Inert-sterile neutrino: Cold or warm dark matter candidate

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(Received 28 January 2010; published 25 March 2010)

In usual particle models, sterile neutrinos can account for the dark matter of the Universe only if they have masses in the keV range and are warm dark matter. Stringent cosmological and astrophysical bounds, in particular, imposed by x-ray observations, apply to them. We point out that in a particular variation of the inert doublet model, sterile neutrinos can account for the dark matter in the Universe and may be either cold or warm dark matter candidates, even for masses much above the keV range. These inert-sterile neutrinos, produced nonthermally in the early Universe, would be stable and have very small couplings to standard model particles, rendering very difficult their detection in either direct or indirect dark matter searches. Their existence could be revealed only by discovering other particles of the model in collider experiments.

DOI: 10.1103/PhysRevD.81.063529 PACS numbers: 95.35.+d

I. INTRODUCTION

Since the first indications more than seven decades ago [1], many different strong pieces of evidence supporting the existence of dark matter have been accumulated (see e.g. Refs. [2–4]). The presence of dark matter has been revealed so far through its gravitational effects. Much effort is being devoted to the detection of dark matter annihilation or decay products or the scattering of dark matter particles off nuclei. However dark matter may consist of particles which will not be revealed (at least in the near future) in this type of searches. We provide here an example of a dark matter candidate found in a simple extension of the standard model (SM), whose nature could be indirectly proven only through the discovery and study in colliders of other nonstandard particles predicted within the model. The dark matter particle candidate we study here is a sterile neutrino with mass in the tens of keV to the tens of GeV range and produced nonthermally in the early Universe, which can be either warm dark matter (WDM) or cold dark matter (CDM).

One or more gauge singlet right-handed (sterile) neutrinos are included in simple extensions of the SM which can easily accommodate neutrino oscillation data [5–8]. These data show that at least two of the active neutrinos have a nonzero mass. In many models sterile neutrinos are the right-handed Dirac mass partners of the active neutrinos. In some seesaw-inspired models, sterile neutrinos have large Majorana masses, which lead to three light (mostly active) neutrinos and several heavier (mostly sterile) neutrinos, the lightest of which is an attractive dark matter candidate [8]. Since this candidate necessarily decays into a light neutrino and a photon, to constitute the dark matter its lifetime must be much longer than the age of the Universe. Thus, this dark matter candidate might be detected through the photons produced in its decay in the dark halos of galaxies. Moreover, to have the required dark matter relic density, the lightest sterile neutrino must usually have a mass in the keV range, although this depends on the mechanism through which sterile neutrinos are produced in the early Universe.

Relic sterile neutrinos with only standard model interactions are produced in the early Universe through active-sterile neutrino oscillations. Sterile neutrinos produced through nonresonant oscillations [6–8] must have masses M_s in the keV range to account for the whole of the dark matter and are WDM. Through a combination of x-ray and structure formation constraints, an upper bound $M_s \leq$ 3–4 keV has been obtained [7–11] (see, however, Ref. [12] for a very recent weak hint of a possible signal). Lyman-α forest data have been used to impose the lower bound of $M_s \geq 5.6$ keV [13] (see also Refs. [14,15] for previous bounds) on nonresonantly produced sterile neutrinos, or the revised limit of $M_s \geq 8$ keV obtained by a new analysis [16], which combined with the previous upper bounds would exclude nonresonantly produced dark matter sterile neutrinos. Even disregarding the controversial Lyman-α bounds, the mass range allowed for these neutrinos is very restricted because there is an independent lower bound $M_s \geq 1.8$ keV [17,18] derived from the analysis of phase space density evolution of dwarf spheroidal galaxies. In general, these bounds do not consider the possibility of a very large lepton asymmetry. In the presence of a large lepton asymmetry $\mathcal{L} \equiv (n_{\nu_e} - n_{\bar{\nu}_e})/s > 10^{-6}$, where n_{ν_e} and $n_{\bar{\nu}_e}$ are the number densities of neutrinos and antineutrinos and s is the entropy density in the Universe, sterile neutrinos may be produced in the early Universe through resonant oscillations [19,20].

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Considering the upper limit of the lepton asymmetry imposed by big bang nucleosynthesis (BBN), $\mathcal{L} < 2.5 \times 10^{-3}$ [20], the range 1 keV $\leq M_s \leq$ 50 keV is in principle allowed for sterile neutrino dark matter [17,20,21]. In slightly more complicated models, sterile neutrino dark matter may be produced as decay products of, for example, a heavy singlet scalar [22,23], or may not completely thermalize as in low reheating temperature scenarios [24]. Yet, in all these models the x-ray constraints are important.

Here, we consider a small variation of the SM in which the lightest sterile neutrino is stable (hence it does not produce photons as decay products) and may constitute all of the dark matter. We study a variation of the inert doublet model [25,26] (in itself an extension of the model in Ref. [27]). We present the model in Sec. II. In Sec. III, we describe the scenario we consider. We show the conditions for the lightest sterile neutrino to constitute the dark matter in Sec. IV and conclude in Sec. V.

II. THE MODEL

In this variation of the inert doublet model [25,26], one scalar doublet $\eta(\eta^+, \eta_0)$ and three sterile neutrinos, which we call inert-sterile neutrinos, N_i with $i = 1, 2, 3$ odd under a new parity Z_2, are added to the SM. All the particles in the SM are even under the additional Z_2 symmetry. These assignments make the new particles ''inert'' because their couplings to the SM particles are very limited. The leptonic Yukawa couplings in this model are

$$\mathcal{L}_Y = f_{ij}(\phi^- \nu_i + \bar{\phi}^0 l_i) l_j^c + h_{ij}(\nu_i \eta^0 - l_j \eta^+) N_j + \text{H.c.} \tag{2.1}$$

Here $\phi = (\phi^+, \phi^0)$ is the SM scalar doublet field, and $L(\nu_i, l_i)$ are the SM lepton fields. Under the extended electroweak symmetry $SU(2)_L \times U(1)_Y \times Z_2$, the fields η, N, ϕ, and L are in the $(2, 1/2; -)$, $(1, 0; -)$, $(2, 1/2; +)$, and $(2, -1/2; +)$ representations, respectively. The inert and the standard doublet scalar also couple through the scalar potential [26],

$$\begin{aligned} V = {}& \mu_1^2 \Phi^\dagger \Phi + \mu_2^2 \eta^\dagger \eta + \lambda_1 (\Phi^\dagger \Phi)^2 + \lambda_2 (\eta^\dagger \eta)^2 \\ & + \lambda_3 (\Phi^\dagger \Phi)(\eta^\dagger \eta) + \lambda_4 (\Phi^\dagger \eta)(\eta^\dagger \Phi) \\ & + \tfrac{1}{2}\lambda_5 [(\Phi^\dagger \eta)^2 + \text{H.c.}]. \end{aligned} \tag{2.2}$$

In particular the last quartic coupling provides the mass splitting between the two physical inert neutral scalar particles $\eta_H = \sqrt{2}\,\text{Im}(\eta^0)$ and $\eta_L = \sqrt{2}\,\text{Re}(\eta^0)$ [25,26], which are the heaviest and the lightest for $\lambda_5 < 0$ (otherwise the two would be exchanged)

$$m_{\eta_H}^2 - m_{\eta_L}^2 = |\lambda_5| v^2. \tag{2.3}$$

The masses of the inert scalar bosons are

$$\begin{aligned} m_{\eta^+}^2 &= \mu_2^2 + \lambda_3 v^2/2, \\ m_{\eta_H}^2 &= \mu_2^2 + (\lambda_3 + \lambda_4 - \lambda_5) v^2/2\mu_2^2 + (\lambda_L - 2\lambda_5) v^2/2, \\ m_{\eta_L}^2 &= \mu_2^2 + (\lambda_3 + \lambda_4 + \lambda_5) v^2/2\mu_2^2 + \lambda_L v^2/2. \end{aligned} \tag{2.4}$$

Here $v/\sqrt{2} = 174$ GeV is the vacuum expectation value (VEV) of the SM Higgs field (the inert scalar does not acquire a VEV), $m_{\eta^\pm}$ is the mass of the charged scalars, λ_5 has been chosen to be real, and we define $\lambda_L = \lambda_3 + \lambda_4 + \lambda_5$. The only other terms in the Lagrangian allowed by the Z_2 symmetry are Majorana mass terms for the inert-sterile neutrinos,

$$\tfrac{1}{2} M_i N_i N_i + \text{H.c.} \tag{2.5}$$

The Z_2 symmetry forbids sterile-active neutrino mixings. The N_i's are not the Dirac mass partner of the ν_i as in usual extensions of the SM and active neutrino Majorana masses are generated at one-loop level. The active neutrino mass matrix elements are [25]

$$\begin{aligned} (\mathcal{M}_\nu)_{ij} = \sum_k h_{ik} h_{jk} \frac{M_k}{16\pi^2} \Bigg[& \frac{m_{\eta_H}^2}{m_{\eta_H}^2 - M_k^2} \ln\left(\frac{m_{\eta_H}^2}{M_k^2}\right) \\ & - \frac{m_{\eta_L}^2}{m_{\eta_L}^2 - M_k^2} \ln\left(\frac{m_{\eta_L}^2}{M_k^2}\right)\Bigg]. \end{aligned} \tag{2.6}$$

We will assume in what follows that m_{η_H} is of the order of 100 GeV and m_{η_L} of the order of tens of GeV; thus the first term in Eq. (2.6) is dominant.

The Z_2 parity implies that the lightest inert particle is stable and thus a good dark matter candidate. Both the lightest inert scalar [25,26,28–33] and the lightest sterile neutrino [25,34,35] could be dark matter candidates. We will assume the second possibility.

In Refs. [34,35] it was assumed that the mass difference between η_L and η_H is small, i.e. the coupling λ_5 is very small. In this case, in order to generate the observed active neutrino masses, the h_{ij} couplings cannot be very small. In addition, it was assumed that $m_0 = (m_{\eta_H}^2 + m_{\eta_L}^2)/2 > M_1$, M_2, M_3 and the lightest N_i is produced thermally. Under these assumptions, Ref. [34] found that the lightest inert-sterile neutrino can be CDM and account for the whole of the dark matter if its mass is in the range 7–300 GeV.

III. PARAMETER CONSTRAINTS

Here we will explore a range of values of the coupling constants different from those previously considered, namely, λ_5 not very small and h_{ij} Yukawa couplings small enough to ensure that the sterile neutrinos N_i are never in equilibrium in the early Universe. We will not study the flavor structure of the couplings h_{ij}, but only their order of magnitude. We call generically h_1, h_2, and h_3 the couplings of N_1, N_2, and N_3, respectively. We assume a hierarchy in the couplings, with $h_1 < h_2 \simeq h_3$. We also assume that only the lightest sterile neutrino, which we take to be N_1, is lighter than the lightest inert scalar η_L and hence, it is the dark matter candidate. The η_L particles are produced thermally in the early Universe and decouple when they are nonrelativistic. The subsequent late decay of the η_L produces the inert-sterile N_1 relic particles that now constitute the dark matter. In this scenario, depending on the mass, abundance, and lifetime of η_L, the N_1 can be either CDM or WDM and account for the whole of the dark matter with mass in the range of ~few keV to tens of GeV. We will show that all requirements on the model are fulfilled: active neutrino masses of the right order of magnitude are obtained, the upper bound on the $h_{i,j}$ from $\mu \to e\gamma$ is easy to fulfill, all N_i producing reactions in the early Universe are out of equilibrium, and the necessary relic density and decay rate of η_L for different values of the η_L and N_1 masses are obtained, while respecting all the collider and other bounds imposed on the model.

Let us see first how large the Yukawa couplings h must be to get reasonable values for the active neutrino masses, i.e. $(M_\nu)_{i,j} \simeq 10^{-1}$ eV. Using Eq. (2.6), and assuming that η_L is significantly lighter than η_H, that $M_{2,3}$ is of the same order of magnitude but larger than m_{η_H}, and that the contributions of N_2 and N_3 are dominant, we get

$$h_{2,3} \simeq 0.7 \times 10^{-5}\left(\frac{M_{2,3}}{100\ \mathrm{GeV}}\right)^{1/2}\left(\frac{100\ \mathrm{GeV}}{m_{\eta_H}}\right) \times \left[\ln\left(\frac{M_{2,3}^2}{m_{\eta_H}^2}\right)\right]^{-1/2}. \tag{3.1}$$

Notice that when m_{η_H} is large with respect to m_{η_L}, Eq. (2.3) implies that $m_{\eta_H} \simeq \sqrt{|\lambda_5|}v$. Moreover, when $M_{2,3}$ are larger than but similar to m_{η_H}, the logarithm in Eq. (3.1) is close to 1, thus

$$h_{2,3} \simeq \frac{3 \times 10^{-6}}{\sqrt{|\lambda_5|}}\left(\frac{M_{2,3}}{100\ \mathrm{GeV}}\right)^{1/2}. \tag{3.2}$$

Lepton flavor transitions like the $\mu \to e\gamma$ process in Fig. 1 occur in this model. The branching ratio, $B_{\mu\to e\gamma} = \Gamma_{\mu\to e\gamma}/\Gamma_{\mu\to e\nu\nu}$ in the inert doublet model is [34,36]

$$B_{\mu\to e\gamma} = \frac{192\pi^3\alpha}{G_F^2}\left(\left|\sum_j \frac{h_{\mu j}h_{ej}}{4(4\pi)^2 m_{\eta^-}^2} F_2\left(\frac{M_j^2}{m_{\eta^-}^2}\right)\right|\right)^2, \tag{3.3}$$

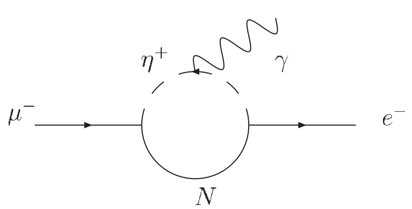

FIG. 1. Diagram for $\mu \to e\gamma$ transition in the variation of the inert doublet model we consider here.

where α is the fine structure constant and G_F is the Fermi constant. For $M_{2,3} \simeq m_{\eta^\pm}$ the function $F_2(x) \times [1 - 6x + 3x^2 + 2x^3 - 6x^2\ln(x)][6(1-x)^4]^{-1}$ is $F_2(1) \simeq 1/12$, whereas for $M_1 < m_{\eta^\pm}$ it is $F_2(0) \simeq 1/6$ [36]. The experimental upper bound on the branching ratio, $B(\mu \to e\gamma) \leq 1.2 \times 10^{-11}$ [37], implies

$$h_{2,3} \leq 2 \times 10^{-2}\left(\frac{m_{\eta^\pm}}{100\ \mathrm{GeV}}\right) \tag{3.4}$$

for $h_1 \ll h_{2,3}$.

Let us now see how small the couplings h_{ij} must be in order for the N_i to never be in equilibrium in the early Universe. The upper bounds are particularly important for N_2 and N_3, whose generic couplings, h_2 and h_3, are larger than the coupling h_1 of N_1. The N_i can be produced through the reactions in Fig. 2, i.e. two to two reactions $L\bar{L} \to N_iN_i$ mediated by any of the four physical inert particles η_H, η_L, η^+, and η^-, which we call now generically η, or $\eta\eta \to N_iN_i$ mediated by L. The N_i could in principle be produced through the decay $\eta \to N_iL$ and the inverse decay of $\eta L \to N_i$. The production rate for N_2, for example, is

$$\Gamma_{N_2} = \sum_L (2\langle\sigma v\rangle_{L\bar{L}\to N_2N_2} + \langle\sigma v\rangle_{L\bar{L}\to N_2N_3})n_L^2/n_{N_2}^{\mathrm{eq}} + \sum_\eta (2\langle\sigma v\rangle_{\eta\eta\to N_2N_2} + \langle\sigma v\rangle_{\eta\eta\to N_2N_3})n_\eta^2/n_{N_2}^{\mathrm{eq}}, \tag{3.5}$$

where $n_{N_2}^{\mathrm{eq}}$ is the N_2 equilibrium number density which appears in the equation as a normalization factor, and n_L and n_η are the number densities of the standard leptons and the inert scalars, respectively, at the temperature considered.

Equation (3.5) is derived from the Boltzmann equation for the production of N_i ($i = 1, 2, 3$) in the process $ab \to N_ic$, where a, b, and c are particles and we assume that only the initial particles, a and b, have initially a nonzero particle density. If the particles a and b are in equilibrium, assuming Maxwell-Botzmann density distributions, the time evolution of the number density n_{N_i} depends on the number densities of particles a and b in the following manner [see e.g. Eqs. (5.8) and (5.23) of Chap. 5 of Ref. [38]]:

$$\frac{dn_{N_i}}{dt} + 3Hn_{N_i} = \sum_{a,b,c} n_a n_b \langle\sigma_{ab\to N_ic}|v|\rangle. \tag{3.6}$$

As usual, it is convenient to change variables to $Y \equiv n_{N_i}/s$

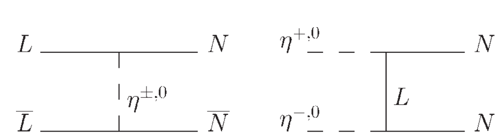

FIG. 2. Production processes of inert-sterile neutrinos in the early Universe.

and $x \equiv m_{N_i}/T$ [see Eq. (5.16) of Ref. [38]] and obtain

$$\frac{dY}{dx} = \frac{1}{Hxs} \sum_{a,b,c} n_a n_b \langle \sigma_{ab\to N_i c} |v| \rangle. \tag{3.7}$$

Now, dividing and multiplying the right-hand side of Eq. (3.7) by $n_{N_i}^{\rm eq}$ as a normalizing function, one gets

$$\frac{x}{Y^{\rm eq}} \frac{dY}{dx} = \frac{\Gamma_{N_i}}{H}, \tag{3.8}$$

where Γ_{N_i} is defined as in Eq. (3.5) above. Equation (3.7) is equivalent in this case to Eq. (5.26) of Ref. [38] and shows that Y_{N_i} is never significantly different from zero if $\Gamma_{N_i}/H < 1$.

Following Refs. [39,40] and making use of Refs. [41,42], for relativistic η, N_i, and L the thermal averaged cross sections of $L\bar{L} \to N_i N_i$ and of $\eta\eta \to N_i N_i$ are approximately given by

$$\langle \sigma v \rangle_{L\bar{L}\to N_i N_i} \simeq 0.7 \times 10^{-1} \frac{h_i^4}{T^2}, \tag{3.9}$$

$$\langle \sigma v \rangle_{\eta\eta\to N_i N_i} \simeq 3 \times 10^{-1} \frac{h_i^4}{T^2} \ln\left(\frac{4T^2}{M_i^2 + m_\eta^2}\right), \tag{3.10}$$

which show that the process $\eta\eta \to N_i N_i$ is dominant and

$$\Gamma_{N_i} \simeq \Gamma_{\eta\eta\to N_i N_i} \simeq 0.7 \times 10^{-1} h_i^4 T \ln\left(\frac{4T^2}{M_i^2 + m_\eta^2}\right). \tag{3.11}$$

The production is out of equilibrium if the rate is smaller than the expansion rate of the Universe, H,

$$\Gamma_{N_i} < H = 1.66\sqrt{g_*}\frac{T^2}{M_{\rm Pl}}, \tag{3.12}$$

where g_* is the number of degrees of freedom and $M_{\rm Pl}$ is the Planck mass. Since the right-hand side of Eq. (3.12) decreases faster than the left-hand side for decreasing T, if the condition is fulfilled for the smallest T value in the range considered, i.e. the smallest T for which all the particles involved in the production are relativistic, then it is fulfilled for all larger T.

At high temperatures $T > M_{2,3} \simeq m_{\eta_H}$ we need to write the condition in Eq. (3.12) at $T \simeq M_k \simeq m_{\eta_H}$. Thus, the production of relativistic $N_{2,3}$ is out of equilibrium at $T > M_{2,3} \simeq m_{\eta_H}$ if

$$h_{2,3} < 2 \times 10^{-4}\left(\frac{g_*}{106.75}\right)^{1/8}\left(\frac{M_{2,3}}{100\ \mathrm{GeV}}\right)^{1/4}. \tag{3.13}$$

Since we are assuming $M_1 < m_{\eta_L} \ll M_{2,3}$, the condition in Eq. (3.12) for relativistic N_1 and η_L must be taken at $T \simeq m_{\eta_L}$; thus the production of relativistic N_1 from relativistic η_L is out of equilibrium if

$$h_1 < 2 \times 10^{-4}\left(\frac{g_*}{106.75}\right)^{1/8}\left(\frac{m_{\eta_L}}{10\ \mathrm{GeV}}\right)^{1/4}. \tag{3.14}$$

At temperatures in the range $M_{2,3} > T > m_{\eta_L}$, in which the $N_{2,3}$ are nonrelativistic (but η_L and L are relativistic), the relevant thermal average cross sections for $N_{2,3}$ production are approximately

$$\langle \sigma v \rangle_{L\bar{L}\to N_i N_i} \simeq 0.8 \times 10^{-2} \frac{h_i^4}{T^2} \exp\left(-\frac{2M_i}{T}\right), \tag{3.15}$$

$$\langle \sigma v \rangle_{\eta\eta\to N_i N_i} \simeq 2 \times 10^{-1} \frac{h_i^4}{T^2} \exp\left(-\frac{2M_i}{T}\right). \tag{3.16}$$

The production is again dominated by the $\eta_L\eta_L \to N_i N_i$ process, thus

$$\Gamma_{N_i} \simeq \Gamma_{\eta_L\eta_L\to N_i N_i} \simeq 0.7 \times 10^{-1} h_i^4 \frac{T^{5/2}}{M_i^{3/2}} \exp\left(-\frac{M_i}{T}\right). \tag{3.17}$$

Because this rate decreases faster than H, if $\Gamma_{N_i} < H$ is fulfilled at $T = M_{2,3}$ where Γ is maximum within the T interval, the process will be out of equilibrium for lower values of T, and thus we obtain

$$h_{2,3} < 3 \times 10^{-4}\left(\frac{g_*}{106.75}\right)^{1/8}\left(\frac{M_{2,3}}{100\ \mathrm{GeV}}\right)^{1/4}. \tag{3.18}$$

For still lower temperatures $T < m_{\eta^{\pm,0}}$, for which all the inert bosons are nonrelativistic but the N_1 are relativistic, we need to verify that the N_1 are not produced thermally (recall we are assuming that $m_\eta^{\pm,0} > M_1$). In this case

$$\langle \sigma v \rangle_{\eta\eta\to N_1 N_1} \simeq \frac{3}{4\pi} h_1^4 \frac{M_1^2}{m_\eta^4}, \tag{3.19}$$

and

$$\Gamma_{N_1} \simeq \Gamma_{\eta\eta\to N_1 N_1} = 10^{-2} h_1^4 \frac{M_1^2}{m_\eta} \exp\left(-\frac{2m_\eta}{T}\right). \tag{3.20}$$

This rate decreases faster than H as T decreases; thus if $\Gamma_{N_1}/H < 1$ at the highest temperature in the range considered, $T = m_\eta$, the condition is fulfilled at any lower T. Thus,

$$h_1 < 3 \times 10^{-2}\left(\frac{g_*}{106.75}\right)^{1/8}\left(\frac{m_{\eta_L}}{10\ \mathrm{GeV}}\right)^{3/4}\left(\frac{\mathrm{MeV}}{M_1}\right)^{1/2}. \tag{3.21}$$

After considering all the required upper bounds on the h_{ij} Yukawa couplings, we conclude that Eq. (3.13) provides the most restrictive upper bound on the Yukawa couplings of the heaviest inert-sterile neutrinos, $h_{2,3}$, and they are compatible with the value assigned to $h_{2,3}$ in Eq. (3.1), which is necessary to account for the active neutrino masses. The most restrictive bound on h_1, the Yukawa couplings of the lightest inert-sterile neutrino N_1, will be given in Eq. (4.2) below and is derived from our requirement of a long enough lifetime of the lightest inert bosons η_L into N_1.

Let us now consider the decays of the $\eta^\pm$ and η_H into inert-sterile neutrinos. If $m_\eta^\pm > m_\eta^0 + m_W$, then the process $\eta^\pm \to \eta^0 + W$ can occur. The branching ratio of the decay mode $\eta^\pm \to N_i L$ with respect to the dominant $\eta^\pm \to \eta^0 + W$ mode is proportional to the ratio of the couplings h_i^2/g_W^2, where g_W is the weak coupling. Using the value of $h_{2,3}$ necessary to produce the active neutrino masses, given in Eq. (3.2) with $|\lambda_5| \simeq 0.2$, for example, $h_i^2/g_W^2 \simeq 10^{-10}$ $(M_i/100\ \text{GeV})$, which is negligible. Thus, the heavier inert-sterile neutrinos $N_{2,3}$ are not produced in the decays of the inert charged bosons. Neither is the lightest inert-sterile neutrino produced in these decays, since $h_1 \ll h_{2,3}$. If, instead $m_{\eta^\pm} < m_{\eta^0} + m_W$, the 3-body decay $\eta^\pm \to \eta^0 + L + \bar{L}$ dominates the decay of $\eta^\pm$; the branching ratio of $\eta^\pm \to N_i L$ then goes as $h_i^2/g_W^4 \simeq 10^{-11}$ $(M_i/100\ \text{GeV})$ for the heavier inert-sterile neutrinos. The branching ratio is even smaller for N_1. Again, the decay of the charged inert bosons into the inert-sterile neutrinos N_i is negligible. For the decays of the heavier neutral inert boson η_H the same arguments apply but changing the W's by Z's. Thus the inert-sterile neutrinos are not produced in the decays of $\eta^\pm$ and η_H.

IV. RELIC DENSITY ANALYSIS

We need to insure that η_L, the lightest inert scalar particle, is produced thermally in the early Universe and that it is in equilibrium before decoupling while it is already nonrelativistic, at freeze-out, $T_{\text{f.o.}} < m_{\eta_L}$. The dominant processes that maintain the η_L particles in equilibrium depend on the couplings of η_L with the SM particles. The η_L gauge couplings and its couplings in the scalar potential are the same that occur in the inert doublet model in the absence of sterile neutrinos. Using the same couplings, in Refs. [26,29,32] η_L with mass in the GeV range are found to be good dark matter candidates. We want instead for the η_L decay into the lightest inert-sterile neutrino N_1, which constitutes the dark matter now. After the η_L particles decay through the process $\eta_L \to N_1 \nu_i$, there is one N_1 per each η_L. In order for N_1 to account for the whole of the dark matter, the number density of η_L at their decoupling must be larger for the case considered here than in the scenarios in which they constitute the dark matter [26,29,32]. The number density n_{N_1} that is needed for nonrelativistic N_1 to be the dark matter at present must be the same relic number density n_{η_L} the η_L should have at present had they not decayed. Thus the relic density of N_1 is now $n_{N_1} M_1 n_{\eta_L} M_1$ and

$$\Omega_{N_1} h^2 = \Omega_{\eta_L} h^2 \left(\frac{M_1}{m_{\eta_L}}\right), \tag{4.1}$$

where $\Omega_{\eta_L} h^2$ is the relic density the η_L would have at present if they were stable. When the N_1 can be either

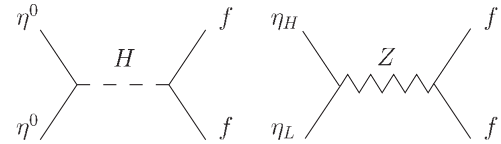

FIG. 3. Dominant η_L annihilation channels into standard model fermions f.

CDM or WDM we require the N_1 density to be that of the observed relic density of dark matter $\Omega_{\text{DM}} h^2 = 0.1099 \pm 0.0062$ [43]. If the N_1 are instead hot dark matter (HDM) we should impose the upper bound $\Omega_{N_1} h^2 \le 0.014 \equiv \Omega_{\text{HDM-max}} h^2$ (the 95% C.L. on the relic density of light neutrinos) [43].

If $m_{\eta_L} > m_W$, the η_L annihilate efficiently into two W bosons and their relic density is too small even to constitute the bulk of the dark matter; thus we are not interested in this mass range. When $m_{\eta_L} < m_W$, the processes in Fig. 3 and their inverse processes keep η_L in equilibrium. The lightest scalar η_L coannihilates with the heaviest inert scalar partner η_H. The coannihilation $\eta_H \eta_L \to f\bar{f}$ into SM fermions f is mediated by the Z boson and its cross section depends on the mass splitting $\Delta = m_{\eta_H} - m_{\eta_L}$ which in turn, depends on λ_5 [see Eq. (2.3)]. The η_L also coannihilates with $\eta^\pm$, via $W^\pm$ exchange, with a cross section which depends on the mass split between them. The process $\eta_L \eta_L \to f\bar{f}$ via Higgs exchange also keeps η_L in equilibrium, and in the particular range of masses we explore below is the dominant process. We use the public code MICROMEGAS [44] to compute the η_L relic density.

The decay $\eta_L \to N_1 L$ must happen after the η_L freeze-out at $T_{\text{f.o.}} = m_{\eta_L}/x_f$, where x_f is in the 20–30 range. Thus, the decay rate must be $\Gamma_{\eta \to N_1 L} \simeq h_1^2 m_{\eta_L}/16\pi < H$ for $T > T_{\text{decay}}$ and $\Gamma_{\eta \to N_1 L} \simeq H$ for $T = T_{\text{decay}}$ with $T_{\text{decay}} < T_{\text{f.o.}}$. These conditions lead to the most stringent bound on h_1

$$h_1 < 2 \times 10^{-9} \left(\frac{20}{x_f}\right) \left(\frac{m_{\eta_L}}{10\ \text{GeV}}\right)^{1/2} \left(\frac{g_*}{10.75}\right)^{1/4}. \tag{4.2}$$

Note that this bound on h_1 is consistent with the previous requirements.

We can now show that the inert-sterile neutrinos produced in this model may be either WDM or CDM, which are characterized by the freestreaming length λ_{fs} [45,46]

$$\lambda_{\text{fs}} = 2 r t_{\text{EQ}} (1 + z_{\text{EQ}})^2 \times \ln\left(\sqrt{1 + \frac{1}{r^2 (1 + z_{\text{EQ}})^2}} + \frac{1}{r(1 + z_{\text{EQ}})^2}\right). \tag{4.3}$$

Here the subscript EQ denotes matter-radiation equality and $r = a(t) p(t)/M_1$, where $a(t)$ and $p(t)$ are the scale factor of the Universe and the dark matter particle characteristic momentum at time t, respectively. As the Universe expands, the ratio r remains constant. At the time of matter-radiation equality, $\lambda_{\rm fs}$ must be 0.1 Mpc [47] for WDM, which fixes $r \simeq 10^{-7}$. At the moment of decay of the η_L (we make the approximation of instantaneous decays) the scale factor of the Universe is $a \simeq T_o/T_{\rm decay}$, where T_o is the photon temperature today, and the momentum of the relativistic N_1 decay products is $m_{\eta_L}/2$. Thus, $r \simeq T_o m_{\eta_L}/(2T_{\rm decay}M_1)$. Therefore, $r = 10^{-7}$ fixes the mass of N_1 to be

$$(M_1)_{\rm WDM} \simeq 2.4\ {\rm MeV}\left(\frac{m_{\eta_L}}{10\ {\rm GeV}}\right)\left(\frac{5\ {\rm MeV}}{T_{\rm decay}}\right). \qquad (4.4)$$

Given a particular $T_{\rm decay}$, Eq. (4.4) provides the N_1 mass for which the N_1 would constitute WDM. Heavier N_1 (smaller $\lambda_{\rm fs}$) would be CDM and lighter ones (larger $\lambda_{\rm fs}$) HDM.

We require the decay temperature to be $T_{\rm decay} \gtrsim 5$ MeV, in order not to affect the success of BBN predictions, and $T_{\rm decay} < m_{\eta_L}/x_f$, because the decays of η_L happen after they decouple. Thus, the range of masses for which N_1 could be a good WDM candidate is

$$24\ {\rm keV}\left(\frac{x_f}{20}\right) < (M_1)_{\rm WDM} < 2.4\ {\rm MeV}\left(\frac{m_{\eta_L}}{10\ {\rm GeV}}\right). \qquad (4.5)$$

Finally, in order to choose suitable sets of parameters there are a number of constraints that need to be considered. The null result for the process $e^+e^- \to Z^* \to \eta_H\eta_L$ in LEP II searches for neutralinos, and imposes the bound $m_{\eta_H} > 120$ GeV when $m_{\eta_L} < 80$ GeV [48]. Alternatively, in a range of parameters we will not explore, the neutral inert boson mass difference must be $m_{\eta_H} - m_{\eta_L} < 8$ GeV [48] for $m_{\eta_L} + m_{\eta_H} > m_Z$ due to the LEP I measurement of the Z width, which implies $m_{\eta_L} > 40$ GeV. In addition, the suitable set of parameters should also be within the allowed range provided by electroweak precision measurements [26,32].

There are also constraints on the λ couplings in the scalar potential, Eq. (2.2). Vacuum stability of the scalar potential imposes [26]

$$\lambda_{1,2} > 0, \qquad \lambda_2 < 1, \\ \lambda_3,\ \lambda_L - \lambda_5 - |\lambda_5| > -2\sqrt{\lambda_1\lambda_2}, \qquad (4.6)$$

and perturbativity of the scalar potential imposes [26]

$$\lambda_3{}^2 + (\lambda_L - \lambda_5)^2 + \lambda_5{}^2 < 12\lambda_1{}^2. \qquad (4.7)$$

In Figs. 4 and 5, we show regions of the $m_{\eta_L} - M_1$ plane in which N_1 has the right dark matter density for two different sets of parameters. The Higgs mass is $M_{\rm Higgs} = 160$ GeV, $m_{\eta_H} = 125$ GeV, and $m_{\eta^\pm}$130 GeV in Fig. 4 and the Higgs mass is $M_{\rm Higgs} = 500$ GeV, $m_{\eta_H} = 150$ GeV, and $m_{\eta^\pm}$300 GeV in Fig. 5. The upper panels of the figures show the bounds on λ_L obtained from vaccum stability (cross-hatched violet regions) and perturbativity (shaded

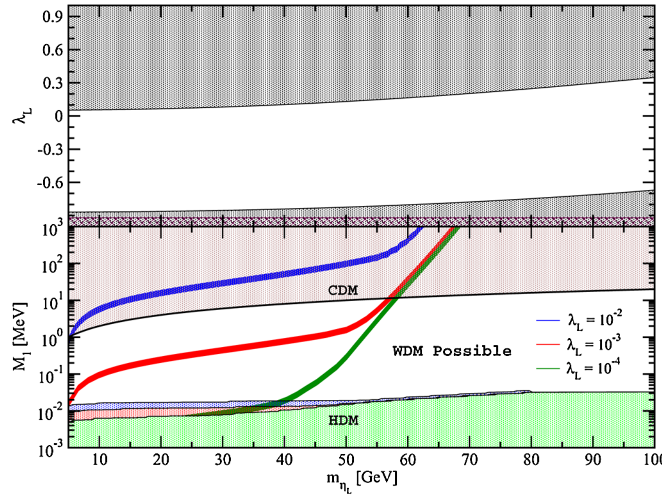

FIG. 4 (color online). In both panels, $M_{\rm Higgs} = 160$ GeV, $m_{\eta_H} = 125$ GeV, and $m_{\eta^\pm} = 130$ GeV. Upper panel: The shaded areas correspond to forbidden values of λ_L from vacuum stability (cross-hatched violet region) and perturbativity (shaded gray region) arguments. Lower panel: From top to bottom, the (blue, red, and green colored) narrow strips show the regions where N_1 would have the right dark matter density for the corresponding top to bottom values of λ_L given in the panel. The unshaded (middle) background region labeled "WDM Possible" corresponds to the range in Eq. (4.5), where the N_1 may constitute WDM (above it, N_1 can only be CDM and below it, only HDM). For any particular value of $T_{\rm decay}$ between 5 MeV (upper boundary of unshaded region) and m_{η_L}/x_f (lower boundary of unshaded region, which depends slightly on λ_L through x_f and is higher for larger λ_L) there is one value of M_1 given by Eq. (4.4), within the unshaded background region for which N_1 would be WDM (and it would be CDM for all larger values of M_1 and HDM for all smaller ones). In order for N_1 to be allowed as HDM, its mass must be at least a factor of $\Omega_{\rm DM}h^2/\Omega_{\rm HDM\text{-}max}h^2 0.1099/0.014 \simeq 8$ smaller than that corresponding to the center of the narrow (colored) bands for a given m_{η_L}.

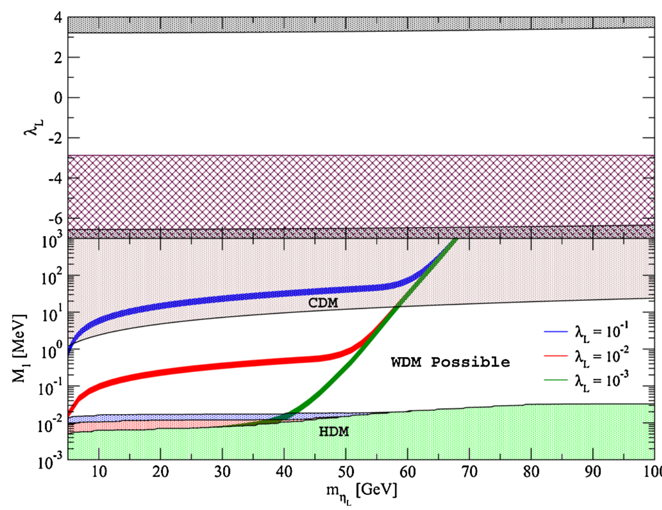

FIG. 5 (color online). Same as in Fig. 4 but for $M_{\rm Higgs}$500 GeV, $m_{\eta_H} = 150$ GeV, and $m_{\eta^\pm} = 300$ GeV.

gray region) arguments. From top to bottom, the (blue, red, and green colored) narrow strips in the lower panels of Figs. 4 and 5 show the regions in the $m_{\eta_L} - M_1$ plane in which $\Omega_{N_1} h^2$ in Eq. (4.1) is within the 3σ measured range for the dark matter (either CDM or WDM). The top, middle, and bottom strips corresponds, respectively, to the top, middle, and bottom values of λ_L shown in the panels. The unshaded background region labeled "WDM Possible" corresponds to the range in Eq. (4.4), where the N_1 may constitute WDM (above it, it can only be CDM and below it, only HDM). For any particular value of $T_{\rm decay}$ between 5 MeV (which defines the upper boundary of the unshaded region) and m_{η_L}/x_f (which defines the lower boundary of the unshaded region) there is one value of M_1 given by Eq. (4.4), within the unshaded background region for which N_1 would be WDM (N_1 would be CDM for all larger values of M_1 and HDM for all smaller ones). Notice that the lower boundary of the unshaded region depends on λ_L through x_f; it is slightly higher for higher values of λ_L (thus the blue, red, and green colors of the lower regions, for which the N_1 can only be HDM). Therefore, within the unshaded background region N_1 could be WDM or CDM, depending on $T_{\rm decay}$. For a given set of parameters defining the model (and hence a given $T_{\rm decay}$), in order for N_1 to be allowed as HDM, its mass M_1 must be, at least, a factor of $\Omega_{\rm DM} h^2/\Omega_{\rm HDM\text{-}max} h^2 = 0.1099/0.014 \simeq 8$ smaller than the value at the center of the narrow strips defined by Eq. (4.1) for a given m_{η_L} (for the corresponding values of λ_L). The figures show that the lightest inert-sterile neutrino could be HDM even for masses as large as $\sim$1 keV.

V. CONCLUSIONS

In conclusion, we have shown that inert-sterile neutrinos, produced nonthermally in the early Universe, could be a viable WDM or CDM candidate. They are virtually nondetectable in either direct or indirect dark matter searches because of their extremely weak couplings to SM particles. Thus, their existence could be revealed only by discovering other particles of the model in collider experiments. We should keep in mind that the dark matter may consist of an admixture of different types of particles and that particles undetectable in dark matter searches may be part of it. The existence of these particles could only be inferred from collider data, supplemented by the null results from dark matter searches or with results from these searches which find other detectable dark matter components with a density smaller than required to constitute the whole of the dark matter. Unveiling the nature of the dark matter necessarily requires the combination of collider and direct and indirect searches.

ACKNOWLEDGMENTS

We thank E. Dolle and S. Su for helpful discussions. This work was supported in part by the U.S. Department of Energy Grant No. DE-FG03-91ER40662, Task C at UCLA. S. P. R. is partially supported by the Portuguese FCT through CERN/FP/83503/2008 and CFTP-FCT UNIT 777, which are partially funded through POCTI (FEDER), and by the Spanish Grant No. FPA2008-02878 of the MCT by MICINN. G. G. and S. P. R. would like to thank the Aspen Center for Physics and CERN where part of this work took place. E. O. would also like to thank CERN for hospitality.

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