

Probing a New Regime of Neutrino Self-Interactions with Astrophysical Neutrinos and the Relativistic Cosmic Neutrino Background

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(Received 23 January 2025; revised 28 June 2025; accepted 10 September 2025; published 28 October 2025)

Neutrino self-interactions beyond the standard model have profound implications in astrophysics and cosmology. In this Letter, we study an uncharted scenario in which one of the three neutrino species has a mass smaller than the temperature of the cosmic neutrino background. This results in a relativistic component that significantly broadens the absorption feature on the astrophysical neutrino spectra, in contrast to the sharply peaked absorption expected in the extensively studied scenarios assuming a fully nonrelativistic cosmic neutrino background. By solving the Boltzmann equations for neutrino absorption and regeneration, we demonstrate that this mechanism provides novel sensitivity to sub-keV mediator masses, well below the traditional $\sim 1\text{--}100$ MeV range. Future observations of the diffuse supernova neutrino background with Hyper-Kamiokande could probe coupling strengths down to $g \sim 10^{-8}$, surpassing existing constraints by orders of magnitude. These findings open new directions for discoveries and offer crucial insights into the interplay between neutrinos and the dark sector.

DOI: [10.1103/9ddp-j1z9](https://doi.org/10.1103/9ddp-j1z9)

Introduction—Neutrinos, interacting feebly with matter, could strongly interact among themselves. This intriguing possibility, which is not only allowed by laboratory constraints [1–6] but also well motivated from various beyond-the-standard-model theories [7–14], has received growing interest in recent years—see, e.g., Ref. [15] for a review. Moreover, the cosmological and astrophysical implications of neutrino self-interactions are rich [16–41], with the most prominent ones related to cosmic microwave background (CMB) data interpretations [16–19], supernova dynamics [20–24], and astrophysical neutrino propagation [25–33].

Astrophysical neutrinos propagating through cosmic distances could be attenuated due to scattering with the cosmic neutrino background (CNB) via neutrino self-interactions [25–33]. To achieve observable attenuation, the mediator of neutrino self-interactions is often assumed to be light, typically around $m_\phi \sim 1\text{--}100$ MeV, with the coupling $g_\nu \sim 10^{-3} - 10^{-1}$. The corresponding four-Fermi effective interaction strength is $G_X \equiv g^2/m_\phi^2 \sim 10^5 G_F$ with G_F being the Fermi constant. Below the MeV scale,

cosmological observables from big bang nucleosynthesis (BBN) and the effective number of neutrino species (N_{eff}) typically set much stronger constraints than astrophysical neutrinos and other probes.

In this Letter, we propose an exceptionally interesting effect that could be exploited to probe a new regime of neutrino self-interactions *well below the MeV scale*.

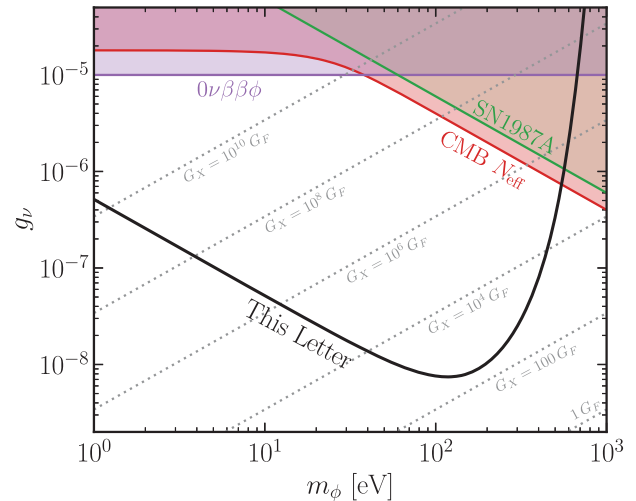


FIG. 1. The parameter space of neutrino self-interactions in the low-mass regime. The observation of DSNB in upcoming neutrino detectors such as Hyper-Kamiokande can substantially improve current known bounds (colored regions) [37,44,45] by a few orders of magnitude (the black line).

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Compared to previous studies, our Letter takes into account a crucial factor that was neglected in the past but could substantially enhance the sensitivity reach, much stronger than the cosmological constraints. That is, the lightest neutrino species in the CNB today can still be relativistic, i.e., its mass being lower than the CNB temperature, which is allowed by all experimental data.

Figure 1 presents a preview of our results. We demonstrate that future observations of the diffuse supernova neutrino background (DSNB) [42,43] can probe m_ϕ of eV to sub-keV and g_ν down to 10^{-8} . This surpasses existing bounds by a few orders of magnitude. In terms of G_X , the best scenario may even reach $\mathcal{O}(100)G_F$.

Relativistic neutrinos may significantly enhance the absorption of astrophysical neutrinos during propagation. This occurs at the resonance of the s -channel scattering, $\nu\nu \rightarrow \phi \rightarrow \nu\nu$, where the mediator ϕ can be either on or off shell. In previous studies [25–33], the CNB being scattered is assumed to be fully nonrelativistic. Thus, the resonance occurs only in a very narrow energy range of astrophysical neutrinos. In contrast, for relativistic CNB neutrinos that follow a thermal momentum distribution, the resonance can be reached at much wider energies. This is because each astrophysical neutrino with a certain momentum can always find a CNB neutrino with appropriate momentum such that their Mandelstam variable s matches m_ϕ^2 . In other words, *relativistic CNB neutrinos widen the resonance*, which is the key observation made in this Letter.

The relativistic scenario becomes even more motivated after the recent DESI publication reporting the latest cosmological bound on the sum of neutrino masses [46,47], as the upper bound is approaching the minimal value compatible with neutrino oscillation data [48,49].

In the subsequent analysis, we estimate the absorption rate of DSNB neutrinos by the relativistic CNB to provide a qualitative discussion, highlighting the advantage of using such a scenario to probe neutrino self-interactions. To quantitatively study the impact on the DSNB, we perform a numerical computation by solving, for the first time, the Boltzmann equation for the coupled ν - ϕ system instead of that for ν only.

Interaction and mean-free path—We consider the following Lagrangian for neutrino self-interactions:

$$\mathcal{L} \supset \frac{1}{2}(\partial\phi)^2 - \frac{1}{2}m_\phi^2\phi^2 + g_\nu\phi(\nu\nu + \nu^\dagger\nu^\dagger), \quad (1)$$

where ϕ is a real scalar and ν represents the two-component Weyl–Van der Waerden spinor of a neutrino. Hence, Majorana neutrinos are implied by the Lagrangian. Like the Majorana mass terms of neutrinos, the above interaction violates the lepton number, implying that antineutrinos can be produced from neutrinos via, e.g., $\nu\nu \rightarrow \phi \rightarrow \bar{\nu}\bar{\nu}$. Since both the source (DSNB) and the target (CNB) are approximately symmetric under $\nu \leftrightarrow \bar{\nu}$, our analysis below is also ν - $\bar{\nu}$ symmetric, leading to ν - $\bar{\nu}$ symmetric results, unaffected

by this lepton-number-violating feature. For simplicity, we do not consider the flavor structure in this Letter. A full flavor analysis including the ν - $\bar{\nu}$ difference will be conducted in our upcoming work.

With the interaction in Eq. (1), an astrophysical neutrino ν_A propagating through the CNB may scatter off a neutrino in the CNB ν_C via

$$\nu_A + \nu_C \rightarrow X, \quad X \in \{2\nu_A, \phi, 2\phi, \dots\}. \quad (2)$$

Throughout this Letter, we use the term “astrophysical neutrinos” for neutrinos that acquire energies directly or indirectly from astrophysical sources. This includes neutrinos both emitted directly from such sources and regenerated through self-interactions.

It is useful to introduce the mean-free path of the propagation, L_{mean} . In the ultrarelativistic (UR) and nonrelativistic (NR) scattering regimes, L_{mean} can be approximated as

$$L_{\text{mean}} \simeq \begin{cases} \left(\frac{g_\nu^2 m_\phi^2 T}{16\pi E_\nu^2} \exp\left[-\frac{m_\phi^2}{4TE_\nu}\right] + \dots \right)^{-1} & (\text{UR}) \\ (n_C \sigma_{\text{NR}})^{-1} & (\text{NR}) \end{cases}, \quad (3)$$

where E_ν denotes the energy of the incoming ν_A , T is the temperature of the CNB, n_C is the number density of ν_C , and σ_{NR} denotes the NR cross section, which is almost independent of the momentum of ν_C . The UR and NR regimes correspond to the neutrino mass $m_\nu \ll T$ or $m_\nu \gg T$, respectively. For more general cases, L_{mean} can be derived from collision terms of Boltzmann equations—see Supplemental Material [56]. In the UR regime of Eq. (3), we only show the dominant contribution of $\nu_A + \nu_C \rightarrow \phi$, with “...” representing the remaining contributions. While in the NR regime, the contribution of $\nu_A + \nu_C \rightarrow \phi$ to σ_{NR} reads

$$\sigma_{\text{NR}}^{(\text{res})} \simeq \pi g_\nu^2 \delta(s - m_\phi^2), \quad (4)$$

where $s \simeq 2E_\nu m_\nu$. Because of the δ function, Eq. (4) gives rise to a very sharp and narrow absorption of ν_A by the CNB—see, e.g., Fig. 1 of Ref. [25]. In practice, Eq. (4) is usually integrated into the resonance of s -channel scattering via the Breit-Wigner formalism [cf. Eq. (25) in [57]], rather than being calculated separately.

In contrast to the narrow absorption in the NR regime, the absorption in the UR regime is much wider. As can be seen from Eq. (3), L_{mean}^{-1} in the UR regime peaks at $E_\nu \sim m_\phi^2/4T$, implying that ν_A of this energy can be absorbed most efficiently. If E_ν varies around this value by a factor of a few, the absorption is still effective.

Therefore, by comparing the NR and UR regimes, we see that *the most crucial difference is a very sharp and narrow absorption versus a much wider one*.

Taking the UR result with the present CNB temperature $T \simeq 1.9$ K and a few benchmark values indicated below, we find

$$\frac{L_{\text{mean}}}{1 \text{ Gpc}} \simeq 0.8 \left(\frac{10^{-8}}{g_\nu} \times \frac{E_\nu}{20 \text{ MeV}} \times \frac{0.1 \text{ keV}}{m_\phi} \right)^2 e^\lambda, \quad (5)$$

where $\lambda \equiv (m_\phi^2/4TE_\nu)$ and for the above benchmark values it gives $e^\lambda \sim \mathcal{O}(1)$. Obviously, the obtained L_{mean} for this benchmark is much shorter than the radius of the observable Universe, $R_{\text{Universe}} \simeq 14$ Gpc. This implies that the relativistic CNB over the entire Universe would be opaque to a 20 MeV neutrino in the benchmark scenario. Therefore, observations of astrophysical neutrinos at such energies from sources that are cosmologically distant should be able to probe a sub-keV neutrino self-interaction mediator with $g_\nu \sim 10^{-8}$.

The above estimate of the mean-free path serves as a qualitative study and has neglected the cosmological redshift. If ν_A is produced from high-redshift sources (e.g., redshift $z \sim 4-5$), it would propagate through a much denser and more energetic (hence more likely to be relativistic) neutrino background than the present CNB. A quantitative study, taking the cosmological redshift and other effects such as the production of ϕ and secondary ν_A into account, requires solving the Boltzmann equation, as we discuss below.

Solving the Boltzmann equation—The Boltzmann equation that governs the evolution of the phase space distributions of ν and ϕ in the expanding Universe reads

$$(\partial_t - H p \partial_p) \begin{bmatrix} f_\nu \\ f_\phi \end{bmatrix} = \begin{bmatrix} -f_\nu \Gamma_\nu^- + (1 - f_\nu)(\Gamma_\nu^+ + \mathcal{S}_\nu) \\ -f_\phi \Gamma_\phi^- + (1 + f_\phi)\Gamma_\phi^+ \end{bmatrix}, \quad (6)$$

where f_x is the phase space distribution of species $x = \nu$ or ϕ , $H = a^{-1} da/dt$ is the Hubble parameter with a the scale factor, Γ_x^+ and Γ_x^- are the production and depletion rates of x via particle physics processes such as those in Eq. (2), and $\mathcal{S}_\nu$ accounts for the production of ν from astrophysical sources.

In principle, f_ν could include both astrophysical and CNB neutrinos, with the former being viewed as a high-energy nonthermal addition to the thermal spectrum of the latter. However, the very hierarchical energy regimes between them and the low number density of the former make this treatment impractical. Hence, we separate the high-energy part from the thermal part and consider f_ν as the phase space distribution of ν_A only. Under this convention, we have $f_{\nu,\phi} \ll 1$, which implies that the Pauli-blocking and Bose-enhancement factors $1 \mp f_{\nu,\phi}$ can be neglected.

In the UR regime, the dominant processes contributing to $\Gamma_{\nu,\phi}^\pm$ are $\nu_A + \nu_C \rightarrow \phi$ and $\phi \rightarrow \nu_A + \nu_A$, with their contributions proportional to g_ν^2 . Other processes such as $\nu_A +$

$\nu_C \rightarrow \nu_A + \nu_A$ (without resonance) and $\nu_A + \nu_C \rightarrow \phi + \phi$ are subdominant in the regime of our interest because their contributions are suppressed by g_ν^4 . In this Letter, we focus on the $\mathcal{O}(g_\nu^2)$ contributions, which can be calculated analytically (see Sec. S1 in Supplemental Material [56] for the derivations),

$$\Gamma_\nu^- = \frac{g_\nu^2 m_\phi^2 T}{16\pi E_\nu^2} \exp\left[-\frac{m_\phi^2}{4TE_\nu}\right], \quad (7)$$

$$\Gamma_\nu^+ = \frac{2g_\nu^2 m_\phi^2}{16\pi E_\nu^2} \int_{E_\nu^{\min}}^\infty dE_\phi f_\phi, \quad (8)$$

$$\Gamma_\phi^+ = \frac{g_\nu^2 m_\phi^2}{16\pi E_\phi p_\phi} \int_{E_\nu^-}^{E_\nu^+} dE_\nu f_\nu \exp\left[-\frac{E_\phi - p_\nu}{T}\right], \quad (9)$$

$$\Gamma_\phi^- = \frac{g_\nu^2 m_\phi^2}{16\pi E_\phi}, \quad (10)$$

where E_ϕ and p_ϕ denote the energy and momentum of ϕ , $E_\phi^{\min} = (m_\phi^2/4E_\nu) + E_\nu$, and $E_\nu^\mp = (E_\phi \mp p_\phi)/2$. With the analytical expressions of $\Gamma_{\nu,\phi}^\pm$ and a given source $\mathcal{S}_\nu$, it is straightforward to solve the Boltzmann equation numerically.

Figure 2 shows an example of a monochromatic and instantaneous source emitting neutrinos with $E_\nu = 20$ MeV at $z = 4$. In addition, we set $m_\phi = 100$ eV, $g_\nu = 6 \times 10^{-9}$ for neutrino propagation. The left and right panels show how the particle number and energy in each given momentum bin evolve, respectively. Here the particle number and energy in a bin are defined as $N_{\text{bin}} \equiv \int_{\text{bin}} \tilde{f}_{\nu+\phi}(\tilde{p}) [d^3\tilde{p}/(2\pi)^3]$ and $E_{\text{bin}} \equiv \int_{\text{bin}} \tilde{f}_{\nu+\phi}(\tilde{p}) \tilde{p} [d^3\tilde{p}/(2\pi)^3]$, where $\tilde{p} = ap$ is the comoving momentum and $\tilde{f}_{\nu+\phi}(\tilde{p}) \equiv f_\nu(p) + f_\phi(p)$. Note that our numerical computation throughout the Letter uses much finer momentum binning than what is shown in the figure.

As is shown in the left panel, the total number of astrophysical neutrinos (black dashed line) increases rapidly after the initial production from the source. This is because $\nu_A + \nu_C \rightarrow \phi$ and $\phi \rightarrow \nu_A + \nu_A$ proceed efficiently at this stage. The two processes together lead to an exponential growth of ν_A with degraded energies and meanwhile deplete ν_A at the initial energy. Despite the proliferation of the particle number, the total energy is approximately conserved (up to proper comoving factors), as is shown by the black dashed line in the right panel. The energy conservation relies on two approximations: (i) the energy contribution from ν_C is negligible, and (ii) the involved species are all relativistic [58]. The validity of (i) is guaranteed by the low temperature of CNB; and the validity of (ii) is implied by $E_\nu \gg m_\phi$. From the energy conservation and the particle number proliferation, we anticipate that the most prominent effect of neutrino

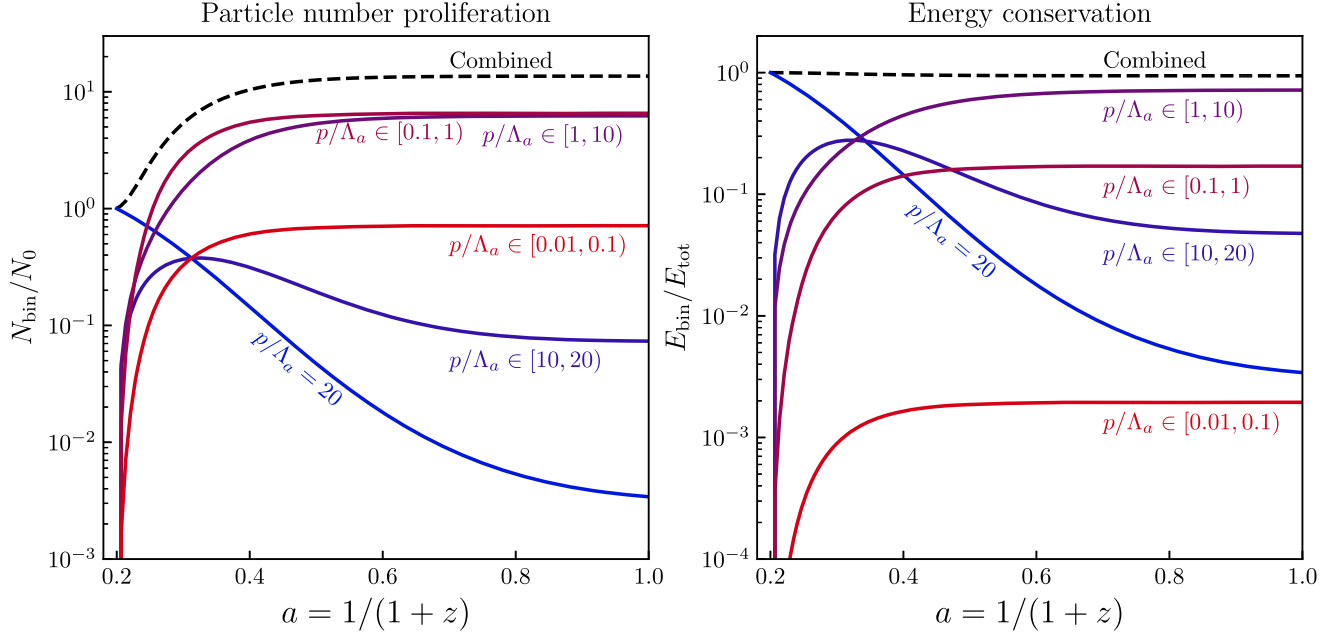


FIG. 2. Evolution of particle numbers (left) and energy (right), assuming $m_\phi = 100$ eV, $g_\nu = 6 \times 10^{-9}$, and a monochromatic and instantaneous source emitting neutrinos with $E_\nu = 20$ MeV at $z = 4$. Here, the monotonically decreasing curves represent neutrinos that retain their initial comoving momentum, corresponding to $p/\Lambda_a = 20$ where $\Lambda_a \equiv \text{MeV} \times a_i/a$ and $a_i = 0.2$. Other colored curves represent particles that are subsequently produced. The black dashed lines represent the total number and energy of these particles.

self-interactions on the DSNB should be a substantial enhancement of its low-energy spectrum in combination with significant absorptions at high energies.

Impact on the DSNB spectrum—To reach the resonance in the UR regime, E_ν needs to be around 1–100 MeV, which roughly matches the energy range of supernova neutrinos. Hence, the most suitable neutrino source to be considered in this Letter would be supernovae, provided that their distances are at the cosmological scale (gigaparsec). Although the neutrino flux of a single supernova at such a large distance is far from being observable in current and future experiments, the combined neutrino flux of all supernovae in the entire Universe (i.e., DSNB) is likely to be detected in the foreseeable future [43]. Therefore, it is tempting to investigate whether and how the DSNB might be altered in the new physics scenario considered here.

The source term of the DSNB can be computed by [42,43]

$$S_\nu(E_\nu, a) = \frac{2\pi^2}{E_\nu^2} \times \frac{dN}{dE_\nu} \times \frac{R_{\text{SN}}}{a^3}, \quad (11)$$

where dN/dE_ν denotes the neutrino emission per supernova and R_{SN} is the rate of core-collapse supernova. Both can be approximated by analytical expressions and have been provided by Ref. [42]. Following the same setup in Ref. [42], we calculate Eq. (11) and substitute it into the Boltzmann equation. Solving the equation, we obtain the phase space distribution f_ν , which can be recast into the neutrino flux Φ_ν by

$$\Phi_\nu = a^3 \frac{E_\nu^2}{2\pi^2} f_\nu. \quad (12)$$

Figure 3 shows the impact on the DSNB spectrum from neutrino self-interaction. In the upper panel, the black dashed curve represents the DSNB neutrino flux without neutrino self-interactions, assuming a temperature of 6 MeV. This curve is confirmed to align with Fig. 4 in Refs. [42,43] and is also approximately consistent with Fig. 1 in [59]. The colored curves depict the results of neutrino self-interaction. We show the result with three different couplings, $g = 10^{-8}$, $g = 2 \times 10^{-8}$, and $g = 4 \times 10^{-8}$ for $m_\phi = 100$ eV, while also showing the result of $g = 10^{-8}$ for $m_\phi = 10$ eV to demonstrate the impact of the mediator mass. The lower panel illustrates the normalized results compared to free propagation, highlighting both creation and absorption effects. In all cases with the interaction, the neutrino flux in the MeV range is absorbed by $\nu_A + \nu_C \rightarrow \phi$. One can see that the absorption in the lower panel behaves as a “wide dip,” indicating the result of the “widened resonance.” Next, the produced ϕ particles decay into two neutrinos, with each roughly carrying half of the energy of ν_A , given that CNB temperature is lower than ν_A energy by a few orders of magnitude. These created neutrinos further scatter with CNB to create more ϕ particles, which then decay into pairs of neutrinos with halved energies. Consequently, there is an accumulation of neutrinos in the low-energy spectrum, leading to a blowup behavior in flux in the lower-energy range, as shown in Fig. 3.

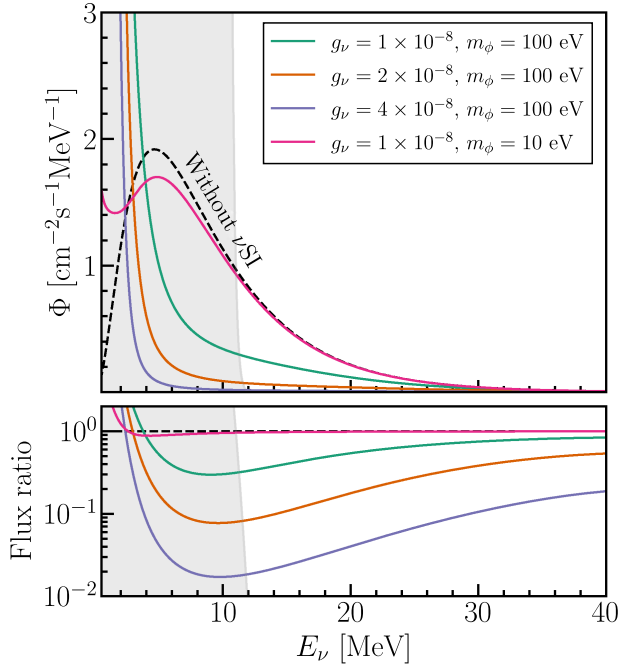


FIG. 3. The impact of neutrino self-interactions (ν SI) on the DSNB energy spectrum. Upper: neutrino flux. Lower: neutrino flux normalized by the result from the propagation of DSNB neutrinos without ν SI. The gray-shaded region indicates the background from reactor neutrinos [59,60]. The black dashed curves indicate the results without ν SI. The DSNB temperature is assumed to be 6 MeV here.

Detection and χ^2 analysis—For DSNB detection, we focus on the dominant inverse β decay (IBD) channel [61] in the upcoming Hyper-Kamiokande (HK) detector [62,63] that is expected to detect the DSNB with robust statistics due to its unprecedentedly large volume and excellent background reduction if it is gadolinium doped. With a 3740 ktyr exposure, HK is expected to detect ~ 800 DSNB events in the standard scenario above 12 MeV, while in the new physics scenario, the events above a certain energy could be substantially reduced. The calculation of DSNB event rates at HK has been well established in the literature—see, e.g., Ref. [32]. We detail it together with a discussion on the detection backgrounds in Sec. S4 of Supplemental Material [56].

To quantify the observability of the new physics effect, we adopt the following χ^2 function [64]:

$$\chi^2 = \left(\frac{y}{\sigma_y}\right)^2 + \sum_i \left(\frac{(1+y)N_i - N_{\text{st},i}}{\sigma_i}\right)^2, \quad (13)$$

where N_i and $N_{\text{st},i}$ denote the numbers of events in the i th energy bin in the new physics and the standard model scenarios, respectively; y is a normalization factor, and σ_i and σ_y denote the uncertainties of N_i and y , respectively. With sufficient statistics ($N_{\text{st},i} \gg 1$), one can take

$\sigma_i = \sqrt{N_{\text{st},i} + N_{\text{bkg},i}}$, where $N_{\text{bkg},i}$ is the number of background events. For the normalization uncertainty, we take $\sigma_y = 0.5$, which can be justified from DSNB calculations over the past ten years—see Supplemental Material [56] for detailed discussions. This uncertainty is partially caused by the star formation rate [42] and could be improved by future astrophysical observations. In our analysis, we adopt the frequentist treatment to marginalize y [64]. In addition, the DSNB temperature T_{SN} is uncertain and may affect the spectral shape. We vary T_{SN} in [4, 8] MeV and marginalize it as well. Figure 1 presents our projected sensitivity reach. Our results indicate that observations of DSNB at HK can probe neutrino self-interactions with m_ϕ from eV to sub-keV and g_ν down to 10^{-8} .

In the shown parameter space, there are a few existing bounds derived from cosmology [37,65], the detection of SN1987A neutrinos [44,66–69], and neutrinoless double β decay experiments searching for Majoron emissions ($0\nu\beta\beta\phi$) [5,45,70–72]. The cosmological bound relies on the production of ϕ in the early Universe via $\nu\nu \rightarrow \phi$ and $\nu\nu \rightarrow \phi\phi$, with the production rates proportional to g_ν^2 and g_ν^4 , respectively. In the low-mass regime, the former is suppressed by the mass of ϕ , causing the latter to dominate the production of very light m_ϕ . After neutrino decoupling, the produced ϕ and their subsequent decay products may affect the Hubble expansion rate, thereby altering BBN and CMB predictions. This effect can be quantified by the cosmological parameter N_{eff} . Here we take the N_{eff} bound on neutrinophilic scalars from Ref. [37], assuming the allowed deviation of N_{eff} is less than 0.285 [73]. For the SN1987A bound, we adopt the most stringent one from Ref. [44]. As for the $0\nu\beta\beta\phi$ bound, since m_ϕ here is well below the Q value of the process, it is almost independent of m_ϕ . Experimental searches for $0\nu\beta\beta\phi$ typically set an upper bound on g_ν around 10^{-5} [45,72].

Comparing our result with the existing bounds in Fig. 1, we see that a substantial part of the unexplored parameter space can be probed by DSNB observations. In this Letter, we focus on supernova neutrinos mainly because the energy scale at $\mathcal{O}(10)$ MeV is ideal for the resonant production of sub-keV ϕ . Nevertheless, our calculation can be readily applied to other astrophysical sources at much higher energies. For instance, PeV or EeV neutrinos would allow us to explore the widened resonance for $m_\phi \sim 1$ or 30 MeV, respectively. We leave these interesting scenarios to future work.

Summary and conclusions—Strong beyond-the-standard-model neutrino self-interactions are well motivated and exhibit rich phenomena in astrophysics and cosmology. An intriguing consequence is that astrophysical neutrinos propagating through cosmic distances could be attenuated due to scattering with the CNB, typically resulting in a sharp absorption feature on the spectrum due to the resonance production of the

mediator from the self-interaction. This mechanism provides stringent constraints on light mediators with masses of $\sim 1\text{--}100$ MeV and coupling strengths in the range of $10^{-3} - 10^{-1}$.

In this Letter, we consider an uncharted scenario in which the absorption is broadened significantly due to a relativistic component of the CNB. This occurs if the mass of the lightest neutrino species is below the CNB temperature, which is allowed by neutrino oscillation measurements and motivated by recent DESI data; thus, the CNB could follow a wide thermal momentum distribution, enabling a much wider range of astrophysical neutrino energies to satisfy the resonance condition.

We demonstrate that this broadened absorption feature offers a novel sensitivity to sub-keV mediator masses, well below the traditional 1–100 MeV range. By solving the Boltzmann equations for the coupled neutrino-mediator system, we quantitatively evaluate the absorption and regeneration of DSNB neutrinos during propagation. Our results show that, for mediator masses in the eV to sub-keV range, future observations of the DSNB by Hyper-Kamiokande would have a sensitivity of the coupling strength up to $g \sim 10^{-8}$, surpassing existing constraints by a few orders of magnitude.

This scenario warrants various studies in future work. For example, the proliferation of sub-10-MeV neutrinos could exceed the fluxes of neutrinos from reactors and other sources. Furthermore, rich phenomenology could manifest in TeV–EeV astrophysical neutrinos observed by IceCube and other observatories. Altogether, they will significantly extend our reach in the parameter space of neutrino self-interactions, offering valuable insights into the mystery of neutrinos and their possible interplay with the dark sector.

Acknowledgments—We are grateful to John Beacom and Pedro Machado for their helpful discussions. I. R. W. and B. Z. were supported by Fermi Research Alliance, LLC under Contract No. DE-AC02-07CH11359 with the U.S. Department of Energy, Office of Science, Office of High Energy Physics and are currently supported by Fermi Forward Discovery Group, LLC under Contract No. 89243024CSC000002 with the U.S. Department of Energy, Office of Science, Office of High Energy Physics. I. R. W. is also supported by DOE Distinguished Scientist Fellowship Grant No. FNAL 22-33. The work of X. J. X. is supported in part by the National Natural Science Foundation of China (NSFC) under Grant No. 12141501 and also by the CAS Project for Young Scientists in Basic Research (YSBR-099).

Data availability—The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be

prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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