# Avoiding BBN constraints on mirror models for sterile neutrinos

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We point out that in models that explain the Los Alamos neutrino scintillation detector result for neutrino oscillations using the mirror neutrinos, the big bang nucleosynthesis constraints on the number of sterile neutrinos can be avoided by using the late time phase transition that only helps to mix the active and the sterile neutrinos. The main idea is to have a standard model singlet scalar field that mixes the visible and the mirror sector and has a low vacuum expectation value ( $\sim 100$ keV) so that the active and sterile neutrinos remain unmixed at the big bang nucleosynthesis epoch. The model predicts different effective number of neutrinos at the epoch of recombination compared to the standard model and is therefore testable in experiments that measure cosmic microwave spectrum with a higher precision.

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## I. INTRODUCTION

The existence of neutrino oscillations has now been confirmed for solar and atmospheric neutrinos as well as for reactor and accelerator neutrinos. It is remarkable that all the data from many different experiments can be well understood in terms of only three neutrinos that mix among themselves. They imply very narrow ranges of both the mass difference squares as well as their respective mixings.

There is however another piece of evidence for oscillations, which if confirmed will require severe departure from the successful three neutrino scheme just mentioned. It is the apparent observation of the muon antineutrino oscillating to the electron type antineutrino in the Los Alamos neutrino scintillation detector (LSND) [1] experiment. An attempt was made to confirm this result by the KARMEN [2] collaboration, which eliminated a large fraction of the parameter space allowed by LSND. It is hoped that the Mini-BOONE experiment at FERMILAB, currently under way will settle the LSND anomaly in the near future [3].

If the LSND experiment is confirmed, one straightforward way to understand the results would be to postulate the existence of one or more extra neutrinos with mass in the eV range, the so-called sterile neutrinos (with no/negligible coupling to $W, Z$ bosons), which mix with the known active neutrinos. There are various models of sterile neutrinos, the details of which depend on the neutrino mass hierarchy. They are known in the literature as the $2 + 2$ [4], $3 + 1$ , as well as $3 + 2$ [6] models. Of the three, the $3 + 1$ model seems less disfavored than the $2 + 2$ by the null results of other oscillation experiments. However, the more recently proposed $3 + 2$ scenario [6] that involves two sterile neutrinos is apparently in better agreement with all data, more so than the other models.

The major challenge posed by the sterile neutrino for theory is to understand its ultra-lightness despite its being a standard model singlet. A class of particle physics models that successfully answers this challenge are the mirror matter models. The basic assumption of these models is that there is an identical copy of the standard model (both constituents and forces) in nature [7,8] that coexists with the familiar standard model matter and forces. The neutrinos of the mirror sector do not feel the standard model forces and can therefore be identified as sterile neutrinos [8,9] in order to explain the LSND anomaly. We will be focusing on the asymmetric mirror models of Ref. [9], which assumes that the weak scale in the visible sector is much smaller than that of the mirror sector. In such models the mirror neutrinos, like all mirror fermions, will be heavier than the corresponding particles in the visible sector.

We will assume that the two lightest mirror neutrinos have masses in the eV range and mix with the known active neutrinos in order to explain the LSND anomaly. The ultralightness of sterile neutrinos in these models is not a problem since the same mechanism that keeps the active neutrinos light will have its mirror analog and keep the mirror neutrinos light. This resolves a major conceptual difficulty with light sterile neutrinos. The mixing angles however cannot be predicted; so we adjust them to agree with experiment. More detailed models of this kind have been widely discussed [9,10]. Phenomenological and astrophysical constraints on these models have also been extensively discussed [11,12].

Sterile neutrino models for LSND face two cosmological hurdles that we would like to address in this paper. The issues are: how to make them consistent with (i) our understanding of big bang nucleosynthesis (BBN) and (ii) the recent bounds on neutrino masses from Wilkinson microwave anisotropy probe (WMAP) observations. The first problem is that BBN allows the number of neutrinos $N_\nu$, in equilibrium when the temperature of the Universe is 1 MeV, to be restricted by $He^4$ and $D_2$ observations to be very close to three [13].[1] On the other hand for $\nu_s$ mass in

[1] A recent reanalysis of the $He^4$ data [14] seems to allow a large range for $N_\nu$ i.e. $N_\nu \leq 4.44$ at 95% confidence level assuming $N_\nu > 3$. This is also problematic for the $3 + 2$ scenario.

 

the eV range and mixing in the few percent range required to explain the LSND data, rapid $\nu_e - \nu_s$ oscillations while the active neutrinos are still coupled to the primeval plasma can thermally populate sterile neutrinos leading to $N_\nu = 4$ for the 3 + 1 and 2 + 2 scenarios and $N_\nu = 5$ for the 3 + 2 scenario. Same $\nu_e - \nu_s$ oscillations after neutrino decoupling but before neutron to proton freeze-out can also affect $He^4$ predictions making the constraints on the number of extra sterile neutrinos much more tight.

The WMAP [15] constraints are on the sum of all neutrino masses in equilibrium at the epoch of structure formation which corresponds to a temperature around an eV. According to [16], $\sum m_\nu \leq 1.38$ eV for one sterile and $\sum m_\nu \leq 2.12$ eV for two extra ones assuming that they went into equilibrium at the BBN epoch. These constraints are also quite important since taken at face value, they would seem to rule out the 3 + 2 model for LSND.

It is therefore important to look for scenarios that may allow one to avoid both the above constraints while at the same time providing an explanation of the LSND experiment. Recently, it has been suggested [17] that late time phase transition can generate the masses and mixings of both the active and sterile neutrinos.[2] In these models one can avoid both these constraints. In Ref. [17], it is shown that this can be achieved by endowing a scalar field $\phi$ with vacuum expectation values (vevs) in the 100 keV range so that at the BBN time the sterile as well as the active neutrinos are massless. As a result there is no oscillation among them that can bring the sterile neutrinos into equilibrium. Since the sterile neutrinos decouple from Hubble expansion at very high temperatures, their abundance at the BBN epoch is suppressed leading to concordance with the BBN constraints. Cosmological signatures of generic models of this type have been given in Ref. [17].

In this paper we propose an alternative way to avoid the cosmological constraints using the same idea of late time phase transitions. We show that if the sterile neutrinos are the mirror neutrinos, we need only generate the mixing between the active and the sterile neutrino (and not masses) by the late time phase transition to avoid the BBN and WMAP constraints. An advantage of this model is that the contribution of the sterile neutrinos to the energy density of the universe at the BBN epoch is governed by a free parameter which is the ratio of the reheat temperature of the mirror sector to the visible sector (unlike the model of Ref. [17]). We further find a convenient realization of mirror model with the seesaw scale in the TeV range which implements this scenario. We find that we must employ the double seesaw mechanism [19] to get small neutrino masses, which in turn implies that there must be an extended gauge symmetry visible at the TeV scale. We consider the extended gauge symmetry as $SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$ and construct an explicit scenario with late phase transition and discuss its cosmological and astrophysical implications. The detailed field theoretical models for them can be worked out but we do not discuss them here.

To summarize, the two main ingredients in our work are (i) sterile neutrinos are mirror neutrinos; (ii) their mixing is induced in the post BBN regime.

An important cosmological consequence of the model is that even though there are sterile neutrinos in the early phase of the universe's evolution, the present universe has only active neutrinos since the sterile neutrinos decay or annihilate just before the recombination epoch. The cosmological imprint of their existence is a shifted value for the neutrino temperature which in principle could be detected by high precision study of the matter power spectrum of the universe.

## II. AN EXTENDED MIRROR MODEL

In the mirror model, one assumes that the universe consists not only of the observed standard model particles and forces but also coexisitng with it is an identical but different set of constituents experiencing analogous but different gauge forces. Gravity is common to all the particles. The forces are dictated by the gauge group $G \otimes G$ where one of the gauge groups $G$ acts in the standard model sector and on its fermions whereas the other acts in the other and on the mirror fermions [8]. Mirror symmetry keeps the gauge couplings equal but the effective strength of various forces in both sectors may be different due to different patterns of symmetry breaking. We assume that the weak scale in the mirror sector is about 20–30 times larger than that in the visible sector [9]. This is not a firm number but dictated by the fact that mirror neutrinos have masses in the eV range or higher whereas the heaviest active neutrino mass is known to be 0.05 eV from atmospheric neutrino data.

The masses of the known and mirror neutrinos are assumed to arise from the double seesaw mechanism in each sector. However, their mixing arises from a particle that can mix both sectors as we see below.

In what follows, we denote all particles and parameters of the mirror sector by a prime over the corresponding familiar sector symbol—e.g. mirror quarks are $u', d', s'$, , etc., and mirror Higgs field as $H'_{u,d}$, etc. We will assume that theory is supersymmetric so that the smallness of some of the couplings in the theory can be maintained naturally. Whether R-parity is exact or broken is of no consequence to the neutrino sector discussion; but for simplicity, we will keep R-parity exact.

We now give some details of the model we use. It is an extension of the standard model gauge group and is given by $SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$ (in each sector), which is anomaly free in the presence of the right handed neutrino

[2]It has been brought to the attention of the authors that the idea of late time phase transition as a solution to the BBN problem was first mentioned in [18].

$\nu^c$. All the fermions have obvious quantum numbers under the gauge group. We add a gauge singlet chiral fermion, $S$ in each sector, one per family. In order to get the standard model gauge group from the extended group in each sector, we need to add a pair of new Higgs bosons $\Delta(1, +\frac{1}{2}, -1)$ and a conjugate field $\bar{\Delta}(1, -\frac{1}{2}, +1)$ in the visible sector and two similar fields in the mirror sector. We then add a gauge singlet Higgs field $\chi$ ( and a mirror $\chi'$), whose vacuum expectation value gives Majorana mass for the singlet fermions $S$ by an interaction of the form $\lambda'_{\alpha\beta}(S_\alpha S_\beta \chi + S'_\alpha S'_\beta \chi')$. The particle content of the model is given in Table I.

The full superpotential relevant for neutrinos in each sector can be written as:

$$W = h_\nu L H_u \nu^c + \lambda_1 \nu^c \Delta S + \lambda' S S \chi; \qquad (1)$$

There is an identical set of terms for the superpotential in the mirror sector; we have suppressed the generation index. For the three generation case that we will be interested in, $h_\nu$, $\lambda_1$ and $\lambda'$ are $3 \times 3$ matrices.

The $U(1)_{I_{3R}} \times U(1)_{B-L}$ part of the gauge symmetry is broken by the vev of field $\langle\Delta\rangle$ assumed to be in the multi-TeV range. We choose the vev of the $\chi$ field to be in the GeV range. It is then clear that this leads to the double seesaw form [19] for the $(\nu, \nu^c, S)$ mass matrix:

$$M_\nu = \begin{pmatrix} 0 & h_\nu v & 0 \\ h_\nu^T v & 0 & \lambda_1 v_R \\ 0 & \lambda_1^T v_R & \lambda' \langle\chi\rangle \end{pmatrix}. \qquad (2)$$

There is also a similar matrix for the mirror neutrinos. This leads to the light neutrino mass matrix of the form:

$$\mathcal{M}_\nu = h_\nu M_R^{-1} \lambda' \langle\chi\rangle M_R^{-1T} h_\nu^T v^2 \qquad (3)$$

where $M_R = \lambda_1 \langle\Delta\rangle = \lambda_1 v_R$. It follows that if we choose $\langle\chi\rangle$ about a GeV, $h_\nu \sim 10^{-1}$ and $\lambda' \sim 10^{-4}$, then for $\frac{v}{M_R} \sim \frac{v'}{M'_R} \sim 10^{-2}$, the neutrino masses are in the 0.1 eV range as required by observations. This puts the vev of $\langle\Delta\rangle \simeq 10$ TeV and $\langle\Delta'\rangle \simeq 100$ TeV or so. Also, typical neutrino mass textures can be built into the coupling matrix $\lambda'$.

A similar situation will occur in the mirror sector, where we can choose $\langle\chi'\rangle$ about a factor of 10 higher to get $m_{\nu_s}$ in the eV range to fit LSND (henceforth, we will call the sterile neutrinos $\nu'$ as $\nu_s$).

TABLE I.

| Particles | Visible sector | Mirror sector |
|---|---|---|
| Gauge bosons | $W, Z_1, Z_2, \gamma$ | $W', Z'_1, Z'_2, \gamma'$ |
| Matter | $Q \equiv (u, d)$ | $Q' \equiv (u', d')$ |
| | $u^c, d^c, e^c, \nu^c$ | $u^{c'}, d^{c'}, e^{c'}, \nu^{c'}$ |
| Singlet matter | $S$ | $S'$ |
| MSSM Higgs fields | $H_u, H_d$ | $H'_u, H'_d$ |
| New Higgs fields | $\Delta, \bar{\Delta}, \chi$ | $\Delta', \bar{\Delta}', \chi'$ |
| Connector field | $\phi$ | $\phi$ |

The masses of the singlet fermions $\nu^c$, $\nu^{c'}$ and $S, S'$ are in the multi-TeV range and decouple from the standard model. This places a lower limit on the coupling $\lambda \sim 0.1$–$1$.

It is worth emphasizing at this point that the double seesaw form for the neutrinos could not have been implemented in the context of the standard model. Extra gauge degrees of freedom are needed. A natural and convenient choice for this gauge group is $SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$.

In order to generate mixing between the active and sterile neutrinos, we postulate the existence of a scalar field $\phi$ that connects the two sectors. This can only be done through an interaction of the form $\beta S S' \phi$ [20]. A simple tree level diagram via the exchange of $\nu^c, S$ and $\nu^{c\prime}, S'$ then leads to an effective coupling of the form: $\lambda'' \frac{L H_u L' H'_u \phi}{M_R M'_R}$, where $\lambda'' \simeq \beta h_\nu^2 \sim 10^{-2}$. We assume as in Ref. [17] that the vev of the field $\langle\phi\rangle \sim 100$ keV, so that at the BBN epoch the active and sterile neutrinos are unmixed.[3] In our model this is the only interaction that connects the visible and the mirror worlds. We will see below that the product of the scales $M_R M'_R$ is constrained by the cosmological requirement that the two worlds are not in thermal equilibrium with each other prior to the BBN epoch.

The effective superpotential for our theory below 1 MeV is:

$$\mathcal{W}_{\text{eff}} = g_{\nu\nu_s\phi} \nu \nu_s \phi + \kappa \phi^3 + m_\nu \nu\nu + m_{\nu_s} \nu_s \nu_s \qquad (4)$$

where $g_{\nu\nu_s\phi} \simeq \lambda'' \frac{v^2 \langle\phi\rangle}{M_R M'_R} \sim 10^{-6}$ for $\lambda'' \simeq 10^{-2}$. The resulting $\nu - \nu_s$ mixing is then given by: $m_{\nu-\nu_s} \simeq g_{\nu\nu_s\phi}\langle\phi\rangle \sim 10^{-6}\langle\phi\rangle$. This gives the right order of mixing for the LSND experiment. One difference between our model and that of Ref. [17] is that, late phase transition was used in [17] to generate both masses and mixings whereas in our case, only the mixing needs to be generated at a late stage. As we will see, the mirror model has the advantage that contribution of the sterile neutrinos to the energy density of the Universe at the BBN epoch is given by an arbitrary parameter. This is due to the fact that we assume asymmetric inflation which implies that the reheat temperature in the mirror sector ($T'_R$) of the universe is lower than that of the visible sector ($T_R$). The extra parameter here is the ratio $T'_R/T_R$. The model can therefore work even if the BBN constraints on the number of extra neutrino species tightened further.

[3]The smallness of the $\phi$ vev can be justified if we embed our theory into a brane bulk scenario and have the $\langle\phi\rangle$ vev occur in a distant brane and get transmitted to the our brane via a bulk scalar field.

## III. BBN, ASYMMETRIC INFLATION AND NEUTRINO MIXINGS

Before we address the issue of neutrino mixings and BBN, we note that in the mirror model, we have three light neutrinos, a mirror photon and a mirror electron that could be potential contributors to the energy density at the BBN epoch and affect the success of BBN. In order to reduce their contribution to a negligible level, the idea of asymmetric inflation [21] needs to be invoked as discussed in the context of the mirror model in the second paper of Ref. [9]. In this scenario, it is assumed that the reheat temperature after inflation in the mirror sector is lower than that in the visible sector by a factor of 10 or so i.e. $T'_R \simeq T_R/10$. $T'_R/T_R$ is a new parameter necessary for the consistency of the mirror model. If the interactions linking the two sectors are such that they are not in thermal contact for $T \leq T_R$, then Hubble expansion will roughly maintain the ratio of the two temperatures until the BBN epoch apart from minor corrections arising from particle annihilation in both sectors. Thus at the BBN epoch, the total contribution of the light mirror particles to $\rho_{\rm tot}$ is at the level of about $10^{-3}\rho_\gamma$ which therefore keeps the predictions of standard BBN unchanged. This will also keep the number densities of the light particles such as mirror neutrinos and mirror photons suppressed.

We now proceed to discuss the constraints on this model from cosmology. We consider successive epochs. In principle, there is an epoch for which $T \geq \langle\Delta\rangle, \langle\Delta'\rangle$ when the visible and the mirror sectors are in equilibrium. We will assume that this has to be above the reheat temperature in order for the light mirror particles not to populate the universe and spoil the successes of BBN. We therefore do not consider this separately.

The first constraints on the parameters of the model come from the following epoch. We assume the usual scenario of inflation followed by reheating. The constraint on reheating temperature comes from the fact that below the reheating temperature, the interaction rates for $L + H \rightarrow L' + H' + \phi$ should be out-of-equilibrium, otherwise, the two sectors will be in the same thermal bath and at BBN, the effective $N_\nu$ will far exceed the allowed limit due to the fact that the population of the sterile neutrinos will build up their density to the level of ordinary active neutrinos. The only exception is if the reheat temperature after inflation is below an MeV which we do not invoke here. The relevant epoch for this is when $T \geq \langle H\rangle, \langle H'\rangle$, which we consider now.

### A. $T \geq \langle H\rangle, \langle H'\rangle$

In this region, the condition for being out-of-equilibrium is

$$T \leq T_D \sim \left(g_*^{1/2}\frac{M_R^4}{M_{P\ell}}\right)^{1/3} \tag{5}$$

This inequality implies that for $T \gg$ TeV (note that the mirror Higgs mass is expected to be in the TeV range), the visible and the mirror sector in our model are in equilibrium. So to be consistent with BBN requirements, we must require that the reheat temperature after inflation be less than about a TeV or so. Once this condition is satisfied, the mirror and the visible worlds will be out of thermal contact until after the BBN epoch, since there are no other interactions that can connect the two worlds. It also follows that, below a TeV, there is no production mechanism for the connector field $\phi$ if we forbid its coupling to the inflaton field (which is assumed to couple to the Higgs fields of both sectors (see second paper in Ref. [9]). Therefore, the density of $\phi$ fields is negligible down to $T \sim 0.1$ MeV.

### B. 0.1 MeV $\leq T \leq \langle H\rangle, \langle H'\rangle$

In this regime, the effective interaction connecting the visible to the mirror sector is the coupling $\nu\nu_s\phi$ with a strength given by $g_{\nu\nu_s\phi} \sim \lambda''(\frac{v}{M_R}\frac{v'}{M'_R}) \approx 10^{-6}$. In discussing whether this interaction is in equilibrium, it has been noted in Ref. [17] that the process $\nu\nu_s \rightarrow \phi$, vanishes in the limit of $m_\phi = 0$ by energy momentum conservation. The rate for this process must therefore be proportional to $m_\phi^2(T)$. This can be seen as follows. Using the notation, $d\Phi_{123} = \frac{d^3p_\nu}{(2\pi)^3 2E_\nu}\frac{d^3p_{\nu_s}}{(2\pi)^3 2E_{\nu_s}}\frac{d^3p_\phi}{(2\pi)^3 2E_\phi}(2\pi)^4\delta^4(p_\nu + p_{\nu_s} - p_\phi)$, the thermally averaged cross section for the process $\nu + \nu_s \rightarrow \phi$ goes like [22],

$$\langle\sigma v\rangle \propto \frac{1}{T^6}\int d\Phi_{123}|M(\nu + \nu_s \rightarrow \phi)|^2 f_\nu^{\rm eq} f_{\nu_s}^{\rm eq}, \tag{6}$$

where $f_\nu^{\rm eq}$ and $f_{\nu_s}^{\rm eq}$ are the thermal phase space distribution (since $\langle\sigma v\rangle$ is the thermally averaged) of the particles $\nu$ and $\nu_s$ respectively. Using the energy conservation enforced by the $\delta$-function in $d\Phi_{123}$ we get

$$\Gamma(\nu + \nu_s \rightarrow \phi) \sim T^3\frac{1}{T^6}\int\frac{d^3p_\phi}{(2\pi)^3}f_\phi^{\rm eq}\int d\Phi_{23}|M(\nu + \nu_s \rightarrow \phi)|^2 \tag{7}$$

Since the process $\nu + \nu_s \rightarrow$ is $CP$ conserving we obtain

$$\Gamma(\nu + \nu_s \rightarrow \phi) \sim \Gamma_{\rm rest}(\phi \rightarrow \nu + \nu_s)\left\langle\frac{m_\phi}{E_\phi}\right\rangle \tag{8}$$

where $\Gamma_{\rm rest}(\phi \rightarrow \nu + \nu_s)$ is the decay width of $\phi$ at rest. For $T \gg m_\phi$ we get

$$\Gamma(\nu + \nu_s \rightarrow \phi) \sim \frac{g_{\nu\nu_s\phi}^2}{4\pi}\kappa^2 T \sim 10^{-21}T \tag{9}$$

as noted, where $\kappa$ is the self coupling of the connector field $\phi$ (defined by the superpotential $W_\phi \sim \kappa\phi^3$). If we choose the scalar self coupling $\kappa$ to be of order $10^{-4}$, and use the fact that the number density of $\nu_s$ is down by a factor of $10^{-3}$ relative to the $\nu$'s, we expect the rate

$\Gamma(\nu\nu_s \to \phi) \simeq \frac{\kappa T}{4\pi} \frac{g^2_{\nu\nu_s\phi}}{4\pi} \frac{n_{\nu_s}}{n_\nu} \sim 10^{-21}T$. This leads to the conclusion that above $T = T_\phi \sim 100$ keV, the interactions connecting the visible with the mirror sector are out of equilibrium.[4]

As has already been emphasized in Sec. III A, as long as the reheat temperature $T_R \leq 1$ TeV, above $T_\phi$, there is no thermal contact between the visible and the mirror world. Asymmetric inflation (i.e. $T'R \ll T_R$) then guarantees that the standard big bang cosmology receives only slight perturbation until $T \simeq T_\phi$ and our understanding big bang nucleosynthesis remains unchanged. When the two worlds come into thermal equilibrium after $T_\phi$, the perturbation to the standard big bang picture is very slight and as we show below, it manifests only in a slight shift of the cosmic neutrino background temperature.

### C. $T \leq 0.1$ MeV

This regime is below the scale of $\langle\phi\rangle$. Therefore, $\nu - \nu_s \to \phi$ is kinematically forbidden since Im$\phi$ becomes a pseudo-Goldstone boson. The only process that can lead to production of sterile neutrinos is $\nu\nu \to \nu_s\nu_s$. The rate for this process is given by $10^{-25}T$. This process is in equilibrium below $T \leq 100$ eV.

Below this temperature the mirror sector and the standard model neutrinos will thermalize without a significant transfer of energy. The thermalization can be seen by looking at the rate for $\Gamma(\nu + \nu_s \to \phi)$ in Eq. (7) which is bigger than the Hubble expansion rate $H = \sqrt{g_*}T^2/M_P$.

The second process which can contribute to thermalization of $\nu_s$ is $\nu + \nu \to \nu_s + \nu_s$, via virtual phi exchange. This interaction rate goes like $\Gamma(\nu + \nu \to \nu_s + \nu_s) \sim (g_{\text{eff}}/\kappa)^4 T^5/(T^2 + m_\phi^2)^2$ with $g_{\text{eff}} \sim 10^{-6}$ this interaction becomes faster than the Hubble expansion rate and can dominate over the first process if kappa is much smaller than $1/3$. In any case for the parameter range of our interest, the $\nu_s$ are in thermal equilibrium.

The next set of events depends on the mass of the field $\phi$. If $m_\phi < m_{\nu_s}$, then $\nu_s$ would be unstable. Since our interest is to explain LSND results, it is simpler to assume a stable $\nu_s$, which therefore leads us to assume that $m_\phi > m_{\nu_s}$ (which incidentally implies that $\kappa \geq 10^{-5}$ which is consistent with our assumption that $\kappa \sim 10^{-4}$). In this case, the field $\phi$ will first decay to $\nu + \nu_s$ and below $T \simeq m_{\nu_s}$, the sterile neutrinos will annihilate via the process $\nu_s\nu_s \to \nu\nu$ and disappear.

To calculate the final temperature of neutrino bath in terms of the photon temperature, we proceed as follows: we first remember that the boson $\phi$ is part of a supersymmetric multiplet whose scalar field part has only one surviving imaginary part and the real part has decoupled at very high temperature. We assume that the imaginary part is also superlight. Then we proceed through the following steps: Just below $T \simeq m_e$, electron-positron annihilation heats up the photons leading to the relation $T_\nu^0 = (\frac{4}{11})^{1/3}T_\gamma$. The $\nu_s$ and $\phi$ at this stage are not in thermal contact with the active neutrinos. Once the temperature of the universe cools below 100 eV, the $\nu - \nu_s\phi$ system comes into full thermal equilibrium. Using energy conservation [23] and taking the effect of the incomplete $\phi$ supermultiplet into account, we get (if $\phi$ is brought into thermal contact)

$$\left(3 + \frac{4}{7} + n'\right)T^4_{\nu+\phi+\nu_s} = 3T^4_{\nu_0} \tag{10}$$

where $n'$ is the number of sterile neutrinos and the factor $\frac{4}{7}$ takes into account the contribution of the bosonic part of singlet Higgs superfield $\phi$. If the field $\phi$ is heavier than the sterile neutrinos (which can be the case without fine tuning), as the universe cools below $m_\phi$, the $\phi$ decay to $\nu + \nu_s$. Using entropy conservation at this stage, we get

$$(3 + n')T^3_{\nu+\nu_s} = \left(3 + \frac{4}{7} + n'\right)T^3_{\nu+\phi+\nu_s} \tag{11}$$

Using the above two equations, we get for the temperature of the $\nu + \nu_s$ system

$$T^4_{\nu+\nu_s} = \frac{1}{(1 + \frac{n'}{3})^{4/3}}\left[1 + \frac{n' + 4/7}{3}\right]^{1/3}\left(\frac{4}{11}\right)^{4/3}T^4_\gamma \tag{12}$$

Noting that $\rho_\nu \propto (3 + n')T^4_{\nu+\nu_s}$, we find the effective number of neutrinos at matter radiation equality is

$$N_\nu = 3\left[1 + \frac{4/7}{3 + n'}\right]^{1/3} \tag{13}$$

As an example if $n' = 1$ the effective number of neutrinos visible to cosmic microwave background (CMB) experiments is $N_\nu = 3.14$ which is still consistent with present CMB data [16]. Future CMB experiments like PLANCK [24] and CMBpol [25] will be able to improve the limit on $N_\nu$ and can provide a test of this model. The contribution of the active neutrinos to the critical energy density for $n' = 1$ is

$$\Omega_\nu = \left(\frac{32}{21}\right)^{1/4}\frac{\sum m_\nu}{92h^2} \tag{14}$$

Using the upper recent bound on the neutrino energy density [26] one finds

$$\sum m_\nu < 0.41 \text{ eV} \tag{15}$$

Typical values of $N_\nu$ for this case are given in Table II.

Incidentally, the same steps can be repeated for theories without supersymmetry. In this case, we will assume that Re$\phi$ has a mass of order of the $\phi$-vev or about 100 keV,

[4]The suppression factor $n_{\nu_s}/n_\nu$ arises because since the Universe does not have any $\phi$ particles until after $T_\phi$, the process $\nu + \nu_s \to \phi$ can go only via collisions between $\nu$ and $\nu_s$.

TABLE II.

| $n'$ | $m_\phi > m_{\nu_s}$ |
|---|---|
| 1 | 3.14 |
| 2 | 3.11 |
| 3 | 3.10 |

whereas Im$\phi$ has a mass above an eV. The same steps as in the supersymmetric case can now be repeated and one gets a very small change in the upper limit on the sum of neutrino masses given in Eq. (11).

Our scenario for sterile neutrinos has also interesting astrophysical implications. The first point is to look for any new mechanism for energy loss from the supernova core via emission of $\nu_s$ or $\phi$. Since $\nu - \nu_s$ mixing arises from spontaneous symmetry breaking at scale $\ll$ MeV, inside hot astrophysical environments such as a supernova, the active and sterile neutrinos remain unmixed. As a result, there is no energy loss via the emission of $\nu_s$. However, there could be energy loss due to the processes $\nu\nu \rightarrow \phi\phi, \nu_s\nu_s$. The rates for these processes are estimated to be: $\sim 10^{-25}T \sim (50 \text{ sec})^{-1}$. Comparing this with typical supernova explosion time scale, we expect this energy loss mechanism not to be significant. Also due to zero mixing between $\nu - \nu_s$, all supernova results based on three active neutrinos [27] remain unaffected. Only in the very outer layers of the supernova explosion when the temperature drops below 100 keV, will these mixings become operative.

## IV. DISCUSSION AND SUMMARY

In summary, we have presented a mirror model for the sterile neutrinos that can explain the LSND results and yet be consistent with stringent constraints from big bang nucleosynthesis and cosmic microwave background as well as structure formation bounds on neutrino properties. We make predictions for the effective neutrino number to which the next generation CMB measurements are sensitive. An important requirement of this model is that the reheat temperature after inflation must be less than a TeV. The model has also other interesting properties discussed earlier such as the mirror hydrogen being a dark matter [28], which remain unaffected by our modification. Similarly, suggestions that the ultra high energy neutrinos could be originating from topological defects in the mirror sector [29] remain unchanged by our extension.

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[1] LSND Collaboration, Phys. Rev. Lett. **77**, 3082 (1996).
[2] B. Armbruster *et al.*, Phys. Rev. D **65**, 112001 (2002).
[3] BooNe Collaboration, A. O. Bazarko, hep-ex/9906003.
[4] D. Caldwell and R. N. Mohapatra, Phys. Rev. D **46**, 3259 (1993); J. Peltoniemi and J. W. F. Valle, Nucl. Phys. **B406**, 409 (1993); J. Peltoniemi, D. Tommasini, and J. W. F. Valle, Phys. Lett. B **298**, 383 (1993).
[5] S. Bilenky, W. Grimus, C. Giunti, and T. Schwetz, hep-ph/9904316; V. Barger, B. Kayser, J. Learned, T. Weiler, and K. Whisnant, Phys. Lett. B **489**, 345 (2000); for a review, see S. Bilenky, C. Giunti, and W. Grimus, Prog. Part. Nucl. Phys. **43**, 1 (1999).
[6] M. Sorel, J. Conrad, and M. Shavitz, Phys. Rev. D **70**, 073004 (2004).
[7] T. D. Lee and C. N. Yang, Phys. Rev. **104**, 254 (1956); K. Nishijima, private communication; Y. Kobzarev, L. Okun, and I. Ya Pomeranchuk, Yad. Fiz. **3**, 1154 (1966); M. Pavsic, Int. J. Theor. Phys. **9**, 229 (1974); S. I. Blinnikov and M. Y. Khlopov, Astron. Zh. **60**, 632 (1983).
[8] R. Foot, H. Lew, and R. R. Volkas, Phys. Lett. B **272**, 67 (1991); Mod. Phys. Lett. A **7**, 2567 (1992); R. Foot and R. Volkas, Phys. Rev. D **52**, 6595 (1995).
[9] Z. Berezhiani and R. N. Mohapatra, Phys. Rev. D **52**, 6607 (1995); Z. Berezhiani, A. Dolgov, and R. N. Mohapatra, Phys. Lett. B **375**, 26 (1996).
[10] B. Brahmachari and R. N. Mohapatra, Phys. Lett. B **437**, 100 (1998); K. S. Babu and R. N. Mohapatra, Phys. Lett. B **522**, 287 (2001); K. S. Babu and G. Seidl, Phys. Lett. B **591**, 127 (2004).
[11] M. Cirelli, G. Marandella, A. Strumia, and F. Vissani, Nucl. Phys. **B708**, 215 (2005).
[12] V. Berezinsky, M. Narayan, and F. Vissani, Nucl. Phys. **658**, 254 (2003).
[13] For a review and references, see Richard H. Cyburt, Brian D. Fields, and Keith A. Olive, Phys. Lett. B **567**, 227 (2003); A. Dolgov and F. Villante, Nucl. Phys. **B679**, 261 (2004); K. Abajazian, Astropart. Phys. **19**, 303 (2003).
[14] R. Cyburt, B. Fields, K. Olive, and E. Skillman, astro-ph/0408033.
[15] D. Spergel *et al.*, Astrophys. J. Suppl. Ser. **148**, 175 (2003).
[16] For a review, see S. Hannestand, hep-ph/0312122; J. Cosmol. Astropart. Phys. 05 (2003) 004.
[17] Z. Chacko, L. J. Hall, S. J. Oliver, and M. Perelstein, hep-ph/0405067 [Phys. Rev. Lett. (to be published)]; Z. Chacko, L. J. Hall, T. Okui, and S. J. Oliver, Phys. Rev. D **70**, 085008 (2004).
[18] R. Foot and R. R. Volkas, Phys. Rev. D **52**, 6595 (1995).
[19] R. N. Mohapatra and J. W. F. Valle, Phys. Rev. D **34**, 1634 (1986).
[20] M. Collie and R. Foot, Phys. Lett. B **432**, 134 (1998); B. Brahmachari and R. N. Mohapatra, in Ref.[10].
[21] E. W. Kolb, D. Seckel, and M. Turner, Nature (London) **314**, 415 (1985); H. Hodges, Phys. Rev. D **47**, 456 (1993).

[22] E. Kolb and M. Turner, *Early Universe*, (Addison-Wesley, Reading, MA, 1990), Chap. 6, Fig, 6.4, p. 167.
[23] K. Enquist *et al.*, Nucl. Phys. **B373**, 498 (1992); X. Shi, D. N. Schramm, and B. D. Fields, Phys. Rev. D **48**, 2563 (1993); K. Abajazian, Astropart. Phys. **19**, 303 (2003); A. D. Dolgov and F. L. Villante, Nucl. Phys. **B679**, 261 (2004).
[24] http://www.rssd.esa.int/index.php?project=PLANCK
[25] http://spacescience.nasa.gov/missions/concepts.htm
[26] U. Seljak, ICTP Summer School, 2004 (unpublished).
[27] See for instance A. S. Dighe and A. Y. Smirnov, Phys. Rev. D **62**, 033007 (2000); C. Lunardini and A. Y. Smirnov, Astropart. Phys. **21**, 703 (2004) and references therein.
[28] R. N. Mohapatra and V. L. Teplitz, Phys. Rev. D **62**, 063506 (2000); Z. Berezhiani, D. Comelli, and F. L. Villante, Phys. Lett. B **503**, 362 (2001).
[29] V. S. Berezinsky and A. Vilenkin, Phys. Rev. D **62**, 083512 (2000).