Supernova neutrino oscillations:
Adiabaticity improvement by Majoron fields

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Abstract

If the lepton numbers are associated with global symmetries spontaneously broken at a scale below 1 TeV, neutrino oscillations in supernovae produce classic Majoron fields that perturb the neutrino propagation itself and may change the oscillation patterns in the periods of largest $\nu$ fluxes. The impact of the Majoron fields on the same transitions as $\nu_e \rightarrow \nu_X$ that presumably occur in the Sun is studied in the case of the non-adiabatic MSW solution. It is shown how the back reaction of the Majoron fields may improve the adiabaticity of these oscillations in a supernova environment which has implications on the outgoing $\nu_e$ spectrum.

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I. INTRODUCTION

Several experiments from solar to atmospheric neutrinos and laboratory oscillation experiments indicate that neutrinos oscillate and leptonic flavors are not conserved. This is in contrast with the minimal standard model (SM). It is quite possible that the energy scale of breaking of the three lepton numbers is comparable or even smaller than the Fermi scale. Furthermore, if they are spontaneously broken by the expectation values of some scalar fields then, a few bosons with zero mass should exist - the Nambu-Goldstone bosons (NG) - one per broken global symmetry.

It was recently pointed out that these NG bosons (called Majorons or familons when associated with lepton numbers) couple to the time rate of creation of the respective lepton numbers carried by the matter particles and therefore coherent NG fields are produced whenever lepton number violating processes occur simultaneously. That is the case if neutrinos change flavor on their way out from stars as seems to happen in the solar system. Once the NG fields are generated (the triggering process may be a normal Mikheyev-Smirnov-Wolfenstein (MSW) resonant conversion), they change in turn the relative potentials of the different neutrino species and so get a life of their own.

The numbers show that if the scale of symmetry breaking is below 1 TeV then, the Majoron fields are important enough to play a role in supernova neutrino oscillations. They are however too small in the case of the Sun unless the scale of symmetry breaking lies below 1 KeV. As a result, supernova neutrinos may exhibit oscillation patterns in contradiction with the observations of solar, atmospheric and terrestrial neutrinos. We ought to be prepared, in the event of a close by supernova explosion, for the possible kind of effects caused by NG fields.

In the previous paper, the example studied was that Majoron field generated by the conversion $\nu_e \rightarrow \nu_\tau$, assumed to take place in a certain resonance shell, yield neutrino potentials which become competitive with the standard electroweak potentials at larger radii and therefore affect the other flavor transitions characterized by smaller $\Delta m^2$. The effects can be so dramatic as the resonant oscillation $\bar{\nu}_e \leftrightarrow \bar{\nu}_\mu$ in a context of $m_{\nu_e} < m_{\nu_\mu}$ hierarchy, where the resonance is otherwise possible for $\nu_e \leftrightarrow \nu_\mu$ but not for the anti-neutrinos, if they only interact via standard $W$ and $Z$ bosons. In the present paper, I want to discuss what happens if the Majoron potentials are already significant in the very region where the oscillations they are generated from occur. Then, a back reaction effect takes place yielding an interesting flavor dynamics.

Suppose that in the Sun the electron-neutrino oscillates into the muon-neutrino with the parameters of the non-adiabatic, small mixing angle solution. Then, the $\nu_e \leftrightarrow \nu_\mu$ transitions are also non-adiabatic in a supernova and only a fraction of each neutrino species is converted into the other. Furthermore, since the level crossing probability is an increasing function of the energy, the hotter $\nu_\mu$s have larger survival probabilities than the cooler $\nu_e$s. It will be shown that the back reaction of the NG fields improves the adiabaticity of the neutrino transitions, thus yielding a hotter energy spectrum for the outgoing $\nu_e$s. That kind of effect can in principle be traced in those detectors such as Super-Kamiokande and SNO, capable of detecting supernova electron-neutrinos.
In the following it is assumed that the partial lepton number $L_e$ is conserved at the Lagrangian level but the global symmetry associated with it, $U(1)_{L_e}$, is spontaneously broken by the expectation values of one or more scalar iso-singlets $\sigma_i$. Then, a NG boson $\xi_e$ exists with zero mass. The neutrino mass matrix violates in principle the three lepton numbers but for simplicity I will ignore the other possibly existing Majorons. It may be interpreted as meaning that the respective scales of symmetry breaking are slightly higher.

It is well established that the Nambu-Goldstone bosons only interact through derivative couplings [7,8] (related to the soft pion low energy theorems). This is clarified [9,2] by changing variables from the original fields with definite $L_e$ charges namely, fermions $\chi^a$ and scalars $\sigma_i$, as follows:

$$\chi^a = \exp(-i\xi_e L_e^a) \psi^a ,$$  \hspace{1cm} (1)

$$\sigma_i = \exp(-i\xi_e L_e^i) (\langle\sigma_i\rangle + \rho_i) .$$  \hspace{1cm} (2)

The so defined physical weak eigenstates, the fermions $\psi^a = e, \nu_e, ...$ and massive bosons $\rho_i$, are invariant under the group $U(1)_{L_e}$. In terms of these fields the symmetry is simply realized as translations of the Majoron field, $\xi_e \rightarrow \xi_e + \alpha$, and consequently the Lagrangian can only depend on the derivatives $\partial_\mu \xi_e$.

The non-standard interactions relevant to this work are contained in the expression

$$\mathcal{L} = -\nu_L^a T C m_{ab} \nu_L^b + \mathrm{h.c.} + \frac{1}{2} \Omega_{e}^{2} \partial_\mu \xi_e \partial^\mu \xi_e + \mathcal{L}_{\mathrm{int}}(\partial_\mu \xi_e) ,$$  \hspace{1cm} (3)

where $\Omega_{e}^{2} = 2 \sum_i | L_i \langle\sigma_i\rangle |^2$ and $\nu_L^a$ denote the standard $\nu_e^L, \nu_{\mu}^L, \nu_{\tau}^L$ or any extra neutrino singlets. The $\xi_e$ equation of motion,

$$\partial_\mu \partial^\mu \xi_e = -\partial_\mu J_\mu^e / \Omega_{e}^{2} ,$$  \hspace{1cm} (4)

identifies with the conservation law of the Nöther current associated with the symmetry $\xi_e \rightarrow \xi_e + \alpha$. All the Majoron interactions are cast in the current $J_\mu^e$ obtained from $\mathcal{L}_{\mathrm{int}}(\partial_\mu \xi_e)$. Its leading terms do not depend on the particular model and are derived from the $\chi^a, \sigma_i$ kinetic Lagrangian using Eqs. (1), (2). The result is the following:

$$\mathcal{L}_{\mathrm{int}} = (\partial_\mu \xi_e) (\bar{\nu}_e \gamma^\mu e + \bar{\nu}_e \gamma^\mu \nu_e + \cdots) ,$$  \hspace{1cm} (5a)

$$J_\mu^e = \bar{e} \gamma^\mu e + \bar{\nu}_e \gamma^\mu \nu_e + \cdots .$$  \hspace{1cm} (5b)

The dots represent scalar boson terms and model dependent radiative corrections that are not relevant for this work and will ignored in the following.

The expressions above are typical of one kind of models, the Abelian-singlet Majorons [10-12]. Singlet means that the scalar fields that spontaneously break the lepton number symmetries are all singlets under the SM gauge group $SU(2) \times U(1)$, thus complying with
the LEP results on the $Z^0$ invisible width, unlike the triplet-Majoron model \cite{13}. By Abelian I mean U(1) symmetry groups, not horizontal in flavor space, associated to $L_e$, $L_\mu$, $L_\tau$ or any linear combinations of them. As the respective currents do not change flavor these models are not bound by the laboratory limits on the familon models \cite{14}. In addition, of all interactions involving SM particles the neutrino masses are the most important source of lepton numbers violation. Thus, in single collision or decay reactions, the effective strength of the Majoron couplings, resulting from $\partial_\mu J_\mu^\nu / \langle \sigma \rangle$, is proportional to the neutrino masses and suppressed by the symmetry breaking scale: $g \sim m_\nu / \langle \sigma \rangle$. For that reason, one does not expect observable zero neutrino double beta decays accompanied by Majoron emission as shown in \cite{15}. With a sensitivity to neutrino-Majoron couplings of the order of $10^{-5}$ \cite{16} they cannot even probe relatively low symmetry breaking scales. This kind of models also escape present astrophysical bounds on the couplings of Nambu-Goldstone bosons. Neutrino-Majoron couplings with strengths $g \sim m_\nu / \langle \sigma \rangle$ are too far from $10^{-4}$ to change supernova collapse dynamics \cite{17}, and even below the threshold of $\sim 10^{-8.5}$ for supernova cooling through singlet Majoron emission \cite{18}. Finally, the pseudo-scalar couplings to electrons, that could be responsible for energy loss in stars \cite{19}, only arise through radiative corrections and are so further suppressed.

In the $\xi_e$ equation of motion, the source term $\partial_\mu J_\mu^e$ is nothing but the time rate of creation of $L_e$-number carried by matter particles per unity of volume. If the neutrinos $\nu_e$ and $\nu_\mu$ oscillate into each other outside the supernova neutrinospheres, but not their anti-particles, the net variation of $L_e$ is given by the difference between the numbers of converted neutrinos $N(\nu_\mu \rightarrow \nu_e)$ and $N(\nu_e \rightarrow \nu_\mu)$. In a stationary regime the fluxes are constant in time and Eq. \eqref{4} reduces to a Poisson equation with a Coulombian solution for $\xi_e$ \cite{2}. The gradient $\vec{A}_e = -\vec{\nabla} \xi_e$ obeys a Gauss law. In a spherical symmetric configuration it only has a radial component,

$$A_e(r) = -\frac{1}{\Omega_\xi_e} \frac{\hat{L}_e(r)}{4\pi r^2}, \quad \text{(6)}$$

that is determined by the integral of the source term over the volume of radius $r$, in this case, $L_e$-number created per unity of time in that volume, $\hat{L}_e(r) = \int d^3 x \partial_\mu J_\mu^e$. It can also be expressed as \cite{3}

$$A_e(r) = -\frac{1}{\Omega_\xi_e} \left[ j(\nu_\mu \rightarrow \nu_e) - j(\nu_e \rightarrow \nu_\mu) \right], \quad \text{(7)}$$

where $j(\nu_e \rightarrow \nu_\mu)$ denotes the flux of $e$-neutrinos converted to $\mu$ flavor and $j(\nu_\mu \rightarrow \nu_e)$ the reciprocal. Both these quantities are functions of the radius $r$.

Taking in consideration the interactions specified by Eqs. \eqref{3}, \eqref{5a}, the equation of motion of the neutrino wave function is, including flavor space,

$$\left( i \partial - \gamma^0 V_0 - \vec{\gamma} \cdot \vec{A}_e \hat{L}_e \right) \psi_L = m \psi_R, \quad \text{(8)}$$

where $m$ is the $\nu$ mass matrix (real for simplicity), $\hat{L}_e$ is the flavor-valued quantum number (1 for $\nu_e$ and 0 for $\nu_\mu, \nu_\tau$) and $V_0$ designates the flavor conserving SM potential in a medium at rest. One derives in the same way as in the case of a scalar potential $V_0$ the equations governing flavor oscillations \cite{20} (see also \cite{3}),
\[ i \frac{\partial}{\partial r} \nu = \left( \frac{m^2}{2E} + V_0 + A_e \frac{d}{dr} \right) \nu , \]  
which give (after absorbing a flavor universal term in the wave function)

\[ i \frac{\partial}{\partial r} \left( \begin{array}{c} \nu_e \\ \nu_\mu \end{array} \right) = \frac{1}{2E} \left( \begin{array}{cc} 2E(V_W + A_e) & \frac{1}{2} \Delta m^2 \sin 2\theta \\ \frac{1}{2} \Delta m^2 \sin 2\theta & \Delta m^2 \cos 2\theta \end{array} \right) \left( \begin{array}{c} \nu_e \\ \nu_\mu \end{array} \right) . \]  

\[ V_W = \sqrt{2} G_F n_e \]  
and \( \theta \) is the mixing angle. It is worth to notice that the results do not change if one considers alternatively a NG boson associated with breaking of \( L_\mu \) or \( L_e - L_\mu \). The reason is, \( \nu_e \leftrightarrow \nu_\mu \) oscillations only violate \( L_e - L_\mu \) (the non-conservation of \( L_e + L_\mu \) is suppressed by \( m_\nu^2/E^2 \)) and only care about the difference between the \( \nu_e \) and \( \nu_\mu \) potentials.

In order to calculate \( A_e(r) \) one needs to know the fluxes of \( \nu_\mu \) and \( \nu_e \) as functions of the radius with and without oscillations. I make the simplification of neglecting the oscillation length by saying that a neutrino with energy \( E_R \) is converted to the other flavor at the position where the resonance condition,

\[ E_R = \frac{\Delta m^2 \cos 2\theta}{2(V_W + A_e)} , \]  
is fulfilled. In addition, having in mind that in a non-adiabatic regime only a fraction \( 1 - P_c \) changes flavor (small mixing angle), the level crossing probability \( P_c \) is calculated using the Landau-Zener approximation [11,20]:

\[ P_c(E) = \exp \left\{ -\frac{\pi}{4} \left| \frac{\Delta m^2}{dE_R/dr} \sin^2 2\theta \right| \right\} . \]  

I believe that these approximations change the numbers but not the lesson taken by comparing the results with and without a Majoron field.

Let the number of emitted particles per unity of time and energy be specified by the distribution functions \( f_{\nu_e}(E) = d\hat{N}_{\nu_e}/dE \) and \( f_{\nu_\mu}(E) = d\hat{N}_{\nu_\mu}/dE \). They are normalized by the relation between the number luminosity \( \hat{N}_\nu \), the energy luminosity \( L_\nu \) and the average energy \( \bar{E}_\nu \) of each \( \nu \) flavor namely, \( \hat{N}_\nu = L_\nu/\bar{E}_\nu \). \( L_\nu \) and \( \hat{N}_\nu \) will be given below in units of ergs/s and ergs/s/MeV respectively. Equation (11) gives the resonance energy as a function of the radius: particles with lower energies reach the resonance position at higher density regions and the hottest neutrinos oscillate at the largest radii. The statement that the number of converted neutrinos is the fraction \( 1 - P_c \) of the number of particles with resonance energy translates into an equation for the time rate \( d\hat{L}_e = d\hat{N}(\nu_\mu \to \nu_e) - d\hat{N}(\nu_e \to \nu_\mu) \) of \( L_e \)-number creation in a shell with depth \( dr \):

\[ d\hat{L}_e = (1 - P_c) \left[ f_{\nu_\mu}(E_R) - f_{\nu_e}(E_R) \right] \frac{dE_R}{dr} dr . \]  

This establishes a differential equation for \( \hat{L}_e(r) \). Notice that the derivative of \( E_R \) is not independent from the derivative of \( \hat{L}_e \) because the Majoron potential \( A_e(r) \) that enters in
FIG. 1. Level crossing probability $P_c$ as a function of the $\nu$ energy. Dashed curve holds for the SM with constants $\bar{M} = 4 \times 10^{31}$ g and $Y_e = 1/2$. In the dotted and bold curves the Nambu-Goldstone field $\xi_e$ exists with $G_e = G_F$ and $G_e = 4G_F$ respectively, and the luminosities are $10^{52}$ ergs/s for $\nu_e$ and $7 \times 10^{51}$ ergs/s for $\nu_\mu$.

$E_R$ depends also on $\dot{L}_e(r)$, as Eq. (3) shows. In addition, the probability $P_c$ depends on the $E_R$ and $\dot{L}_e$ derivatives as well, and that makes a non-linear differential equation for $\dot{L}_e(r)$ specified by Eqs. (6), (11) - (13).

It remains to tell the density profile of the medium. In the regions of a supernova star with densities typical of the Sun the mass density goes as $1/r^3$, the constant $\bar{M} = \rho r^3$ lying between $10^{31}$ g and $15 \times 10^{31}$ g depending on the star [22]. Then, in terms of the electron abundance $Y_e \approx 1/2$, the electroweak potential ($\sqrt{2}G_F n_e$) is

$$V_W = 0.76 \ Y_e \frac{\bar{M}}{10^{31} \text{g}} \ r_{10}^{-3} \times 10^{-12} \text{eV},$$

where $r_{10} = r/10^{10}$ cm. This is to be compared with

$$A_e = 1.48 \ \frac{G_e}{G_F} \frac{\bar{L}_e}{10^{51} \text{ergs/s}/\text{MeV}} \ r_{10}^{-2} \times 10^{-12} \text{eV},$$

where $G_F = 11.66$ TeV$^{-2}$ is the Fermi constant and $G_e = 1/\Omega_e^2$. Clearly, if the neutrino luminosities are sufficiently high say, $10^{52}$ ergs/s for 10 MeV neutrinos, and the scale of lepton symmetry breaking is around or below the Fermi scale, the Majoron potential $A_e$ becomes competitive with $V_W$ at radii where the resonance occurs for $\Delta m^2$ values interesting for solar neutrinos ($10^{-5} - 10^{-4}$ eV$^2$). More generally, at large enough distances the Majoron potentials decay as the inverse square radius and necessarily dominate over the local interactions.
Let us examine the $\nu_e \leftrightarrow \nu_\mu$ oscillations with the mixing parameters of the non-adiabatic solar neutrino solution (for a recent update see [23]), choosing in particular the values $\Delta m^2 = 7 \times 10^{-6}$ eV$^2$, $\Delta m^2 \sin^2 2\theta = 4 \times 10^{-8}$ eV$^2$. In a supernova the resonance is non-adiabatic as well and, as Eq. (12) indicates, the survival probability $P_c$ increases with the $\nu$ energy. This was studied in detail in the framework of the SM [24]. In Fig. 1, $P_c$ in the Landau-Zener approximation is plotted against the energy. The dashed curve holds for the SM potential with a constant $\tilde{M} = 4 \times 10^{31}$ g. It is manifest the aggravation of the non-adiabaticity with the energy.

To study the Majoron case one has to specify the energy spectra and luminosities. I used Fermi-Dirac distributions with the following values of temperature and chemical potential [30]: for $\nu_e$, $T = 2.4$ MeV and $\mu = 3.2 T$; for $\nu_\mu$, $T = 5.1$ MeV, $\mu = 4.1 T$. This gives average energies of 10 and 23 MeV respectively. The luminosity intensities are in turn $10^{52}$ ergs/s for $\nu_e$ and $7 \times 10^{51}$ ergs/s for $\nu_\mu$ which amount to particle emission rates of $10^{51}$ and $3 \times 10^{50}$ ergs/s/MeV respectively. Because the $e$-neutrinos are more numerous, the $\nu_e \leftrightarrow \nu_\mu$ oscillations produce a net destruction of $L_e$-number and a positive Majoron potential $A_e$. The dynamics is the following: the less energetic $\nu_e$ oscillate to $\nu_\mu$ at the smallest radii; this conversion produces a positive $A_e(r)$ which attenuates the fall of the total potential $V_e = V_W + A_e$ with the radius; as a consequence, the adiabaticity improves at larger $r$ and the most energetic neutrinos change flavor with higher probabilities.

Figure 2 shows the total potential $V_e$ as a function of the radius. The dashed curve stands for $V_W$, the SM potential, with $\tilde{M} = 4 \times 10^{31}$ g and $Y_e = 1/2$. In the dotted and bold curves the Nambu-Goldstone field associated to $L_e$ symmetry breaking operates with a constant $G_e = G_F$ and $G_e = 4 G_F$, respectively. The potential falls more slowly as the field $\xi_e$ grows. The level crossing probability is plotted in Fig. 1 for both cases (dotted and
FIG. 3. The dotted curves are the assumed luminosity distributions for $\nu_e$ and $\nu_\mu$ as emitted from the neutrinospheres. The dashed curve represents the $\nu_e$ luminosity after electroweak neutrino oscillations whereas in the bold curve a $\xi_e$ field exists with $G_e = 4G_F$.

The improvement of adiabaticity makes more $\nu_\mu$s to convert into $\nu_e$ and less $\nu_e$ to survive, and because the $\mu$-neutrinos are more energetic, the outgoing $\nu_e$ spectrum is harder than if there was no Majoron field. The average energy of the outgoing $e$-neutrinos is 13 MeV if $G_e = 0$ but rises to 17 MeV if $G_e = G_F$ and 21 MeV if $G_e = 4G_F$.

### III. CONCLUSIONS AND DISCUSSION

To summarize, if the explanation of the solar neutrino deficit is the MSW non-adiabatic oscillation $\nu_e \rightarrow \nu_\mu$ (or $\nu_e \rightarrow \nu_\tau$) then, the standard model of electroweak interactions predicts that in a supernova the $\nu_e \leftrightarrow \nu_\mu$ transitions are also non-adiabatic. It means that, to a large extent, the $e$-neutrinos preserve their lower energy spectrum, unless $\nu_e$ also mixes to another flavor with a too high or too low $\Delta m^2$ to show up in solar neutrinos. If
however, $L_e$ is a spontaneously broken quantum number, the associated Nambu-Goldstone boson, $\xi_e$, will acquire a classic field configuration which may be strong enough to produce a back reaction with the net effect of improving the adiabaticity of the $\nu_e \leftrightarrow \nu_\mu$ transitions. The final result is a $\nu_e$ energy spectrum harder than expected.

In 1987, the existing detectors were only able to detect electron anti-neutrinos but the now operating Super-Kamiokande and SNO experiments will be capable of detecting supernova $\nu_e$ events. The analysis of the energy distribution can in principle reveal or put limits on that kind of effect.

The scenarios of neutrino mixing change considerably if one considers the evidences from atmospheric and terrestrial neutrino experiments (for a review see [1]). The atmospheric neutrino anomaly and the zenith angle dependence observed by Super-Kamiokande can be explained by $\nu_\mu \to \nu_\tau$ oscillations, best fit $^{[2]} \Delta m^2 = 5 \times 10^{-3} \text{eV}^2, \sin^2 2\theta = 1$. The alternative $\nu_\mu \to \nu_e$ is excluded by the CHOOZ limits $^{[26]}$. This can still accommodate $\nu_e \to \nu_\mu$ or $\nu_e \to \nu_\tau$ as solar neutrino solutions but that is no longer true if one takes in consideration the evidence from the Liquid Scintillation Neutrino Detector (LSND) for $\bar{\nu}_\mu \to \bar{\nu}_e$ $^{[27]}$ and $\nu_\mu \to \nu_\tau$ $^{[28]}$ oscillations. The very different $\Delta m^2$ scales involved in LSND ($\Delta m^2_{e\mu} > 0.2 \text{eV}^2$), atmospheric and solar neutrinos call for a fourth flavor - a sterile neutrino. In that picture solar $\nu_s$ oscillate into the sterile $\nu_s$.

We now examine the consequences for supernovae neutrinos always assuming the non-adiabatic solar neutrino solution, ignoring for definiteness the possible mixing between $\nu_s$ and $\nu_\mu$ or $\nu_\tau$. The oscillation pattern is the following: 1) MSW conversion of $\nu_e \to \nu_\mu$ with LSND $\Delta m^2$, 2) sequential oscillation $\nu_\mu \to \nu_e \to \nu_s$, the first a LSND transition, the second a solar $\nu$ process. The outcome is a hard spectrum for $\nu_e$ depleted by $\nu_e \to \nu_s$, but only partially because of the non-adiabaticity of this transition. If alternatively, a NG field $\xi_e$ exists (created by $\nu_e \to \nu_\mu, \nu_\mu \to \nu_e$ and $\nu_e \to \nu_s$), it improves the adiabaticity of $\nu_e \to \nu_s$ causing a $\nu_e$ depletion stronger than predicted by SM interactions.

If one repeats the analysis with other scenarios of $\nu$ mixing the effects will be different in detail but with one thing in common: the signature of NG fields is a surprise i.e., an oscillation pattern not consistent with the $\nu$ mixing derived from solar, atmospheric and terrestrial $\nu$ experiments. It should be kept in mind that the situation turns more complex and rich if there is mixing between the three NG bosons associated with the three lepton flavors, a very natural feature if they are all spontaneously broken. This was explored in $^{[2]}$.

A point that cannot be overstressed is that the Nambu-Goldstone fields are proportional to the rate of charge violation processes and therefore to the very reaction rates. In the case of neutrino oscillation this manifests as a strong dependence of the Majoron fields on the magnitudes of the neutrino luminosities. In the numeric simulation I chose $10^{32}$ ergs/s for $\nu_e$ and $7 \times 10^{51}$ ergs/s for $\nu_\mu$, values produced and even exceeded in the about half a second that lasts between the neutronisation $\nu_e$ burst and the supernova explosion $^{[8,29,30]}$. The neutrino luminosities decay afterwards in a time scale of 1 second, or rather 4 seconds $^{[30,31]}$, as indicated by SN 1987A events $^{[32]}$. The highest luminosity happens during the first $\nu_e$ burst - above $10^{53}$ ergs/s in the peak $^{[8,29,30]}$ - and the Majoron field may be even stronger then. However, the time scale of the rise and fall of the $\nu_e$ signal is about 5 ms, too short to authorize a stationary approximation in the calculation of $\xi_e$. In fact, distances of the order of $10^{10}$ cm in such a period of time are
beyond the light cone and a special study is required.

The effects of the Majoron fields on the neutrino spectra, if any, will be observed in a shorter or longer interval of time depending on the actual scale of lepton number symmetry breaking. The observation of such a correlation with the flux magnitudes, by itself a signature of the NG fields, would thus provide a measurement of the scale of spontaneous symmetry breaking.

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REFERENCES

[1] D. O. Caldwell, in Proceedings of "COSMO 97: First International Workshop on Particle Physics and the Early Universe", Ambleside, England, 1997 (edited by L. Roszkowski, World Scientific, 1998), hep-ph/9804367. Y. Totsuka, in Proceedings of the "XVIII International Symposium on Lepton-Photon Interactions", Hamburg, 1997 (edited by A. De Roeck and A. Wagner, World Scientific, 1998).

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

[3] S. P. Mikheyev and Y. Yu. Smirnov, Yad. Fiz. 42, 1441 (1985) [Sov. J. Nucl. Phys. 42, 913 (1985)] and Nuovo Cimento 9C, 17 (1986).

[4] W. C. Haxton, Phys. Rev. Lett. 57, 1271 (1986); S. J. Parke, ibid, 1275 (1986).

[5] S. P. Rosen and J. M. Gelb, Phys. Rev. D 34, 969 (1986).

[6] A. Burrows, D. Klein, and R. Gandhi, Phys. Rev. D 45, 3361 (1992).

[7] T.-P. Cheng and L.-F. Li, "Gauge Theory of Elementary Particle Physics" (Oxford University Press, New York, 1989).

[8] S. Weinberg, "The Quantum Theory of Fields: II" (Cambridge University Press, Cambridge, England, 1996).

[9] G. B. Gelmini, S. Nussinov and T. Yanagida, Nucl. Phys. B 219, 31 (1983).

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

[11] J. Schechter and J. W. F. Valle, Phys. Rev. D 25, 774 (1982).

[12] B. Grinstein, J. Preskill, and M. B. Wise, Phys. Lett. 159 B, 57 (1985).

[13] G. B. Gelmini and M. Roncadelli, Phys. Lett. 99B, 411 (1981).

[14] Particle Data Group, R. M. Barnett et. al., Phys. Rev. D 54, 1 (1996).

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

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

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

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

[19] G. C. Raffelt, in Proceedings of "Beyond the Desert 1997", Tegernsee, Germany, 1997 (edited by H. V. Klapdor-Kleingrothaus, and H. Paes, IOP Publishing, 1998) (and references therein), astro-ph/9707268. M. Fukugita, S. Wata mura, and M. Yoshimura, Phys. Rev. Lett. 48, 1522 (1982).

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

[21] L. Wolfenstein, Phys. Rev. D 17, 2369 (1978).

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

[23] N. Hata and P. Langacker, Phys. Rev. D 56, 6107 (1997).

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

[25] Super-Kamiokande Collaboration, A. Suzuki, talk given at the "International Symposium on Lepton and Baryon Number Violation", Trento, 1998.

[26] M. Appollonio et al., Phys. Lett. B. 420, 397 (1998).

[27] LSND Collaboration, C. Athanassopoulos et al., Phys. Rev C 54, 2685 (1996); Phys. Rev. Lett. 77, 3082 (1996).

[28] LSND Collaboration, C. Athanassopoulos et al., Phys. Rev C 58, 2489 (1998); Phys. Rev. Lett. 81, 1774 (1998).

[29] R. Mayle, J. R. Wilson, and D. N. Schramm, Astrophys. J. 318, 288 (1987).

[30] A. Burrows, Annu. Rev. Nucl. Sci. 40, 181 (1990).
[31] S. A. Bludman and P. J. Schinder, Astrophys. J. 326, 265 (1988).
[32] K. Hirata et. al., Phys. Rev. Lett. 58, 1490 (1987); R. M. Bionta, ibid, 1494.