

Opening up new parameter space for sterile neutrino dark matter

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Sterile neutrinos are compelling dark matter (DM) candidates, yet the minimal production mechanism solely based on active-sterile ($\nu_a - \nu_s$) oscillations is robustly excluded by astrophysical observations. Nonstandard self-interactions in either the active ($\nu_a - \nu_a$) or sterile ($\nu_s - \nu_s$) sector are known to alter the sterile neutrino DM production in the early Universe, which could alleviate the tension with astrophysical constraints. Here, we propose a new solution where scalar-mediated nonstandard interactions between active and sterile neutrinos ($\nu_a - \nu_s$) lead to new production channels for sterile neutrino DM, independent of the active-sterile mixing and without the need for any fine-tuned resonance or primordial lepton asymmetries. This framework enables efficient sterile neutrino DM production even at vanishingly small mixing angles and opens up new viable regions of parameter space that can be tested with future x-ray and gamma-ray observations. It also provides a new mechanism for observable neutrino-DM interactions while being consistent with DM relic density and has broad phenomenological implications for astrophysical neutrinos and beyond.

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Introduction. The nature of dark matter (DM) remains an open question of fundamental importance in physics. As the so-called weakly interacting massive particle (WIMP) miracle with GeV-TeV-scale thermal freeze-out DM loses its appeal due to null experimental searches over the past several decades [1], there is growing interest in alternatives to the WIMP paradigm [2], such as WIMP-less DM [3], freeze-in DM [4], axions [5], sterile neutrinos [6], wavelike DM [7], and primordial black holes [8]. Here, we focus on keV-scale sterile neutrinos, which are particularly well-motivated candidates for (warm) DM due to their potential connection to other outstanding puzzles like neutrino mass and matter-antimatter asymmetry [9–12].

In a minimal extension of the Standard Model (SM), sterile neutrinos are produced nonthermally in the early Universe through active-sterile ($\nu_a - \nu_s$) oscillations. This nonresonant production, known as the *Dodelson-Widrow (DW) mechanism* [13], is theoretically attractive since it relies solely on neutrino mixing and requires no new particles beyond sterile neutrinos. However, achieving the correct DM relic abundance in this scenario requires relatively “large” mixing angles that are strongly constrained by observational data, mostly from x-ray and gamma-ray line searches for the radiative decay $\nu_s \rightarrow \nu_a + \gamma$ [14–22], effectively excluding the entire DW parameter space; see Fig. 1.

A possible alternative to the DW mechanism is the *Shi-Fuller mechanism* [23], in which a primordial lepton asymmetry ($\eta_L \sim 10^{-3}$ to 10^{-1}) resonantly enhances the sterile neutrino production. However, the required lepton asymmetry to yield efficient DM production is several orders of magnitude larger than the observed baryon asymmetry ($\eta_B \sim 10^{-10}$), challenging its consistency with standard cosmology and big bang nucleosynthesis (BBN) [24–26], although flavor-specific lepton asymmetries can mitigate the BBN bounds to some extent [27].

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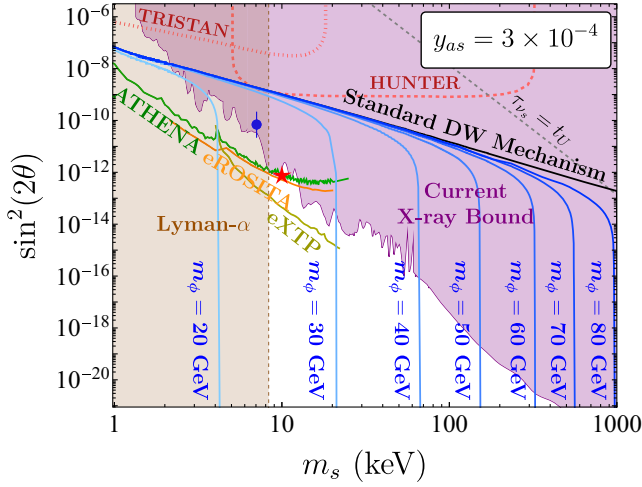


FIG. 1. Sterile neutrino mass versus mixing angle for a fixed active-sterile neutrino coupling $y_{as} = 3 \times 10^{-4}$. Shaded regions indicate astrophysical constraints. The blue contours satisfy the correct DM relic abundance $\Omega_s h^2 = 0.12$ for different values of the mediator mass m_ϕ . The black line corresponds to the correct relic density using the DW mechanism. The red star marks a benchmark point (BP) used later to illustrate the yield behavior. The blue point corresponds to the purported 3.5 keV x-ray line signal. See text for details.

These challenges motivate the need for alternative sterile neutrino DM production mechanisms, involving nonstandard interactions (NSIs) in the neutrino sector [28] or using heavy particle decays [29–31]. Prior studies using NSI have explored the effect of self-interactions between the active neutrinos ($\nu_a - \nu_a$) [32–35] or between the sterile neutrinos ($\nu_s - \nu_s$) [36–38] on the DM production. In this Letter, we explore an intriguing new possibility, i.e., NSI between active and sterile neutrinos ($\nu_a - \nu_s$). While active-sterile NSI has been previously explored in the contexts of BBN [39] and high-energy neutrinos [40,41], its role in sterile neutrino DM production has not been studied before.

As summarized in Fig. 1, active-sterile neutrino interactions lead to unique features very different from the active-active and sterile-sterile self-interaction scenarios. In particular, the active-sterile NSIs introduce new number-changing processes, such as $\nu_a \nu_a \leftrightarrow \nu_s \nu_s$, providing an efficient sterile neutrino production channel that differs fundamentally from the DW mechanism, as it does not rely on the active-sterile mixing. As a result, it extends the accessible parameter space to arbitrarily small mixing angles. It significantly improves upon earlier self-interaction-based production scenarios, which necessarily depend on the mixing and underproduce sterile neutrinos below a threshold mixing angle. Additionally, unlike scenarios involving sterile self-interactions, this mechanism does not rely on any fine-tuned resonance to achieve efficient DM production.

Active-sterile NSI. We consider a keV-scale Majorana sterile neutrino, which mixes with active neutrinos with

strength θ and also participates in NSI via a complex scalar mediator, assumed to have mass $m_\phi \gtrsim \mathcal{O}(\text{GeV})$. The relevant effective interaction Lagrangian is given by

$$-\mathcal{L} \supset y_{as} \bar{\nu}_a \nu_s \phi + \text{H.c.}, \quad (1)$$

where we have taken the coupling constant y_{as} to be real.¹ We have also suppressed the active neutrino flavor index, as the DM production happens above the neutrino decoupling and all flavors lead to qualitatively similar results, although some of the observational constraints are flavor dependent.

The effective active-sterile interaction (1) can be motivated by models of DM-neutrino interactions from a (sterile) neutrino portal [42–53]. While self-interaction terms such as $y_{ss} \bar{\nu}_s \nu_s \phi$ and $y_{aa} \bar{\nu}_a \nu_a \phi$ may also be present in a concrete, gauge-invariant ultraviolet completion of this effective interaction [48], we set these additional Yukawa couplings to zero as a working assumption in order to isolate the effect of active-sterile NSI. Moreover, one can always forbid these extra terms by invoking a Z_2 symmetry under which the SM particles are uncharged, while ν_s and ϕ have -1 charges, so that only the $\bar{\nu}_s \phi$ term is allowed, while $\bar{\nu} \nu \phi$ and $\bar{\nu}_s \nu_s \phi$ terms are not. We further assume that the ϕ field does not acquire a vacuum expectation value, and therefore there is no induced active-sterile neutrino mixing due to the interaction (1).

For concreteness, we adopt a minimal setup with no additional interactions for ϕ or ν_s and assume vanishing primordial abundances so that both are populated entirely through their coupling to active neutrinos. For $m_\phi \gtrsim \mathcal{O}(\text{GeV})$, the decay rate of $\phi \rightarrow \nu_a \nu_s$ is much faster than the relevant production timescales at $\mathcal{O}(\text{MeV})$, provided the coupling satisfies $y_{as} \gtrsim 10^{-8}$. Under this condition, the evolution of the mediator ϕ can be safely neglected in the “heavy mediator” limit, which we consider here.

Since sterile neutrinos never reach full equilibrium with the active sector, we model their distribution using a modified Fermi-Dirac form,

$$f_s(p) = \frac{\alpha}{e^{p/T_s} + 1}, \quad (2)$$

where p is the magnitude of the 3-momentum, T_s is an effective temperature encoding the deviation from thermal equilibrium, and $0 < \alpha \lesssim 1$ is a normalization factor encoding the departure from chemical equilibrium. Both T_s and α are treated as functions of the SM temperature T . In the dilute regime $f_s(p) \ll 1$, where quantum statistical effects are negligible, Eq. (2) approaches the Maxwell-Boltzmann limit. On the other hand, the fully thermalized case is recovered in the *large* coupling limit

¹A pseudoscalar mediator with an imaginary coupling will lead to the same physical effects in the relativistic limit, as shown in Ref. [41] for the high-energy neutrino case.

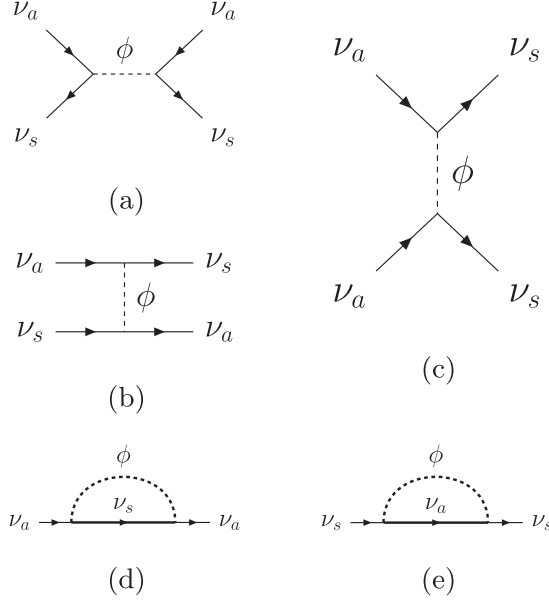


FIG. 2. Feynman diagrams illustrating active-sterile NSI contributions to scattering, number-changing interactions, and thermal potentials.

($\alpha \rightarrow 1, T_s \rightarrow T$), and $f_s(p)$ coincides with the active neutrino distribution $f_a(p) = 1/(e^{p/T} + 1)$.

The new NSI (1) opens up additional channels for active-sterile neutrino scattering [see Figs. 2(a) and 2(b)] and introduces new contributions to the neutrino self-energy [see Figs. 2(d) and 2(e)]. It also enables a number-changing process, $\nu_a \nu_a \rightarrow \nu_s \nu_s$ [see Fig. 2(c)], which plays a central role in the production of sterile neutrinos. Unlike the DW mechanism, this does not rely on active-sterile mixing and can remain efficient even when mixing angles vanish. These reactions are especially important at high temperatures, when production from mixing is negligible.

Evolution of sterile neutrino DM. The time evolution of the sterile neutrino distribution function $f_s(p, t)$ can be described by a semiclassical Boltzmann equation derived from the quantum kinetic equation (QKE) [13,54]. This equation serves as a classical approximation to the full QKE framework, incorporating quantum corrections through an effective active-sterile transition probability. The approximation remains valid in the small mixing angle limit [55–57].

In the presence of active-sterile NSI, the Boltzmann equation is modified to

$$\frac{\partial f_s(t, p)}{\partial t} - H p \frac{\partial f_s(t, p)}{\partial p} = \mathcal{C}_{\nu_a \leftrightarrow \nu_s} + \mathcal{C}_{\nu_a \nu_a \leftrightarrow \nu_s \nu_s}, \quad (3)$$

where $H(t)$ is the Hubble rate. The first collision term on the right-hand side describes sterile neutrino production via active-sterile oscillations. It depends on both the total interaction rate $\Gamma_{\text{tot}} = \Gamma_{\text{SM}} + \Gamma_{\text{NSI}}$ and effective potential

$\mathcal{V}_{\text{tot}} = \mathcal{V}_{\text{SM}} + \mathcal{V}_{\text{NSI}}$, which include contributions from the SM as well as NSI. The second collision term in Eq. (3) accounts for number-changing interactions induced by NSI (see supplemental material [58], Sec. I).

The SM contribution to the thermally averaged scattering rate of ν_a , arising from interactions with charged leptons, is given by [54,61]

$$\Gamma_{\text{SM}}(p) = C_a(T) G_F^2 p T^4, \quad (4)$$

where G_F the Fermi constant and $C_a(T)$ is a temperature-dependent function defined in Refs. [54,61,62]. The quantity Γ_{NSI} denotes the interaction rate arising from additional scattering channels induced by NSI. In this scenario, the only relevant process is $\nu_a \nu_s \leftrightarrow \nu_a \nu_s$, which contributes to the interaction rates of both active and sterile neutrinos. Specifically, this rate is given by the sum

$$\Gamma_{\text{NSI}} = \Gamma_{\nu_a \nu_s \leftrightarrow \nu_a \nu_s}^{(a)} + \Gamma_{\nu_a \nu_s \leftrightarrow \nu_a \nu_s}^{(s)}. \quad (5)$$

In the heavy-mediator limit $m_\phi^2 \gg pT$, these rates simplify to (see supplemental material [58], Sec. II, for the derivation)

$$\Gamma_{\nu_a \nu_s \leftrightarrow \nu_a \nu_s}^{(i)}(p) \simeq \frac{7\pi y_{as}^4}{108 m_\phi^4} p \times \begin{cases} \alpha T_s^4 & \text{for } i = a \\ T^4 & \text{for } i = s \end{cases}. \quad (6)$$

The SM effective potential for ν_a can be expressed as [63–65]

$$\mathcal{V}_{\text{SM}}(p) = \frac{2\sqrt{2}\zeta(3)G_F\eta_B T^3}{4\pi^2} - \frac{8\sqrt{2}G_F p}{3m_Z^2}(\rho_{\nu_a} + \rho_{\bar{\nu}_a}) - \frac{8\sqrt{2}G_F p}{3m_W^2}(\rho_a + \rho_{\bar{a}}), \quad (7)$$

where the first term is the baryonic contribution and the rest are the thermal contributions, with ρ_a, ρ_{ν_a} being the energy densities of corresponding charged leptons and active neutrinos, respectively. The NSI contribution to the thermal potential arises from the self-energy diagrams shown in Figs. 2(d) and 2(e) and is given by the difference between the active and sterile neutrino effective potentials [38,66],

$$\mathcal{V}_{\text{NSI}} = \mathcal{V}_a - \mathcal{V}_s. \quad (8)$$

Assuming the sterile neutrino mass is negligible compared to the relevant temperatures, the thermal potential for a neutrino species ν_i (with $i = a$ or s) induced by the other species ν_j (with $j \neq i$) in the heavy-mediator limit takes the form [38]

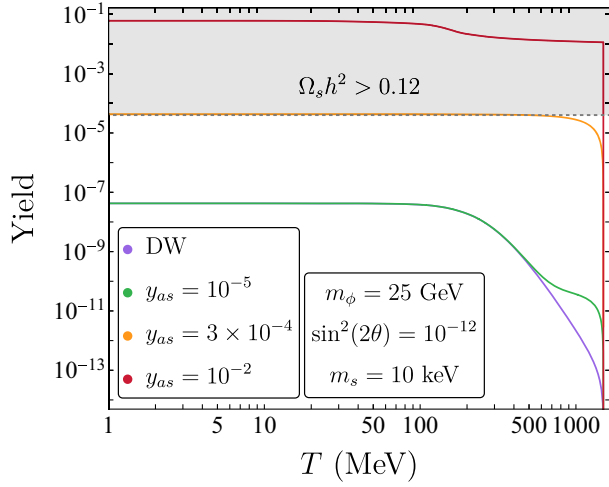


FIG. 3. Evolution of the comoving yield for sterile neutrino DM production for the chosen BP as a function of the SM temperature T .

$$\mathcal{V}_i(p) \simeq -\frac{7\pi^2 y_{as}^2}{45m_\phi^4} p \times \begin{cases} \alpha T_s^4 & \text{for } i = a \\ T^4 & \text{for } i = s \end{cases}. \quad (9)$$

In our framework, sterile neutrinos can also be produced directly via the number-changing process $\nu_a \nu_a \rightarrow \nu_s \nu_s$, as shown in Fig. 2(c). This contribution is captured by the second collision term in Eq. (3). At early times, when production from the modified DW mechanism is negligible, these interactions play a crucial role in initiating the sterile neutrino population and driving DM production. For a heavy mediator, the corresponding rates simplify to

$$\Gamma_{\nu_i \nu_i \rightarrow \nu_j \nu_j}(p) \simeq \frac{7\pi y_{as}^4}{216m_\phi^4} p \times \begin{cases} T^4 & \text{for } i = a, j = s \\ \alpha T_s^4 & \text{for } i = s, j = a \end{cases}. \quad (10)$$

See Ref. [58], Sec. III, for a comparison of the relevant interaction rates with the Hubble rate.

We solve the Boltzmann equation (3) by first rewriting it in terms of the evolution equations for the comoving number density n_s and energy density ρ_s of sterile neutrinos,

$$\dot{n}_s + 3Hn_s = C_{n_s}, \quad \dot{\rho}_s + 4H\rho_s = C_{\rho_s}, \quad (11)$$

where C_{n_s} and C_{ρ_s} are the respective moments of the collision term in Eq. (3); see Ref. [58], Sec. I. These equations are then solved numerically from an initial temperature of $T_i = 1.5$ GeV up to the active neutrino decoupling temperature, $T \approx 1$ MeV.² In Fig. 3, we show the evolution of the sterile neutrino comoving yield for a representative benchmark point (BP) with $m_s = 10$ keV,

²Our final result is insensitive to the choice of the initial temperature, as long as it is well below the mediator mass and well above the neutrino decoupling temperature.

$\sin^2 2\theta = 10^{-12}$, $m_\phi = 25$ GeV shown by the red star in Fig. 1, which satisfies current astrophysical constraints and lies within the projected reach of future x-ray telescopes. We notice the following features:

- (i) For very small couplings (e.g., $y_{as} \lesssim 10^{-5}$), number-changing interactions are only briefly active and quickly become inefficient. In this regime, modifications to the DW mechanism from the additional scatterings and thermal potentials are negligible, and the yield closely tracks the DW baseline.
- (ii) For slightly larger couplings (e.g., $y_{as} \sim 10^{-4}$), number-changing processes remain efficient for a longer duration, sustaining production down to $T \sim 100$ MeV. As a result, the yield exceeds the DW curve before freezing in. After number-changing production ceases, the spectrum gradually returns to a DW-like shape.
- (iii) For large couplings (e.g., $y_{as} \gtrsim 10^{-2}$), number-changing interactions remain efficient throughout the entire production window, keeping the sterile sector in thermal contact with the active neutrinos. The yield quickly rises to the overproduction region (gray shaded), well above the DW line. This scenario is cosmologically excluded as it would indicate overproduction of DM beyond the observed value $\Omega_s h^2 = 0.120 \pm 0.001$ [67], as well as of extra number of relativistic degrees of freedom at BBN.

With our chosen parametrization of the distribution function (2), the final relic abundance is determined by the suppression factor α , while the effective temperature T_s controls the spectral shape of the sterile neutrino population. The evolution of these quantities is given in Ref. [58], Sec. IV. However, we emphasize that the final relic density remains the same irrespective of the choice of the distribution function, since Ω_s (or ρ_s) can be directly obtained by solving Eq. (11) without the need to specify the exact form of $f_s(p)$.

Results. Figure 1 presents the viable parameter space in the plane of sterile neutrino mass and active-sterile mixing angle for a fixed value of $y_{as} = 3 \times 10^{-4}$. We find that the presence of NSI significantly enlarges the allowed region beyond the traditional DW scenario. In particular, viable production can now happen at much smaller mixing angles and even remains efficient in the limit of vanishing mixing. This behavior differs substantially from both the standard DW mechanism and earlier works on both active and sterile neutrino self-interaction scenarios [32,33,35–38], where underproduction necessarily occurs at small mixing due to the lack of efficient production channels.

Moreover, in the previously discussed active-active and sterile-sterile NSI scenarios, for heavy mediators, a substantial departure from DW typically requires couplings $\gtrsim 10^{-2}$ [32,38]. However, in our case, such large couplings lead to overproduction of sterile neutrinos. In fact, efficient production can be achieved at smaller couplings without

relying on active-sterile mixing. Figure 1 indicates a transition in the parameter space: for a given coupling and mediator mass, there exists a minimum sterile neutrino mass above which number-changing interactions alone can generate the observed relic abundance irrespective of the mixing angle. As the mediator mass increases, the resulting abundance approaches that of the standard DW scenario, as expected in the heavy mediator regime.

The viable parameter space now extends into regions where sterile neutrino DM can be produced efficiently without violating existing astrophysical constraints from x-ray telescopes such as Chandra [17], XMM-Newton [14], NuSTAR [21], and XRISM [22] as well as from gamma-ray telescopes like INTEGRAL/SPI [20] at higher energies. We also include projected sensitivities from future missions such as ATHENA [68], eROSITA [69], and eXTP [70]. Upcoming MeV gamma-ray telescopes can further extend the sensitivity to sterile neutrino DM in the (0.2–100) MeV mass range (not shown in Fig. 1) [71]. The blue circle in Fig. 1 represents the reported detection of an unidentified x-ray line near 3.5 keV [72–74] which would correspond to a 7.1 keV sterile neutrino, but the existence of the 3.5 keV line signal remains debatable [75]. In addition to astrophysical probes, sterile neutrino DM coupling to electron flavor can also be tested via precision β -decay experiments, as shown by the dotted and dashed red lines in Fig. 1 for TRISTAN [76] and HUNTER [77] projected sensitivities.

Furthermore, cosmological constraints from structure formation impose additional restrictions on warm DM candidates like sterile neutrinos, which, due to their relatively long free-streaming length, suppress the formation of small-scale structures, a feature that must remain consistent with observations of the matter power spectrum, Lyman- α , and Milky Way satellites [78–82]. Specifically, Lyman- α forest measurements impose a lower bound on the warm DM mass, which corresponds to a present-day root-mean-square velocity of $v_{\text{rms, today}} \lesssim 16$ m/s [83]. From this, we can infer that the sterile neutrino mass should satisfy a lower bound of $m_s \gtrsim 8$ keV (shown by the vertical brown-shaded region in Fig. 1), which is stronger than the generic lower bound of $m_s \gtrsim 2$ keV for any fermionic DM from phase-space density arguments based on observations of dwarf galaxies [61,84–86].

Additionally, any viable DM candidate must have a lifetime longer than the age of the Universe, $t_U \sim 13.8$ Gyr [67], since particles with shorter lifetimes would decay before contributing significantly to structure formation [10,54]. The gray dashed curve in Fig. 1 marks the parameter values for which the sterile neutrino lifetime equals the age of the Universe.

With the active-sterile interaction of the form (1), there are further constraints from the loop-induced coupling of the sterile neutrino to the SM Z-boson [46,87], including invisible Z decay and direct detection constraints. However, these contributions are suppressed by $\sim y_{as}^2/16\pi^2$ and

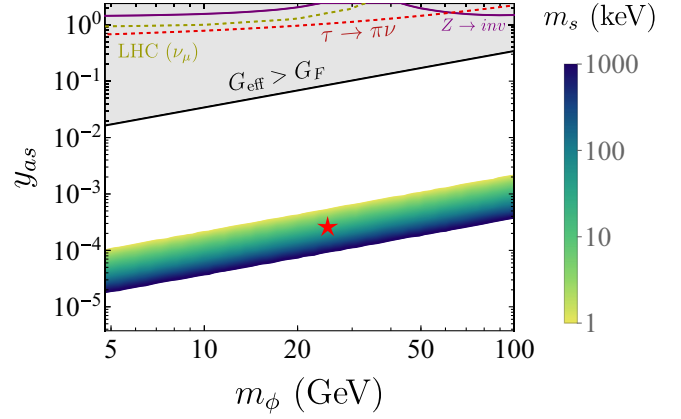


FIG. 4. Mediator mass versus coupling strength for allowed values of mixing angle. The colored band denotes the region with the correct DM abundance ($\Omega_s h^2 = 0.12$). The laboratory constraints shown here are all within the cosmologically disfavored regime with $G_{\text{eff}} > G_F$ (gray shaded). The red star corresponds to the BP shown in Fig. 1 that gives the correct yield in Fig. 3.

therefore are negligibly small for the BP value of $y_{as} = 3 \times 10^{-4}$ assumed in Fig. 1. Even for larger values of y_{as} , the direct detection of keV-scale sterile neutrino DM is rather challenging with the current detector thresholds, but electron recoil signals from inelastic scattering of sterile neutrinos from bound electrons could potentially provide competitive bounds in the future [88].

Figure 4 shows the allowed parameter space in the plane of mediator mass m_ϕ and coupling y_{as} , assuming sterile neutrino production is dominated by NSI. These results hold for any mixing angle allowed by astrophysical bounds. We find that sterile neutrinos with masses in the range 1–100 keV can account for the observed relic abundance even in the absence of mixing, for couplings of order $y_{as} \sim 10^{-4}$ – 10^{-3} . At smaller couplings ($y_{as} \lesssim 10^{-5}$), the number-changing process $\nu_a \nu_a \rightarrow \nu_s \nu_s$ becomes inefficient, and the combined effect of scattering and oscillations is also insufficient to generate the observed relic abundance. The range of couplings considered in our analysis remains consistent with current experimental constraints. For larger values ($y_{as} \gtrsim 10^{-3}$), sterile neutrinos are typically overproduced unless the mediator is sufficiently heavy to suppress the production rate. The gray shaded regions in Fig. 4 indicate parameter choices where the effective neutrino coupling $G_{\text{eff}} \equiv y_{as}^2/m_\phi^2$ is greater than G_F and is therefore cosmologically excluded, because it would hinder neutrino decoupling at MeV-scale prior to BBN. Notably, we find that laboratory constraints from $Z \rightarrow \text{invisible}$ decays, tau decays, and LHC searches [89,90] are already encompassed within the cosmologically disfavored region. Complementary views of the model parameter space are provided in Ref. [58], Sec. V.

The lowest mediator mass considered in Fig. 4 is 5 GeV because below this, for our choice of the initial temperature T_i , we enter a regime where the mediator can become

kinematically accessible in the thermal bath during the initial stage of DM production, potentially modifying the sterile sector evolution and enabling additional scattering processes involving on-shell ϕ exchange. A detailed treatment of this “light” mediator regime is left for future work.

Conclusion. The freeze-in mechanism discussed here for the efficient production of ν_s DM via the interaction term (1) provides a novel way to realize observable neutrino-DM interactions while being consistent with the DM relic density requirement—a feat otherwise deemed extremely challenging, if not impossible, in the thermal freeze-out scenario [52]. This result could have major implications for the future observations of supernova neutrinos (either galactic or diffuse) at upcoming large-volume neutrino detectors like JUNO [91], Hyper-K [92], and DUNE [93] as well as high-energy astrophysical neutrinos at neutrino telescopes like IceCube [94] and KM3NeT [95]. In particular, it introduces an energy-dependent opacity via $\nu_a\nu_s \rightarrow \nu_a\nu_s$ scattering mediated by ϕ that would attenuate the astrophysical neutrino flux passing through a DM halo, with detectable effects in the observed neutrino energy spectrum [40,41,52,96–102].

The presence of keV sterile neutrinos can also influence the evolution of core-collapse supernovae [103–107], but after taking into account the dynamical feedback effect [108–110], it did not show up in Fig. 1. However, all these analyses were done in the context of $\nu_a - \nu_s$ oscillations only. Including NSI effects could change the picture significantly. For instance, the NSI process $\nu_a\nu_a \rightarrow \nu_s\nu_s$ mediated by ϕ can potentially affect the core-collapse supernova dynamics by depleting the active neutrino population, thereby reducing the neutrino pressure and altering the cooling rate of the protoneutron star. A detailed treatment of this effect is beyond the scope of the present work, but recent results from Refs. [111,112] show that the ensuing constraints from SN1987A might become relevant in the parameter space shown in Fig. 4.

In addition, the $\nu_a\nu_a \rightarrow \nu_s\nu_s$ process would reduce the number of active neutrinos escaping from core-collapse supernova events throughout the cosmic history. This leads to a suppressed diffuse supernova neutrino background (DSNB) flux, as compared to the standard expectation [113–116]. Given the recent hint of a potential DSNB detection in Super-K Gd run data [117], this could become another important probe of the parameter space shown in Fig. 4.

The NSI considered here could also have interesting cosmological implications. In particular, the $\nu_a\nu_a \rightarrow \nu_s\nu_s$ process with sub-MeV sterile neutrinos serves as an efficient neutrino cooling channel after the standard neutrino decoupling [118]. This could in fact mimic self-interacting neutrinos and resolve some persistent tensions in current cosmological datasets [119].

UV-complete models of the effective interaction term (1) offer additional interesting phenomenology in laboratory experiments. For instance, the model proposed in Ref. [48] was originally motivated by the MiniBooNE, KOTO, and muon $g - 2$ anomalies. Although the latter two anomalies have since disappeared, the MiniBooNE excess still remains unresolved.

In conclusion, our new mechanism for the freeze-in production of sterile neutrino DM using nonstandard interaction with active neutrinos without relying on the active-sterile neutrino mixing provides a new way to realize potentially observable neutrino-DM interactions while being consistent with DM relic density and has broad phenomenological implications.

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