Constraining the Self-Interacting Neutrino Interpretation of the Hubble Tension

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Large, nonstandard 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 particles, and that a UV-complete model requires a nonminimal mechanism to simultaneously generate neutrino masses and appreciable self-interactions.

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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 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. This tension cannot be addressed by modifying ΛCDM at low redshift [6–9]; adding new physics before recombination seems more promising [10–17]. The solutions in Refs. [10–17] operate at temperatures $\gtrsim 1$ eV to modify the sound horizon and the inferred value of H_0. 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 Refs. [11–16].

One resolution to both issues is nonstandard neutrino self-interactions [18–22]

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

where $G_{\rm eff}$ is a dimensionful coupling with flavor indices suppressed. If $G_{\rm eff}$ is much larger than the standard model (SM) Fermi constant, G_F, neutrinos remain tightly coupled to each other until relatively late times. This inhibits their free-streaming, resulting in enhanced power on small scales and a shift in the acoustic peaks of the CMB spectrum relative to ΛCDM [23].

The effect of self-interactions is degenerate with other parameters in the CMB fit, including the angular scale of the sound horizon, the spectral index and amplitude of primordial fluctuations, and extra radiation. These degeneracies enable a *preference* for $G_{\rm eff} \gg G_F$ in cosmological data [18,20–22] while relaxing the H_0 tension [20–22]. Reference [22] extended previous analyses, allowing for finite neutrino masses and extra radiation at CMB times. They found that $G_{\rm eff}$ in the "strongly" interacting (SIν) or "moderately" interacting (MIν) regimes

$$G_{\rm eff} = \begin{cases} (4.7^{+0.4}_{-0.6}\,\mathrm{MeV})^{-2} & (\mathrm{SI}\nu) \\ (89^{+171}_{-61}\,\mathrm{MeV})^{-2} & (\mathrm{MI}\nu) \end{cases} \tag{2}$$

could simultaneously reduce the H_0 and σ_8 tensions. (These regions correspond to the Planck TT + lens + BAO + H_0 datasets. Other dataset combinations considered in Ref. [22] prefer similar values of $G_{\rm eff}$.) 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 range of $G_{\rm eff}$ in Eq. (2) vastly exceeds the strength of weak interactions, whose coupling is $G_F \simeq (2.9\times 10^5\ \mathrm{MeV})^{-2}$. We show that this interaction can only arise from the virtual exchange of a force carrier ("mediator") with $\mathcal{O}(\mathrm{MeV})$ mass and appreciable couplings to neutrinos. For this mass scale, the effective interaction in Eq. (1) is valid at energies of order $\lesssim 100$ eV, which prevail during the CMB era. However, at higher energies, this mediator is easy to produce on shell, and is subject to stringent cosmological and laboratory bounds.

We find that if strong neutrino self-interactions resolve the H_0 tension, then, (i) *Flavor-universal $G_{\rm eff}$ excluded*: If $G_{\rm eff}$ is neutrino flavor universal, both SIν and MIν regimes in Eq. (2) are excluded by laboratory searches for rare K decays and neutrinoless double-beta decay. (ii) MIν *interactions with ν_τ favored*: Couplings to ν_e, ν_μ with $G_{\rm eff}$ in the range of Eq. (2) are also excluded, except for a small island for ν_μ coupling. The only viable scenario involves neutrinos interacting through their ν_τ components in the MIν regime. (iii) *Vector forces excluded*: Constraints from big bang nucleosynthesis (BBN) exclude most self-consistent vector mediators. (iv) *Dirac neutrinos disfavored*:Mediator-neutrino interactions thermalize the right-handed components of Dirac neutrinos, significantly increasing the number of neutrino species at BBN. This excludes nearly all scenarios except the MIν regime with couplings to ν_τ. (v) *Minimal seesaw models disfavored*: Achieving the necessary interaction strength from a gauge-invariant, UV-complete model, while simultaneously accounting for neutrino masses is challenging in minimal seesaw models.

This work is organized as follows: We first demonstrate that a light new particle is required to generate the interaction in Eq. (1) with appropriate strength; we then present cosmological bounds on this scenario. Next we discuss corresponding laboratory constraints and show how Eq. (1) can arise in UV complete models before offering some concluding remarks.

The necessity of a light mediator.—Refs. [20–22,24] assume that all left-handed (LH) neutrinos undergo $2 \to 2$ flavor-universal scattering described by the interaction in Eq. (1). The largest detected CMB multipoles correspond to modes that entered the horizon when the neutrino temperature was < 100 eV. This sets the characteristic energy scale of scattering reactions during this epoch: it is important that the form of the Lagrangian in Eq. (1) is valid at this temperature. At higher energies, however, this description breaks down. As previously noted in Refs. [18,20–22], the operator in Eq. (1) is nonrenormalizable, and thus is necessarily replaced by a different interaction with new degree(s) of freedom at a scale higher than the $\sim\mathcal{O}(100\ \text{eV})$ energies probed by the CMB (see Ref. [25] for a review).

The interaction in Eq. (1) can be mediated by a particle ϕ with mass m_ϕ and coupling to neutrinos g_ϕ:

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

where ν_α are two-component left-handed neutrinos, and we allow for generic flavor structure $g_\phi^{\alpha\beta}$ of the interaction. In Eq. (3) we have assumed that ϕ is a real scalar; our conclusions are unchanged if ϕ is *CP* odd or complex. Vector forces face stronger constraints than scalars, as discussed below.

Using Eq. (3), we see that the $\nu\nu \to \nu\nu$ scattering amplitude is $\mathcal{M} \propto g_\phi^2/(m_\phi^2 - q^2)$. If the momentum transfer q satisfies $|q^2| \ll m_\phi^2$, then $\mathcal{M} \propto G_{\text{eff}}(1 + q^2/m_\phi^2 + \cdots)$, where

$$G_{\text{eff}} \equiv \frac{g_\phi^2}{m_\phi^2} = (10\ \text{MeV})^{-2} \left(\frac{g_\phi}{10^{-1}}\right)^2 \left(\frac{\text{MeV}}{m_\phi}\right)^2. \tag{4}$$

If $m_\phi^2 \ll |q^2|$, $\mathcal{M} \propto g_\phi^2/q^2$, leading to qualitatively different energy dependence for neutrino self-interactions; this regime was investigated in Refs. [26,27], which found no improvement in the H_0 tension. (Unlike Ref. [22], Refs. [26,27] fixed N_{eff} and $\sum m_\nu$, but we note that a light mediator would affect multipoles between the first acoustic peak and the diffusion scale. This should be contrasted with the massive mediator case where the self-interaction effects are larger at higher multipoles, allowing for nonstandard values of N_{eff} and $\sum m_\nu$ to compensate. A strongly interacting mode could exist here, but is unlikely to result in a larger value of H_0 after accounting for N_{eff} and $\sum m_\nu$ effects, since these impact higher-ℓ modes of the CMB spectrum.) Thus, we focus on models with a new particle ϕ for which $m_\phi^2 \gg |q^2|$ at energy scales relevant to the CMB.

Throughout this epoch, neutrinos are relativistic, so the typical momentum transfer is $|q^2| \sim T_\nu^2$. Equation (4) is valid if $m_\phi \gg T_\nu$. Comparing the values in Eq. (2) to G_{eff} in Eq. (4),

$$m_\phi \simeq (4\text{–}200) \times |g_\phi| \text{MeV}. \tag{5}$$

Since perturbativity requires $g_\phi \lesssim 4\pi$, a new sub-GeV state is required to realize this self-interacting-neutrino solution. Since $T_\nu < 100$ eV at horizon entry of the highest observed CMB multipoles, the validity of Eq. (1) in the analyses of Refs. [18,20–22,24] requires $m_\phi \gtrsim \text{keV}$ (as noted in Ref. [22]). From Eq. (5), this translates to

$$m_\phi \gtrsim \text{keV} \Rightarrow |g_\phi| \gtrsim 10^{-4}. \tag{6}$$

This bounds the allowed ranges of m_ϕ and g_ϕ. Note that Eq. (5) precludes the new self-interactions from being described within standard model effective theory with no light states below the weak scale [28].

Finally, we note that Eq. (3) is not gauge-invariant at energies above the scale of electroweak symmetry breaking (EWSB). We explore UV completions later.

Cosmological bounds.—Successful predictions of BBN provide a powerful probe of additional light species. New particles in thermal equilibrium with neutrinos increase the 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% C.L. for the SIν-(MIν-) preferred values of the baryon density, as detailed in the Supplemental Material [29].

We emphasize that large $\Delta N_{\text{eff}} \simeq 1$ at CMB times is crucial for the MIν and SIν results [22]. Since BBN does not prefer large N_{eff}, the self-interacting neutrino framework requires an injection of energy between nucleosynthesis and recombination, e.g., via late equilibration of a dark sector [49]. Such scenarios may face additional constraints. To remain model independent, we only consider the implications of BBN for the mediator (and right-handed neutrinos if they are Dirac particles) needed to implement strong neutrino self-interactions.

Mediators and ΔN_{eff}: Eq. (3) induces $\phi \leftrightarrow \nu\nu$ decays and inverse decays, which can equilibrate ϕ with neutrinos before neutrino-photon decoupling at $T_{\rm dec} \sim 1$–2 MeV. Here we show that this necessarily happens for mediators that realize $G_{\rm eff}$ in Eq. (2). Annihilation and scattering processes also contribute, but the corresponding rates are suppressed by additional powers of g_ϕ.

Vector mediators:—If Eq. (1) arises from a vector particle ϕ_μ with mass m_ϕ, then at energies above m_ϕ $\mathcal{L} \to \frac{1}{2} m_\phi^2 \phi^\mu \phi_\mu + (g_\phi \phi_\mu \nu^\dagger \bar{\sigma}^\mu \nu + \text{H.c.})$, where g_ϕ is the gauge coupling. ϕ_μ equilibrates before $T_{\rm dec}$ via $\nu\nu \leftrightarrow \phi$ if the corresponding thermally averaged rate $\Gamma_{\nu\nu\to\phi}$ exceeds Hubble when $T = \max(T_{\rm dec}, m_\phi)$:

$$\frac{\Gamma_{\nu\nu\to\phi}}{H} \sim \frac{g_\phi^2 m_\phi^2 M_{\rm Pl}}{\max(T_{\rm dec}, m_\phi)^3} > 10^8 \frac{G_{\rm eff}}{(10\ \text{MeV})^{-2}}, \qquad (7)$$

where $M_{\rm Pl} = 1.22 \times 10^{19}$ GeV and we have used Eqs. (4) and (6). This reaction is in equilibrium for all values of couplings and masses of interest. As a result, ϕ_μ has a thermal number density at $T_{\rm dec}$ in both MIν and SIν scenarios. Counting degrees of freedom, we find $\Delta N_{\rm eff} = (8/7)(3/2) \simeq 1.7$ assuming ϕ_μ remains relativistic throughout BBN; if ϕ_μ becomes nonrelativistic between $T_{\rm dec}$ and the end of BBN, then $\Delta N_{\rm eff} \approx 2.5$. Thus, ϕ_μ must become nonrelativistic well before $T_{\rm dec}$. Reference [16] found that Boltzmann suppression for massive vectors is effective for $m_\phi > 10$ MeV (95% CL). Using Eq. (4), this requires $g_\phi \gtrsim \mathcal{O}(0.1)$, which is excluded in all theoretically consistent (or anomaly-free) vector models with neutrino couplings [50,51]. Anomaly-free vectors, such as those coupled to lepton-family-number currents, would introduce large $\bar{\nu}\nu\bar{e}e$ interactions which would likely spoil the CMB fit.

Scalar mediators:—Similarly, any scalar mediator ϕ that realizes $G_{\rm eff}$ from Eq. (3) with $g_\phi \gtrsim 10^{-4}$ [required by Eq. (6)] 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 1 (2) degree(s) of freedom. The ϕ density must become Boltzmann suppressed before neutrino-photon decoupling, leading to a lower limit on m_ϕ. We use AlterBBN 2.1 [52,53] as described in the Supplemental Material [29] to obtain lower bounds (95% C.L.)

$$m_\phi > \begin{cases} 1.3\ \text{MeV} & \text{(real scalar)} \\ 5.2\ \text{MeV} & \text{(complex scalar)} \end{cases}, \qquad (8)$$

for the SIν preferred values of the baryon density (corresponding MIν bounds are somewhat weaker—see Supplemental Material [29]). SIν and MIν BBN bounds are presented in Fig. 1 as thick and thin red vertical lines, respectively.

Constraining dirac neutrinos:—If neutrinos are Dirac all neutrino masses arise from the interaction $\mathcal{L}_{\rm Dirac} \supset y_\nu H L \nu_R \to m_\nu \nu \nu_R$, where $m_\nu \equiv y_\nu v/\sqrt{2}$, 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 SM alone, the Yukawa coupling $y_\nu \sim 10^{-12}(m_\nu/0.1\ \text{eV})$ is insufficient to thermalize right-handed states, so relic neutrinos consist of left-handed neutrinos and right-handed antineutrinos [61].

The 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

$$\frac{\Gamma_{\phi\to\nu\nu_R}}{H} \simeq \frac{g_\phi^2 m_\nu^2 M_{\rm Pl}}{m_\phi^3} = 10^6 \frac{G_{\rm eff}}{(10\ \text{MeV})^{-2}} \frac{\text{MeV}}{m_\phi}, \qquad (9)$$

where $T = m_\phi \gtrsim T_{\rm dec}$ is the temperature at which RHN production is maximized relative to H. See the Supplemental Material for more details [29].

Neutrino oscillation results 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 [62]. For all values of m_ϕ we consider in the SIν range, at least one RHN will thermalize before BBN, leading to $\Delta N_{\rm eff} \gtrsim 1$. We therefore assume that neutrinos are Majorana particles for the remainder of this work.

Secret neutrino interactions: The $N_{\rm eff}$ bounds considered here can, in principle, be evaded by "secret" interactions which are communicated to active neutrinos via mixing with a light sterile neutrino, which couples directly to a mediator. In these scenarios the active-sterile mixing angle is suppressed at early times by plasma effects, but can become large at later times when the universe is cooler [63–66]. The mixing angle may be smaller than $\sim 10^{-9}$ for $T \gtrsim 50$ keV when BBN ends (to avoid thermalization) and subsequently grow to $\sim \mathcal{O}(1)$ by $T \sim 100$ eV (to enable a large active neutrino self-interaction during the CMB era, thereby resolving the H_0 tension). This sharp transition over a narrow temperature range requires significant fine-tuning of the active-sterile mass splitting and a large lepton asymmetry. See Supplemental Material for a discussion [29].

Laboratory bounds.—Because terrestrial experiments routinely reach energies above the MeV scale, the model of Eq. (3) is well constrained. We focus on scalar mediators, commenting on pseudoscalars later. Laboratory constraints arise from the following:

Double beta decay:—If $g_\phi^{ee} \neq 0$ 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 constrain $|g_\phi^{ee}| \lesssim 10^{-4}$ if $m_\phi \lesssim 2$ MeV [58–60], shown in the top row of Fig. 1.

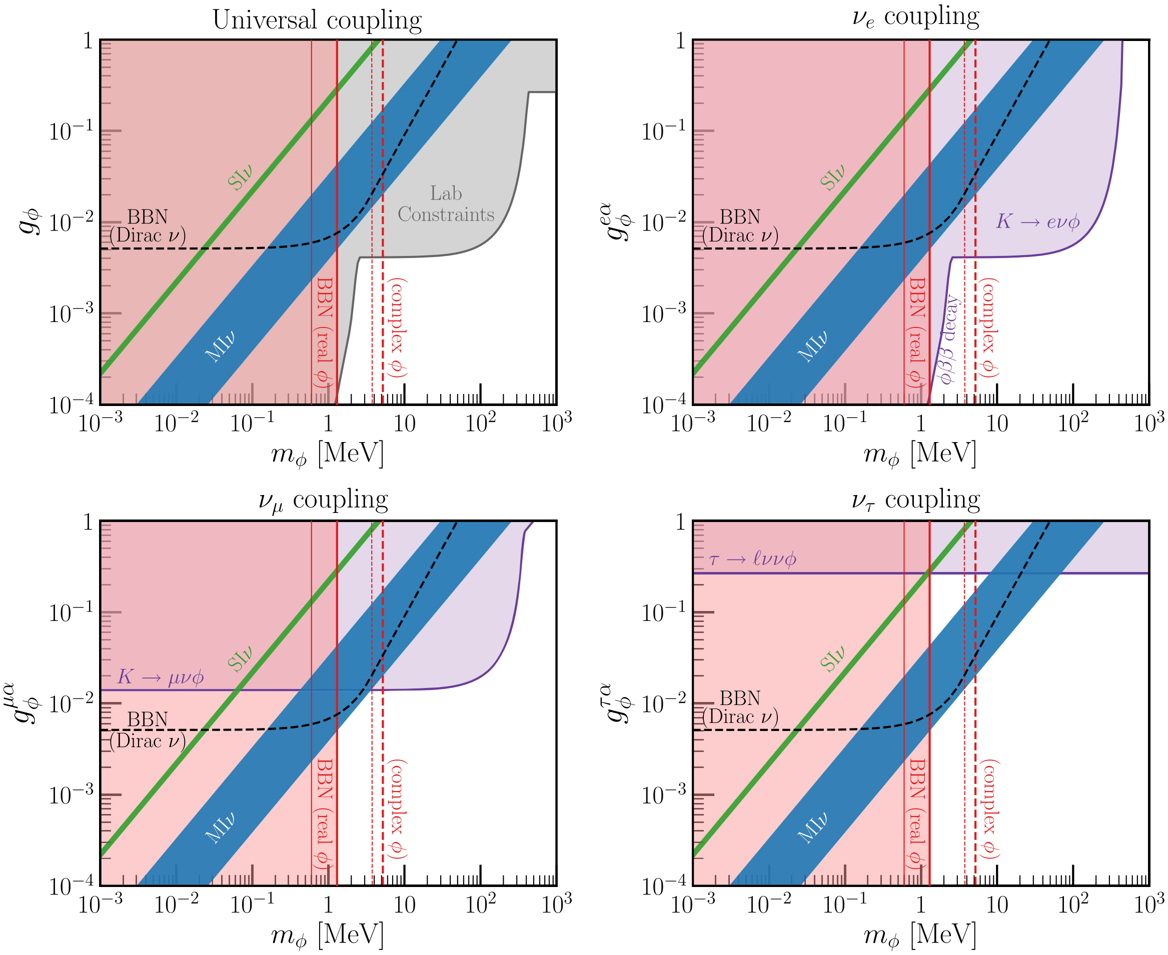

FIG. 1. Bounds (shaded regions) on light neutrino-coupled mediators with flavor-universal couplings (top left), and flavor-specific couplings to ν_e (top-right), ν_μ (bottom-left), and ν_τ (bottom-right). The bands labeled MIν and SIν are the preferred regions from Eq. (2) [22] translated into the g_ϕ-m_ϕ plane. Also shown are constraints from τ and rare meson decays [54–57], double-beta decay experiments [58–60] (purple), and BBN (red). We combine the τ/meson decay and double-beta decay constraints as "lab constraints" in the upper-left panel. BBN yields depend on the baryon density η_b; thick (thin) lines correspond to the SIν (MIν) preferred values of η_b. Nucleosynthesis constraints are stronger for complex scalar mediators (dashed red) than for real scalars (solid red). If neutrinos are Dirac, their right-handed components equilibrate before BBN above the dashed black line.

Meson decays:—Nonzero $g_\phi^{\alpha\beta}$ can allow for meson decays $\mathfrak{m}^\pm \to \ell_\alpha^\pm \nu_\beta \phi$ if $m_\phi < m_\mathfrak{m} - m_{\ell_\alpha}$ [54–57,67,68]. $\mathrm{Br}(K^+ \to e^+\nu_e)/\mathrm{Br}(K^+ \to \mu^+\nu_\mu) = (2.416 \pm 0.043) \times 10^{-5}$ constrains $g_\phi^{e\beta}$ as shown in the top row of Fig. 1 [69,70]. $\mathrm{Br}(K^+ \to \mu^+\nu_\mu\nu\bar{\nu}) < 2.4 \times 10^{-6}$ [71] constrains $g_\phi^{\mu\alpha}$, shown by the purple region in the bottom-left panel of Fig. 1.

τ decays: The decay $\tau^- \to \ell_\beta \bar{\nu}_\beta \bar{\nu}_\tau \phi$ constrains $g_\phi^{\tau\tau}$. Reference [69] found $g_\phi^{\tau\tau} \lesssim 0.3$ for light ϕ, depicted as a purple band in the bottom-right panel of Fig. 1.

Figure 1 summarizes our findings: values of G_{eff} from Eq. (2) favored by the H_0 tension are excluded if ϕ couples universally to all neutrinos (top-left), which was explicitly considered in Refs. [18,20–22,24], or (in the SIν solution) if ϕ couples predominantly to ν_e or ν_μ (top-right and bottom-left panels, respectively). Similarly, we can exclude 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 (different eigenstates will start to free-stream at different times), and the results of Refs. [18,20–22,24] may not apply.

However, a *flavor-restricted* coupling leads to approximately the same neutrino mass-eigenstate interactions as in Refs. [18,20–22,24], since the flavor eigenstates are well mixed in the mass basis. A ν_τ-only coupling, in which the matrix $g_{\alpha\beta}$ is zero except for $g_{\tau\tau}$, is potentially viable since τ decays are less constraining than meson decays. Thus, we are unable to fully exclude an interaction $G_{\mathrm{eff}}^\tau \bar{\nu}_\tau \nu_\tau \bar{\nu}_\tau \nu_\tau$.

In this case, $G_{\mathrm{eff}}^\tau = A \times G_{\mathrm{eff}}$ for G_{eff} defined in Eq. (4) and $A \sim \mathcal{O}(1)$ is a constant that accounts for the reduced scattering probability of each mass eigenstate. 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. [18,20–22,24]. Nonetheless, we expect that the preferred coupling range should shift slightly *up* relative to the flavor-universal case; a complete study is necessary to know how this affects the full SIν range. The MIν range is still allowed in a τ-flavor-only scenario, though a dedicated study is needed.

Ultraviolet completions.—Here we demonstrate the difficulty of realizing the operator in Eq. (1) from renormalizable, gauge-invariant interactions that also accommodate neutrino masses and mixings. We consider models of Majorana neutrinos with an additional particle ϕ, specifically the type-I and II seesaw mechanisms. In both, we find the resulting $\phi\nu\nu$ 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 $G_{\rm eff}$ to address the H_0 tension.

We note that the coupling of ϕ to LH neutrinos in Eq. (3) violates lepton number in analogy to neutrino masses, so it is a compelling possibility to relate these phenomena. The SM Lagrangian preserves lepton number, so the scale f of lepton-number violation must arise from new interactions. In type-I models, f is related to the RH neutrino mass, while in type-II it is proportional to the Higgs-triplet mixing parameter [72]. The interaction of ϕ with the neutrino sector occurs through the combination $f + \lambda\phi$, where λ is a coupling constant. While the relation of λ to the neutrino-masses is model-dependent, the interaction with ϕ takes on the universal form $g_\phi \approx \lambda m_\nu/f \Rightarrow G_{\rm eff} \sim \lambda^2 m_\nu^2/(m_\phi^2 f^2)$ in both type-I and type-II seesaw scenarios. Realizing $G_{\rm eff} \approx (4\text{–}300\ \text{MeV})^{-2}$ requires $f \sim 10^3 m_\nu \sim 10$ eV. In the type-I model, this scale sets the mass of the RHNs, which thermalize before BBN and contribute to $\Delta N_{\rm eff}$ as in the Dirac case discussed earlier. In type-II models this scale is bounded by nonobservation of rare lepton-number-violating processes [72–75]. Therefore, minimal scenarios where the same seesaw generates neutrino masses and the operator in Eq. (1) with the magnitude in Eq. (2) are not possible.

These arguments also apply to the Majoron, the Nambu-Goldstone boson of lepton-number breaking [76–78]. In these models, ϕ is a pseudoscalar particle, but its coupling to neutrinos is still suppressed by m_ν/f. However, the bounds we considered still apply, because all limits derive from relativistic neutrinos, for which there is no distinction between scalar and pseudoscalar.

Finally, we note that large $G_{\rm eff}$ can be obtained using *separate* seesaw mechanisms to generate the neutrino masses and the $\phi\nu\nu$ interaction—we can use the type-I seesaw for the light neutrino masses and the type-II seesaw mechanism can produce large $g_\phi^{\tau\tau}$ (as long as it does not contribute to m_ν). The size of $g_\phi^{\tau\tau}$ decouples from the neutrino masses.

Concluding remarks.—We have shown that the self-interacting neutrino explanation of the H_0 tension requires the existence of a light ~MeV-scale mediator, subject to stringent cosmological and laboratory bounds. Consequently, for both the SIν and MIν regimes in Eq. (2), the flavor-universal interactions considered in Refs. [18,20–22] are robustly excluded by BBN-only bounds on $\Delta N_{\rm eff}$ and by laboratory searches for rare K decays and neutrinoless double-beta decay; the SIν regime is excluded for all flavor structures.

Intriguingly, we find that flavor-dependent variations of the MIν regime may viably resolve the H_0 tension if a ~10 MeV scalar mediator with large coupling interacts almost exclusively with ν_τ or ν_μ (though there is little parameter space for ν_μ coupling). A dedicated exploration of the τ-only scenario is necessary to determine if the preferred region to resolve the H_0 tension persists without 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.

However, realizing such strong, flavor specific 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/f, where $f \gg m_\nu$ is the appropriate seesaw scale. A compelling and viable model remains to be found.

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[1] A. G. Riess *et al.*, Astrophys. J. **826**, 56 (2016).
[2] T. Shanks, L. Hogarth, and N. Metcalfe, Mon. Not. R. Astron. Soc. **484**, L64 (2019).
[3] A. G. Riess, S. Casertano, D. Kenworthy, D. Scolnic, and L. Macri, arXiv:1810.03526.
[4] N. Aghanim *et al.* (Planck Collaboration), arXiv:1807.06209.
[5] A. G. Riess, S. Casertano, W. Yuan, L. M. Macri, and D. Scolnic, Astrophys. J. **876**, 85 (2019).
[6] M. Vonlanthen, S. Rsnen, and R. Durrer, J. Cosmol. Astropart. Phys. 08 (2010) 023.
[7] L. Verde, E. Bellini, C. Pigozzo, A. F. Heavens, and R. Jimenez, J. Cosmol. Astropart. Phys. 04 (2017) 023.

[8] J. Evslin, A. A. Sen, and Ruchika, Phys. Rev. D **97**, 103511 (2018).
[9] K. Aylor, M. Joy, L. Knox, M. Millea, S. Raghunathan, and W. L. K. Wu, Astrophys. J. **874**, 4 (2019).
[10] J. Lesgourgues, G. Marques-Tavares, and M. Schmaltz, J. Cosmol. Astropart. Phys. 02 (2016) 037.
[11] E. Di Valentino, C. Bohm, E. Hivon, and F. R. Bouchet, Phys. Rev. D **97**, 043513 (2018).
[12] V. Poulin, T. L. Smith, D. Grin, T. Karwal, and M. Kamionkowski, Phys. Rev. D **98**, 083525 (2018).
[13] F. D'Eramo, R. Z. Ferreira, A. Notari, and J. L. Bernal, J. Cosmol. Astropart. Phys. 11 (2018) 014.
[14] V. Poulin, T. L. Smith, T. Karwal, and M. Kamionkowski, Phys. Rev. Lett. **122**, 221301 (2019).
[15] K. L. Pandey, T. Karwal, and S. Das, arXiv:1902.10636.
[16] M. Escudero, D. Hooper, G. Krnjaic, and M. Pierre, J. High Energy Phys. 03 (2019) 071.
[17] P. Agrawal, F.-Y. Cyr-Racine, D. Pinner, and L. Randall, arXiv:1904.01016.
[18] F.-Y. Cyr-Racine and K. Sigurdson, Phys. Rev. D **90**, 123533 (2014).
[19] M. Archidiacono and S. Hannestad, J. Cosmol. Astropart. Phys. 07 (2014) 046.
[20] L. Lancaster, F.-Y. Cyr-Racine, L. Knox, and Z. Pan, J. Cosmol. Astropart. Phys. 07 (2017) 033.
[21] I. M. Oldengott, T. Tram, C. Rampf, and Y. Y. Y. Wong, J. Cosmol. Astropart. Phys. 11 (2017) 027.
[22] C. D. Kreisch, F.-Y. Cyr-Racine, and O. Dor, arXiv: 1902.00534.
[23] S. Bashinsky and U. Seljak, Phys. Rev. D **69**, 083002 (2004).
[24] G. A. Barenboim, P. B. Denton, and I. M. Oldengott, Phys. Rev. D **99**, 083515 (2019).
[25] T. Cohen, *Proc. Sci.*, TASI2018 (2019) 011 [arXiv: 1903.03622].
[26] F. Forastieri, M. Lattanzi, and P. Natoli, J. Cosmol. Astropart. Phys. 07 (2015) 014.
[27] F. Forastieri, M. Lattanzi, and P. Natoli, arXiv:1904.07810 [Phys. Rev. D (to be published)].
[28] M. B. Gavela, D. Hernandez, T. Ota, and W. Winter, Phys. Rev. D **79**, 013007 (2009).
[29] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.123.191102 for more detail on BBN constraints, dirac neutrino thermalization, Supernova 1987A constraints, and strong neutrino self-interactions via sterile neutrino interactions, which includes Refs. [30–48].
[30] P. D. Serpico and G. G. Raffelt, Phys. Rev. D **70**, 043526 (2004).
[31] C. Boehm, M. J. Dolan, and C. McCabe, J. Cosmol. Astropart. Phys. 08 (2013) 041.
[32] K. M. Nollett and G. Steigman, Phys. Rev. D **91**, 083505 (2015).
[33] E. Aver, K. A. Olive, and E. D. Skillman, J. Cosmol. Astropart. Phys. 07 (2015) 011.
[34] R. J. Cooke, M. Pettini, and C. C. Steidel, Astrophys. J. **855**, 102 (2018).
[35] R. H. Cyburt, B. D. Fields, K. A. Olive, and T.-H. Yeh, Rev. Mod. Phys. **88**, 015004 (2016).
[36] A. Berlin, N. Blinov, and S. W. Li, Phys. Rev. D **100**, 015038 (2019).
[37] G. Steigman, J. Cosmol. Astropart. Phys. 10 (2006) 016.
[38] C. Boehm, M. J. Dolan, and C. McCabe, J. Cosmol. Astropart. Phys. 12 (2012) 027.
[39] G. Mangano, G. Miele, S. Pastor, T. Pinto, O. Pisanti, and P. D. Serpico, Nucl. Phys. **B729**, 221 (2005).
[40] P. F. de Salas and S. Pastor, J. Cosmol. Astropart. Phys. 07 (2016) 051.
[41] H. K. Dreiner, H. E. Haber, and S. P. Martin, Phys. Rep. **494**, 1 (2010).
[42] A. Burrows and J. M. Lattimer, Astrophys. J. **318**, L63 (1987).
[43] A. Burrows, M. S. Turner, and R. P. Brinkmann, Phys. Rev. D **39**, 1020 (1989).
[44] A. Burrows, M. T. Ressell, and M. S. Turner, Phys. Rev. D **42**, 3297 (1990).
[45] G. G. Raffelt, *Stars as Laboratories for Fundamental Physics* (University of Chicago Press, Chicago, 1996).
[46] J. H. Chang, R. Essig, and S. D. McDermott, J. High Energy Phys. 01 (2017) 107.
[47] E. W. Kolb, D. L. Tubbs, and D. A. Dicus, Astrophys. J. **255**, L57 (1982).
[48] G. M. Fuller, R. Mayle, and J. R. Wilson, Astrophys. J. **332**, 826 (1988).
[49] A. Berlin and N. Blinov, Phys. Rev. Lett. **120**, 021801 (2018).
[50] M. Bauer, P. Foldenauer, and J. Jaeckel, J. High Energy Phys. 07 (2018) 094.
[51] Y. Farzan and I. M. Shoemaker, J. High Energy Phys. 07 (2016) 033.
[52] A. Arbey, Comput. Phys. Commun. **183**, 1822 (2012).
[53] A. Arbey, J. Auffinger, K. P. Hickerson, and E. S. Jenssen, arXiv:1806.11095.
[54] K. Blum, A. Hook, and K. Murase, arXiv:1408.3799.
[55] J. M. Berryman, A. de Gouvêa, K. J. Kelly, and Y. Zhang, Phys. Rev. D **97**, 075030 (2018).
[56] K. J. Kelly and Y. Zhang, Phys. Rev. D **99**, 055034 (2019).
[57] G. Krnjaic, G. Marques-Tavares, D. Redigolo, and K. Tobioka, arXiv:1902.07715.
[58] M. Agostini *et al.*, Eur. Phys. J. C **75**, 416 (2015).
[59] K. Blum, Y. Nir, and M. Shavit, Phys. Lett. B **785**, 354 (2018).
[60] T. Brune and H. Ps, Phys. Rev. D **99**, 096005 (2019).
[61] A. D. Dolgov, K. Kainulainen, and I. Z. Rothstein, Phys. Rev. D **51**, 4129 (1995).
[62] I. Esteban, M. C. Gonzalez-Garcia, A. Hernandez-Cabezudo, M. Maltoni, and T. Schwetz, J. High Energy Phys. 01 (2019) 106.
[63] K. Abazajian, G. M. Fuller, and M. Patel, Phys. Rev. D **64**, 023501 (2001).
[64] S. Hannestad, R. S. Hansen, and T. Tram, Phys. Rev. Lett. **112**, 031802 (2014).
[65] S. Hannestad, I. Tamborra, and T. Tram, J. Cosmol. Astropart. Phys. 07 (2012) 025.
[66] N. Saviano, A. Mirizzi, O. Pisanti, P. D. Serpico, G. Mangano, and G. Miele, Phys. Rev. D **87**, 073006 (2013).

[67] K. C. Y. Ng and J. F. Beacom, Phys. Rev. D **90**, 065035 (2014); **90**, 089904(E) (2014).

[68] K. Ioka and K. Murase, Prog. Theor. Exp. Phys. (**2014**), 61E01.

[69] A. P. Lessa and O. L. G. Peres, Phys. Rev. D **75**, 094001 (2007).

[70] L. Fiorini (NA48/2 Collaboration), Nucl. Phys. B, Proc. Suppl. **169**, 205 (2007).

[71] M. Tanabashi *et al.* (Particle Data Group), Phys. Rev. D **98**, 030001 (2018).

[72] Y. Cai, J. Herrero-García, M. A. Schmidt, A. Vicente, and R. R. Volkas, Front. Phys. **5**, 63 (2017).

[73] A. G. Akeroyd, M. Aoki, and H. Sugiyama, Phys. Rev. D **79**, 113010 (2009).

[74] D. N. Dinh, A. Ibarra, E. Molinaro, and S. T. Petcov, J. High Energy Phys. 08 (2012) 125; 09 (2013) 23.

[75] P. S. B. Dev, C. M. Vila, and W. Rodejohann, Nucl. Phys. **B921**, 436 (2017).

[76] Y. Chikashige, R. N. Mohapatra, and R. D. Peccei, Phys. Lett. **98B**, 265 (1981).

[77] Y. Chikashige, R. N. Mohapatra, and R. D. Peccei, Phys. Rev. Lett. **45**, 1926 (1980).

[78] J. Schechter and J. W. F. Valle, Phys. Rev. D **25**, 774 (1982).