

Addressing the Hubble tension with sterile neutrino dark matter

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One of the promising dark matter (DM) candidates is a keV-scale sterile neutrino. In the early Universe, the observed relic of the sterile neutrino DM is generated via the Dodelson-Widrow mechanism. However, this production scenario is severely constrained by various astrophysical observations. Many nonstandard interactions between active (ν_a) and sterile (ν_s) neutrinos have been proposed to evade these astrophysical bounds. Here, we study sterile neutrinos in the context of a mass-varying scenario by coupling both active and sterile neutrinos to a scalar field. This novel mechanism opens up a new parameter space that generates the observed DM relic and addresses the Hubble tension. We find that the resulting parameter space can be fully probed by future x-ray missions.

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Introduction. Understanding the fundamental nature of dark matter (DM) remains one of the central open problems in astroparticle physics. Among the proposed candidates, sterile neutrinos with masses in the keV range have received significant attention [1–18]. In the simplest scenario, sterile neutrinos are produced in the early Universe through oscillations with active neutrinos prior to their decoupling and before the onset of big bang nucleosynthesis (BBN). This nonresonant production process is commonly referred to as the Dodelson-Widrow (DW) mechanism [1]. Sterile neutrinos are unstable and can undergo radiative decay into an active neutrino accompanied by a photon [19]. Such decays would generate monochromatic x-ray or γ -ray lines, depending on the sterile neutrino mass, from regions where dark matter is concentrated, such as galaxies and galaxy clusters. However, results from dedicated x-ray and γ -ray line searches place strong constraints on this scenario, creating significant tension with the DW production mechanism and excluding most of its viable parameter space [20–23].

To evade the experimental bound, thus opening up new viable parameter space, many nonstandard interactions (NSIs) have been proposed in the DW framework. These NSIs can be between active-active ($\nu_a - \nu_a$) neutrino [24–26], active-sterile ($\nu_a - \nu_s$) neutrino [27], or sterile-sterile ($\nu_s - \nu_s$) [28,29] neutrino sectors. In the presence of an

NSI, new scattering or annihilation channels open up in the active (ν_a) and/or sterile (ν_s) neutrino sector. So, the efficient production of the sterile neutrino from the active neutrino happens even for a small mixing angle.

Another unsolved problem in modern cosmology is Hubble tension (see, e.g., [30–32] for reviews). The tension arises from the persistent disagreement between determinations of the present-day expansion rate of the Universe, H_0 , inferred from early-Universe observations and those obtained from late-Universe measurements. *Planck* satellite measurement of cosmic microwave background (CMB) reported $H_0 = 67.4 \pm 0.5 \text{ km s}^{-1} \text{ Mpc}^{-1}$ [33]. In contrast, local distance-ladder measurements based on Cepheid-calibrated type Ia supernovae from the SH0ES Collaboration report a significantly higher value, $H_0 = 73.30 \pm 1.04 \text{ km s}^{-1} \text{ Mpc}^{-1}$ [34,35]. So, there is a $5\text{-}\sigma$ difference between the measured value of H_0 by *Planck* and SH0ES. To alleviate this tension, many modifications to the Λ CDM cosmology have been proposed [36–40]. One such solution is early dark energy (EDE) [41–45]. References [43,46] studied EDE by quadratically coupling the Standard Model (SM) neutrinos to a background scalar field. Because of this coupling, the ensemble of SM neutrinos, i.e., active neutrinos, creates an effective potential for the scalar, but the vacuum expectation value (VEV) of the scalar backreacts and gives active neutrinos an effective mass. These effects generate the required dynamics and can address the Hubble tension. This tension can also be eased by increasing ΔN_{eff} [47] via the decay of a sterile neutrino of mass $\mathcal{O}(30) \text{ MeV}$ [48] or $\mathcal{O}(150\text{--}450) \text{ MeV}$ [49,50] at BBN.

In this study, we assume that the background scalar field is coupled to both the sterile and active neutrinos. Hence, the sterile neutrino also gets an effective mass due to the VEV of the background scalar field. This feature is unique

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compared to other NSI or lepton-asymmetry models [2] for sterile neutrinos. In those models, the mass of the sterile neutrino remains the same throughout its cosmological history, from production until today. Whereas, in our case, the sterile neutrino mass varies as the background scalar field evolves over time. We have shown that, due to an increased effective sterile neutrino mass in the early Universe, a viable parameter space in the mixing-angle–mass plane opens up, evading existing experimental constraints. We further show that, in the allowed region of the sterile neutrino parameter space, the inferred value of H_0 is modified compared to the standard Λ CDM scenario and can reproduce the SH0ES-measured value.

We organize the rest of the Letter as follows. We describe our model in the next section and discuss the cosmological evolution of the scalar field and the sterile neutrino in the third section. The results are presented in the penultimate section, and we conclude in the final section.

Model setup. In this model, an active neutrino and a sterile neutrino are quadratically coupled with a real homogeneous scalar field ϕ . The interaction Lagrangian is given by

$$-\mathcal{L}_{\text{int}} \supset \epsilon \frac{m_{\nu_1}}{M_{\text{pl}}^2} \bar{\nu}_1 \nu_1 \phi^2 + \epsilon \frac{m_{\nu_4}}{M_{\text{pl}}^2} \bar{\nu}_4 \nu_4 \phi^2. \quad (1)$$

Here, ϵ is the coupling strength, M_{pl} is the Planck mass, and m_{ν_1} and m_{ν_4} are Dirac masses for active and sterile neutrinos, respectively. As we impose a $\mathbb{Z}_2$ symmetry ($\phi \rightarrow -\phi$) on the Lagrangian, the above dimension-five operator is the leading-order interaction term between the neutrinos and the scalar field ϕ ¹ [52]. In the absence of any NSIs in the neutrino sector, sterile and active neutrino flavors mix through the vacuum mixing angle θ .

Active and sterile neutrinos are produced on shell during their cosmological evolution. Therefore, using the equation of motion, effective neutrino masses can be written as [46]

$$\begin{aligned} m_{\nu_1, \text{eff}} &= m_{\nu_1} \left(1 + \epsilon \frac{\phi^2}{M_{\text{pl}}^2} \right), \\ m_{\nu_4, \text{eff}} &= m_{\nu_4} \left(1 + \epsilon \frac{\phi^2}{M_{\text{pl}}^2} \right). \end{aligned} \quad (2)$$

As seen from Eq. (2), a universal coupling ϵ rescales both the active and sterile neutrino masses by the same factor, leaving the vacuum mixing angle θ unchanged. Motivated by this observation, and to keep the model as minimal as possible, we assume the coupling ϵ to be the same for both the active and sterile neutrino sectors.

¹We assume that the scalar field self-interactions are negligible. For the effect of the quartic self-interaction term of the scalar field on the Hubble tension, see Ref. [51].

In an isotropic and homogeneous Friedmann-Lemaître-Robertson-Walker background, the Euler-Lagrange equation corresponding to the interaction Lagrangian in Eq. (1) yields the following equation of motion for the homogeneous scalar field:

$$\frac{d^2 \phi}{dt^2} + 3H(t) \frac{d\phi}{dt} + \left(m_\phi^2 + 2\epsilon \frac{m_{\nu_1}}{M_{\text{pl}}^2} \bar{\nu}_1 \nu_1 + 2\epsilon \frac{m_{\nu_4}}{M_{\text{pl}}^2} \bar{\nu}_4 \nu_4 \right) \phi = 0, \quad (3)$$

where m_ϕ is the bare mass of the scalar field ϕ , t is the proper time, and $H(t)$ is the total Hubble rate including contributions from active and sterile neutrinos and the scalar field. At the quantum level, the terms proportional to $\bar{\nu}_1 \nu_1$ and $\bar{\nu}_4 \nu_4$ are substituted by their corresponding expectation values evaluated in their respective neutrino backgrounds. This expectation value is given by (see, for example, Appendix A of Ref. [53])

$$\langle \bar{\nu}_i \nu_i \rangle = \int \frac{d^3 p}{(2\pi)^3} \frac{m_{\nu_i, \text{eff}}}{E_p} f(p) = \frac{\rho_{\nu_i} - 3P_{\nu_i}}{m_{\nu_i, \text{eff}}}, \quad (4)$$

with $i = 1, 4$ corresponding to active and sterile neutrinos, respectively. Here, ρ_{ν_i} and P_{ν_i} are the energy density and pressure of the i th neutrino species, respectively. Then, using Eq. (4), we can rewrite the effective potential in Eq. (3) as [43]

$$\begin{aligned} \frac{1}{\phi} \frac{\partial V_{\text{eff}}}{\partial \phi} &= m_\phi^2 + 2\epsilon \frac{m_{\nu_1}}{M_{\text{pl}}^2} \frac{\rho_{\nu_1} - 3P_{\nu_1}}{m_{\nu_1, \text{eff}}} \\ &\quad + 2\epsilon \frac{m_{\nu_4}}{M_{\text{pl}}^2} \frac{\rho_{\nu_4} - 3P_{\nu_4}}{m_{\nu_4, \text{eff}}}. \end{aligned} \quad (5)$$

As a result, the pressure and energy density of both neutrinos affect the evolution of the background scalar field through their contribution to the effective potential V_{eff} .

At any time t , the Hubble rate receives a contribution from the standard Λ CDM part plus an extra contribution due to the changes of effective masses of ν_1 and ν_4 for their interaction with the background scalar field ϕ , i.e.,

$$H^2 = H_{\Lambda\text{CDM}}^2 + \frac{8\pi}{3M_{\text{pl}}^2} \rho_{\text{extra}}, \quad (6)$$

where ρ_{extra} is

$$\begin{aligned} \rho_{\text{extra}} &= \rho_{\nu_1}(m_{\nu_1, \text{eff}}, T_{\nu_1}) - \rho_{\nu_1}(m_{\nu_1}, T_{\nu_1}) \\ &\quad + \rho_{\nu_4}(m_{\nu_4, \text{eff}}) - \rho_{\text{CDM}} + \frac{1}{2} \left(\frac{d\phi}{dt} \right)^2 + \frac{1}{2} m_\phi^2 \phi^2. \end{aligned} \quad (7)$$

Here, $\rho_{\nu_1}(m_{\nu_1, \text{eff}}, T_{\nu_1})$ and $\rho_{\nu_4}(m_{\nu_4, \text{eff}})$ are the modified active and sterile neutrino energy density due to their

effective masses, respectively. Also, $\rho_{\nu_1}(m_{\nu_1}, T_{\nu_1})$ is the active neutrino energy density in Λ CDM for $m_{\nu_1} = 0.05$ eV and ρ_{CDM} is the cold dark matter energy density. As $\rho_{\nu_1}(m_{\nu_1}, T_{\nu_1})$ and ρ_{CDM} are already included in $H_{\Lambda\text{CDM}}$, we subtracted those terms from Eq. (7).

$$\frac{\partial}{\partial t} f_s(p, t) - H p \frac{\partial}{\partial p} f_s(p, t) \approx \frac{1}{4} \frac{\Gamma_a(p) \Delta^2(p) \sin^2 2\theta}{\Delta^2(p) \sin^2 2\theta + (\Gamma_a(p))^2/4 + [\Delta(p) \cos 2\theta - V_T]^2} \times [f_a(p, t) - f_s(p, t)], \quad (8)$$

where $f_s(p, t)$ and $f_a(p, t)$ denote the sterile neutrino momentum distribution function and the active neutrino Fermi-Dirac distribution function, respectively. Here, $\Gamma_a(p)$ is the active neutrino forward interaction rate and $\Delta(p) = \frac{m_{\nu_4, \text{eff}}^2 - m_{\nu_1, \text{eff}}^2}{2p}$. The effective matter-mixing angle is given by

$$\sin^2 2\theta_m = \frac{\Delta^2(p) \sin^2 2\theta}{\Delta^2(p) \sin^2 2\theta + [\Delta(p) \cos 2\theta - V_T]^2}. \quad (9)$$

In the SM, $\Gamma_a(p)$ is obtained by summing over all leptonic and hadronic channels, and it is expressed as $\Gamma_a(p) = C_a(T) G_F^2 p T^4$ [4,58,59], where $C_a(T)$ is a temperature-dependent function. We have numerically evaluated $\Gamma_a(p)$ from Ref. [60]. The SM thermal potential V_T is given by [61–63]

$$V_T(p) = -\frac{8\sqrt{2}G_F p}{3m_W^2} \left[(\rho_{l_a} + \rho_{\bar{l}_a}) + \frac{m_W^2}{m_Z^2} (\rho_{\nu_a} + \rho_{\bar{\nu}_a}) \right], \quad (10)$$

where $\rho_{l_a}(\rho_{\bar{l}_a})$ is the energy density of the a th flavor charged lepton (l_a) [antilepton ($\bar{l}_a$)], $\rho_{\nu_a}(\rho_{\bar{\nu}_a})$ is the energy density of the a th flavor active neutrino, and p is the active

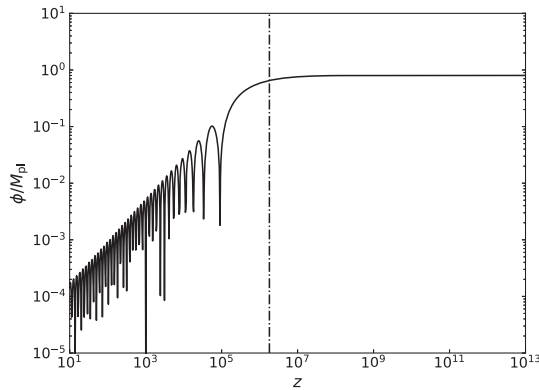


FIG. 1. Evolution of the scalar field with respect to the redshift z for $\phi_0/M_{\text{pl}} = 0.8$ and coupling, $\epsilon = 1.56 \times 10^3$. Here, the sterile neutrino mass is set to 10 keV, and the effective mass is 10 MeV. The vertical dash-dotted line corresponds to $z_{\text{osc}} = 1.81 \times 10^6$.

The production of sterile neutrinos proceeds via oscillations between active and sterile neutrino states in the early Universe, prior to BBN, through the so-called Dodelson-Widrow mechanism [1]. The evolution of the sterile neutrino phase-space distribution is described by a semi-classical Boltzmann equation of the form [3,4,54–57]

neutrino momentum. As the NSI terms are suppressed by M_{pl}^2 , the NSI scattering or annihilation channel will be highly suppressed compared to the SM weak interaction channels. Hence, the contribution from the NSI terms in $\Gamma_a(p)$ or V_T is strongly suppressed.

Cosmological evolution of the sterile neutrino. In the early Universe at some high temperature ($T \gtrsim 100$ GeV), the scalar field ϕ is displaced from the minimum of its potential to some value ϕ_0 , close to the Planck scale. This displacement of the scalar field will come from some UV dynamics. Since during this period $H^2 \gg \frac{1}{\phi} \frac{\partial V_{\text{eff}}}{\partial \phi}$, the evolution of ϕ , governed by Eq. (3), is highly overdamped leaving the scalar field frozen at ϕ_0 (see Fig. 1). So, the effective mass of the sterile neutrino is constant initially and is given by $m_{\nu_4, \text{eff}}^{\text{max}} = m_{\nu_4} \left(1 + \epsilon \frac{\phi_0^2}{M_{\text{pl}}^2} \right)$. For a given ϵ and ϕ_0 , $m_{\nu_4, \text{eff}}^{\text{max}}$ is the maximum effective mass of the sterile neutrino. With this enhanced effective mass, the sterile neutrino (ν_4) is produced from the active neutrino (ν_1) via the DW mechanism. We calculate the number density of the sterile neutrino (n_{ν_4}) by numerically solving Eq. (8) after modifying the publicly available STERILE-DM package [60]. In Fig. 2, we show the evolution of the sterile neutrino yield,

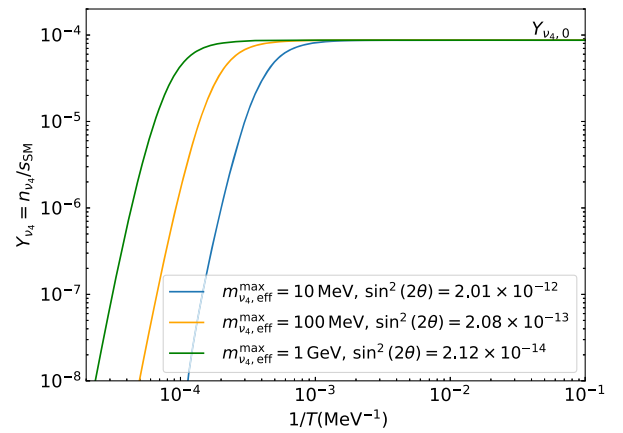


FIG. 2. Yield of sterile neutrino for present-day sterile mass of $m_{\nu_4} = 10$ keV. Increasing $m_{\nu_4, \text{eff}}^{\text{max}}$ by one order reduces $\sin^2 2\theta$ by one order.

$Y_{\nu_4} \equiv n_{\nu_4}/s_{\text{SM}}$, where s_{SM} is the entropy density of the background radiation, for temperatures ranging from 50 GeV to 10 MeV for present-day sterile neutrino mass $m_{\nu_4} = 10$ keV. Blue, orange, and green contours correspond to early-Universe sterile neutrino masses, $m_{\nu_4, \text{eff}}^{\text{max}} = 10$ MeV, 100 MeV, and 1 GeV, respectively. From Fig. 2 we observe that increasing $m_{\nu_4, \text{eff}}^{\text{max}}$ triggers the production of the sterile neutrino at increasingly high temperatures. This happens because $\Delta(p)$ is proportional to the effective sterile neutrino mass ($m_{\nu_4, \text{eff}}$). As a result, increasing $m_{\nu_4, \text{eff}}$ extends the temperature interval over which $\Delta(p)$ dominates over the thermal potential (V_T), relative to the standard DW scenario. For a fixed m_{ν_4} , this yields a larger effective matter-mixing angle (θ_m), see Eq. (9), compared to its value in the standard DW scenario. This effect leads to an enhanced production of sterile neutrinos at high temperatures. The sterile neutrino production reaches a maximum at temperature [1]

$$T_{\text{max}} \approx 133 \text{ MeV} \left(\frac{m_{\nu_4, \text{eff}}^{\text{max}}}{1 \text{ keV}} \right)^{1/3}. \quad (11)$$

For $T \gg T_{\text{max}}$, thermal potential V_T dominates over $\Delta(p)$. Hence, using Eq. (8) we deduce that $Y_{\nu_4} \propto (m_{\nu_4, \text{eff}}^{\text{max}})^4 \sin^2(2\theta)/T^9$. At $T \ll T_{\text{max}}$, sterile comoving number density freezes out and the yield Y_{ν_4} becomes constant until today. The freeze-out occurs before the BBN epoch in our parameter space. The sterile neutrino relic density can then be computed using the relation $\Omega_{\text{DM}} = Y_{\nu_4, 0} s_0 m_{\nu_4} / \rho_0$, where $Y_{\nu_4, 0}$, s_0 , and ρ_0 are the sterile neutrino yield, entropy density, and critical energy density today, respectively. Utilizing Eq. (11), the sterile neutrino relic density is

$$\Omega_{\text{DM}} \propto \left(\frac{m_{\nu_4, \text{eff}}^{\text{max}}}{\text{keV}} \right) \left(\frac{m_{\nu_4}}{\text{keV}} \right) \sin^2(2\theta). \quad (12)$$

Therefore, to satisfy the observed DM relic density, $\omega_{\text{DM}} = \Omega_{\text{DM}} h^2 = 0.12$ [33], we can accommodate a smaller mixing between the active and sterile neutrino than in the standard DW scenario, by increasing $m_{\nu_4, \text{eff}}^{\text{max}}$. From Fig. 2, we see that, for a fixed m_{ν_4} , a change in $m_{\nu_4, \text{eff}}^{\text{max}}$ can be compensated by inversely tuning $\sin^2 2\theta$.

The ϕ starts oscillating when the effective potential of the scalar field dominates over the Hubble friction term, i.e., $\frac{1}{\phi} \frac{\partial V_{\text{eff}}}{\partial \phi} \gg H^2$. As the active neutrino (ν_1) is relativistic during the entire period that is relevant to our analysis, it behaves as radiationlike, i.e., $P_{\nu_1} \approx \frac{1}{3} \rho_{\nu_1}$. The mass of the scalar field should be $m_\phi \lesssim H(z_{\text{CMB}}) \approx 10^{-29}$ eV such that the ϕ decays by the time of CMB and it has a negligible contribution to the Hubble parameter. Therefore, we set $m_\phi = 0$ in our numerical analysis. So, from Eq. (5) we see that $\frac{\partial V_{\text{eff}}}{\partial \phi}$ is dominated by the contribution from the sterile

neutrino (ν_4). The sterile neutrino becomes nonrelativistic near the BBN epoch in our region of interest. Hence, ignoring the pressure term for the sterile neutrino in Eq. (5), the redshift at which the scalar field starts oscillating, determined by equating $\frac{1}{\phi} \frac{\partial V_{\text{eff}}}{\partial \phi} = H^2$, is given by

$$z_{\text{osc}} \simeq \frac{3\epsilon \omega_{\text{DM}}}{4\pi \omega_{\text{rad}}}, \quad (13)$$

where ω_{DM} and ω_{rad} are the present-day relic density of DM and radiation, respectively. In deriving Eq. (13) we have assumed oscillation starts during the radiation-dominated epoch. From Fig. 1, we see that the value of z_{osc} , represented by the vertical dash-dotted line, is of the order of 10^6 . Amplitude of the oscillation of ϕ , calculated using the WKB approximation [52], is

$$\phi_{\text{amp}} \propto \left(\frac{1}{\phi} \frac{\partial V_{\text{eff}}}{\partial \phi} \right)^{-1/4} (1+z)^{3/2} \propto (1+z)^{3/4}. \quad (14)$$

Hence, the kinetic energy density of the scalar field redshifts away as $\frac{1}{2} \left(\frac{d\phi}{dt} \right)^2 \propto (1+z)^{9/2}$ during matter domination (see Supplemental Material for details [64]), which is faster than any other components of the Universe. So, the scalar field will not have any relic in the present day.

The sterile neutrino energy density is greater than the cold dark matter energy density in the Λ CDM scenario by a factor of $m_{\nu_4, \text{eff}}^{\text{max}}/m_{\nu_4}$ in the early Universe [see Eq. (S1) in the Supplemental Material [64]]. This extra energy density before photon decoupling changes the sound horizon at the surface of last scattering, altering the inferred value of H_0 . This change can be estimated as [41,46]

$$1 + \frac{\delta H_0}{H_0} \approx \frac{\int_{z_{\text{CMB}}}^{\infty} dz (\rho_{\Lambda\text{CDM}} (1 + R(z)))^{-1/2}}{\int_{z_{\text{CMB}}}^{\infty} dz ((\rho_{\Lambda\text{CDM}} + \rho_{\text{extra}}) (1 + R(z)))^{-1/2}} \times \frac{\int_0^{z_{\text{CMB}}} dz (\rho_{\Lambda\text{CDM}} + \rho_{\text{extra}})^{-1/2}}{\int_0^{z_{\text{CMB}}} dz (\rho_{\Lambda\text{CDM}})^{-1/2}}, \quad (15)$$

where $R(z) = \frac{3}{4} \frac{\omega_b}{\omega_\gamma} (1+z)^{-1}$ and $\omega_b = 2.24 \times 10^{-2}$, $\omega_\gamma = 2.47 \times 10^{-5}$ are the present-day baryon and photon relic densities in Λ CDM, respectively. We have taken the redshift of recombination to be $z_{\text{CMB}} = 1080$ and the central value of H_0 in Λ CDM to be $67.4 \text{ km s}^{-1} \text{ Mpc}^{-1}$.

Results. In Fig. 3, we delineate the parameter space for the scalar field ϕ that reproduces the SH0ES measurement of the Hubble parameter. In each contour line, $m_{\nu_4, \text{eff}}^{\text{max}}/m_{\nu_4}$ is constant with $m_{\nu_4} = 10$ keV and satisfies the DM relic for a specific value of $\sin^2 2\theta$. On each contour, the mixing angle (θ) required to satisfy the DM relic density is smaller as we increase the ratio of $m_{\nu_4, \text{eff}}^{\text{max}}/m_{\nu_4}$. For example, increasing $m_{\nu_4, \text{eff}}^{\text{max}}/m_{\nu_4}$ from 10^3 to 10^4 requires reducing the $\sin^2 2\theta$

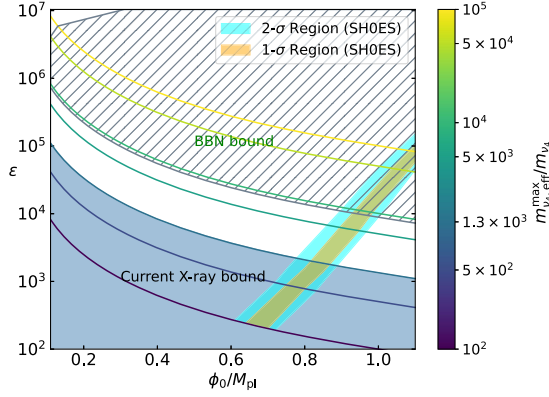


FIG. 3. Scalar field parameter space for $m_{\nu_4} = 10$ keV. Each contour line corresponds to different $m_{\nu_4,eff}^{max}/m_{\nu_4}$ and satisfies the DM relic, $\Omega_{DM}h^2 = 0.12$, for specific values of $\sin^2 2\theta$. The H_0 value increases as we move from left to right along each contour. We have marked the SH0ES 1- σ and 2- σ regions in orange and cyan, respectively. The region below $m_{\nu_4,eff}^{max}/m_{\nu_4} = 1.3 \times 10^3$ (blue-shaded area) is excluded by x-ray observations. The hatched region above $m_{\nu_4,eff}^{max}/m_{\nu_4} = 8.8 \times 10^3$ is excluded by BBN constraints.

from 2.01×10^{-12} to 2.08×10^{-13} , in order to satisfy relic density. As a result, in our scenario, for a fixed sterile neutrino mass, the x-ray bounds on the sterile neutrino can be evaded as we require a smaller mixing angle compared to the standard DW case.

Along each contour, moving from left to right, the initial scalar field value ϕ_0 increases while the coupling ϵ decreases. Consequently, according to Eq. (13), the oscillation redshift z_{osc} decreases, implying that the scalar field begins to oscillate at a later epoch. As a result, the extra energy density ρ_{extra} contributes more to the total energy density at later times, especially close to matter-radiation equality, leading to a larger expansion rate prior to the recombination. This reduces the sound horizon and, through Eq. (15), yields a larger inferred value of the present-day Hubble parameter H_0 . For every point in the (ϵ, ϕ_0) parameter space, we compute the corresponding value of H_0 using Eq. (15). The orange and cyan regions in Fig. 3 denote the parameter space for which the predicted H_0 lies within the SH0ES 1 σ ($72.26 \leq H_0 \leq 74.34$ km s $^{-1}$ Mpc $^{-1}$) and 2 σ ($71.22 \leq H_0 \leq 75.38$ km s $^{-1}$ Mpc $^{-1}$) intervals, respectively, corresponding to the SH0ES determination $H_0 = 73.30 \pm 1.04$ km s $^{-1}$ Mpc $^{-1}$.

Using the publicly available package ALTERBBN [65,66], we determine that the maximum allowed extra DM energy density at the time of BBN ($T = 1$ MeV) is 10^{-2} times the photon density at BBN. DM densities larger than this value are excluded at the 2 σ (95% confidence level) level by the BBN constraints from the primordial helium mass fraction Y_p and the deuterium-to-hydrogen ratio ($\frac{D}{H}$) $_p$. The values of Y_p and ($\frac{D}{H}$) $_p$ are taken to be 0.245 ± 0.003 [67–73] and $(25.47 \pm 0.29) \times 10^{-6}$ [67,74–80], respectively. In the

Λ CDM scenario, the ratio of DM density to photon density at BBN ($T = 1$ MeV) is of the order of 10^{-6} . In our scenario, at BBN, the sterile neutrino DM density is $m_{\nu_4,eff}^{max}/m_{\nu_4}$ times larger than its value in the Λ CDM scenario (see Supplemental Material for details [64]). So, in our case, this sets a maximum allowed value of $m_{\nu_4,eff}^{max}/m_{\nu_4}$ at BBN to be 8.8×10^3 . We represent this bound by the hatched region in Fig. 3. One can evade this BBN bound by decaying the scalar field before BBN by increasing the coupling ϵ such that z_{osc} becomes greater than the redshift of BBN. In Fig. 3, in the white region at the top left corner, the decay of ϕ happens before BBN. However, for these values of ϵ and ϕ_0 , the Hubble tension cannot be alleviated for any m_{ν_4} .

We depict the resulting sterile neutrino parameter space in the $\sin^2 2\theta - m_{\nu_4}$ plane in Fig. 4. The solid black line is the Dodelson-Widrow line corresponding to the absence of new nonstandard interactions. In the presence of a new interaction given in Eq. (1) viable parameter space is opened up without violating current x-ray bounds from Chandra [20], XMM-Newton [21], NuSTAR [22], and XRISM [23] telescopes. The x-ray constraint is marked by the blue region. The black dotted line and gray solid line correspond to the relic lines for $m_{\nu_4,eff}^{max}/m_{\nu_4} = 10^3$ and 8.8×10^3 , respectively. As noted earlier, the more we increase the $m_{\nu_4,eff}^{max}/m_{\nu_4}$ ratio, the less mixing we need between the active and sterile neutrino sector. The hatched region in Fig. 4 is excluded by BBN. The white region of Fig. 4 corresponds to the parameter space in which the modified background expansion history yields the present-day value of the Hubble parameter, H_0 , as measured by the SH0ES Collaboration. The light yellow region is excluded

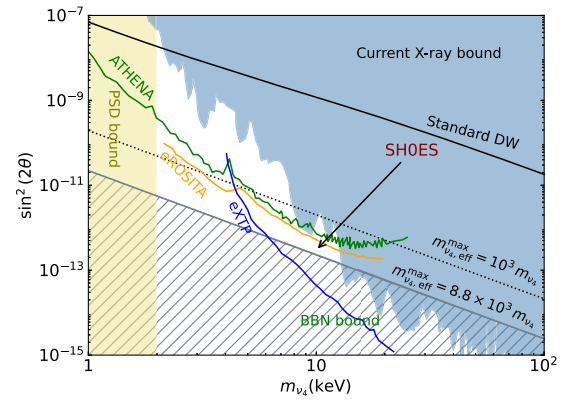


FIG. 4. Sterile neutrino parameter space in the presence of our new interaction term. In the white region, DM relic and Hubble tension both can be satisfied. The black dotted and gray solid lines correspond to relic line for $m_{\nu_4,eff}^{max}/m_{\nu_4} = 10^3$ and 8.8×10^3 , respectively. X-ray and phase-space density (PSD) bounds are denoted by blue and light yellow regions, respectively. The sensitivity curves of future x-ray missions [27] are shown in green (ATHENA), in blue (eXTP), and in orange (eROSITA). The black solid line is the standard Dodelson-Widrow line.

by the phase-space density constraint applicable to fermionic dark matter candidates, derived from observations of dwarf galaxies [81–83]. Following Ref. [27], we also add future sensitivities of upcoming missions of ATHENA (green line) [84], eROSITA (orange line) [85], and eXTP (blue line) [86]. The parameter region of interest lies in the sweet spot of the sensitivities of future missions.

Conclusion. We have studied the production of sterile neutrino dark matter before BBN via the DW mechanism in the presence of a new interaction term given in Eq. (1). This interaction universally couples sterile and active neutrinos to a real homogeneous scalar field (ϕ) quadratically and generates the observed relic. The scalar field gives an effective mass ($m_{\nu_4, \text{eff}}$) to the sterile neutrino. Because of this enhanced effective mass, production of the sterile neutrino from the active neutrino happens at high temperature. This results in a smaller vacuum mixing angle ($\sin^2 2\theta$) needed for the DM relic generation compared to the standard DW scenario. We have used this feature to allow a smaller $\sin^2 2\theta$ by increasing $m_{\nu_4, \text{eff}}^{\text{max}}$ in the parameter space of the sterile neutrino. In the standard DW scenario, this region of smaller $\sin^2 2\theta$ yields an underabundant DM relic.

In addition to that, we have demonstrated that the modified background cosmological evolution in this region of the sterile neutrino parameter space can yield the present-day value of the Hubble parameter, H_0 , consistent with the SH0ES measurement for suitable choices of the coupling ϵ and the initial scalar field value ϕ_0 . This enhancement arises because the sterile neutrino dark matter contributes an energy density larger than that in the standard Λ CDM cosmology by a factor of $m_{\nu_4, \text{eff}}/m_{\nu_4}$ during the early Universe. Consequently, the Hubble expansion rate prior to the recombination epoch increases, reducing the sound horizon at photon decoupling and leading to a larger inferred present-day Hubble parameter.

Moreover, an increased Hubble rate during BBN leads to a larger neutron-to-proton ratio at freeze-out, thereby

enhancing the primordial helium and deuterium abundances beyond the observed limits [67]. The corresponding constraint on $\sin^2 2\theta$ translates into a relic-density contour defined by $m_{\nu_4, \text{eff}}^{\text{max}}/m_{\nu_4} = 8.8 \times 10^3$. After imposing the astrophysical constraints, the phase-space density bound for fermionic dark matter, and the BBN constraint, we identify a parameter region where sterile neutrino masses in the range 2–13 keV and vacuum mixing angle $\sin^2 2\theta \sim 10^{-8} - 10^{-13}$, yield the background value of H_0 compatible with the SH0ES measurement. Interestingly, this entire region will be probed by upcoming x-ray missions such as ATHENA [84], eROSITA [85], and eXTP [86].

We emphasize, however, that the present analysis is restricted to the background cosmological evolution. A quantitative assessment of whether the model alleviates the Hubble tension requires a rigorous cosmological likelihood analysis that includes CMB anisotropies, baryon acoustic oscillations, type Ia supernovae, and large-scale structure data. Such an analysis necessitates implementing the model in a Boltzmann equation solver, such as CLASS [87] or CAMB [88], and is left for future work.

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