New Dark Matter Candidate: Nonthermal Sterile Neutrinos

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We propose a new and unique dark matter candidate: ~100 eV to ~10 keV sterile neutrinos produced via lepton-number-driven resonant Mikheyev-Smirnov-Wolfenstein conversion of active neutrinos. The requisite lepton number asymmetries in any of the active neutrino flavors range from 10^{-3} to 10^{-1} of the photon number. The unique feature here is that the adiabaticity condition of the resonance strongly favors the production of lower energy sterile neutrinos. The resulting nonthermal (cold) energy spectrum can cause these sterile neutrinos to revert to nonrelativistic kinematics at an early epoch, so that free-streaming lengths at or below the dwarf galaxy scale are possible. Therefore, the main problem associated with light neutrino dark matter can be circumvented in our model. [S0031-9007(99)08792-X]

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Ample evidence shows that the dominant component of the matter content of the Universe is nonbaryonic dark matter contributing a fraction of the critical density $\Omega_m \sim 0.2$ to 1 [1,2]. In this Letter we suggest a novel means to allow sterile neutrinos with masses ~100 eV to ~10 keV to be the dark matter. We note that the existence of light sterile neutrinos may be an inevitable conclusion if the neutrino oscillation interpretations of the Super Kamiokande atmospheric neutrino data, solar neutrino data, and the data from the Los Alamos Liquid Scintillator Neutrino Detector are correct [3].

We envision production of our sterile neutrinos from active neutrinos via a net-lepton-number-driven MSW (Mikheyev-Smirnov-Wolfenstein [4]) resonant conversion process occurring near or during the big bang nucleosynthesis (BBN) epoch. Only neutrinos which evolve through resonances adiabatically are efficiently converted. In turn, lower energy neutrinos have larger ratios of resonance width to oscillation length and, hence, evolve more adiabatically. Therefore, the resonant process can produce a final energy distribution for the sterile neutrinos which is grossly nonthermal and skewed to very low energies. These neutrinos have energy distribution functions which are "cold." This, coupled with their relatively large rest mass, implies that they go nonrelativistic at quite early epochs, therefore essentially becoming cold dark matter (CDM), albeit with a cutoff at or below the dwarf galaxy scale. Below this cutoff scale, free streaming of these neutrinos can effectively damp out density fluctuations. Since this cutoff scale is much lower than any previous light neutrino dark matter models (e.g., hot dark matter (HDM) [5], warm dark matter (WDM) [6], and mixed cold plus hot (CHDM) [7] models), we deem these sterile neutrinos to be "cool dark matter" (CoolDM).

The contribution to the matter density of our Universe from an active neutrino species (ν_e, ν_μ, or ν_τ) with mass m_ν is

$$\Omega_\nu \approx \frac{m_\nu}{91.5h^2\ \mathrm{eV}}\,. \tag{1}$$

Here $h \equiv H_0/(100\ \mathrm{km/sec\,Mpc})$ where H_0 is the Hubble constant. Therefore, active neutrinos with masses in the ~10 eV range naturally serve as dark matter candidates. They are "hot" because their kinetic temperature is relatively high as a result of their light masses. It has been shown that models with HDM as the sole dark matter ingredient of the Universe fail to form galaxies at an early enough epoch. This is due to the large free streaming of light neutrinos which damp out density fluctuations below the free streaming scale [5]

$$\lambda_{\mathrm{f.s.}} \sim 40\left(\frac{m_\nu}{30\ \mathrm{eV}}\right)^{-1}\ \mathrm{Mpc}. \tag{2}$$

This generic problem of the large free streaming length of light neutrinos is difficult to circumvent. One possibility is to have non-Gaussian "seeds" in the primordial density fields [8]. Another is to have light neutrinos comprise only a small fraction of the dark matter ($\Omega_\nu/\Omega_m \lesssim 0.2$) [7,9].

WDM models consisting of $\sim 10^2$ eV sterile neutrinos have a free streaming length much shorter than that of HDM [10]. Because of this free streaming, WDM models have a reduced amplitude of density fluctuations at the galaxy cluster scale, which fits observations of structures at this scale better than the COBE-normalized standard CDM (sCDM) model (with $\Omega_m = 1$ and $h = 0.5$) [10]. However, density fluctuations at smaller galaxy scales ($\sim 1h^{-1}$ Mpc) in WDM models may be too small [10] to accommodate observations of damped Lyman-α systems at high redshifts, and observations of "Lyman-break" galaxies which are highly clustered galaxies at redshifts $z \sim 3$ [11].

The neutrino cool dark matter candidate proposed here is "colder" (i.e., with a much smaller free-streaming length) than all three classes of models mentioned above. In particular, the cold nonthermal energy distribution of our candidate sterile neutrinos causes them to move much more slowly than WDM sterile neutrinos with similar masses. The small free-streaming lengths in CoolDM

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models enable a circumvention of the generic problem associated with light neutrino dark matter candidates.

The production mechanism of our sterile neutrino CoolDM candidate is very different from that of its closest kin—the sterile neutrino WDM candidate. In the model proposed by Dodelson and Widrow [6], sterile neutrinos were produced by active-sterile neutrino oscillation at the BBN epoch. The oscillation was driven by finite-temperature matter effects [12]. The end result was a sterile neutrino population with an energy spectrum resembling that of an active neutrino species but with an overall suppression in normalization [6]. The average energy of these sterile neutrinos is therefore the same as the active ones, $\approx 3.151T$, where T is the temperature of the active neutrino species. In our model, the sterile neutrinos are produced by a *resonant* active-sterile neutrino transformation driven by a preexisting lepton number asymmetry. The resonant transformation is only adiabatic at the low energy portion of the neutrino energy spectrum. This feature creates a nonthermal sterile neutrino population, with only the low energy states populated.

We quantify this production mechanism in the density matrix formalism [13]. In this formalism, the $\nu_\alpha \leftrightarrow \nu_s$ ($\alpha = e$, μ, or τ, and ν_s is the sterile neutrino) system is described by a four vector $(P_0, \mathbf{P})$. The number densities of the neutrinos are $n_{\nu_\alpha} \equiv (P_0 + P_z)/2$ and $n_{\nu_s} \equiv (P_0 - P_z)/2$. The evolution of the mixing system in the BBN epoch satisfies the equation

$$\dot{\mathbf{P}} = \mathbf{V} \times \mathbf{P} + \dot{P}_0 \hat{\mathbf{z}} - D\mathbf{P}_\perp , \qquad (3)$$

where $\mathbf{P}_\perp = P_x \hat{\mathbf{x}} + P_y \hat{\mathbf{y}}$. The D term is the quantum damping term that acts to reduce the mixing system into flavor eigenstates and suppress the oscillation. Numerically $D \sim G_F^2 T^5$, where G_F is the Fermi constant. The vector $\mathbf{V}$ is the effective potential of the oscillation. Its components are

$$V_x = \frac{\delta m^2}{2E} \sin 2\theta, \qquad V_y = 0,$$
$$V_z = -\frac{\delta m^2}{2E} \cos 2\theta + V_\alpha^L + V_\alpha^T , \qquad (4)$$

where $\delta m^2 \equiv m_{\nu_s}^2 - m_{\nu_\alpha}^2 \approx m_{\nu_s}^2$, θ is the vacuum mixing angle, and E is the neutrino energy. The matter-antimatter asymmetry contribution to the effective potential is [12]

$$V_\alpha^L \approx 0.35 G_F T^3 \left[L_0 + 2L_{\nu_\alpha} + \sum_{\beta \neq \alpha} L_{\nu_\beta} \right], \qquad (5)$$

where L_0 represents the contributions from the baryonic asymmetry and electron-positron asymmetry, $\sim 10^{-10}$; and L_{ν_β} is the asymmetry in the other active neutrino species ν_β. For convenience, we will denote $\mathcal{L} \equiv 2L_{\nu_\alpha} + \sum_{\alpha \neq \beta} L_{\nu_\beta}$. The contribution to $\mathbf{V}$ from a thermal neutrino background is $V_\alpha^T \sim -10^2 G_F^2 E T^4$ [12].

Similar quantities for the antineutrino channel $\overline{\nu}_\alpha \leftrightarrow \overline{\nu}_s$ can be defined and they satisfy

$$\dot{\bar{\mathbf{P}}} = \bar{\mathbf{V}} \times \bar{\mathbf{P}} + \dot{\bar{P}}_0 \hat{\mathbf{z}} - \bar{D} \bar{\mathbf{P}}_\perp , \qquad (6)$$

where $\bar{V}_x = V_x$, $\bar{V}_y = V_y$, $\bar{V}_z = -\delta m^2 \cos 2\theta / 2E - V_\alpha^L + V_\alpha^T$, and $\bar{D} \approx D$.

The sign of $\mathcal{L}$ does not matter in the dark matter problem. For illustrative purposes we assume that it is positive, so that V_α^L is positive. Then the $\nu_\alpha \leftrightarrow \nu_s$ oscillation system encounters a resonance ($V_z = 0$, or, equivalently, a maximally matter-enhanced mixing) as a result of the nonzero $\mathcal{L}$ at a temperature

$$T_{\mathrm{res}} \approx 9 \left(\frac{m_{\nu_s}}{10^2 \ \mathrm{eV}} \right)^{1/2} \left(\frac{\mathcal{L}}{0.1} \right)^{-1/4} \epsilon^{-1/4} \ \mathrm{MeV}, \qquad (7)$$

where $\epsilon \equiv E/T$. At $T \sim T_{\mathrm{res}}$, V_α^T is completely negligible. We have implicitly taken $\cos\theta \approx 1$ since the mixing angle must be small. The antineutrino $\overline{\nu}_\alpha \leftrightarrow \overline{\nu}_s$ oscillation, on the other hand, has $|V_z| > |- \delta m^2/2E|$. It is therefore suppressed relative to a vacuum oscillation situation. The net effect is destruction of $\mathcal{L}$, because ν_α is resonantly transformed into sterile neutrinos while $\overline{\nu}_\alpha$ is not.

Quite similar to the putative matter-enhanced resonance transition in the solar neutrino problem, the adiabaticity condition of the $\nu_\alpha \to \nu_s$ resonant transformation at the resonant energy bin ϵ_{res} is

$$V_x^2 \left| \frac{\mathrm{d}\epsilon_{\mathrm{res}}}{\mathrm{d}V_z} \right| \left| \frac{\mathrm{d}\epsilon_{\mathrm{res}}}{\mathrm{d}t} \right|^{-1} > 1 . \qquad (8)$$

In the equation, $|V_x|$ is the transformation rate, $|V_x(\mathrm{d}\epsilon_{\mathrm{res}}/\mathrm{d}V_z)|$ is the energy width of the resonance, and $|\mathrm{d}\epsilon_{\mathrm{res}}/\mathrm{d}t|$ is the speed of movement of the resonance energy bin across the neutrino spectrum as ν_α neutrinos of different energies encounter resonance at different temperatures. (Strictly speaking, our expressions for the rate and the energy width of the resonances are valid only when collisions are not important, i.e., $D \ll |V_x|$. However, the adiabaticity condition is nonetheless the same when $D \gg |V_x|$ because the suppression in rate due to collisions is compensated by the increase in the energy width of the resonance.) From Eq. (4) we have

$$\left| V_x \frac{\mathrm{d}\epsilon_{\mathrm{res}}}{\mathrm{d}V_z} \right| \approx \epsilon_{\mathrm{res}} \sin 2\theta . \qquad (9)$$

From Eq. (7) we have

$$\frac{\mathrm{d}\epsilon_{\mathrm{res}}}{\mathrm{d}t} \approx \epsilon_{\mathrm{res}} \left[4H - \frac{\mathrm{d}\mathcal{L}/\mathrm{d}t}{\mathcal{L}} \right], \qquad (10)$$

where $H \approx 5.5 T^2/m_{\mathrm{Pl}}$ is the Hubble expansion rate and $m_{\mathrm{Pl}} \approx 1.22 \times 10^{28}$ eV is the Planck mass. The adiabaticity condition is therefore satisfied if

$$4 \times 10^9 \left(\frac{m_{\nu_s}}{10^2 \ \mathrm{eV}} \right)^{1/2} \left(\frac{\mathcal{L}}{0.1} \right)^{3/4} \epsilon_{\mathrm{res}}^{-1/4} \times \sin^2 2\theta \left[\frac{1}{1 - (\mathrm{d}\mathcal{L}/\mathrm{d}t)/4H\mathcal{L}} \right] > 1 . \qquad (11)$$

Assuming adiabatic neutrino evolution through resonances, we have $\mathrm{d}\mathcal{L}/\mathrm{d}t = f(\epsilon_{\mathrm{res}}) \mathrm{d}\epsilon_{\mathrm{res}}/\mathrm{d}t = f(\epsilon_{\mathrm{res}}) \epsilon_{\mathrm{res}} [4H - (\mathrm{d}\mathcal{L}/\mathrm{d}t)/\mathcal{L}]$ where $f(\epsilon)$ is the neutrino distribution function. We therefore derive $\mathcal{L} = \mathcal{L}^{\text{init}} - \int_0^{\epsilon_{\text{res}}} f(\epsilon)\, d\epsilon$ where superscript "init" indicates initial values. As a result, $|(d\mathcal{L}/dt)/\mathcal{L}| \lesssim H$ is always true unless $|\mathcal{L}| \ll |\mathcal{L}^{\text{init}}|$. This implies that the adiabaticity condition Eq. (11) holds true for $\nu_\alpha \leftrightarrow \nu_s$ vacuum mixing that is not too small ($\sin^2 2\theta \gtrsim 10^{-9}$), until most of $\mathcal{L}$ is destroyed by the resonant $\nu_\alpha \leftrightarrow \nu_s$ transformation. The final $\mathcal{L}$ is $\sim$0 after the resonant conversion process. In this limit, the total change of the $\nu_\alpha\overline{\nu}_\alpha$ asymmetry is $\Delta L_{\nu_\alpha} \approx \mathcal{L}^{\text{init}}/2$. This change is entirely due to the ν_α to ν_s transformation, implying that the ν_s sea produced has a number density that is a fraction

$$F \approx \frac{4}{3}\,\Delta L_{\nu_\alpha} \tag{12}$$

of the number density of an active neutrino species. The sterile neutrino contribution to the matter density today is then

$$\Omega_\nu \approx F\left(\frac{m_{\nu_s}}{91.5h^2\ \text{eV}}\right) \approx \left(\frac{m_{\nu_s}}{343\ \text{eV}}\right)\left(\frac{h}{0.5}\right)^{-2} \times \left(\frac{2L_{\nu_\alpha} + \sum_{\beta\neq\alpha} L_{\nu_\beta}}{0.1}\right), \tag{13}$$

with $\alpha, \beta = e, \mu, \tau$. Here we have assumed that all neutrino asymmetries are their initial values and we have dropped the superscript "init" (L_{ν_β} will not change anyway).

From Eqs. (4) and (13) we note that if the sign of $2L_{\nu_\alpha} + \sum_{\beta\neq\alpha} L_{\nu_\beta}$ is negative, the antineutrino counterpart of ν_s will be resonantly produced from $\overline{\nu}_\alpha$, and may therefore become the dark matter candidate if the counterpart of Eq. (13) is satisfied.

As the transition is only adiabatic when $\mathcal{L}$ is a significant fraction of its initial value, only the low energy sterile neutrinos [which encounter resonances first while $\mathcal{L}$ is still relatively large, see Eq. (7)] are resonantly produced. Higher energy neutrinos go through the resonance later, when most of $\mathcal{L}$ is damped. Evolution through resonances for these high energy neutrinos is therefore nonadiabatic, producing no significant $\nu_\alpha \to \nu_s$ conversion. The end result is a sterile neutrino energy spectrum as illustrated in Fig. 1, with a characteristic mean neutrino energy E approximately satisfying $\int_0^{E/T} f(E'/T)d(E'/T) \approx 8\Delta L_{\nu_\alpha}/3 \approx 4\mathcal{L}/3$. As an example, for $\mathcal{L} = 0.02$, we have E/T is 0.7, i.e., less than one-quarter of the average E/T for active species. The comoving free-streaming length of these sterile neutrinos becomes

$$\lambda_{\text{f.s.}} \sim 40\left(\frac{m_\nu}{30\ \text{eV}}\right)^{-1}\left(\frac{E/T}{3.15}\right)\ \text{Mpc}\,, \tag{14}$$

which for a given sterile neutrino mass can be more than a factor of several shorter than that calculated from Eq. (2).

In this Letter we limit our discussion of the sterile neutrino production to a temperature range below the quark-hardron phase transition, $\sim$150 MeV. The reason for this is that the neutrino transformation process can be modeled confidently below this temperature. With this restriction, Eq. (7) shows that the preexisting $\mathcal{L}$ has to be $\gtrsim 10^{-3}$. The mass of the sterile neutrino dark matter candidate is then $\lesssim$10 keV from Eq. (13). ($\mathcal{L}$ and the sterile neutrino mass are limited simultaneously because their product must yield $\Omega_\nu \sim 1$.) If we also restrict the sterile neutrino mass to be $m_{\nu_s} \gtrsim 10^2$ eV in order to have potentially successful structure formation at galaxy scales, Eq. (13) implies that $\mathcal{L} \lesssim 10^{-1}$. For example, if the preexisting $2L_{\nu_\mu} + L_{\nu_e} + L_{\nu_\tau}$ is 0.02, a sterile neutrino with $m_{\nu_s} = 1160$ eV and mixing with ν_μ will yield $\Omega_\nu = 0.4$ for $h = 0.65$. The comoving free-streaming length of this dark matter is only $\lambda_{\text{f.s.}} \sim$ 0.4 Mpc, giving rise to a mass-scale cutoff of $M \sim 10^{10} M_\odot$, roughly the scale of dwarf galaxies. In this case, structure formation in CoolDM models differs from that in CDM models only at subgalactic scales, such as in halo structure and formation history at high redshift. These differences may potentially be tested by observations of galaxy rotation curves [14] and high redshift galaxies [11,15]. However, if the mass of the sterile neutrinos is as high as 10 keV, resulting from $\mathcal{L} \sim 10^{-3}$, then CoolDM essentially becomes CDM. It is intriguing that oscillation between active neutrinos and keV sterile neutrinos in supernovas in the presence of strong magnetic fields may potentially give rise to "pulsar kicks" [16].

An $\mathcal{L}$ between $\sim 10^{-3}$ and $\sim 10^{-1}$ necessarily implies that the asymmetries in one or more active neutrino species are of the same order (or larger if there is cancellation between asymmetries). Lepton asymmetries between $\sim 10^{-3}$ and $\sim 10^{-1}$ are not ruled out by big bang nucleosynthesis [17]. In fact, if this asymmetry resides in the ν_μ or ν_τ sector, its impact on the primordial ^{4}He yield (which gives the constraint on lepton asymmetry) is well below the current detection level. If some of the asymmetry is in the ν_e sector, with a magnitude $\gtrsim 10^{-2}$, its impact on the primordial ^{4}He yield is quite appreciable. Just where the limit on L_{ν_e} stands is hard to quantify precisely because of the large uncertainty in the primordial ^{4}He abundance measurements [18,19]. In particular, Kohri *et al.* [20] has argued that an $L_{\nu_e} \approx 0.015$ is what is needed to bring a low set of primordial ^{4}He abundance values [18] into agreement with the measured low primordial deuterium abundance [21].

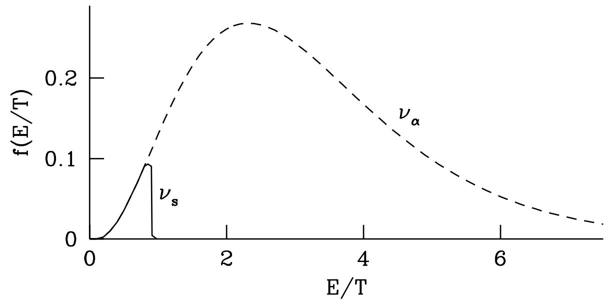

FIG. 1. Illustration of the energy spectrum of the sterile neutrinos produced vs the thermal spectrum of the active neutrino species.

On the other hand, in individual neutrino sectors, lepton asymmetries of order $\sim 10^{-3}$ to $\sim 10^{-1}$ are very large compared to the baryon asymmetry, $\sim 10^{-10}$. But this is not to say that their sum, the total lepton number asymmetry, cannot be as small as $\sim 10^{-10}$. In the absence of such a cancellation, however, the resultant large lepton number must be generated below the electroweak phase transition to avoid the transfer between lepton number and baryon number. Models that generate a large lepton number despite a small baryon number asymmetry indeed exist [22,23].

In summary, we have proposed that $\sim 10^2$ eV to ~ 10 keV sterile neutrinos ν_s will be a cool or even a cold dark matter candidate if they were produced in the early Universe via active-sterile neutrino resonant oscillation in the presence of a lepton number asymmetry. The amplitude of this lepton number asymmetry is $\sim 10^{-3}$ to $\sim 10^{-1}$ for at least one active neutrino species, which is consistent with constraints from big bang nucleosynthesis. The minimal set of free parameters in our cool dark matter model, and its potentially testable predictions for halo structures and protogalaxy formation, all make the cool dark matter scenario interesting and worthy of further invistigations.

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