. (2-3). They interact with fermions through the vector potentials $V_a^\mu = -\Lambda_a \partial^\mu \phi / V_\Lambda$ which manifests in the flavor evolution equations like (13) in the form of flavor dependent potentials ($\mathbf{v}$ is the neutrino velocity)

$$V_a = -\Lambda_a (\dot{\phi} + \mathbf{v} \cdot \nabla \phi)/V_\Lambda \; . \tag{17}$$

In the two-flavor oscillation case analyzed in the previous section, the important quantity is the difference between potentials, $V_e - V_\mu$. In a stationary, spherical symmetric regime, the majoron field is given by Eq. (9) and induces the potential

$$(V_e - V_\mu)_{\rm NG} = \frac{4}{V_\Lambda^2} \frac{1}{4\pi r^2} \left[\dot{N}(\nu_e \to \nu_\mu) - \dot{N}(\nu_\mu \to \nu_e) \right] \; . \tag{18}$$

This shows that the typical order of magnitude of the majoron potentials ($V_{\rm NG}$) is the neutrino flux, j_ν, divided by V_Λ^2. In the Sun, $j_\nu \, V_\Lambda^{-2} \sim 10^{-2} R_\odot / r^2$ for a scale V_Λ as low as 1 KeV. The phase shift produced by $V_{\rm NG}$, $\int V_{\rm NG} dr$, seems too small to have any effect on solar neutrino oscillations. In a supernova however the neutrino fluxes are many orders of magnitude higher than in the Sun and the situation changes completely. To compare the majoron with the electroweak potential [8],

$$(V_e - V_\mu)_{\rm SM} = \sqrt{2}\, G_{\rm F}\, n_e \; , \tag{19}$$

one has to compare $V_\Lambda^{-2} j_\nu \simeq V_\Lambda^{-2} n_\nu$ with $G_{\rm F}\, n_e$. It turns out [3, 6] that the neutrino density n_ν surpasses the electron density n_e at a radius of the order of 10^{10} cm, where the matter density is similar to the density in the Sun. Furthemore, the neutrino density decreases with r^{-2} whereas the electron density goes as r^{-3} in a supernova envelope star [10]. Hence, the majoron potentials are not only comparable with the electroweak potentials for scales V_Λ as high as $G_{\rm F}^{-1/2}$ but even overcome at large enough radius.

In principle, a majoron field produced by a certain flavor transition can affect the propagation of other neutrino particles or flavors outside the region it was created or even outside the star. More exciting, the magnitude of the majoron potentials can be significative in regions where the MSW resonance takes place for values of Δm^2 that are relevant for solar or atmospheric neutrinos. It has been shown that majoron fields may improve the adiabaticity of $\nu_e \leftrightarrow \nu_X$ transitions [6] or cause *resonant* anti-neutrino oscillations [3] otherwise impossible with solely standard weak interactions.

In the Early Universe, the system is not stationary but is not very different in the sense that time takes the role of the radius in a star. This will be clear in section 3.3. The majoron fields are functions of time but the majoron induced potentials are again proportional to $V_\Lambda^{-2} n_\nu$ or, to be more exact, $V_\Lambda^{-2}(n_\nu - n_{\bar{\nu}})$. Clearly, the neutrino densities are not small in comparison with the electron or other weakly interacting particles.

3.1 Improvement of adiabaticity: majoron back reaction

Suppose that ν_e oscillates in the Sun into an active neutrino $\nu_X = \alpha_\mu \nu_\mu + \alpha_\tau \nu_\tau$ in accordance with the non-adiabatic, small mixing angle solution [11]-[14]. Then, the same $\nu_e \to \nu_X$ transitions must occur in a supernova envelope, non-adiabatic as well, accompanied by $\nu_X \to \nu_e$ since $\nu_{\mu,\tau}$ are also produced in a supernova. First, as shows the resonance condition,

$$2E(V_e - V_X) = \Delta m^2 \cos 2\theta \; , \tag{20}$$

the neutrinos with higher energies are converted at larger radii. Second, the transitions are less adiabatic and less frequent for the more energetic neutrinos [15] which has important consequences: since the $\nu_{\mu,\tau}$ flavors are emitted with an hotter spectrum than ν_e, the final observable ν_e spectrum is softer than if the transitions were completely adiabatic.

The situation changes [6] if there is a majoron field associated with a partial lepton number, let it be L_e or $L_e - L_X$, as considered in the previous sections. The potential $V_e - V_X$ departs from the charged current potential as shows Eq. (18). There, the rates $\dot{N}(\nu_e \to \nu_X)$ and $\dot{N}(\nu_X \to \nu_e)$ are functions of the radius, in other words, the gradient of the majoron field and the potential induced by it at a certain radius is determined by the total rates of neutrino conversions inside the sphere of that radius.

Since a supernova emits more ν_e than $\nu_{\mu,\tau}$ the majoron potential is positive which has the effect of slowing down the fall of $V_e - V_X$ with the radius. That improves the adiabaticity of the oscillations and conversion efficiency of the more energetic neutrinos. The implication, observable in detectors like Super-Kamiokande and SNO [16], is a ν_e spectrum harder than expected from a non-adiabatic solar neutrino solution (and softer ν_μ, ν_τ spectra) within the framework of electroweak interactions.

3.2 Resonant conversion of anti-neutrinos

If not only one but all three partial lepton numbers are spontaneously broken one expects to exist mixing between the three NG bosons ξ_e, ξ_μ, ξ_τ [3]. By going to the unitary basis as in Eqs. (4-5),

$$\chi^a = \exp\{-i(\xi_e L_e^a + \xi_\mu L_\mu^a + \xi_\tau L_\tau^a)\}\, \psi^a \ , \tag{21}$$

$$\sigma_i = \exp\{-i(\xi_e L_e^i + \xi_\mu L_\mu^i + \xi_\tau L_\tau^i)\}\, (\langle\sigma_i\rangle + \rho_i) \ , \tag{22}$$

the χ^a fermion and σ_i boson kinetic terms deliver the following interaction Lagrangian

$$\mathcal{L}_{\rm int} = \frac{1}{2} V_{\ell m}^2\, \partial_\mu \xi_\ell\, \partial^\mu \xi_m + \partial_\mu \xi_\ell (\bar{\ell}\gamma^\mu\, \ell + \bar{\nu}_\ell \gamma^\mu \nu_\ell) \ , \qquad \ell, m = e, \mu, \tau \ . \tag{23}$$

The kinetic matrix $V_{\ell m}^2$ and its inverse $G_{\ell m}$ (same dimension as $G_{\rm F}$) are not diagonal in general [3]. As a result, long-range majoron fields couple different flavors to each other. Instead of an equation like (18), a neutrino with flavor ν_ℓ feels a NG potential

$$V_\ell = -(G_{\ell e}\, \dot{L}_e + G_{\ell\mu}\, \dot{L}_\mu + G_{\ell\tau}\, \dot{L}_\tau)/4\pi r^2 \ . \tag{24}$$

It means that the violation of one flavor in one place creates majoron fields that act at a distance on any of the other flavors. This may have interesting consequences.

Suppose that resonant $\nu_e \leftrightarrow \nu_H$ oscillations take place in a supernova environment (H stands for Heavy, it may be the heaviest of the ν mass eigenstates involved in atmospheric ν oscillations [17]). In addition ν_e also mixes with ν_L (L means Light, possibly the neutrino involved in solar ν oscillations) and, in absence of other than weak interactions, ν_e oscillates resonantly into ν_L, but not $\bar{\nu}_e \to \bar{\nu}_L$, in a region of lower density than where $\nu_e \leftrightarrow \nu_H$ take place. If majoron fields come into play, the ν_e and ν_L propagation is modified by the potentials in Eq. (24) that are created by the $\nu_e \leftrightarrow \nu_H$ oscillations, with $\dot{L}_e = -\dot{L}_\mu < 0$ because a supernova emits more ν_es than ν_Hs. If the following hierarchy, $|G_{e\ell}| << |G_{LH}|, G_{HH}$, exists between these constants, the relevant potential for $\bar{\nu}_e \leftrightarrow \bar{\nu}_L$ oscillations is approximately

$$V_{\bar{\nu}_e} - V_{\bar{\nu}_L} = -\sqrt{2}\, G_{\rm F}\, n_e \ - G_{LH}\, \dot{L}_H/4\pi r^2 = -(V_{\nu_e} - V_{\nu_L}) \ . \tag{25}$$

In absence of NG fields this potential is negative and, by hypothesis, the resonance condition (20) cannot be satisfied for $\bar{\nu}_e \leftrightarrow \bar{\nu}_L$ contrary to $\nu_e \leftrightarrow \nu_L$. However, if the majoron fields exist and G_{LH} is negative, the NG potential is positive and necessarily overcomes at large enough radius because n_e falls faster than $1/r^2$. Therefore, if Δm_{eL}^2 is sufficiently low and the ν fluxes high enough resonant anti-neutrino oscillations become possible. Such effect can in principle be observed because the predicted energy spectrum of $\bar{\nu}_\mu$ and $\bar{\nu}_\tau$ is harder than the $\bar{\nu}_e$ spectrum.

The SN1987A data [18] seem to disfavor $\bar{\nu}_e \leftrightarrow \bar{\nu}_{\mu,\tau}$ oscillations [19] (see however [20]) but do not exclude that they occur with a probability below a certain limit. As stressed in [3], they may occur in the first instants of ν emission, when the ν luminosity and hypothetical majoron fields are stronger, to disappear later at some ν flux threshold.

3.3 Neutrino Oscillations in the Early Universe

The observed baryon number asymmetry was supposed to be produced in the Early Universe out of an initial state with a zero net baryon number. It is natural that a lepton number asymmetry was also dynamically generated. In fact, some of the proposed mechanisms of baryogenesis go through a primary process of leptogenesis [21] that is converted into baryon asymmetry in spharelon mediated reactions [22]. It is also quite natural that the partial lepton numbers were not conserved and the respective asymmetries developped in the Early Universe. Suppose that the linear combination $\Lambda = \Lambda_e L_e + \Lambda_\mu L_\mu + \Lambda_\tau L_\tau + \Lambda_s L_s$ (s stands for a possible extra sterile neutrino) is a spontaneously broken quantum number associated with a NG boson ϕ. When a Λ asymmetry is generated a majoron field is also produced because the total current $J^\mu = J^\mu_\Lambda + V_\Lambda \partial^\mu \phi$ is conserved. In an homogeneous and isotropic Universe it remains zero everywhere, J^0_Λ and ϕ only depend on time and are related as $\dot\phi = -J^0_\Lambda / V_\Lambda$. This means that $\phi = 0$ before the Λ-genesis (up to an additive constant, meaningless due to the global symmetry) but $\dot\phi$ remains different from zero after the Λ violating processes become ineffective because the Λ number per unity of comoving volume is then conserved and finite. At temperatures below the μ mass the Λ-number density is

$$J^0_\Lambda = \sum_a \Lambda_a (n_{\nu_a} - \bar n_{\nu_a}) + \Lambda_e (n_e - n_{e^+}) \, , \qquad a = e, \mu, \tau, s \, . \tag{26}$$

The electron-positron asymmetry is set to be the same as the baryon asymmetry and therefore very small but the asymmetries in the neutrino sector can in principle be much larger than that. It is well known [23, 24] that lepton asymmetries,

$$\hat L_a = (n_{\nu_a} - \bar n_{\nu_a}) / n_\gamma \, , \tag{27}$$

as large as 10^{-5} or more help preventing the oscillation of standard neutrinos into an extra sterile neutrino at the time of nucleonsynthesis and thus relax the limits on the neutrino masses and mixings with ν_s. The standard model neutrino potentials are given as [25] ($n_e - n_{e^+} = n_p - n_{\bar p}$ and $n_n - n_{\bar n}$ are negligible)

$$\begin{aligned} V_{\rm SM}[\nu_a] &= \sqrt2 G_{\rm F} n_\gamma (\hat L_a + \hat L_e + \hat L_\mu + \hat L_\tau - A_a T^2 M_W^{-2}) \, , \qquad a = e, \mu, \tau \, , \\ V_{\rm SM}[\bar\nu_a] &= -\sqrt2 G_{\rm F} n_\gamma (\hat L_a + \hat L_e + \hat L_\mu + \hat L_\tau + A_a T^2 M_W^{-2}) \, , \qquad a = e, \mu, \tau \, , \end{aligned} \tag{28}$$

where $A_e = 55$ and $A_{\mu,\tau} = 15.3$. In turn, a Λ majoron field contributes with the potentials

$$V_{\rm NG}[\nu_a] = -V_{\rm NG}[\bar\nu_a] = -\Lambda_a \dot\phi / V_\Lambda = V_\Lambda^{-2} n_\gamma \sum_b \Lambda_a \Lambda_b \hat L_b \, , \qquad a, b = e, \mu, \tau, s \, . \tag{29}$$

Besides its different flavor dependence, this majoron potential may actually be dominant because the scale V_Λ of Λ spontaneous breaking can be significantly smaller than the Fermi scale of weak interactions. If that is the case the limits on the mixing with a sterile neutrino or the minimum required lepton asymmetries to conciliate ν_s with the Nucleosynthesis limits [26] on the number of extra light degrees of freedom are further relaxed [27].

4 Conclusions

Nambu-Goldstone bosons only have derivative couplings but long-range NG fields can still be obtained. Their source is the rate of variation of the quantum number associated with the NG boson per unity of time. If one or more of the lepton numbers are spontaneously broken quantum numbers, the respective NG bosons, called majorons, acquire long range fields in several systems. Neutrino oscillations in stars produce majoron fields as a result of the constructive interference between the wave functions of the oscillating flavors. In the Early Universe, the reactions that may generate global lepton asymmetries necessarily generate coherent majoron fields too. Such fields have significant effects on neutrino oscillations in supernovae and in the Early Universe, competitive with the standard model weak interactions, if the scale of spontaneous breaking of the lepton nummbers is around or under the Fermi scale. This implies a sensitivity to effective neutrino-majoron couplings, $\sim m_\nu / V_\Lambda$, of the order of $m_\nu G_{\rm F}^{1/2}$, orders of magnitude better than is attainable with scattering processes considerations, at most $\sim 10^{-7.5}$ [28]-[33].

In the Early Universe majoron fields may prevent the oscillation into a sterile neutrino flavor at the time of nucleosynthesis. In the case of supernovae, they have a surprise effect. Because in solar, atmospheric and laboratory neutrino experiments the fluxes are most certainly too low to produce significant majoron fields, their effects on supernova neutrinos cannot be anticipated from those kind of experiments which may appear incompatible with future supernova neutrino observations. The other possible signature is the correlation with the time dependent magnitudes of the supernova neutrino fluxes which, if observed, will tell exactly what is the scale of lepton number symmetry breaking.

This work was supported in part by the grant PESO/P/PRO/1250/98. This talk was also presented, with the exception of section 3.3, in the *X International Baksan School: Particles and Cosmology* [4].

References

[1] E. K. Akhmedov, hep-ph/0001264 and hep-ph/0001264; S. M. Bilenky, C. Giunti and W. Grimus, Prog. Part. Nucl. Phys. **43**, 1 (1999).

[2] G. B. Gelmini, S. Nussinov and T. Yanagida, Nucl. Phys. B **219**, 31 (1983); A. A. Anselm and N. G. Uraltsev, Sov. Phys. JETP **57**, 1142 (1983).

[3] L. Bento, Phys. Rev. D **57**, 583 (1998).

[4] L. Bento, in: *Proceedings of the X International School: Particles and Cosmology*, Baksan Valley, Kabardino-Balkaria, Russia, 1999 (hep-ph/9908506).

[5] Y. Chikashige, R. N. Mohapatra and R. D. Peccei, Phys. Lett B **98**, 265 (1981).

[6] L. Bento, Phys. Rev. D **59**, 015013 (1999).

[7] L. Bento and Z. Berezhiani, Phys. Rev. D **62**, 055003 (2000).

[8] L. Wolfenstein, Phys. Rev. D **17**, 2369 (1978); S. P. Mikheyev and A. Yu. Smirnov, Yad. Fiz. **42**, 1441 (1985) [Sov. J. Nucl. Phys. **42**, 913 (1985); Nuovo Cimento **9C**, 17 (1986).

[9] T. K. Kuo and J. Pantaleone, Rev. Mod. Phys. **61**, 937 (1989).

[10] J. R. Wilson *et. al.*, Ann. N. Y. Acad. Sci. **470**, 267 (1986).

[11] W. C. Haxton, Phys. Rev. Lett. **57**, 1271 (1986); S. J. Parke, *ibid*, 1275; S. P. Rosen and J. M. Gelb, Phys. Rev. D **34** , 969 (1986).

[12] N. Hata and P. Langacker, Phys. Rev. D **56**, 6107 (1997); J. N. Bahcall, P. I. Krastev and A. Yu. Smirnov, Phys. Rev. D **58**, 096016 (1998).

[13] C. Giunti, M. C. Gonzalez-Garcia, C. Peña-Garay, Phys.Rev. D **62**, 013005 (2000).

[14] G. L. Fogli, E. Lisi, D. Montanino, A. Palazzo, Phys.Rev. D **62**, 113003 (2000); M. C. Gonzalez-Garcia, M. Maltoni, C. Peña-Garay, J. W. F. Valle, (hep-ph/0009350).

[15] H. Minakata and H. Nunokawa, Phys. Rev. D **38**, 3605 (1988).

[16] A. Burrows, Annu. Rev. Nucl. Sci. **40**, 181 (1990); A. Burrows, D. Klein, and R. Gandhi, Phys. Rev. D **45**, 3361 (1992).

[17] (Super-Kamiokande Collaboration) Y. Fukuda *et al.*, Phys. Rev. Lett. **81**, 1562 (1998); S. Fukuda *et al.*, Phys. Rev. Lett. **85**, 3999 (2000).

[18] K. Hirata *et. al.*, Phys. Rev. Lett. **58**, 1490 (1987); R. M. Bionta, *ibid*, 1494.

[19] B. Jegerlehner, F. Neubig and G. Raffelt, Phys. Rev. D **54**, 1194 (1996); A. Yu. Smirnov, D. N. Spergel and J. N. Bahcall, Phys. Rev. D **49**, 1389 (1994).

[20] P. J. Kernan and L. M. Krauss, Nucl. Phys. B **437**, 243 (1995).

[21] M. Fukujita and T. Yanagida, Phys. Lett. B **174** 45 (1986); W. Buchmüller and M. Plümacher, Phys. Rep. **320**, 329 (1999) and Int. J. Mod. Phys. A **15**, 5047 (2000).

[22] V. A. Kuzmin, V. A. Rubakov and M. E. Shaposhnikov, Phys. Lett. **155B**, 36 (1985); J. A. Harvey and M. S. Turner, Phys. Rev. D **42**, 3344 (1990).

[23] R. Foot and R. R. Volkas, Phys. Rev. Lett. **75**, 4350 (1995).

[24] R. Foot, M. Thomson and R. R. Volkas, Phys. Rev. D **53**, 5349 (1996); R. Foot and R. R. Volkas, Phys. Rev. D **55**, 5147 (1997); X. Shi and G. M. Fuller, Phys. Rev. D **59**, 063006 (1999); P. Di Bari, P. Lipari and M. Lusignoli, Int. J. Mod. Phys. A **15**, 2289 (2000).

[25] D. Notzold and G. Raffelt, Nucl. Phys. B **307**, 924 (1988); K. Enqvist, K. Kainulainen and J. Maalampi, Phys. Lett. B **249**, 531 (1990); Nucl. Phys. B **349**, 754 (1991).

[26] S. Burles, K. M. Nollet, J. W. Truran and M. S. Turner, Phys. Rev. Lett. **82**, 4176 (1999); K. A. Olive, G. Steigman and T. P. Walker, Phys. Rep. **333-334**, 389 (2000).

[27] L. Bento, in preparation.

[28] M. Hirsch *et al.*, Phys. Lett. B **372**, 8 (1996).

[29] T. Bernatowicz *et al.*, Phys. Rev. Lett. B **69**, 2341 (1992).

[30] E. W. Kolb, D. L. Tubbs, D. A. Dicus, Astrophys. J. (Letters) **255**, L57, (1982); G. M. Fuller, R. Mayler, and J. R. Wilson, Astrophys. J. **332**, 826 (1988).

[31] Z. G. Berezhiani and A. Y. Smirnov, Phys. Lett. B **220**, 279 (1989).

[32] K. Choi and A. Santamaria, Phys. Rev. D **42**, 293 (1990).

[33] M. Kachelriess, R. Tomàs and J. W. F. Valle, Phys. Rev. D **62**, 023004 (2000).