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Sterile Neutrinos, Dark Matter, and Pulsar Velocities in Models with a Higgs Singlet

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We identify the range of parameters for which the sterile neutrinos can simultaneously explain the cosmological dark matter and the observed velocities of pulsars. To satisfy all cosmological bounds, the relic sterile neutrinos must be produced sufficiently cold. This is possible in a class of models with a gauge-singlet Higgs boson coupled to the neutrinos. Sterile dark matter can be detected by the x-ray telescopes. The presence of the singlet in the Higgs sector can be tested at the CERN Large Hadron Collider.

Sterile neutrinos with masses of the order of a few keV can explain the observed velocities of pulsars [1–4], can play a role in the star formation and reionization of the Universe [5], and can be cosmological dark matter [6–9] if their relic population is sufficiently "cold."

While it is possible that the keV mass scale is a new fundamental constant of nature, it is of interest to consider the possibility that sterile neutrino masses arise from the Higgs mechanism, just like the masses of the other fermions. Then the keV scale is generated by the Higgs vacuum expectation value (VEV), and it depends on the coupling. The sterile neutrinos can couple to the $SU(3) \times SU(2) \times U(1)$ singlet Higgs boson S, whose vacuum expectation value can give them the Majorana masses. Models of this kind have been proposed [10,11], and their implications for the electroweak phase transition, baryogenesis, and collider searches have been studied in detail [12]. The production of relic sterile neutrinos has also been studied in one specific limit of parameters, in which the singlet S with a sub-GeV mass can play the role of the inflaton [11]. Here we will consider a more general range of parameters, focusing on the region in which the sterile neutrinos can simultaneously explain pulsar kicks and dark matter, while satisfying the Lyman-α, x-ray, and other bounds. We will see that, if the scalar Higgs boson mass is of the order of the electroweak scale, all these constraints can be satisfied simultaneously, and, in particular, the momentum distribution of the relic sterile neutrinos can be cold enough for dark matter. The mixing between the Higgs doublet and the singlet can be probed by the upcoming experiments at the LHC.

The neutrino masses can be introduced by means of the following addition to the standard model Lagrangian:

$$\mathcal{L} = \mathcal{L}_0 + \bar{N}_a (i \not\partial) N_a - y_{\alpha a} H \bar{L}_\alpha N_a - \frac{h_a}{2} S \bar{N}_a^c N_a$$
$$+ V(H, S) + \text{H.c.}, \tag{1}$$

where $\mathcal{L}_0$ includes the gauge and kinetic terms of the standard model, H is the Higgs doublet, S is the real singlet, L_α ($\alpha = e, \mu, \tau$) are the lepton doublets, and N_a

($a = 1, \ldots, n$) are the additional singlet neutrinos. Let us consider the following scalar potential:

$$V(H, S) = m_H^2 |H|^2 + m_S^2 S^2 + \lambda_{HS} |H|^2 S^2$$
$$+ \lambda_S S^4 + \lambda_H |H|^4. \tag{2}$$

After the electroweak symmetry breaking, the Higgs doublet and singlet fields each develop a VEV, $\langle H \rangle = v_0 = 247$ GeV, $\langle S \rangle = v_1$, and the singlet neutrinos acquire Majorana masses $M_a = h_a v_1$. The masses of the Higgs doublet and singlet at zero temperature are

$$\tilde{m}_{H,S}^2 = \lambda_H v_0^2 + \lambda_S v_1^2 \pm \sqrt{D}, \tag{3}$$

where $D = (\lambda_H v_0^2 - \lambda_S v_1^2)^2 + \lambda_{HS}^2 v_0^2 v_1^2$. Below the scale of this symmetry breaking, the low-energy effective Lagrangian is

$$\mathcal{L} = \mathcal{L}_{SM} + \bar{N}_a (i \not\partial) N_a - y_{\alpha a} H \bar{L}_\alpha N_a$$
$$- \frac{M_a}{2} \bar{N}_a^c N_a + \text{H.c.}, \tag{4}$$

where $\mathcal{L}_{SM}$ is the standard model Lagrangian. This is the usual seesaw Lagrangian [13]. The number n of singlet neutrinos is not limited by the anomaly constraints or any other theoretical considerations, and the experimental limits exist only for larger mixing angles [14]. Supernova 1987A provides a constraint [7,15], which depends on the mixing angle and the sterile neutrino mass [7]. To explain the neutrino masses inferred from the atmospheric and solar neutrino experiments, $n = 2$ singlets are sufficient [16], but a greater number is called for if the Lagrangian (4) is to explain the Liquid Scintillator Neutrino Detector experiment [17], the r-process nucleosynthesis [18], the pulsar kicks [2–4], and dark matter [6–9,19]. The scale of the right-handed Majorana masses M_a is unknown; it can be much greater than the electroweak scale [13], or it may be as low as 1–10 eV [17,20]. The seesaw mechanism [13] can explain the smallness of the neutrino masses in the presence of the Yukawa couplings of the order of 1 if the Majorana masses are much larger than the electroweak scale; then the light neutrino masses

are suppressed by the ratio $\langle H \rangle / M$. However, the origin of the Yukawa couplings remains unknown, and there is no evidence that these couplings must be of order 1. In fact, the Yukawa couplings of the charged leptons are much smaller than 1. Theoretical *naturalness* arguments in favor of the low-energy seesaw [17] appear to be as compelling as those in favor of the high-scale seesaw [13]. In both limits one can have a successful leptogenesis: in the case of the high-scale seesaw, the baryon asymmetry can be generated from the out-of-equilibrium decays of heavy neutrinos [21], while in the case of the low-energy seesaw, the neutrino oscillations produce the asymmetry [22].

Let us first consider just one singlet neutrino, with Majorana mass of the order of several keV, in light of its possible role in explaining the pulsar kicks and dark matter. The same neutrino can play an important role in star formation and reionization of the Universe [5], as well as other astrophysical phenomena [23].

The range of parameters consistent with the explanation of pulsar kicks is shown in Fig. 1. It contains two regions corresponding to resonance oscillations [2] and off-resonance oscillations [3]. The boundaries of these regions are defined by the requirements of (i) supernova energetics, including SN 1987A constraints, (ii) adiabaticity and weak damping for resonant oscillations, (iii) sufficient anisotropy to explain observed pulsar speeds [1–3].
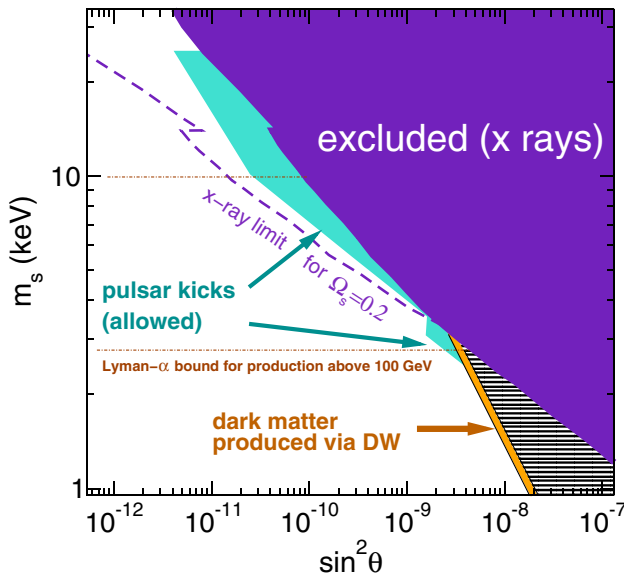

FIG. 1 (color online). The x-ray limits reported in Ref. [27] (dashed line) apply if sterile neutrinos account for all the dark matter ($\Omega_s = 0.2$). The value of Ω_s depends on production mechanisms, but it cannot be lower than the amount produced via the DW mechanism [6] (except for low-reheat scenarios [20,29]). The model-independent exclusion plot [purple (or solid black) region] is obtained by assuming this minimal value. A sterile neutrino with mass 3 keV and $\sin^2\theta \approx 3 \times 10^{-9}$, produced at some temperature above 100 GeV, can explain both pulsar kicks and dark matter.

If the coupling h in Eq. (1) is small enough, the sterile neutrinos are out of thermal equilibrium at any time after inflation. This is the case if the annihilations $NN \rightarrow NN$, $NN \rightarrow$ scalars, etc., are not fast enough to keep the sterile neutrinos in equilibrium for temperatures T in the range $\tilde{m}_S < T < T_{\text{reheat}}$. For $\tilde{m} \sim 1$ TeV and the reheat temperature $T_{\text{reheat}} < 10^{16}$ GeV, the sterile neutrinos are out of equilibrium for $h \ll 10^{-6}$ [12].

The sterile neutrinos can be produced out of equilibrium in two different processes. First, since S is in thermal equilibrium at high temperature, sterile neutrinos are produced through decays $S \rightarrow NN$ [11]. Most of these neutrinos are produced at temperatures of the order of $\tilde{m}_S \sim (0.1 - 1)$ TeV. Second, at much lower temperatures, $T \sim 0.1$ GeV, the sterile neutrinos are produced from oscillations of active neutrinos, as in the Dodelson-Widrow (DW) scenario [6]. If the lepton asymmetry is relatively large, the sterile neutrinos are produced much more efficiently [9]. The lepton asymmetry of the Universe is not know, but strong upper bounds do exist [24].

These two production mechanisms operate sequentially. S decays are governed by the coupling h, while the production via DW mechanism depends on mixing angles. In the limit $y_{\alpha i} v_0 \ll \min_i(M_i)$, these mixing angles are given by the usual seesaw relations: $\theta_{\alpha i} = y_{\alpha i} \langle H \rangle / M_i$. For simplicity we will assume that only one sterile neutrino has mass of several keV and that only one of the mixing angles is nonzero. This mixing angle is a function of both couplings, y and h, as well as $\tan\beta = v_0 / v_1$:

$$\theta = \frac{y\langle H \rangle}{h\langle S \rangle} = \frac{y}{h}\tan\beta. \tag{5}$$

There are two limits which, in combination, appear to stymie the DW proposal [6] for sterile dark matter. We will see that the models with singlet Higgs bosons offer a way to reconcile these seemingly contradictory bounds.

One important limit on the abundance of relic sterile neutrinos comes from x-ray observations. Relic neutrinos can decay into one of the active neutrinos and a photon. Since this is a two-body decay, photon energy is equal to one-half of the particle mass. The flux of expected x-ray photons depends on the decay rate, which is proportional to $(\sin^2\theta)$ [25]. This flux also depends on the relic abundance of sterile neutrinos, which, in turn, depends on the value of h (for the $S \rightarrow NN$ production at temperature $T \sim \tilde{m}_S$) and θ (for the production via neutrino mixing at temperature $T \sim 0.1$ GeV). If the latter is the dominant production mechanism, then the sterile neutrinos can be the dark matter in the range of masses [26] below 3.5 keV, as shown in Fig. 1. For larger masses, the x-ray limits on the decay photons [27] disallow the relic sterile neutrinos, unless their abundance is smaller than what is required for dark matter. However, if S decays at $T \sim (0.1–1)$ TeV are the dominant source of sterile neutrinos, they can still make up the entire dark matter, even if θ is arbitrarily small.

Another limit on dark matter comes from the Lyman-α observations [28] and is based on the requirement that dark matter be cold enough to generate the smallest structures observed in the absorption spectra of distant quasars. This limits the free-streaming length of dark matter, whose relation with the particle mass depends on the production mechanism. If the sterile neutrinos are produced via mixing, then the Lyman-α bound is $m > 10$ keV [28]. Obviously, this limit is in conflict with the x-ray limit mentioned above, namely $m < 3.5$ keV.

If sterile neutrinos account for only a part of dark matter, then Lyman-α bounds do not apply, and x-ray bounds are weaker. X-ray limits [27], shown as a dashed line in Fig. 1, are based on the assumption that sterile neutrinos constitute the entire dark matter, i.e., $\Omega_s = 0.2$. However, if $h \to 0$, this condition is not satisfied for points along the dashed line; in fact, $\Omega_s \ll 0.2$ for most of these points. This is because, in the absence of S decays, the density of sterile neutrinos in the Universe may be too small to explain dark matter. The x-ray signal is correspondingly smaller in this case, and the x-ray bounds are weaker. However, there still exists an x-ray bound on a sterile neutrino with a given mass and mixing angle, although this particle may be only a small part of dark matter.

Production via neutrino mixing [6] provides the lowest possible abundance of sterile neutrinos, as long as the Universe was reheated to a GeV or higher temperature. In low-reheat cosmological scenarios the bounds are relaxed considerably [20,29]. Here we will assume that the Universe has reached temperatures well above TeV after inflation. In Fig. 1 we show both the bounds based on the assumption $\Omega_s = 0.2$ and the model-independent exclusion region [solid purple (or solid black)] based on production only via the DW mechanism [6], in which case Ω_s can be smaller than 0.2. To calculate this production, we used an analytical fit to the numerical calculation of Abazajian [26]. There may be some hadronic uncertainties in this calculation [30].

There is a range of parameters for which the sterile neutrinos can explain the pulsar velocities and can affect the star formation, although they may not be the dominant component of the dark matter (see Fig. 1). However, it is also possible to explain pulsar kicks and dark matter simultaneously. Indeed, if the S-decay mechanism dominates the production of relic neutrinos, they may be redshifted sufficiently as the Universe cools from $T_S \sim (0.1–1)$ TeV to $T < $ MeV.

If all the dark matter is made up of sterile neutrinos, the bound $m_s > 10$ keV [28] applies to the DW scenario, in which the average momentum of a keV neutrino at low temperature T is $\langle p_s \rangle_{\text{DW}} \approx 2.8\ T$ [26]. Any additional redshifting of sterile neutrinos (e.g., due to some entropy production [31]) relaxes the 10 keV limit to a lower value.

Production of sterile neutrinos via decay $S \to NN$ occurs mainly at temperature $T_S \approx \tilde{m}_S$ [11]. Our model has

two mass eigenstates given by Eq. (3), subject to thermal corrections. In the case of different masses and non-negligible mixing, one of the mass eigenstates decouples before the other. Each mass eigenstate has a nonzero S component and, therefore, contributes to the production of the relic population of sterile neutrinos. If we make a simplifying assumption of equal masses (achieved when $\lambda_H \langle H \rangle^2 \approx \lambda_S \langle S \rangle^2$ and when λ_{HS} is small), the results of Ref. [11] are directly applicable to our case, and one can write that the amount of sterile dark matter is

$$\Omega_s = 0.2 \left(\frac{33}{\xi}\right)\left(\frac{h}{1.4 \times 10^{-8}}\right)^3\left(\frac{\langle S \rangle}{\tilde{m}_S}\right) = \tag{6}$$

$$0.2 \left(\frac{33}{\xi}\right)\left(\frac{h}{1.4 \times 10^{-8}}\right)^3\left(\frac{1}{\sqrt{2\lambda_S}}\right), \tag{7}$$

where ξ is the change in the number density of sterile neutrinos relative to T^3 due to the dilution taking place as the Universe cools down from $T_S \sim (0.1–1)$ TeV to a temperature well below MeV. The sterile neutrinos produced in these decays have an almost thermal spectrum at the time of production. More precisely, their average momentum $\langle p_s \rangle = \pi^6/(378\zeta(5))T \approx 2.45\ T$ [11], while the average momentum of the relativistic fermions in equilibrium is $p_T \approx 3.15\ T$. As the Universe cools down, the number of effective degrees of freedom decreases from $g_*(T_S) = 110.5$ to $g_*(0.1\ \text{MeV}) = 3.36$. Then $\xi = g_*(T_S)/g_*(0.1\ \text{MeV}) \approx 33$. This causes the redshifting of $\langle p_s \rangle$ by the factor $\xi^{1/3}$:

$$\langle p_s \rangle_{(T \ll 1\ \text{MeV})} = 0.76\ T\left[\frac{110.5}{g_*(\tilde{m}_S)}\right]^{1/3} \tag{8}$$

Comparing Eq. (8) with the DW case, one concludes that, as long as the population of sterile neutrinos is dominated by those produced at high temperature (large enough h, small θ), the Lyman-α limit changes from 10 keV to

$$m_s > 2.7\ \text{keV} \tag{9}$$

This lower bound is shown in Fig. 1.

In addition, the electroweak phase transition in this model can be first order, and entropy production can further redshift the momentum distribution of sterile neutrinos. Enqvist *et al.* [12] have found that the energy density increase due to the phase transition can be at most $10 T_0^4$, where T_0 is the transition temperature. This changes ξ by at most a factor 1.3.

The presence of the singlet Higgs boson has important implications for Higgs searches. Although we have focused so far on the keV sterile neutrino, the present data requires $n \geq 3$ sterile neutrinos in Eq. (1) [17,19]. This is not in conflict with big-bang nucleosynthesis (BBN): one or two additional thermalized neutrinos are consistent with BBN constraints at the one (two) sigma level [32]; besides some of these sterile neutrinos can be out of equilibrium. The couplings h_a of the additional sterile neutrinos need

P H Y S I C A L   R E V I E W   L E T T E R S

not be small, and the mixing between H and S can also be large. Hence, the Higgs boson can decay invisibly. This possibility has the effect of weakening the LEP bound on the mass of the lightest Higgs [12,33]. However, one can discover the invisible Higgs boson h_{inv} at the LHC in the $Z + h_{inv}$ channel [34]. For Higgs mass of 120 GeV, the discovery is possible at the LHC already with 10 fb^{-1} in the $Z + h_{inv}$ channel, while 100 fb^{-1} of data can afford the discovery in the weak boson fusion channel [34]. For some range of couplings, the singlet S can decay into visible channels and can be discovered via displaced vertices [35].

To summarize, relic sterile neutrinos with mass of a few keV can simultaneously explain pulsar velocities and dark matter if their production in the early Universe occurred above the electroweak scale. Over a broader range of parameters (Fig. 1), sterile neutrinos can explain pulsar kicks and can play a role in star formation, while they may not be the dominant component of dark matter.

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