How Self-Interactions can Reconcile Sterile Neutrinos with Cosmology

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Short baseline neutrino oscillation experiments have shown hints of the existence of additional sterile neutrinos in the eV mass range. However, such neutrinos seem incompatible with cosmology because they have too large of an impact on cosmic structure formation. Here we show that new interactions in the sterile neutrino sector can prevent their production in the early Universe and reconcile short baseline oscillation experiments with cosmology.

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Introduction.—A variety of short baseline neutrino experiments seem to indicate the existence of at least one additional neutrino species with a mass in the eV range (see, e.g., Refs. [1,2]). In order to be compatible with the Large Electron-Positron Collider (LEP) constraint on the number of light neutrinos coupled to Z [3], these additional neutrinos must be sterile; i.e., they must be singlets under the $SU(2)\times U(1)$ electroweak gauge group. However, the fact that they do not couple to any particles in the standard model by no means implies that they are completely noninteracting. In fact, it is entirely possible, even natural, that the sterile neutrinos couple to other vector bosons which can have different properties from those associated with the $SU(2)\times U(1)$ of the standard model. Another option could be that the sterile neutrinos couple to a new pseudoscalar, as would be the case in Majoron-type models [4,5]. Here we will consider the possibility that sterile neutrinos can be strongly self-coupled through a Fermi 4-point interaction similar to the low energy behavior of neutrinos in the standard model, but with a completely different coupling strength. As we will see below, such a new interaction can have profound effects on active-sterile neutrino conversion in the early Universe and completely change cosmological bounds on sterile neutrinos.

Recent data on the anisotropy of the cosmic microwave background from the Planck satellite [6], in combination with auxiliary data on the large-scale distribution of galaxies, have shown that cosmology can accommodate sterile neutrinos of eV mass, but *not* if they are fully thermalized (see, e.g., Ref. [7]), because the suppression of structure formation is too strong (see Ref. [8] for a detailed discussion). The problem is that the masses and mixing angles preferred by terrestrial data inevitably lead to almost complete thermalization of sterile neutrinos. One possible way of circumventing this problem is to introduce a lepton asymmetry which pushes the resonant region in momentum space to very low values (see, e.g., Refs. [9,10]). The problem with this model is that it is far from clear how to produce this lepton asymmetry. Furthermore, the suppression changes very rapidly from zero to maximum suppression as a function of the lepton asymmetry, so partial thermalization requires some fine-tuning. Here we present an alternative scenario for preventing sterile neutrino production: If sterile neutrinos are strongly self-interacting, they provide a significant matter potential for themselves, which in turn completely changes the active-sterile conversion process. We will demonstrate that self-interactions can prevent sterile neutrino production to a point where bounds from cosmic microwave background and large-scale structure completely disappear—making sterile neutrinos with masses in the eV range perfectly compatible with precision cosmological data.

Scenarios.—We are considering a hidden gauge boson with mass M_X, and we take the mass to be $\gtrsim 100$ MeV, such that we can use an effective 4-point interaction for all temperatures of interest. The interaction strength is then written as

$$G_X \equiv \frac{g_X^2}{M_X^2}. \tag{1}$$

We will assume a $1+1$ scenario, specifically, a muon neutrino (or tau neutrino) and one sterile neutrino species, a simplification which does not qualitatively alter any of our findings. The system can then be fully characterized by a momentum-dependent, 2×2 Hermitian density matrix $\rho(p)$. Since we are not assuming any lepton asymmetry, the evolution of the antiparticle density matrix is trivial, since $\rho(p)=\bar{\rho}(p)$. We expand the density matrix in terms of Pauli matrices:

$$\rho = \frac{1}{2} f_0 (P_0 + \boldsymbol{P}\cdot\boldsymbol{\sigma}), \tag{2}$$

where $f_0 = (e^{p/T}+1)^{-1}$ is the Fermi-Dirac distribution and $\boldsymbol{\sigma}$ is a vector consisting of the three Pauli matrices. The evolution equations for P_0 and $\boldsymbol{P}$ are called the quantum kinetic equations, and they were first derived in Refs. [11–14] (for a presentation closer to the present one, see Refs. [9,15]). It is convenient to form the linear combinations

$$P_a \equiv P_0 + P_z = 2\frac{\rho_{aa}}{f_0}, \qquad P_s \equiv P_0 - P_z = 2\frac{\rho_{ss}}{f_0}, \quad (3)$$

which separate the sterile and the active sector. The equations of motions are then given by

$$\dot{P}_a = V_x P_y + \Gamma_a\left[2\frac{f_0}{f_0} - P_a\right], \quad (4a)$$

$$\dot{P}_s = -V_x P_y + \Gamma_s\left[2\frac{f_{\mathrm{eq},s}(T_{\nu_s}, \mu_{\nu_s})}{f_0} - P_s\right], \quad (4b)$$

$$\dot{P}_x = -V_z P_y - D P_x, \quad (4c)$$

$$\dot{P}_y = V_z P_x - \frac{1}{2}V_x(P_a - P_s) - D P_y. \quad (4d)$$

The Γ_s term is an approximation to the full scattering kernel which is valid in the limit of strong coupling. The sterile equilibrium distribution, $f_{\mathrm{eq},s}(T_{\nu_s}, \mu_{\nu_s}) = (e^{(p-\mu_{\nu_s})/T_{\nu_s}} + 1)^{-1}$, where T_{ν_s} and μ_{ν_s} are the sterile neutrino temperature and pseudochemical potential, respectively, is uniquely determined from the requirement that the interaction must respect energy conservation and number conservation. Γ_a and Γ_s are related to the 4-point interaction constants as

$$\Gamma_a = C_\mu G_F^2 p T^4, \qquad \Gamma_s = G_X^2 p T_{\nu_s}^4 n_{\nu_s}, \quad (5)$$

where $C_\mu \simeq 0.92$, while n_{ν_s} is the normalized number density of sterile neutrinos, $n_{\nu_s} = (2/3\zeta(3)T^3)\int p^2 \rho_{ss}(p)dp$. D quantifies the damping of quantum coherence in the system and is approximately half of the scattering rates: $D \simeq \frac{1}{2}(\Gamma_a + \Gamma_s)$. We have chosen to define Γ_s in analogy with Γ_a, and this means that we do not have exact conservation of ΔN_{eff} for the scattering term in Eq. (4b) since Γ_s depends on p. However, none of the results change significantly when we let $p = 3.15T$ in the expression for Γ_s.

In order to include the sterile neutrino self-interaction, we repeat the derivation in [14] for the self-interaction due to the Z boson in the active sector, but now for an X boson in the sterile sector. This gives an additional term in the matter potential V_z. The potentials are now

$$V_x = \frac{\delta m_s^2}{2p} \sin 2\theta, \quad (6a)$$

$$V_z = V_0 + V_a + V_s, \quad (6b)$$

$$V_0 = -\frac{\delta m_s^2}{2p} \cos 2\theta, \quad (6c)$$

$$V_a = -\frac{14\pi^2}{45\sqrt{2}} p\left[\frac{G_F}{M_Z^2} T_\gamma^4 n_{\nu_a}\right], \quad (6d)$$

$$V_s = +\frac{16 G_X}{3\sqrt{2} M_X^2} p u_{\nu_s}. \quad (6e)$$

Here, δm_s^2 is the mass difference, θ is the vacuum mixing angle, M_Z is the mass of the Z boson, M_X is the mass of the boson mediating the new force, and u_{ν_s} is the physical energy density of the sterile neutrino. We solve the system of equations using a modified version of the public code LASAGNA [16,17].

Results.—In Fig. 1 we show the degree of thermalization of the sterile neutrino, quantified in terms of the total energy density in the active plus sterile sector,

$$N_{\mathrm{eff}} \equiv \frac{u_{\nu_a} + u_{\nu_s}}{u_{\nu_0}}, \qquad u_{\nu_0} \equiv \frac{7}{8}\left(\frac{4}{11}\right)^{4/3} u_\gamma. \quad (7)$$

We have chosen $g_X = 0.1$ and a sample of values for G_X, and we show how ΔN_{eff} develops with the decreasing temperature. We can see that the thermalization of the sterile neutrino moves to lower temperatures when the interaction becomes stronger, and this is what we would expect since a strong interaction means that even a small background of sterile neutrinos can prevent further thermalization.

The amount of thermalization depends on both g_X and G_X, and in Fig. 2 we show ΔN_{eff} as a function of both. It shows that thermalization can be almost completely blocked by the presence of the new interaction for high values of G_X and low values of g_X.

Another interesting observation is that the degree of thermalization depends almost entirely on the mass of the

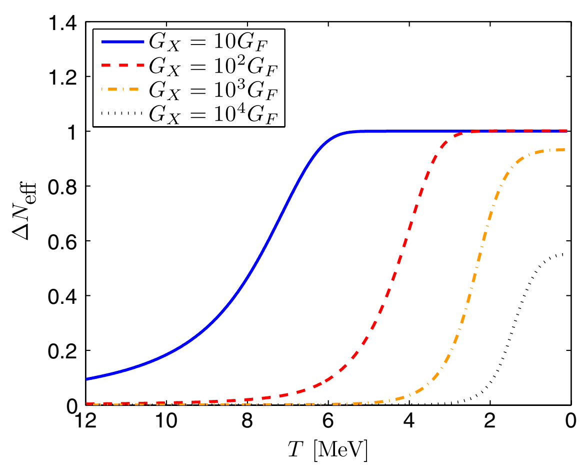

FIG. 1 (color online). The evolution of ΔN_{eff} as the temperature drops for $g_X = 0.1$ and different values of the coupling constant G_X.

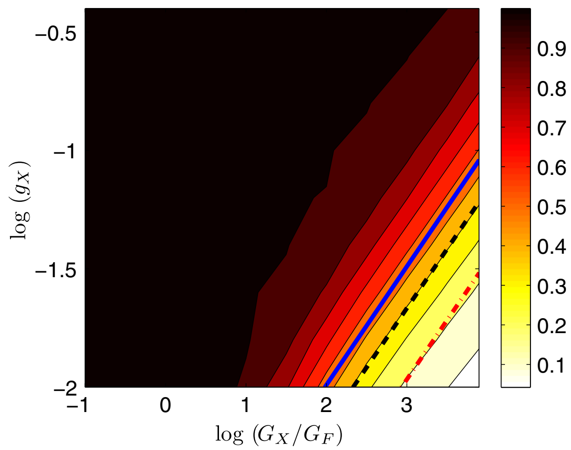

FIG. 2 (color online). Contours of equal thermalization. $\Delta N_{\rm eff}$ is given by the colors. The solid, dashed, and dot-dashed lines correspond to hidden bosons with masses $M_X = 300$, 200, and 100 MeV, respectively.

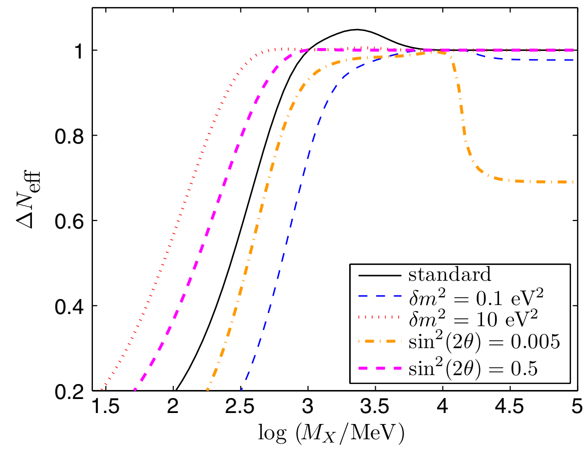

FIG. 3 (color online). Dependence of $\Delta N_{\rm eff}$ on the mixing parameters. $g_X = 0.01$ has been used for all the models, while G_X has been changed to give the variation in mass. The standard parameters are $\sin^2(2\theta) = 0.05$ and $\delta m^2 = 1$ eV^2.

new boson M_X, not on the dimensionless coupling g_X. This can be understood qualitatively from the following simple argument: At high temperature the production of sterile neutrinos is suppressed by rapid scattering (the quantum Zeno effect), but as soon as production commences, the thermalization rate of a sterile neutrino can be approximated by $\Gamma_t \sim \Gamma \sin^2(2\theta_m)$, where Γ is the rate with which "flavor content" (in this context meaning active versus sterile) is measured by the system, and θ_m is the in-medium mixing angle (see, e.g., Refs. [18,19] for a discussion of this in the context of active neutrinos). Γ is entirely dominated by the interaction via X so that $\Gamma \propto G_X^2$ and the in-medium mixing angle is likewise dominated by the potential generated by the new interaction so that $\sin^2(2\theta_m) \propto 1/V_s^2 \propto M_X^4/G_X^2$ leading to the sterile thermalization rate being proportional to M_X^4; i.e., Γ_t does not depend on g_X, only on M_X.

The determination of mixing parameters from accelerator experiments is quite uncertain, and it is therefore interesting to know how our results would be affected if we changed the vacuum mixing angle or the mass difference. The results of such a variation are seen in Fig. 3. Regarding the ability to inhibit thermalization, the results do not change much. A somewhat higher or lower mass will be needed for the hidden boson, but $\Delta N_{\rm eff} = 0.6$ can, for example, be reached by using $M_X = 100$ MeV even at $\delta m^2 = 10$ eV^2. There are, however, two other interesting observations. First, note that when the hidden boson mass is high, $\Delta N_{\rm eff}$ decreases with decreasing $\sin^2(2\theta)$ or δm^2—the well-known limit for noninteracting sterile neutrinos (see, e.g., Refs. [9,20]). As the boson mass is lowered, the new interaction first permits full thermalization of the sterile neutrino before we reach the mass range where the new interaction inhibits the thermalization.

The other interesting observation is that $\Delta N_{\rm eff} > 1$ for some values of M_X. At first this seems very puzzling and counterintuitive. In a model with only oscillations and no new interactions this would be impossible since the number density and energy density of the sterile neutrinos could never exceed the densities of the active neutrinos; the net production of steriles would simply shut off as soon as $\rho_{ss} \sim f_0$. However, in the model presented here there are two effects at play simultaneously: the production of steriles due to oscillations and the redistribution of sterile states due to the new interaction. If the redistribution of energy is sufficiently fast, it can keep $\rho_{ss} < f_0$, allowing for more production of steriles. Figure 4 provides an illustration of the effect by showing a snapshot of the distributions at the point where $\Delta N_{\rm eff}$ crosses 1 for a model with $M_X = 2.3$ GeV. Sterile neutrinos are still being produced in the region close to the resonance at $p/T \approx 5$ since $\rho_{ss} < f_0$ and oscillations therefore populate sterile

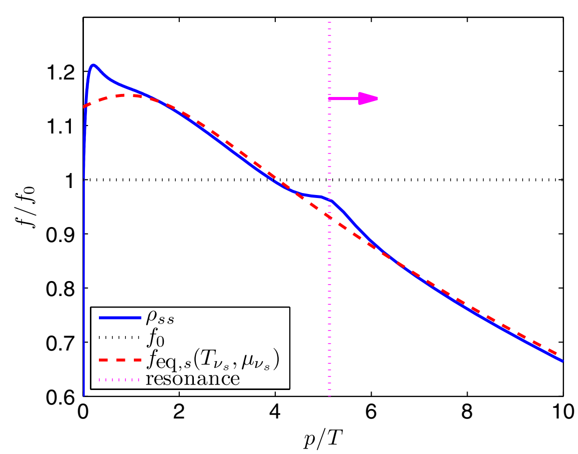

FIG. 4 (color online). The sterile energy distribution relative to f_0 at $T = 4.3$ MeV, where $\Delta N_{\rm eff}$ crosses 1 for $\delta m^2 = 1$ eV^2, $\sin^2(2\theta) = 0.05$, $G_X = G_F$, and $g_X = 0.01$, which corresponds to $M_X = 2.9$ GeV.

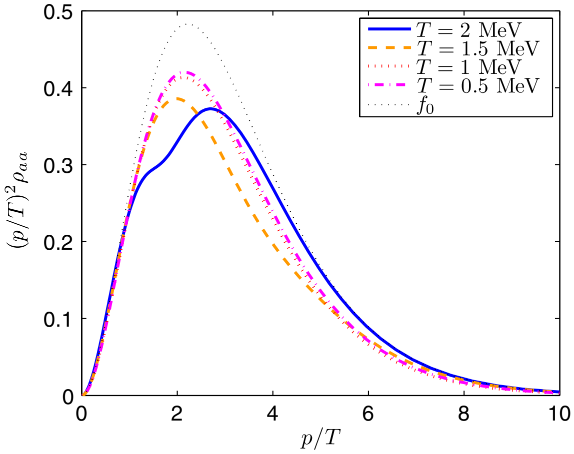

FIG. 5 (color online). The active neutrino distribution for different temperatures. The parameters used are $G_X = 3\times 10^2 G_F$ and $g_X = 0.025$. This corresponds to a hidden boson with the mass $M_X = 424$ MeV.

neutrinos from the active sector. At the same time, ρ_{ss} continues to grow at lower p/T due to the redistribution of states. In total this means that $\Delta N_{\rm eff}$ is still growing and will do so until the resonance has moved to very high p/T, where f_0 becomes very small or the active neutrinos decouple from the electrons. Naively, we would expect $\Delta N_{\rm eff}$ to be highest for low values of M_X because the energy redistribution becomes more efficient. However, when M_X is decreased, the suppression of oscillations due to the effect of M_X on the matter potential quickly wins and $\Delta N_{\rm eff}$ decreases rapidly with decreasing M_X. Therefore, $\Delta N_{\rm eff} > 1$ can only occur in a limited transition region of M_X if it occurs at all (which depends on the mixing parameters δm^2 and $\sin^2(2\theta)$).

Finally, we again stress that our treatment is only consistent if $M_X \gg T$ for *any* temperature relevant to our calculation. For the typical mass differences favored by short baseline measurements, the production of sterile neutrinos takes place at temperatures well below 100 MeV, and we have taken this as a representative minimum mass for the new boson. Note that such a low mass might be excluded for a boson coupling to the active sector [21]. (It should be noted that the bounds quoted in [21] are based on the assumption that the interaction can be treated as a 4-point interaction at $E \sim M_Z$. Since we are looking at very low mass bosons, this assumption does not hold and the bound is therefore expected to be much looser.) However, provided that the coupling is diagonal in "flavor" such that X couples only to the sterile state, such bounds are irrelevant. We also note that, in the case where G_X becomes very large, free-streaming of sterile neutrinos will be inhibited and structure formation bounds changed. However, such effects require G_X to be extremely large, $G_X \gtrsim 10^7 G_F$ (see, e.g., Refs. [22,23]).

*Big bang nucleosynthesis (BBN).—*Apart from the additional energy density in the sterile sector, the oscillations can have another important effect, namely, a distortion of the active neutrino distribution. This can happen even after neutrino decoupling because energy can still be transferred between the active and sterile sectors after the active neutrino decouples from the plasma. In models where the active-sterile conversion is delayed, such as the one presented here or models with a nonzero lepton asymmetry [10], this can in certain cases be the dominant cosmological effect. The reason is that the electron neutrino takes part in the nuclear reaction network relevant for BBN (see, e.g., Ref. [10]). Even if the sterile neutrino mixes primarily with ν_μ or ν_τ, active-active oscillations will transfer part of the distortion to the electron sector. However, a detailed investigation of this effect is beyond the scope of this Letter, and here we simply point out that interesting effects on BBN might occur. For illustration, we show in Fig. 5 how the active distribution can vary as a function of temperature compared to its unperturbed state f_0.

*Discussion.—*We have demonstrated that additional self-interactions of a sterile neutrino can prevent its thermalization in the early Universe and in turn make sterile neutrinos compatible with precision cosmological observations of structure formation. Arguably, the model discussed here is more natural than invoking a nonzero lepton asymmetry, relying only on the sterile sector possessing interactions similar to those in the standard model. In order for the model to work, the new gauge boson mediating the interaction must be significantly lighter than M_Z, but can easily be heavy enough that no significant background of such particles can exist at late times. Finally, we note that if this scenario is indeed realized in nature, future precise measurements of $N_{\rm eff}$ will effectively pinpoint the mass of the hidden gauge boson. In summary, the framework presented here presents a natural way of reconciling short baseline neutrino experiments with precision cosmology.

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*Note added.—*Recently, another Letter [24] on the same topic appeared.

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