FERMILAB-PUB-19-175-A-T

Constraining the Self-Interacting Neutrino Interpretation of the Hubble Tension

Nikita Blinov,* Kevin J. Kelly,† Gordan Krnjaic,‡ and Samuel D. McDermott§
Fermi National Accelerator Laboratory, Batavia, IL, USA

(Dated: May 9, 2019)

Large, non-standard neutrino self-interactions have been shown to resolve the $\sim 4\sigma$ tension in Hubble constant measurements and a milder tension in the amplitude of matter fluctuations. We demonstrate that interactions of the necessary size imply the existence of a force-carrier with a large neutrino coupling ($> 10^{-4}$) and mass in the keV – 100 MeV range. This mediator is subject to stringent cosmological and laboratory bounds, and we find that nearly all realizations of such a particle are excluded by existing data unless it carries spin 0 and couples almost exclusively to τ-flavored neutrinos. Furthermore, we find that the light neutrinos must be Majorana, and that a UV-complete model requires a non-minimal mechanism to simultaneously generate neutrino masses and appreciable self-interactions.

I. INTRODUCTION

The discrepancy between low-redshift and Cosmic Microwave Background (CMB) determinations of the present-day Hubble parameter, H_0, has grown in significance to $\sim 4\sigma$ over the last several years [1–5]. The standard cosmological model, ΛCDM, may need to be augmented if this "H_0 tension" is not resolved by observational systematics. Intriguingly, this tension cannot be addressed by modifying ΛCDM at low redshift [6–9], but adding new physics before recombination seems more promising [10–16]. Furthermore, low redshift measurements of the matter density fluctuation amplitude on 8 Mpc scales, σ_8, also appear to be lower than predicted by ΛCDM from the CMB. This milder σ_8 tension is not ameliorated in the models of [11–16].

A particularly interesting resolution to these issues is a non-standard neutrino self-interaction of the form [17–20]

$$\mathcal{L}_{\rm eff} = G_{\rm eff}(\bar{\nu}\nu)(\bar{\nu}\nu), \tag{1}$$

where $G_{\rm eff}$ is a dimensionful coupling and flavor indices have been suppressed. If $G_{\rm eff}$ is much larger than the Standard Model (SM) Fermi constant, $G_{\rm F}$, neutrinos remain tightly coupled to each other until relatively late times. This inhibits their free-streaming and results in enhanced power on small scales and a shift in the acoustic peaks of the CMB spectrum relative to ΛCDM [21]. The effect of neutrino self-interactions is degenerate with other parameters in the CMB fit, such as the angular scale of the sound horizon, the spectral index and amplitude of primordial fluctuations, and extra free-streaming radiation (as parametrized by $N_{\rm eff}$).

These approximate degeneracies have been shown to prefer $G_{\rm eff} \gg G_{\rm F}$ in cosmological data [17–20] and relax the H_0 tension [18–20]. The latest results of Ref. [20] extended previous analyses by allowing for finite neutrino masses and extra free-streaming radiation at the time of the CMB. They found that $G_{\rm eff}$ in the "strongly interacting" (SIν) or "moderately interacting" (MIν) regimes

$$G_{\rm eff} = \begin{cases} (4.6 \pm 0.5\,{\rm MeV})^{-2} & ({\rm SI}\nu) \\ (90^{+170}_{-60}\,{\rm MeV})^{-2} & ({\rm MI}\nu) \end{cases} \tag{2}$$

could simultaneously reduce the H_0 and σ_8 tensions. Interestingly, the SIν cosmology prefers a value of H_0 compatible with local measurements at the 1σ level, even before including local data in the fit.

The favored range of $G_{\rm eff}$ in Eq. (2) vastly exceeds the strength of weak interactions, whose corresponding coupling is $G_{\rm F} \simeq (2.9 \times 10^5\,{\rm MeV})^{-2}$. As we discuss below, such an interaction can only arise from the virtual exchange of a force carrier ("mediator") with $\mathcal{O}$(MeV) mass and appreciable couplings to neutrinos. This mass range indicates that this scenario faces a variety of stringent cosmological and laboratory constraints.

We find that if the interaction in Eq. (1) resolves the H_0 tension, then:

- **SIν is excluded:** The SIν-range of Eq. (2) cannot be realized in any of the consistent, low-energy, weakly coupled models that we consider.
- **Vector forces are excluded:** Constraints from the epoch of Big Bang Nucleosynthesis (BBN) exclude all self-consistent vector mediators.
- **Dirac neutrinos are excluded:** If neutrinos are Dirac particles, mediator-neutrino interactions efficiently thermalize the right-handed component of neutrinos, thereby significantly increasing the number of neutrino species at BBN.
- **Minimal seesaw models are excluded:** Achieving the necessary interaction strength in Eq. (2) from a gauge-invariant, UV-complete model is challenging with minimal field content in a Type-I or -II seesaw model; a non-minimal mechanism is required to generate neutrino masses.
- **Interactions with ν_τ are favored:** Couplings to ν_e, ν_μ with $G_{\rm eff}$ in the range of Eq. (2) are largely

* ORCID: http://orcid.org/0000-0002-2845-961X
† ORCID: http://orcid.org/0000-0002-4892-2093
‡ ORCID: http://orcid.org/0000-0001-7420-9577
§ ORCID: http://orcid.org/0000-0001-5513-1938

This manuscript has been authored 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.

excluded by laboratory searches for rare K decays and for neutrinoless double-beta decay, except for a small island for the ν_μ coupling. This means that the mass-eigenstate neutrinos only interact via their ν_τ components.

This work is organized as follows: Sec. II demonstrates that a light new particle is required to generate the interaction in Eq. (1) with appropriate coupling strength; Sec. III presents the cosmological bounds on this scenario; Sec. IV discusses the corresponding laboratory constraints; Sec. V shows how Eq. (1) can arise in UV complete models; finally, Sec. VI offers some concluding remarks.

II. THE NECESSITY OF A LIGHT MEDIATOR

The Boltzmann equations used in Refs. [18–20, 22] assume that left-handed (LH), mass-eigenstate neutrinos participate in elastic $2 \to 2$ scattering processes. They also assume that the interactions in Eq. (1) involve *constant* and *flavor-universal* values of $G_{\rm eff}$ during all epochs relevant for the CMB. The largest CMB multipoles observed by Planck correspond to modes that entered the horizon when the universe had a temperature of < 100 eV.This temperature sets the characteristic energy scale of scattering reactions during this epoch, and it is important that the form of the Lagrangian shown in Eq. (1) is valid in this regime. At higher energies, however, this description can break down. In this section, we therefore emphasize the need to introduce new particle content to study laboratory and early universe processes that occur at energies $\sim \mathcal{O}$(MeV).

As noted in Refs. [17–20], the operator in Eq. (1) is non-renormalizable, and thus is necessarily replaced by a different interaction with new degree(s) of freedom at some energy scale higher than the energies probed by the CMB (see Ref. [23] for a review). Since $G_{\rm eff}$ in Eq. (1) is momentum-independent, we will assume this interaction arises from "integrating out" a particle ϕ with mass m_ϕ and a perturbative coupling to neutrinos g_ϕ:

$$\mathcal{L}_{\rm phen} \supset -\frac{1}{2}m_\phi^2\phi^2 + \frac{1}{2}(g_\phi^{\alpha\beta}\nu_\alpha\nu_\beta\phi + {\rm h.c.}), \tag{3}$$

where ν_α are two-component left-handed neutrinos, α and β are flavor indices and the subscript "phen" indicates that we will use this Lagrangian for our phenomenological analysis. In Eq. (3) we have assumed that ϕ is a real scalar without loss of essential generality. In particular, our conclusions remain unaffected if ϕ is a CP-odd or complex scalar. We focus on a scalar mediator here for clarity; introducing a new vector force instead, for example, follows the same reasoning but comes with additional, stronger constraints discussed in more detail below. We also explicitly allow for generic couplings $g_\phi^{\alpha\beta}$ between different neutrino species, and discuss implications of different choices of $g_\phi^{\alpha\beta}$ below.

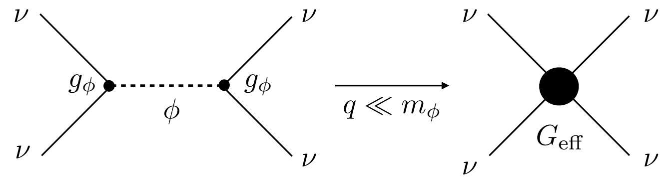

FIG. 1. Cartoon for how renormalizable interactions in Eq. (3) (left diagram) yield the contact interaction in Eq. (1) (right diagram) at low energies; flavor indices are suppressed.

Using Eq. (3), we see that the $\nu\nu \to \nu\nu$ scattering amplitude $\mathcal{M}$ is always proportional to two powers of g_ϕ multiplying the ϕ propagator, as shown schematically in Fig. 1. If the momentum transfer q satisfies $|q^2| \ll m_\phi^2$, we have

$$\mathcal{M} \propto \frac{g_\phi^2}{m_\phi^2 - q^2} \to G_{\rm eff}\left(1 + \frac{q^2}{m_\phi^2} + \cdots\right), \tag{4}$$

where we have suppressed flavor indices and defined

$$G_{\rm eff} \simeq \frac{g_\phi^2}{m_\phi^2} = (10\ \text{MeV})^{-2}\left(\frac{g_\phi}{10^{-1}}\right)^2\left(\frac{\text{MeV}}{m_\phi}\right)^2 . \tag{5}$$

In the opposite limit, $m_\phi^2 \ll |q^2|$, then $\mathcal{M} \propto g_\phi^2/|q^2|$, leading to a qualitatively different energy and temperature dependence of neutrino self-interactions; this regime was investigated in Refs. [24, 25], which found no improvement in the H_0 tension. Thus, for the remainder of this work we focus on models in which $m_\phi^2 \gg |q^2|$ at energy scales relevant to the CMB. The intermediate regime, where the mediator mass is negligible for some CMB wave-numbers and not for others, is beyond the scope of this work. We therefore require a new degree of freedom with $m_\phi^2 \gg |q^2|$.

Throughout this cosmological epoch the neutrinos are relativistic, so the typical momentum transfer is $|q^2| \sim T_\nu^2$. Therefore, the expansion indicated by the arrow in Eq. (4) is valid (and we may neglect the momentum- and, hence, temperature-dependence in $G_{\rm eff}$) only if $m_\phi \gg T_\nu$. Comparing the values in Eq. (2) to the expression for $G_{\rm eff}$ in Eq. (5), we see that

$$m_\phi \simeq (4 - 200) \times |g_\phi|\,\text{MeV} \ , \tag{6}$$

so a new sub-GeV state is generically required to realize the self-interacting-neutrino solution to the H_0 tension. Since $T_\nu < 100\,\text{eV}$ at horizon entry of the highest momentum modes relevant for CMB anisotropies, the validity of the effective interaction in Eq. (1) in the analyses of Refs. [17–20, 22] requires $m_\phi \gtrsim$ keV (as already pointed out in Ref. [20]). From Eq. (6), this condition translates to

$$m_\phi \gtrsim \text{keV} \quad \Longrightarrow \quad |g_\phi| \gtrsim 10^{-4} . \tag{7}$$

Eqs. (6) and (7) bound the range of m_ϕ.

Finally, we note that the interaction in Eq. (3) is not gauge-invariant at energies above the scale of electroweak symmetry breaking (EWSB). This is because ν_α is a component of the $L_\alpha = (\nu_\alpha,\ \ell_\alpha)^T$ doublet of the $SU(2)_W$ group, whose gauge transformations rotate ν_α into the corresponding charged lepton ℓ_α. Thus, some particle with non-trivial electroweak quantum numbers must replace the interaction of Eq. (3) at energies above the Higgs vacuum expectation value $v \simeq 246\,\text{GeV}$. We defer a complete discussion of this issue to Sec V; for now we assume that the interaction of Eq. (1) is valid well below the electroweak scale.

III. COSMOLOGICAL BOUNDS

The successful predictions of BBN in standard cosmology (with only SM particle content) provide a powerful probe of additional light species. New particles in thermal equilibrium with neutrinos increase the Hubble expansion rate during BBN as extra relativistic degrees of freedom or by heating neutrinos relative to photons. Away from mass thresholds, both effects are captured by a constant shift in N_{eff}, the effective number of neutrinos. We find that the observed light element abundances constrain $\Delta N_{\text{eff}} < 0.5$ (0.7) at 95% CL for the SIν- (MIν-) preferred values of the baryon density, as explained in more detail in App. A.

Note that we do not use the Planck limit on ΔN_{eff} [4], since that result assumes free-streaming neutrinos. In fact, a large $\Delta N_{\text{eff}} \simeq 1$ at the time of the formation of the CMB is crucial for reconciling the MIν and SIν cosmologies with the observed CMB power spectrum [20]. The concordance of the CMB and BBN measurements of ΔN_{eff} within the self-interacting neutrino framework thus seems to require an injection of energy between nucleosynthesis and recombination; we remain agnostic on this point. Given this discussion, we conservatively only apply a constraint on ΔN_{eff} from considerations of BBN physics alone.

A. Mediators and ΔN_{eff}

The interaction in Eq. (3) induces $\phi \leftrightarrow \nu\nu$ decays and inverse decays; this process can bring ϕ into into thermal equilibrium with the neutrino bath before BBN.[1] From Eqs. (6) and (7), the mediator that generates sufficiently large neutrino self-interactions must have mass in the keV – few MeV range. Consequently, this particle is at least semi-relativistic when neutrinos decouple from photons at $T_{\text{dec}} \simeq$ MeV. We can estimate the resulting contribution to the expansion rate as parametrized by ΔN_{eff} assuming ϕ is relativistic:

$$\Delta N_{\text{eff}} \simeq \frac{8}{7}\left(\frac{T_\gamma}{T_\nu}\right)^4 \frac{\rho_\phi}{\rho_\gamma}, \tag{8}$$

where ρ_γ is the energy density of photons and T_γ is their temperature. If ϕ becomes non-relativistic before the end of BBN, ΔN_{eff} is increased further as ϕ decays raise the neutrino temperature (see, e.g., Ref. [33]). Here, we show that a mediator whose interactions realize the necessary G_{eff} in Eq. (2) necessarily equilibrates with neutrinos before T_{dec}.

Vector Mediators: If Eq. (1) arises from integrating out a vector particle ϕ_μ with mass m_ϕ, the Lagrangian at energies above m_ϕ is

$$\mathcal{L} \supset \frac{1}{2} m_\phi^2 \phi^\mu \phi_\mu + \left(g_\phi \phi_\mu \nu^\dagger \bar{\sigma}^\mu \nu + \text{h.c.}\right), \tag{9}$$

where g_ϕ is the dimensionless gauge coupling. At low energies, the interaction in Eq. (9) yields

$$G_{\text{eff}} \simeq \frac{g_\phi^2}{m_\phi^2} \simeq (10\ \text{MeV})^{-2}\left(\frac{g_\phi}{10^{-4}}\right)^2\left(\frac{\text{keV}}{m_\phi}\right)^2. \tag{10}$$

The vector equilibrates before T_{dec} via $\nu\nu \leftrightarrow \phi$ if the corresponding thermally averaged rate $\Gamma_{\nu\nu\to\phi}$ exceeds Hubble when $T = \max(T_{\text{dec}}, m_\phi)$:

$$\frac{\Gamma_{\nu\nu\to\phi}}{H} \sim \frac{g_\phi^2 m_\phi^2 M_{\text{Pl}}}{\max(T_{\text{dec}}, m_\phi)^3} > 10^8\,\frac{G_{\text{eff}}}{(10\ \text{MeV})^{-2}}, \tag{11}$$

where $M_{\text{Pl}} = 1.22\times10^{19}$ GeV and we have used Eqs. (10) and (7). Thus, this reaction is in equilibrium for all values of couplings and masses of interest. As a result, ϕ has a thermal number density at T_{dec} in both the MIν and SIν scenarios. Using Eq. (8), we find $\Delta N_{\text{eff}} = (8/7)(3/2) \simeq 1.7$ assuming ϕ remains relativistic throughout BBN. If, instead, ϕ becomes non-relativistic between T_{dec} and $\sim$ 30 keV at the end of BBN, then $\Delta N_{\text{eff}} \approx 2.5$, so ϕ must become non-relativistic well before T_{dec}. In Ref. [16] it was found that the Boltzmann suppression for massive vector particles is effective for $m_\phi > 10$ MeV at 95% CL. Using Eq. (10), this requires $g_\phi \gtrsim \mathcal{O}(0.1)$, which is excluded in all theoretically-consistent (or anomaly-free) vector models with neutrino couplings [34]. Anomaly-free vectors, such as those coupled to $L_e - L_i$ currents, where L_i is a lepton family number, would additionally introduce large $\bar{\nu}\nu\bar{e}e$ interactions which would likely spoil the CMB fit.

Finally, we note that Eq. (9) can arise if active neutrinos mix with heavy sterile-neutrinos, which couple directly to the vector ϕ_μ [35–37]. The mediators in these models are subject to the same BBN constraint above, but because their coupling to active neutrinos is derived from the active-sterile mixing angle, the resulting g_ϕ faces additional bounds from laboratory searches [38] that exclude the interaction strengths in Eq. (2) required to resolve the H_0 tension.

[1] Annihilation and scattering processes (e.g. $\nu\nu \leftrightarrow \phi\phi$ and $\nu\nu \leftrightarrow \nu\nu\phi$) also contribute to ϕ equilibration, but the corresponding rates are suppressed by additional powers of the small coupling $g_\phi < 1$, so their contribution is always subdominant.

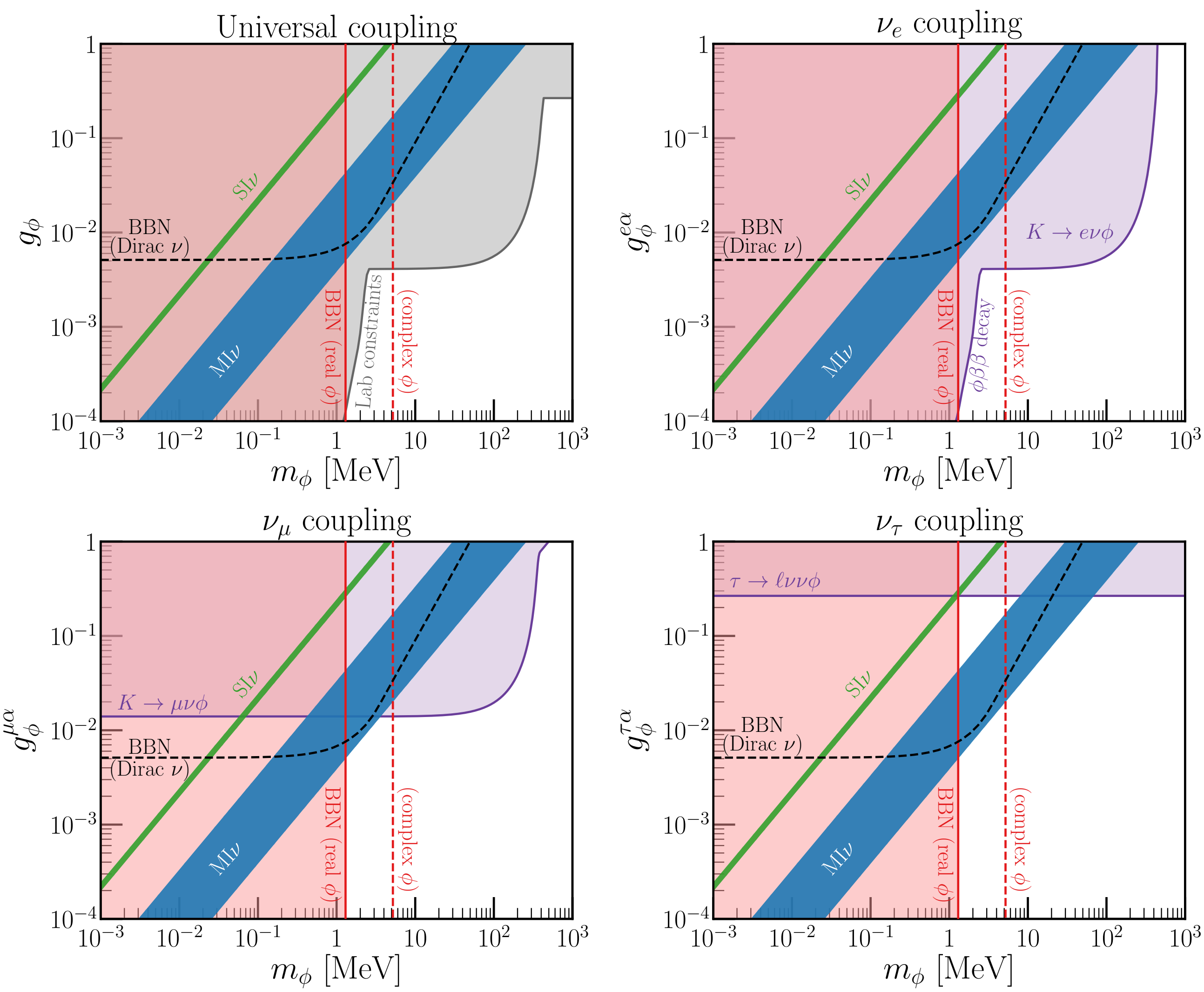

FIG. 2. Bounds on light neutrino-coupled mediators assuming flavor universal couplings (top-left). Bounds on flavor specific couplings to ν_e (top-right), ν_μ (bottom left), and ν_τ (bottom right). The diagonal bands labeled MIν and SIν are the preferred regions from Eq. (2) [20] translated into the g_ϕ-m_ϕ plane. Also shown are constraints from τ and rare meson decays [26–29], double-beta decay experiments [30–32], and BBN.

Scalar Mediators: By a similar argument, any scalar mediator ϕ that realizes $G_{\rm eff}$ from the interaction in Eq. (3) with $g_\phi \gtrsim 10^{-4}$ as required by Eq. (7) also has a thermal abundance at $T_{\rm dec}$. Relativistic scalars in equilibrium with neutrinos contribute $\Delta N_{\rm eff} = 0.57$ (1.1) for a real (complex) ϕ, which has one (two) degree(s) of freedom. As for the vector case, the ϕ density must become Boltzmann-suppressed before neutrino-photon decoupling, leading to a lower limit on m_ϕ. We use `AlterBBN` 2.1 [39, 40] as described in App. A to obtain the 95% CL lower bounds

$$m_\phi > \begin{cases} 1.3 \text{ MeV} & \text{(real scalar)} \\ 5.2 \text{ MeV} & \text{(complex scalar)} \end{cases} , \qquad (12)$$

for the SIν preferred values of the baryon density (the corresponding MIν bounds are somewhat weaker – see App. A). These bounds are are shown as red vertical lines labeled BBN in Fig. 2.

B. Excluding Dirac Neutrinos

If neutrinos are Dirac particles, then all neutrino masses arise from the Higgs-neutrino Yukawa interaction

$$\mathcal{L}_{\rm Dirac} \supset y_\nu H L \nu_R \to m_\nu \nu \nu_R \quad , \quad m_\nu \equiv y_\nu v/\sqrt{2} \quad , \qquad (13)$$

where H is the Higgs doublet, $L = (\nu,\ \ell)^T$ is a lepton doublet, ν_R is a right-handed neutrino (RHN), and flavor indices have been suppressed. The Weyl fermions ν and ν_R become Dirac partners after EWSB and acquire identical masses. In the absence of additional interactions, the Yukawa coupling $y_\nu \sim 10^{-12}(m_\nu/0.1\,\mathrm{eV})$ in Eq. (13) is insufficient to thermalize right handed helicity states, so the relic neutrinos consist exclusively of left-handed neutrinos and right-handed antineutrinos [41].

The light mediator interactions in Eq. (2) are much stronger than the weak force at late times, so ϕ and ν_R can both thermalize. Approximating the RHN production rate as $\Gamma_{\phi\to\nu\nu_R} \simeq (m_\nu/m_\phi)^2\Gamma_{\phi\to\nu\nu}$, for $m_\nu = 0.1$ eV we have

$$\frac{\Gamma_{\phi\to\nu\nu_R}}{H} \simeq \frac{g_\phi^2 m_\nu^2 M_{\rm Pl}}{m_\phi^3} = 10^6 \frac{G_{\rm eff}}{(10\,{\rm MeV})^{-2}}\frac{\rm MeV}{m_\phi}, \quad (14)$$

where $T = m_\phi \gtrsim T_{\rm dec}$ is the temperature at which RHN production is maximized relative to Hubble expansion. For a more detailed discussion, see App. B.

Observations of neutrino oscillations require that at least two of the light neutrinos are massive, with one heavier than $\sim 10^{-2}$ eV and one heavier than $\sim 10^{-1}$ eV [42]. Thus, for all values of m_ϕ we consider in the SIν range, at least one RHN will thermalize before BBN. For this reason, we assume that neutrinos are Majorana particles for the remainder of this work.

C. Supernova 1987A

A light new particle with appreciable couplings to neutrinos can alter the qualitative behavior of neutrino emission that was observed in the aftermath of Supernova 1987A. We find that for the parameter space shown in Fig. 2 the ϕ particle is trapped and does not drain energy from the supernova. See App. C for further discussion.

IV. LABORATORY BOUNDS

Because terrestrial experiments routinely reach energies above the MeV scale, there are many possible probes of the models we consider. Here we specialize to CP-even scalar mediators without loss of essential generality; we comment on CP-odd scalars in Sec. V. The main constraints are as follows:

Double Beta Decay: If $g_\phi^{ee} \neq 0$ in Eq. (3) and ϕ is lighter than the Q-value of a double-beta-decaying nucleus, the process $(Z, A) \to (Z+2, A)\, e^- e^- \phi$ may occur, contributing to measured $2\nu\beta\beta$ rates. Measurements of this process constrain $|g_\phi^{ee}| \lesssim 10^{-4}$ if $m_\phi \lesssim 2$ MeV [30–32]. This constraint is shown in the top row of Fig. 2.

Meson Decays: Nonzero $g_\phi^{\alpha\beta}$ can allow for the meson decays $\mathfrak{m}^\pm \to \ell_\alpha^\pm \nu_\beta \phi$ if $m_\phi < m_{\mathfrak{m}} - m_{\ell_\alpha}$ [26–29]. Note that, unlike $\mathfrak{m} \to \ell\nu$ decays, this process is not helicity suppressed. The strongest constraints come from kaon decay. The measured ratio of ${\rm Br}(K^+ \to e^+\nu_e)/{\rm Br}(K^+ \to \mu^+\nu_\mu) = (2.416 \pm 0.043) \times 10^{-5}$ constrains $g_\phi^{e\beta}$ as shown in the top row of Fig. 2 [43, 44]. The limit on ${\rm Br}(K^+ \to \mu^+\nu_\mu\nu\bar{\nu}) < 2.4 \times 10^{-6}$ [45] constrains $g_\phi^{\mu\alpha}$ as shown by the purple-shaded region in bottom-left panel of Fig. 2. Non-zero $g_\phi^{\mu\alpha}$ still contributes when $m_\phi > m_{K^+} - m_\mu$ via off-shell ϕ [28]. Note that, if ϕ decays to visible matter via $\phi \to \gamma\gamma$ or $\phi \to e^+e^-$ through some other interaction (e.g. ϕ-higgs mixing), the branching ratio must be subdominant to $\phi \to \nu\nu$ decays to avoid tight $\nu - \gamma$ coupling below $T \sim$ MeV, which would spoil the solution to the H_0 tension. Thus, the laboratory bounds presented here are unaffected by other possible couplings (see [46] for a review).

τ **Decays:** Similar to meson decays, ϕ that couples to ν_τ with $g_\phi^{\tau\beta}$ allows for decays $\tau^- \to \ell_\beta \bar{\nu}_\beta \bar{\nu}_\tau \phi$. For $m_\phi \ll m_\tau - m_{\ell_\beta}$, this constrains $g_\phi^{\tau\beta} \lesssim 0.3$ [43]. This is depicted as a purple band in the bottom-right panel of Fig. 2.

In Fig. 2 we summarize our main findings: values of $G_{\rm eff}$ from Eq. (2) favored by the H_0 tension are excluded if ϕ couples universally to all flavor- and mass-eigenstates (top left panel), which was explicitly considered in Refs. [17–20, 22], or if ϕ couples predominantly to ν_e or ν_μ (top-right and bottom-left panels, respectively). Similarly, we can rule out the possibility that ϕ couples to any single mass-eigenstate neutrino, since the ν_e- and ν_μ-composition of each mass eigenstate is similar. Moreover, in this case, the collisional Boltzmann equations would be much more complicated to solve (because different eigenstates will start to free-stream at different times), and the results of Refs. [17–20, 22] may not apply.

However, a *flavor-restricted* coupling leads to approximately the same neutrino mass-eigenstate interactions as in Refs. [17–20, 22], since the flavor eigenstates are well-mixed in the mass basis. In particular, a τ-only coupling, in which the matrix $g_{\alpha\beta}$ is zero except for the $g_{\tau\tau}$ entry, is potentially viable since τ decays are less constraining than meson decays. Thus, we are unable to fully exclude an interaction $G_{\rm eff}^\tau \bar{\nu}_\tau \nu_\tau \bar{\nu}_\tau \nu_\tau$. In this case, $G_{\rm eff}^\tau = A \times G_{\rm eff}$ for $G_{\rm eff}$ defined in Eq. (5) and $A \sim \mathcal{O}(1)$ is a constant that accounts for the reduced scattering probability of each mass eigenstate. However, because mixed mass-eigenstate vertices are possible in this scenario, there are additional diagrams compared to the mass-diagonal case. For this reason, we caution that the effect on the CMB anisotropies of flavor-specific neutrino self-interactions can be mildly different than that considered in Refs. [17–20, 22]. Nonetheless, we expect that the preferred coupling range should shift slightly *up* relative to the flavor-universal case. For this reason, the SIν range is ruled out regardless of choice of flavor structure in Eq. (3). The MIν range is still allowed in a τ-flavor-only scenario, though a dedicated study is necessary and well motivated.

V. POSSIBLE ULTRAVIOLET COMPLETIONS

In this section we quantify the difficulty of realizing the operator in Eq. (1) from renormalizable, gauge invariant interactions that also accommodate neutrino masses and mixings. In light of the discussion in Sec III, we only consider models of Majorana neutrinos. We focus on the type-I and type-II seesaw mechanisms of Majorana neutrino mass generation with an additional scalar particle ϕ. In both cases, we find the resulting $\phi\nu\nu$ Yukawa coupling is suppressed by factors of the light neutrino mass. In these minimal models, it is therefore impossible to simultaneously generate neutrino masses and a large enough $G_{\rm eff}$ in Eq. (2) to address the H_0 tension.

A. Type-I Seesaw

We consider the Lagrangian

$$\mathcal{L}_{\rm I} \supset y_N LHN + \frac{1}{2}\left(M_N + \lambda_N \phi\right) NN + \text{h.c.}, \tag{15}$$

where we have suppressed flavor indices. After EWSB, the first term in Eq. (15) and the $M_N NN$ term generate light neutrino masses after N is integrated out:

$$m_\nu \approx -\frac{y_N^2 v^2}{2M_N}. \tag{16}$$

These interactions also generate the coupling of ϕ to LH neutrinos in Eq. 3 with

$$g_\phi = \lambda_N \frac{m}{M_N} \Rightarrow G_{\rm eff} = \frac{\lambda_N^2 m_\nu^2}{m_\phi^2 M_N^2}. \tag{17}$$

Thus, realizing $G_{\rm eff} \approx (4 - 300\ \text{MeV})^{-2}$ requires $M_N \sim 10^3 m_\nu \sim 10$ eV. If this is true, the RHNs thermalize before BBN and spoil $\Delta N_{\rm eff}$ as discussed in analogy to the Dirac case discussed in Sec. III B. Therefore, a minimal scenario where the same Type-I Seesaw generates the neutrino mass and the operator in Eq. (1) with the magnitude in Eq. (2) is not possible.

We also consider the case in which, instead of M_N violating lepton number explicitly, we replace the term in parentheses in Eq. (15) with a complex scalar Φ. This acquires a vacuum expectation value (VEV) of $\langle\Phi\rangle = f\exp(i\phi/f)$, dynamically generating a mass for N [47–49]. The light neutrinos couple to the pseudo-Goldstone boson ϕ, the Majoron, of this symmetry breaking. The coupling is of the form $i(m_i/f)\phi\nu_i\nu_i$, so the Majoron is a pseudoscalar mediator. We note that our results from the previous sections apply to the Majoron, since all processes we have considered involve relativistic neutrinos for which the distinction between scalar and pseudoscalar ϕ is irrelevant. As with the above case, we see that the neutrino self-interaction strength will be suppressed by the light neutrino masses.

B. Type-II Seesaw

We augment the Standard Model with a complex $SU(2)_W$ triplet T. The type-II Seesaw Lagrangian is

$$\begin{aligned}\mathcal{L}_{\rm II} \supset\, & y_T LTL + (\mu_T + \lambda_T\phi) H T^\dagger H + \text{h.c.} \\ & - M_T^2\, \text{Tr}(T^\dagger T),\end{aligned} \tag{18}$$

where the dimensionful coupling μ_T is a soft breaking of Lepton number. After integrating out T and EWSB, neutrinos acquire masses of order

$$m_\nu \approx \frac{y_T \mu_T v^2}{M_T^2}, \tag{19}$$

and a coupling to ϕ of the form of Eq. 3 with

$$g_\phi = \lambda_T \frac{m_\nu}{\mu_T} \Rightarrow G_{\rm eff} = \frac{\lambda_T^2 m_\nu^2}{m_\phi^2 \mu_T^2}. \tag{20}$$

The scale μ_T is bounded by non-observation of rare lepton-number-violating processes such as $\mu \to e\gamma$ or $\tau \to \mu ee$ [50–52]. Thus, for $m_\phi >$ keV, the m_ν^2/μ_T^2 suppression is too severe to permit values of $G_{\rm eff}$ in Eq. (2), as required to resolve the H_0 tension.

As with the Type-I Seesaw, we may consider the Majoron case where $\mu_T + \lambda_T\phi$ is replaced by a field Φ which acquires a VEV. The neutrinos couple to the pseudo-Goldstone boson of this lepton number symmetry breaking, with coupling of the form $i(m_i/f)\phi\nu_i\nu_i$. Again, this coupling is highly suppressed.

We conclude that the ϕ coupling to active neutrinos cannot arise from interactions with the same field that generates neutrino masses (the RH neutrino N in Type-I seesaw and the triplet T in Type-II), without being suppressed by neutrino masses. A large enough $G_{\rm eff}$ can therefore be obtained using *separate* seesaw mechanisms to generate the neutrino masses and the $\phi\nu\nu$ interaction. For instance, we can use the Type-I Seesaw with $\lambda_N = 0$ for the light neutrino masses and add the Type-II Seesaw with $\mu_T = 0$ to allow for large $g_\phi^{\tau\tau}$. In this scenario, the size of $g_\phi^{\tau\tau}$ is independent of the neutrino masses.

VI. CONCLUDING REMARKS

In this work we have shown that the self-interacting neutrino explanation of the H_0 tension requires the existence of a light $\sim$ MeV-scale mediator, which is subject to stringent cosmological and laboratory bounds. Consequently, most realizations of this scenario are robustly excluded by the BBN-only bounds on $\Delta N_{\rm eff}$, accelerator constraints for rare K decays, and precise results from neutrinoless double-beta decay searches. If ν self-interactions resolve the H_0 tension, they cannot viably arise in models with vector-mediated forces or with Dirac neutrinos, and they cannot reside in the strongly interacting parameter space in Eq. (2).

Our analysis identifies viable phenomenological models with a ~ 10 MeV scalar mediator and large $\mathcal{O}(0.1)$ couplings to τ-flavored neutrinos only. However, realizing such interactions in UV-complete, gauge-invariant models is challenging. We find that sufficiently strong interactions cannot arise in models that generate neutrino masses via a single Type-I or -II seesaw mechanism: the resulting neutrino-scalar coupling is suppressed by factors of m_ν/Λ where $\Lambda \gg m_\nu$ is the appropriate seesaw scale.

Looking forward, a dedicated exploration of the τ-only scattering scenario favored by our analysis will be necessary to see if the moderately interacting-scenario is indeed able to resolve the H_0 tension while not running afoul of laboratory measurements. Our results also motivate exploration of the "intermediate" mediator-mass regime, where neutrino scattering is relevant for a partial range of redshifts explored by the CMB.

ACKNOWLEDGMENTS

We thank André de Gouvêa, Francis-Yan Cyr-Racine, Joshua Isaacson, Martina Gerbino, Stefan Höche, Massimiliano Lattanzi, Kohta Murase, Jessica Turner, and Yue Zhang for helpful conversations. This manuscript has been authored 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. The United States Government retains and the publisher, by accepting the article for publication, acknowledges that the United States Government retains a non-exclusive, paid-up, irrevocable, world-wide license to publish or reproduce the published form of this manuscript, or allow others to do so, for United States Government purposes.

Appendix A: Calculating $\Delta N_{\rm eff}$

The Hubble expansion rate at the time of BBN is sensitive to the energy density in neutrinos and other relativistic species when photon temperatures are below an MeV [53–55]. New relativistic particles or an injection of energy into the Standard Model neutrino bath (via the decays or annihilations of a ν-coupled species) can increase the Hubble expansion rate at this time. A larger expansion rate at the time of BBN modifies the neutron-to-proton ratio and the freeze-out of deuterium-burning reactions, leading to larger yields of Helium-4 and Deuterium. The primordial abundances of these elements are measured in pristine gas clouds to be $Y_p = 0.2449 \pm 0.004$ [56] and 10^5 D/H $= 2.527 \pm 0.03$ [57], respectively. These observations are in good agreement with predictions assuming only Standard Model particle content at the time of BBN [58].

We use the measurements of Y_p and D/H to constrain modifications to the expansion rate using the BBN Boltzmann code `AlterBBN` 2.1 [39, 40]. We follow the Monte Carlo procedure outlined in Ref. [59] to estimate theoretical uncertainties from nuclear reaction rates. In deriving limits we marginalize over a Gaussian prior on the baryon density η_b corresponding to the best-fit points of ΛCDM [4], SIν and MIν [20]. Using the results of Refs. [58, 60] to convert $\Omega_b h^2$ to η_b, we find that $10^{10}\eta_b = 6.133 \pm 0.027$, 6.146 ± 0.082, and 6.248 ± 0.082 in ΛCDM, SIν and MIν, respectively. Note that the best fit values of η_b in these cosmologies are all compatible with the range $5.8 < 10^{10}\eta_b < 6.6$ (95 % CL) extracted only from BBN data [45].

If the new particles remain relativistic throughout nucleosynthesis, their modification of the Hubble rate is specified by a constant shift of the number of relativistic degrees of freedom, $N_{\rm eff}$. We find that the observed values of Y_p and D/H favor the values of $N_{\rm eff}$ shown in Tab. I for the different cosmologies. Note that the ΛCDM result is higher than reported in, e.g., Ref. [58], and it has a smaller uncertainty. This is because of updated D burning rates and observed abundances used in our analysis; for further discussion, see Ref. [59]. These improvements actually weaken the BBN constraint compared to Ref. [58] because of the slight preference for $N_{\rm eff} > 3$. The upper limits in SIν and MIν are weaker still, due to their larger central values and uncertainties of the baryon density: a larger baryon density reduces the Deuterium yield, which can be compensated by increasing $N_{\rm eff}$ [58].

The upper bounds in Tab. I apply to light particles that are fully relativistic at the time of nucleosynthesis. If these particles have masses at the MeV scale, then their decays or annihilations can heat the neutrinos relative to photons; if they are much heavier, they transfer their entropy to neutrinos while those are still in equilibrium with photons. The resulting change in neutrino temperature and the corresponding $\Delta N_{\rm eff}$ can be estimated assuming instantaneous neutrino-photon decoupling at $T_{\rm dec} \approx 1-2$ MeV and by using entropy conservation [33]. This crude estimate along with Tab. I suggests that neutrino-coupled scalars in the SIν cosmology with $m_\phi \lesssim 2-7$ MeV should be incompatible with the observed light element abundances, where the range corresponds to varying $T_{\rm dec}$ and the number of scalar degrees of freedom from 1 (real scalar) to 2 (complex scalar). The results of a full calculation using `AlterBBN` (which does not make these approximations), shown in Tab. II, are compatible with this estimate.

We point out that the contribution of a scalar to $\Delta N_{\rm eff}$ at the 95% CL limit from our BBN analysis is also compatible with the extra radiation density at the best fit point in Ref. [20] at only the 2σ level. In order to be in better agreement with $\Delta N_{\rm eff} \simeq 1.02 \pm 0.29$ for the SIν mode, new semi-relativistic degrees of freedom may need to come into equilibrium with the neutrino bath after BBN is complete [59].

Model	$N_{\rm eff}$	95% upper limit
ΛCDM	3.27 ± 0.137	3.54
SIν	3.27 ± 0.14	3.56
MIν	3.43 ± 0.13	3.72

TABLE I. Preferred values and upper limits on the effective number of neutrino species, $N_{\rm eff}$, from primordial nucleosynthesis. The cosmological models differ through their values of the baryon density parameter η_b (and its uncertainty) determined from the CMB power spectrum as described in the text.

Model	95% CL lower bound on m_ϕ (MeV)	
	real ϕ	complex ϕ
ΛCDM	2.0	6.1
SIν	1.3	5.2
MIν	0.6	3.7

TABLE II. Lower limits on mediator mass m_ϕ from primordial nucleosynthesis, assuming ϕ was in thermal equilibrium before BBN. The cosmological models differ through their values of the baryon density parameter η_b (and its uncertainty) as determined from the CMB power spectrum.

Appendix B: Dirac Neutrino Thermalization

If neutrinos are Dirac fermions, their right-handed (RH) components must not come into equilibrium with the Standard Model plasma before ~ 1 MeV, since they would give $\Delta N_{\rm eff} \gtrsim 1$ during nucleosynthesis, which is clearly incompatible with the results of App. A. In this section we derive the conditions for RH neutrino thermalization.

In the Standard Model, RH neutrinos (neutrinos with the "wrong" helicity) can be created in any interaction that produces LH neutrinos, i.e. weak interactions. The characteristic production rate of RH neutrinos in scattering reactions is $\Gamma_{\rm RH} \sim G_{\rm F}^2 T^5 (m_\nu/T)^2$ [41]. If weak interactions are the only interactions producing neutrinos, then RHNs do not thermalize as long as $m_\nu \lesssim \mathcal{O}(100~\text{keV})$, which is comfortably satisfied in the Standard Model. However, the interaction introduced in Eqs. (1) and (2) is many orders of magnitude stronger than its weak counterpart, and the resulting RH neutrino production rate is much larger.

Thermalization is described by a Boltzmann equation of the form

$$\dot{\rho}_{\nu_R} + 4H\rho_{\nu_R} = C, \tag{B1}$$

where ρ_{ν_R} is the energy density in "wrong helicity" neutrinos and C is the collision term encoding the processes that produce these neutrinos. Thermalization/equilibration occurs when $C \sim 4H\rho_R^{\rm eq}$. The dominant process responsible for RH neutrino production is the decay $\phi \to \nu\nu(\lambda = +1/2)$, where $\lambda = +1/2$ $(-1/2)$ is a helicity label corresponding to RH (LH) neutrinos; scattering reactions such as $\nu\nu \to \nu\nu(\lambda)$ are suppressed by an additional factor of g_ϕ^2. We therefore evaluate the collision term for $\phi \to \nu\nu(\lambda = +1/2)$:

$$C \approx 2 \int d\Phi_3 E_{\nu_R} |\mathcal{M}(\lambda = +1/2)|^2 f_\phi (1 - f_\nu), \tag{B2}$$

where $d\Phi_3$ is the Lorentz-invariant phase-space (including the momentum conservation delta function) and f_ϕ, f_ν are the phase-space distributions of ϕ and LH neutrinos. We have neglected the inverse decay contribution and the RH neutrino Pauli-blocking factors. These approximations are adequate for estimating the onset of equilibrium, assuming the initial abundance of RH neutrinos is negligible, i.e. $f_{\nu_R} \ll 1$. We have also multiplied the right-hand side by two to account for ϕ^* decays which can also produce wrong helicity neutrinos.

We use the Lagrangian in Eq. (3) (with a complex ϕ for a Dirac neutrino) to evaluate the matrix element squared in the plasma frame, *without* summing over one of the helicities [61]

$$|\mathcal{M}(\lambda)|^2 \approx g_\phi^2 \begin{cases} \frac{E_\nu}{2E_{\nu_R}} (1 + \cos\theta) m_\nu^2 & \lambda = +1/2 \\ m_\phi^2 & \lambda = -1/2 \end{cases}, \tag{B3}$$

where θ is the angle between the ν_R and ϕ direction of motion, and we have kept only the leading terms in m_ν/m_ϕ. Note that when $\cos\theta = -1$, the "wrong" helicity amplitude vanishes at this order as a result of angular momentum conservation (the next term is $\propto m_\nu^4$). Using this result in Eq. (B2), we find that

$$C \approx \frac{g_\phi^2 m_\nu^2}{32\pi} \left(\frac{2\zeta(3) T^3}{\pi^2} \right) \mathcal{C}_+(m_\phi/T), \tag{B4}$$

where the factor in parentheses is the number density of relativistic ϕ and the function $\mathcal{C}_+$ is ~ 1 for $m_\phi/T < 1$ and becomes Boltzmann-suppressed for $m_\phi/T > 1$.

We can apply the nucleosynthesis bound as computed in Appendix A if thermalization occurs before $T_{\rm dec} \sim 1$ MeV. By using the approximate thermalization criterion below Eq. (B1), we find that RH neutrinos thermalize before BBN if

$$g_\phi \gtrsim 5 \times 10^{-3} \left(\frac{\max(m_\phi, T_{\rm dec})}{\rm MeV} \right)^{3/2} \left(\frac{0.1~\text{eV}}{m_\nu} \right). \tag{B5}$$

This bound (evaluated using the full numerical $\mathcal{C}_+$) is shown in Fig. 2 as a black dashed line.

Appendix C: Supernova 1987A

A new weakly coupled particle can change the behavior of the neutrino emission that was observed from the explosion of Supernova 1987A. The proto-neutron star cooling phase that was observed in large water Cherenkov detectors was qualitatively similar to the Standard Model-only expectation [62]. If a new particle species X carried away too much energy during the proto-neutron star cooling phase, the time over which neutrinos arrived would have been unacceptably reduced [63, 64]. A semi-analytic criterion that the luminosity of this particle should obey is $L_X \leq L_\nu = 3 \times 10^{52}\,\mathrm{erg}\,/\,\mathrm{s}$ at times of order 1 second after the core bounce [65]. Following the procedure described in more detail in [66], we have

$$L_\phi = \int_0^{R_\nu} dV \int \frac{d^3k}{(2\pi)^3}\, \omega \Gamma_\phi^{\mathrm{prod}} \exp\left(- \int_r^{R_{\mathrm{g}}} \Gamma_\phi^{\mathrm{abs}} dr' \right), \quad \text{(C1)}$$

where: the ϕ has four-momentum $(\omega, \vec{k})$; $\Gamma_\phi^{\mathrm{abs}}$ is the ϕ absorptive width; $\Gamma_\phi^{\mathrm{prod}}$ is the ϕ production rate, which is related to the absorptive width in equilibrium by $\Gamma_\phi^{\mathrm{prod}} = \exp(-\omega/T)\Gamma_\phi^{\mathrm{abs}}$; R_ν is the radius of the neutrinosphere, outside of which neutrinos free-stream; and $R_{\mathrm{g}} = 100\,\mathrm{km}$ is the radius inside of which neutrinos gain energy on average in elastic scattering events. We calculate Γ_ϕ including ϕ decay and ϕ annihilation to neutrino pairs, both of which are important for the masses of interest. We have not included contributions from the neutrino effective potentials, which may be significant at small m_ϕ [32]. We have also neglected neutrino Pauli blocking in Γ_ϕ, which is important near the proto-neutron star core, since this will become unimportant between R_ν and R_g.

We find that L_ϕ given by Eq. (C1) exceeds L_ν if g_ϕ is roughly in the range

$$g_\phi^{\mathrm{excl}} \simeq \left(5 \times 10^{-6} - 6 \times 10^{-5}\right) \times \frac{1}{1 + m_\phi/\mathrm{keV}}. \quad \text{(C2)}$$

The sharp change in the shape of the bound at $m_\phi \simeq$ keV is due to the fact that rate of decay and inverse decay, $\Gamma_{\phi \leftrightarrow \nu\nu} \propto g_\phi^2 m_\phi^2 / T$, becomes subdominant to the annihilation rate, $\Gamma_{\phi\phi \leftrightarrow \nu\nu} \propto g_\phi^4 T$, for masses $m_\phi \lesssim g_\phi T_c$, where $T_c \sim \mathcal{O}(30\,\mathrm{MeV})$ is the core temperature. These bounds are approximately compatible with those shown in [32] at masses above 10 keV. From Eq. (C2), we see that bounds arising from the luminosity of ϕ particles from Supernova 1987A are generally below the coupling range of interest in this work.

It is also interesting to understand the constraints on g_ϕ^{ee} arising from deleptonization of the core, which have been obtained in the $\sim \mathcal{O}(\mathrm{MeV})$ mass range in [32] and which approximately overlap the range in Eq. (C2). At lower masses, this likely has an effect on the early phases of collapse and the collapse progenitor. Such a possibility was suggested in [67] and was studied in the aftermath of Supernova 1987A by [68]. This latter study found a constraint $g_\phi^{ee} < 3 \times 10^{-4}$, neglecting any m_ϕ-dependence. Because these bounds are determined by physics at the beginning of the core collapse, when temperatures are $\sim \mathcal{O}(\mathrm{few\,MeV})$, the bound likely cuts off at $\sim 5\,\mathrm{MeV}$, similar to the $0\nu\beta\beta$ bounds cited above. A detailed study is of interest, but beyond the scope of this work.

[1] A. G. Riess *et al.*, Astrophys. J. **826**, 56 (2016), arXiv:1604.01424 [astro-ph.CO].

[2] T. Shanks, L. Hogarth, and N. Metcalfe, Mon. Not. Roy. Astron. Soc. **484**, L64 (2019), arXiv:1810.02595 [astro-ph.CO].

[3] A. G. Riess, S. Casertano, D. Kenworthy, D. Scolnic, and L. Macri, (2018), arXiv:1810.03526 [astro-ph.CO].

[4] N. Aghanim *et al.* (Planck), (2018), arXiv:1807.06209 [astro-ph.CO].

[5] A. G. Riess, S. Casertano, W. Yuan, L. M. Macri, and D. Scolnic, (2019), arXiv:1903.07603 [astro-ph.CO].

[6] M. Vonlanthen, S. Rsnen, and R. Durrer, JCAP **1008**, 023 (2010), arXiv:1003.0810 [astro-ph.CO].

[7] L. Verde, E. Bellini, C. Pigozzo, A. F. Heavens, and R. Jimenez, JCAP **1704**, 023 (2017), arXiv:1611.00376 [astro-ph.CO].

[8] J. Evslin, A. A. Sen, and Ruchika, Phys. Rev. **D97**, 103511 (2018), arXiv:1711.01051 [astro-ph.CO].

[9] K. Aylor, M. Joy, L. Knox, M. Millea, S. Raghunathan, and W. L. K. Wu, Astrophys. J. **874**, 4 (2019), arXiv:1811.00537 [astro-ph.CO].

[10] J. Lesgourgues, G. Marques-Tavares, and M. Schmaltz, JCAP **1602**, 037 (2016), arXiv:1507.04351 [astro-ph.CO].

[11] E. Di Valentino, C. Behm, E. Hivon, and F. R. Bouchet, Phys. Rev. **D97**, 043513 (2018), arXiv:1710.02559 [astro-ph.CO].

[12] V. Poulin, T. L. Smith, D. Grin, T. Karwal, and M. Kamionkowski, Phys. Rev. **D98**, 083525 (2018), arXiv:1806.10608 [astro-ph.CO].

[13] F. D'Eramo, R. Z. Ferreira, A. Notari, and J. L. Bernal, JCAP **1811**, 014 (2018), arXiv:1808.07430 [hep-ph].

[14] V. Poulin, T. L. Smith, T. Karwal, and M. Kamionkowski, (2018), arXiv:1811.04083 [astro-ph.CO].

[15] K. L. Pandey, T. Karwal, and S. Das, (2019), arXiv:1902.10636 [astro-ph.CO].

[16] M. Escudero, D. Hooper, G. Krnjaic, and M. Pierre, JHEP **03**, 071 (2019), arXiv:1901.02010 [hep-ph].

[17] F.-Y. Cyr-Racine and K. Sigurdson, Phys. Rev. **D90**, 123533 (2014), arXiv:1306.1536 [astro-ph.CO].

[18] L. Lancaster, F.-Y. Cyr-Racine, L. Knox, and Z. Pan, JCAP **1707**, 033 (2017), arXiv:1704.06657 [astro-ph.CO].

[19] I. M. Oldengott, T. Tram, C. Rampf, and Y. Y. Y. Wong, JCAP **1711**, 027 (2017), arXiv:1706.02123 [astro-ph.CO].

[20] C. D. Kreisch, F.-Y. Cyr-Racine, and O. Dor, (2019), arXiv:1902.00534 [astro-ph.CO].

[21] S. Bashinsky and U. Seljak, Phys. Rev. **D69**, 083002 (2004), arXiv:astro-ph/0310198 [astro-ph].

[22] G. A. Barenboim, P. B. Denton, and I. M. Oldengott, Phys. Rev. **D99**, 083515 (2019), arXiv:1903.02036 [astro-ph.CO].

[23] T. Cohen, (2019), arXiv:1903.03622 [hep-ph].
[24] F. Forastieri, M. Lattanzi, and P. Natoli, JCAP **1507**, 014 (2015), arXiv:1504.04999 [astro-ph.CO].
[25] F. Forastieri, M. Lattanzi, and P. Natoli, (2019), arXiv:1904.07810 [astro-ph.CO].
[26] K. Blum, A. Hook, and K. Murase, (2014), arXiv:1408.3799 [hep-ph].
[27] J. M. Berryman, A. De Gouvêa, K. J. Kelly, and Y. Zhang, Phys. Rev. **D97**, 075030 (2018), arXiv:1802.00009 [hep-ph].
[28] K. J. Kelly and Y. Zhang, Phys. Rev. **D99**, 055034 (2019), arXiv:1901.01259 [hep-ph].
[29] G. Krnjaic, G. Marques-Tavares, D. Redigolo, and K. Tobioka, (2019), arXiv:1902.07715 [hep-ph].
[30] M. Agostini *et al.*, Eur. Phys. J. **C75**, 416 (2015), arXiv:1501.02345 [nucl-ex].
[31] K. Blum, Y. Nir, and M. Shavit, Phys. Lett. **B785**, 354 (2018), arXiv:1802.08019 [hep-ph].
[32] T. Brune and H. Ps, (2018), arXiv:1808.08158 [hep-ph].
[33] C. Boehm, M. J. Dolan, and C. McCabe, JCAP **1212**, 027 (2012), arXiv:1207.0497 [astro-ph.CO].
[34] M. Bauer, P. Foldenauer, and J. Jaeckel, JHEP **07**, 094 (2018), arXiv:1803.05466 [hep-ph].
[35] M. Schmaltz and N. Weiner, JHEP **02**, 105 (2019), arXiv:1709.09164 [hep-ph].
[36] E. Bertuzzo, S. Jana, P. A. N. Machado, and R. Zukanovich Funchal, Phys. Rev. Lett. **121**, 241801 (2018), arXiv:1807.09877 [hep-ph].
[37] P. Ballett, S. Pascoli, and M. Ross-Lonergan, Phys. Rev. **D99**, 071701 (2019), arXiv:1808.02915 [hep-ph].
[38] A. de Gouvêa and A. Kobach, Phys. Rev. **D93**, 033005 (2016), arXiv:1511.00683 [hep-ph].
[39] A. Arbey, Comput. Phys. Commun. **183**, 1822 (2012), arXiv:1106.1363 [astro-ph.CO].
[40] A. Arbey, J. Auffinger, K. P. Hickerson, and E. S. Jenssen, (2018), arXiv:1806.11095 [astro-ph.CO].
[41] A. D. Dolgov, K. Kainulainen, and I. Z. Rothstein, Phys. Rev. **D51**, 4129 (1995), arXiv:hep-ph/9407395 [hep-ph].
[42] I. Esteban, M. C. Gonzalez-Garcia, A. Hernandez-Cabezudo, M. Maltoni, and T. Schwetz, JHEP **01**, 106 (2019), arXiv:1811.05487 [hep-ph].
[43] A. P. Lessa and O. L. G. Peres, Phys. Rev. **D75**, 094001 (2007), arXiv:hep-ph/0701068 [hep-ph].
[44] L. Fiorini (NA48/2), *Tau lepton physics. Proceedings, 9th International Workshop, TAU'06, Pisa, Italy, September 19-22, 2006*, Nucl. Phys. Proc. Suppl. **169**, 205 (2007), [,205(2007)].
[45] M. Tanabashi *et al.* (Particle Data Group), Phys. Rev. **D98**, 030001 (2018).
[46] K. C. Y. Ng and J. F. Beacom, Phys. Rev. **D90**, 065035 (2014), [Erratum: Phys. Rev.D90,no.8,089904(2014)], arXiv:1404.2288 [astro-ph.HE].
[47] Y. Chikashige, R. N. Mohapatra, and R. D. Peccei, Phys. Lett. **98B**, 265 (1981).
[48] Y. Chikashige, R. N. Mohapatra, and R. D. Peccei, Phys. Rev. Lett. **45**, 1926 (1980), [,921(1980)].
[49] J. Schechter and J. W. F. Valle, Phys. Rev. **D25**, 774 (1982).
[50] A. G. Akeroyd, M. Aoki, and H. Sugiyama, Phys. Rev. **D79**, 113010 (2009), arXiv:0904.3640 [hep-ph].
[51] D. N. Dinh, A. Ibarra, E. Molinaro, and S. T. Petcov, JHEP **08**, 125 (2012), [Erratum: JHEP09,023(2013)], arXiv:1205.4671 [hep-ph].
[52] P. S. B. Dev, C. M. Vila, and W. Rodejohann, Nucl. Phys. **B921**, 436 (2017), arXiv:1703.00828 [hep-ph].
[53] P. D. Serpico and G. G. Raffelt, Phys. Rev. **D70**, 043526 (2004), arXiv:astro-ph/0403417 [astro-ph].
[54] C. Boehm, M. J. Dolan, and C. McCabe, JCAP **1308**, 041 (2013), arXiv:1303.6270 [hep-ph].
[55] K. M. Nollett and G. Steigman, Phys. Rev. **D91**, 083505 (2015), arXiv:1411.6005 [astro-ph.CO].
[56] E. Aver, K. A. Olive, and E. D. Skillman, JCAP **1507**, 011 (2015), arXiv:1503.08146 [astro-ph.CO].
[57] R. J. Cooke, M. Pettini, and C. C. Steidel, Astrophys. J. **855**, 102 (2018), arXiv:1710.11129 [astro-ph.CO].
[58] R. H. Cyburt, B. D. Fields, K. A. Olive, and T.-H. Yeh, Rev. Mod. Phys. **88**, 015004 (2016), arXiv:1505.01076 [astro-ph.CO].
[59] A. Berlin, N. Blinov, and S. W. Li, (2019), arXiv:1904.04256 [hep-ph].
[60] G. Steigman, JCAP **0610**, 016 (2006), arXiv:astro-ph/0606206 [astro-ph].
[61] H. K. Dreiner, H. E. Haber, and S. P. Martin, Phys. Rept. **494**, 1 (2010), arXiv:0812.1594 [hep-ph].
[62] A. Burrows and J. M. Lattimer, Astrophys. J. **318**, L63 (1987).
[63] A. Burrows, M. S. Turner, and R. P. Brinkmann, Phys. Rev. **D39**, 1020 (1989).
[64] A. Burrows, M. T. Ressell, and M. S. Turner, Phys. Rev. **D42**, 3297 (1990).
[65] G. G. Raffelt, *Stars as laboratories for fundamental physics* (1996).
[66] J. H. Chang, R. Essig, and S. D. McDermott, JHEP **01**, 107 (2017), arXiv:1611.03864 [hep-ph].
[67] E. W. Kolb, D. L. Tubbs, and D. A. Dicus, The Astrophysical Journal **255**, L57 (1982).
[68] G. M. Fuller, R. Mayle, and J. R. Wilson, The Astrophysical Journal **332**, 826 (1988).