Bounds on neutrino-scalar non-standard interactions from big bang nucleosynthesis

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Coherent forward scattering processes by neutrino-scalar non-standard interactions (SNSI) induce an effective neutrino mass. In the Early Universe, a large neutrino effective mass restricts the production of neutrinos. The SNSI effect is modulated by two effective couplings, these account for the coupling between neutrinos and electrons/positrons, $G_{\text{eff}}$, and the neutrino self-interaction, $G_{S}$. These parameters are directly related to the effective number of relativistic species and non-zero values imply a smaller than expected $N_{\text{eff}}$. We employ big bang nucleosynthesis to constraint the SNSI effect. We find that $G_{\text{eff}} < 3.8 \text{ MeV}^{-2}$ and $G_{S} < 6.2 \times 10^{7} \text{ MeV}^{-2}$ at 95% CL. For a scalar mass in the range $10^{-12} \text{ eV} < m_{\phi} < 10^{-6} \text{ eV}$, our neutrino-scalar coupling constraint is more restrictive than any previous result.

I. INTRODUCTION

Despite being light, neutrino gravitational interaction plays an essential role in shaping the distribution of matter and energy in the Universe. Several cosmological surveys have led to the strongest bounds on the sum of the neutrino masses [14]. These are one order of magnitude better than those from experimental counterparts [5]. Cosmology now leads the race to determine the neutrino mass hierarchy, and possibly, measure the mass of at least one neutrino throughout this decade [9,17]. Moreover, three standard neutrinos are required to predict accurately the abundance of light elements on the Universe through big bang nucleosynthesis (BBN) [8,10]. This is in concordance with the standard precision computation of the neutrino contribution to radiation density, that can be expressed in terms of the parameter $N_{\text{eff}} \simeq 3.046$ [11,14].

Cosmological model-independent bounds on neutrinos will be more reliable by disentangling the effects of neutrino parameters with the rest of cosmological ones [1]. As an important step, the existence of relativistic species in the Early Universe has been proven by detecting a phase shift on the acoustic oscillations that cannot be mimicked by other cosmological parameters [15,17]. In this sense, cosmology has become a fruitful Lab to test neutrino physics in the outline of the standard model of particle physics (SM) and beyond (BSM).

Neutrino interactions with matter are crucial to study them. For instance, the Mikheyev Smirnov Wolfenstein (MSW) effect [18], which changes the neutrino oscillations in matter, was used to determine the sign of the square-mass splitting $\Delta m_{31}^{2} > 0$ in the solar neutrino experiments (see for instance [19,20]). The same mechanism is being brought out by long-baseline neutrino experiments aiming to determine the sign of $\Delta m_{31}^{2}$ (see experiments [21,22]).

In cosmology, a neutrino non-standard interaction (NSI) may solve some tensions in the standard theory. It has been studied whether an NSI may explain the discrepancy known as the $H_{0}$-tension, where the measurement of $H_{0}$ by the cosmic microwave background (CMB) and local observations are clearly in statistical disagreement [1,23,25]. There are some approaches that try to solve this problem using NSI, including interactions in the sterile [26], or in the active neutrino sectors [27]. In the latter approach, neutrinos are required to be either strongly self-interacting (SI$\nu$) or moderately self-interacting (MI$\nu$).

SI$\nu$ and/or MI$\nu$ are assumed to be mediated by a scalar particle with a mass larger than $\mathcal{O}(\text{keV})$. And, the parameter space, in this approximation, has been cornered by experimental, astrophysical, and BBN constraints [25,22]. On the other hand, the phenomenology of neutrino scalar non-standard interactions (SNSI) mediated by a light particle is rich and has several consequences. For instance, large-scale structure (LSS) data constrains neutrino dispersion mediated by a scalar much lighter than $\mathcal{O}(\text{eV})$ [33,34]. Furthermore, neutrinos may annihilate and decay into lighter bosons, which, interestingly, may relax the bound on $\sum m_{\nu}$ imposed by LSS [35,38].

Although the information on the light mediator mass is lost when studying two-body dispersion in the regime $m_{\phi} \ll T_{\nu}$, loop diagrams such as mass-correction type, a priory, are mass-dependent regardless of the smallness of the scalar mass. Therefore, studying this kind of diagrams within the Early Universe background is convenient if we are to search for mediator mass-dependent SNSI constraints.
In this manuscript, we explore the cosmological consequences of neutrino SNSI mass correction processes mediated by a light scalar particle. We assess the calculations performed by Babu et al. \[39\] in the Early Universe. Mass correction diagrams involving an SNSI have received recent attention because Ge & Park \[10\] found a small solar neutrino data preference for non-vanishing SNSI couplings with ordinary matter. This result has led to further research about neutrino propagation with SNSI in The Earth, The Sun, and supernovae \[39,41,42\].

For our exploration, we identify two effective parameters that modulate the SNSI effect and study its consequences. We solve numerically the mass contribution and the evolution of the neutrino density. Additionally, we notice that large effective SNSI couplings may noticeably change the neutrino contribution to radiation. This information is encoded through a temperature-dependent change on the effective number of relativistic species \(N_{\text{eff}}\). In order to find \(N_{\text{eff}}\), we employ a modified version of the public code NUDEC_BSM \[19\]. A change on \(N_{\text{eff}}\) straightforwardly alters the expansion rate during radiation dominated era, affecting the proton+neutron freeze-out temperature and, hence, the neutron to proton ratio right at the onset of BBN. Thus, the production of primordial nuclei helps us to constrain the SNSI parameter space. We use a modified version of the public BBN code ALTERNBBN \[43,44\] to find the parameter constraints. Finally, we translate these bounds into the scalar mass - couplings parameter space and compare them with other results.

The rest of the paper is organized as follows. In Section II, we review and discuss the properties of the effective mass coherent forward scattering process (CFS) by SNSI at high temperatures. In Section III, we explore the phenomenological consequences of the effective neutrino mass. Then, in Section IV we constrain the parameter space of the SNSI with BBN theory and the abundance of light elements. In Section V we compare our constraints with laboratory, astrophysical and cosmological bounds on the parameters. Our conclusions are summarized in Section VI.

II. NEUTRINO SCALAR NON-STANDARD INTERACTIONS

The outcomes of neutrino NSI depend on the nature of the mediator particle. On the one hand side, vector-mediated NSI has a phenomenology that produces similar effects as the weak interaction. The SNSI instead, appears as a Yukawa term on the effective Lagrangian \[10\] which induces an effective mass. This mass term depends on the properties of the environment where neutrinos propagate. A dense and hot background may produce a large neutrino mass.

We are interested in the effect of the SNSI in the CFS described by the tadpole diagram of Fig. 1. We consider that neutrinos are propagating in a hot plasma when the Universe had a temperature around some MeVs, this plasma is composed of photons, baryons, charged leptons, and the three standard neutrinos. The SNSI effect in the neutrino propagation can be interpreted as a refractive index \[43,46\]. Here, we focus on a generic scalar interaction ignoring the details of an underlying particle physics model, having the cosmological phenomenology as our main approach.

The effective neutrino mass described by the quantum correction would be

\[
m_{\text{eff}} = m_\nu + 2G_\text{eff} \Delta m(m_\nu; T) + 3G_S \Delta m(m_\nu; T),
\]

where \(m_\nu\) is the bare neutrino mass, and the correction is described by \[39\]

\[
\Delta m(m_f; T) = \frac{m_f}{\pi^2 \Gamma} \int_{m_f}^{\infty} dk \sqrt{k^2 - m_f^2} f(k).
\]

Here \(m_f\) is the mass of the fermion and \(f(k_0)\) is the Fermi-Dirac distribution for the the background fermions. Safely neglecting the chemical potential \[47\], \(\mu = 0\), the Fermi-Dirac distribution for both fermions and anti-fermions is the same \((e^{k/T} + 1)^{-1}\), where \(T\) is the temperature of the thermal background. The two free parameters, \(G_\text{eff}\) and \(G_S\), are then given as

\[
G_\text{eff} = \frac{g_\nu g_\phi}{m_\phi^2},
\]

and

\[
G_S = \frac{g_\phi^2}{m_\phi^2},
\]

where \(m_\phi\) is the the mass of the scalar mediator, \(g_\nu\) is the neutrino-scalar coupling and \(g_f\) is the coupling between the scalar and charged leptons. These effective couplings encode the strength of the interaction and are the ones to be constrained by observations. Here we assume universal couplings with both charged lepton and neutrino flavors. Therefore, all complex phases can be absorbed and one can assume neutrino mass corrections to be always positive. Notice that, at the temperatures that we are interested here, there are not muons/taus.
FIG. 2. Neutrino mass-correction induced by an SNSI interaction, depicted as a function of the photon-baryon temperature. Upper panel: SNSI with electrons and positrons. Lower panel: Neutrino self-interaction for three different values of the neutrino bare mass present in the background. Key cosmological events highlighted: Neutrino decoupling, proton to neutron freeze-out (f.o.), synthesis of light elements, and electron-positron annihilation freeze-out.

present in the plasma, since they have already decayed into lighter particles by then \((T_e \ll m_\mu \sim 105.65 \text{ MeV})\). Hence, we only take into account couplings with electrons and positrons.

The numerical solution of the electron/positron SNSI contributing to the neutrino mass is depicted in Fig. 2. At high temperatures, both contributions to the effective mass are the same. Below the electron mass threshold, the contribution decays exponentially as the Universe cools down. But, when the electron-positron annihilation ends, only electrons remain in the background. However, at temperatures much smaller than \(\mathcal{O}\text{MeV}\), the neutrino mass correction contribution induced by leptons becomes negligible.

Unlike some terrestrial and astrophysical scenarios, here we also need to consider the background composed of relic neutrinos. In this self-interacting case, \(\Delta m\) would have another unknown parameter, the bare neutrino mass \(m_\nu\). Notice that, in order to have a \(\Delta m\) of the same order of magnitude than the one induced by charged leptons, \(G_S\) needs to be roughly \(m_\nu/m_\nu\) times larger than \(G_{\text{eff}}\), see equation (2). As the BBN epoch occurs at temperatures much larger than the bare neutrino mass scales, the mass correction does not drop exponentially with the temperature as it occurs with the electron/positron SNSI. By definition, by constraining \(G_S\), we would be able to find a mediator mass-dependent \(g_\nu\)-bound.

By oscillation experiments, we know that at least two neutrinos are massive. Hereafter, we shall take a conservative value for bare neutrino masses, being one-third of the minimum sum of neutrino masses in the normal hierarchy, \((\sum m_\nu)_{\text{min}} \sim 0.059 \text{ eV}\) [3], assuming an almost degenerate scenario of active neutrinos. Given this, we take \(m_\nu = 0.0195 \text{ eV}\) and assume all three neutrino parameters are universal.

### III. COSMOLOGICAL IMPLICATIONS

In the previous section, we have described how the neutrino \(m_{\text{eff}}\) would be affected by CFS with charged leptons and neutrinos at high temperatures. We now focus on the implementation and implications of neutrino SNSI in the Early Universe. In particular, in this section we compute \(N_{\text{eff}}\) as a function of the SNSI parameters.

The particles in the plasma are in local thermal equilibrium when their interaction rate is larger than the rate of the expansion of the Universe, \(\Gamma \gg H(T_e)\). The Universe at high temperatures \((T_e \sim \mathcal{O}\text{MeV})\) is dominated by radiation and the density of any heavy particle, \(m \gtrsim T_e\), gets suppressed. A large neutrino \(m_{\text{eff}}\) will diminish its production by weak interactions and ultimately the Universe will have less radiation than expected. Therefore, by weighting the effect of \(m_{\text{eff}}\) on \(N_{\text{eff}}\), we will estimate the permitted values of the SNSI parameters.

As we stated in the previous section, the mass correction diagram of Fig. 1 is describing a CFS that implies no transfer of energy and momentum with the plasma. Therefore, a priori, the neutrino thermal evolution should remain unchanged. Nonetheless, we carefully explore whether the \(m_{\text{eff}}\) is capable of changing the thermal evolution of neutrinos.

In this scenario, the weak interaction is the one responsible for keeping neutrinos in thermal equilibrium with the plasma. In equilibrium, the neutrino energy and number density are given by [48]

\[
\rho_\nu(G_{\text{eff}}, G_S; T_e, T_\gamma) = \frac{T_e^4}{\pi^2} \int_\alpha^\infty dx \frac{x^2 \sqrt{x^2 - \alpha^2}}{e^x + 1},
\]

\[
n_\nu(G_{\text{eff}}, G_S; T_e, T_\gamma) = \frac{T_e^3}{\pi^2} \int_\alpha^\infty dx \frac{x \sqrt{x^2 - \alpha^2}}{e^x + 1},
\]

where \(x = E_\nu/T_e\) and \(\alpha = m_{\text{eff}}/T_e\). The effective mass \(m_{\text{eff}}\) encodes all the new physics, as given in equation (1). Notice that the neutrino density gets suppressed with a larger \(m_{\text{eff}}\). As the Universe cools down, the effective neutrino mass drops significantly, this permits the neutrino density to approach and possibly recover its standard value. However, after neutrino decoupling, it is not possible to produce abundantly new neutrinos to reach their standard density. Thus, establishing the neutrino decoupling temperature is important to compute the final neutrino density to a good approximation.

We argue that the change in the rate of expansion due to SNSI leaves the neutrino decoupling temperature almost unchanged. The expansion rate is proportional to the density \(H(T) \propto \sqrt{\rho}\), while the weak interaction rate that keeps neutrinos in local thermal equilibrium is \(\Gamma_{\text{ew}} \propto T^3\). As we will show in section 4V, changes up to 10% in the neutrino density are enough to change
substantially the nuclear reactions, and thus, this is measurable by BBN physics. Therefore, the neutrino decoupling temperature ratio respect to the standard case is $(\rho_{\text{SNSI}}/\rho_{\text{SM}})^{1/10}$. Moreover, a 10% change in the neutrino density implies a 1% change in neutrino decoupling temperature. For the purpose of the present discussion, it is a good approximation to assume no change in neutrino decoupling temperature due to SNSI.

We also track the neutrino temperature evolution after decoupling. We employ a modified version of the public available code NUDEC_BSM [13]. This code solves for the ratio of the neutrino and photon temperature in a much simpler approximation than state-of-the-art codes [14]. Unlike other precise computations of $N_{\text{eff}}$, this code does not include neutrino oscillations, yet, it computes a pretty robust value of $N_{\text{eff}} = 3.045$ in the SM case. In Fig 3 (c) we depicted the evolution of the temperatures for the standard case and a scenario denoted as large $G_{\text{eff}}$ and $G_{S}$ values. We observe that the evolution of temperatures differs only within the numerical error values.

The three-flavour neutrino decoupling happens between $3 - 2$ MeVs, hence, we assume that the neutrinos fully decouple from the plasma at exactly 2 MeV [11-14]. That means that the density freezes out at that temperature and no significant amount of neutrinos gets produced after that. This is because weak interactions would be able to produce only a small percentage of the total neutrino density. Lastly, notice that assuming this late instantaneous neutrino decoupling is the most conservative approach. However, we expect almost the same final neutrino density than using a more complex model for neutrino decoupling. This is because neutrinos decouple the earliest at $\sim 3$ MeV.

We move on to model the neutrino density. The energy density becomes a piece-wise function, where the neutrino density freezes out at the threshold $T_{\gamma} = 2$ MeV,

$$\rho_{\nu} = \begin{cases} \rho_{\nu}(G_{\text{eff}}, G_{S}; T_{\nu}, T_{\gamma}) & T_{\gamma} > 2\text{MeV} \\ \left(\frac{T_{\nu}}{2\text{MeV}}\right)^{4} \rho_{\nu}(G_{\text{eff}}, G_{S}; 2\text{MeV}, 2\text{MeV}) & T_{\gamma} \leq 2\text{MeV} \end{cases}$$

(6)

where $\rho_{\nu}(G_{\text{eff}}, G_{S}; T_{\nu}, T_{\gamma})$ is the thermal density described in [5]. After decoupling, the neutrino density falls due to the adiabatic expansion of the Universe. We sketch this in Figs. 3 (a) & (b) for different values of the SNSI effective parameters.

We proceed to numerically compute the effective number of relativistic species $N_{\text{eff}}$. For this purpose, we again employ the code NUDEC_BSM. We observe that there is
a direct relation between the SNSI parameters and $N_{\text{eff}}$, this is given by

$$N_{\text{eff}}(G_{\text{eff}}, G_S) \approx \frac{8}{7} \left( \frac{11}{4} \right)^{4/3} \frac{3\rho_\gamma(G_S, G_{\text{eff}})}{\rho_\gamma} \text{ for } T_\gamma \ll m_e,$$

where $\rho_\gamma = (2\pi^2/30)T_\gamma^4$ is the photon density and we have assumed a full degeneration of neutrino parameters. In Fig. 4 we illustrate the change on $N_{\text{eff}}$ as a function of the effective couplings $G_S$ and $G_{\text{eff}}$. We can observe a strong positive correlation between the SNSI parameters since they both produce the same effect. In addition, we show the CMB bound by Planck collaboration $N_{\text{eff}} = 2.99^{+0.34}_{-0.33}$ (95% CL) [1]. In the next section, we will constrain these parameters with BBN physics.

Lastly, notice that we are neglecting first order SNSI processes, such as three-level scattering. Thus, we assume that the scalar mediator is out-of-equilibrium with the plasma, so it does not acquire a thermal mass, nor is being thermally produced within the plasma at a significant number. Finally, we assume that neutrino decays into the scalar are irrelevant, hence, the scalar density is insignificant.

Dimensional arguments permit us to explore the validity of these approximations. Our region of interest lies in the small scalar mass regime $m_\phi \ll$ keV. In this case, the SNSI cross-section of processes such as $e^- e^- \rightarrow e^- e^-$ would be $\sigma_{\text{SNSI}} \approx g_\phi^2 g_\nu^2/T^2$. While, the SM cross-section is given by $\sigma_{\text{SM}} \approx 1.4T_{\nu}^2/\alpha^2$, where $\alpha \sim 1/137$ is the fine-structure constant and $T_{\nu} \sim 80$ GeV is the W Boson mass. Comparing both cross-sections, we observed that the condition $g_\phi g_\nu < \alpha T_{\nu}^2/M_w^2$, for $T \sim 1$ MeV implies that $g_\phi g_\nu \lesssim 10^{-12}$. Similarly, the scalar would be prevented from reaching thermal equilibrium as long as the condition $g_\phi^2 < \alpha T_{\nu}^2/M_w^2$ is satisfied. Lastly, the scalar would not significantly contribute to $N_{\text{eff}}$, provided the condition $g_\phi < g_\nu \lesssim 10^{-3}$ for $m_\phi \ll$ keV [50] is not violated. As we will argue along section IV these conditions would be satisfied in the ultralight scalar regime (see Figs. 6 and 7).

IV. BBN CONSTRAINTS

In this section, we present the bounds on the effective SNSI parameters $G_S$ and $G_{\text{eff}}$ by BBN theory and observations of light element abundances.

BBN is one of the cornerstones of modern cosmology and the Big Bang Theory. With the interplay of standard nuclear and particle physics and the standard cosmological model, it describes with great accuracy the synthesis of the lighter nuclei during the very first seconds of cosmic time (For a review see [51]). Despite some uncertainties on the predictions for $^7$Li, which may have a diversity of possible sources [51] [52], it predicts the observed relative abundances of H, D, $^3$He, $^4$He and $^7$Li as a function of a single parameter, the baryon-to-photon ratio, $\eta = n_b/n_\gamma$, or equivalently, the present baryon density $\Omega_b h^2$, which determines the end of the deuterium bottleneck, and therefore, the production rate of heavier nuclei.

Aside from the initial condition on $\eta$, which is thought to be associated with an earlier baryogenesis process, for which the SM seems to have not a satisfactory explanation, BBN success is based on well-known physics, which leaves little space for new or exotic physics. This feature is precisely what makes BBN a useful probe for any non-standard physics that may modify the cosmological evolution during those early times. In particular, any physics that could change the expansion rate during BBN [10] [53] [54]. As the effective neutrino mass that we are discussing changes $N_{\text{eff}}$, it does affect the amount of radiation during that epoch, and so BBN should be sensitive to it. We will focus on this in what follows.

To a good approximation, when the deuterium bottleneck breaks up, most of the neutrons present in the primordial Universe are synthesized in $^4$He. Other elements are then produced at much smaller amounts, with a rate of about $10^{-5}$ for D and $^3$He and $10^{-10}$ for $^7$Li per proton. $^4$He mass fraction is well approximated as

$$Y_p \approx \frac{2(n/p)}{1 + (n/p)},$$

where the neutron to proton ratio, $(n/p)$, at BBN is determined by the output ratio at weak interactions freeze out, when the weak interactions rate per baryon $\Gamma_{\text{ew}} \approx \alpha T^2/M_w^4$ becomes smaller than Hubble expansion, and by neutron number depletion due to $\beta$ decay. In the standard cosmological model at freeze out temperature, $T_* \approx 0.8$ MeV, $(n/p)_B \approx e^{-\Delta m} T_* \sim 1/5$, where $\Delta m = m_\nu - m_p$ is the neutron - proton mass difference, and thus one estimates $(n/p)_{BBN} \approx 1/7$ (for a theoretical calculation of this see for instance [55] [56]).

The key feature for our present analysis resides in the fact that setting $N_{\text{eff}}$ as a free parameter compromises the expansion rate during the radiation dominated epoch. A smaller (larger) value of $N_{\text{eff}}$ than the one computed in
the standard case, reduces (increases) expansion rate and lowers (raises) weak interactions decoupling temperature. Even if the change is mild, due to Boltzmann suppression, a smaller (larger) \( T_e \) implies a lower (higher) \( (n/p) \) and thus a smaller (higher) \( Y_p \).

CMB is sensitive to both \( \eta \) and \( Y_p \) and as a matter of fact, Planck data alone provides a determination of \( Y_p \) \([1]\). Although \( Y_p \) is not sensitive to the baryon-to-photon ratio, as we mentioned earlier, \( \eta \) is an important initial condition for BBN and the production of other light elements. Here, we keep our analysis consistent with the standard case, reduces (increases) expansion rate and with the physical bound of the parameters \( G_S \) and \( G_{\text{eff}} \) which are constrained to be positive. The 100\( \gamma \)% credible region will be defined as

\[
\int_{\theta_{1,\text{bound}}}^{\theta_{1}} d\theta_{1} P^{(1)}(\theta_{1}) = \gamma_1. \tag{9}
\]

Using a finite grid the \( \gamma \)-value for the parameter \( \theta_1 \) is

\[
\gamma_1 = \frac{1}{N} \sum_{i=0}^{N_{\text{bound}}} \sum_{j} e^{-\chi_i^2/2} \delta \theta_{1i} \delta \theta_{2j}, \tag{10}
\]

where \( N = \sum_{i,j} e^{-\chi_i^2/2} \delta \theta_{1i} \delta \theta_{2j} \), \( \delta \theta \) is a small constant finite difference in the parameter sampling, and \( \gamma_2 \) is similarly defined.

We obtain the parameter constraints at 95\( \% \) CL by finding the \( \theta \)-values that make \( \gamma_1 = \gamma_2 = 0.95 \). Our final marginalized bounds are

\[
G_{\text{eff}} < 3.8 \text{ MeV}^{-2} \quad (95\% \text{ CL}), \tag{11}
\]

\[
G_S < 6.2 \times 10^7 \text{ MeV}^{-2} \quad (95\% \text{ CL}).
\]

Notice that, although parameter marginalization is the appropriate statistical procedure to obtain robust constraints, the effective SNSI parameters are strongly positively correlated and the degeneracy cannot be entirely broken up. Furthermore, we check our results for consistency for different values of the baryon-to-photon ratio within its 1-\( \sigma \) error \( \eta = 6.11 \pm 0.03 \times 10^{-10} \). We observe no significant fluctuation in our final bounds.

In the next section, we will disentangle the model parameters (the couplings and the mediator mass) by using our constraints in the effective SNSI parameters and present new bounds on the neutrino-scalar coupling.

V. PARAMETER SPACE COMPARISON

We have discussed and computed the bounds on the effective parameters of SNSI. Here we translate the bounds
into the mass-coupling parameter space and compare our constraints with others from terrestrial experiments as well as astrophysical/cosmological observations.

First, we should discuss one important limit in the cosmological approach. As stated by [39] the scalar mass has a lower bound imposed by the size of the Universe at the relevant epochs. This is because the De Broglie wavelength of the particle, \( l \propto m^{-1} \), cannot be larger than the size of the Universe. Otherwise, it will escape the Hubble horizon. For our considerations, we establish the scalar mass lower bound from Hubble radius at 2 MeV, \( H^{-1}(2\text{MeV}) \). Thus, our results are only valid for \( m_\phi \gtrsim 1.46 \times 10^{-15} \text{eV} \). Interestingly, notice that the size of the Hubble horizon at those epochs is smaller than the size of The Sun.

Here, we present a new stringent bound on the scalar-neutrino coupling \( g_\nu \), which is particularly robust for ultralight scalar masses. Note that the bound on \( G_S \) [from Eq. (11)] permits us to find a mass-dependent bound on \( g_\nu \), that goes as

\[
g_\nu < 7.87 \times 10^{-3} \left( \frac{m_\phi}{\text{eV}} \right) \quad (95\% \text{ CL}) \tag{12}
\]

This new bound restricts a large new region in the parameter space \( (m_\phi, g_\nu) \) for masses \( 1.46 \times 10^{-15} \text{eV} \lesssim m_\phi \lesssim 2.54 \times 10^{-3} \text{eV} \).

In the literature, we spot that there have been extensive efforts to impose bounds on the neutrino-scalar coupling. For instance, it has been constrained by coherent elastic neutrino-nucleus scattering (CEνNS) and by scalar emission in neutrinoless double beta decay experiments [28, 31, 32, 59].

Astrophysical and cosmological observations set the strongest constraints on the neutrino-scalar coupling. A neutrino flavor-dependent scalar interaction is responsible for several non-observed effects in supernovas (SN). Such effects include a loss of SN luminosity, loss of leptons in the supernova core (deleptonization), and trapping of neutrinos by dispersion with a (pseudo)scalar. In Fig. 6 we depict the strongest SN bounds, corresponding to a (pseudo)scalar coupled to electron neutrinos \( |g_{ee}| \) [28, 32].

We also revisit a couple of cosmological bounds. The bound imposed by [39] confronts the positive contribution of a light scalar particle to \( N_{\text{eff}} \), namely \( \Delta N_{\text{eff}} \), with BBN physics. On the other hand, in [33, 34], the authors studied the observable effects on the CMB caused by neutrino self-interactions mediated by a very light scalar particle \( m_\phi \ll T_\nu \). As neutrinos become collisional again at small temperatures, this approximation holds for \( m_\phi \ll T_\nu(z = 100) \), roughly \( m_\phi \lesssim 10^{-3} \text{eV} \). They found the bound \( g_{\nu, \text{eff}} < 2 \times 10^{-7} \), where the ratio between \( g_{\nu, \text{eff}} \) and \( g_\nu \) is no larger than one order of magnitude. In Fig. 6 we depicted all these bounds including our new measurement.

We now discuss the bound on the electron-scalar coupling \( g_\nu \). Notice that here we cannot set a direct constraint on \( g_\nu \) because we do not have a direct measurement on \( g_\nu \), we only have an upper bound. We can only estimate where the bound would lie by using the constraint on \( G_S \) together with the bound on \( G_{\nu, \text{eff}} \) from Eq. (11). Taking the \( g_\nu \) upper value given in Eq. (12) we estimate

\[
g_\nu < 4.83 \times 10^{-10} \left( \frac{m_\phi}{\text{eV}} \right), \tag{13}
\]

for \( g_\nu \) fixed to its upper bound.

Supposing an eventual future measurement of \( g_\nu \), we observe that the bound on \( g_\nu \) is weaker in our analysis than those obtained from neutrinos SNSI from The Sun and SN. The Sun bound is particularly interesting, since [10] found a possible preference for a non-vanishing SNSI \( G_{\nu, \text{eff}} \) effective coupling. As a matter of fact, aside from the fluctuation, we can safely take the solar bound as \( \Delta m_\odot < 7.4 \times 10^{-3} \text{eV} \). The solar medium is non-relativistic, therefore, the mass correction in The Sun goes as \( \Delta m_\odot = G_{\nu, \text{eff}} n_\text{e}\odot \), where the number density of electrons at the solar core is \( n_\text{e}\odot \sim 5.2 \times 10^{11} \text{eV} \). With this \( g_{\text{ee}} = m_\phi^2 \Delta m_\odot / g_\nu \). Comparing the solar bound on \( \Delta m_\odot \) with our Early Universe bound, and taking the upper value in eq. (12), we observe that, indeed, the Early Universe bound is weaker than the solar one (see Fig. 7).

In Fig. 7 we show different constraints compared with our results. On the one hand side, there are strong bounds, \( g_\nu < 10^{-15} \), from stellar physics where an electron-scalar coupling would diminish stars to a cooler than expected state. This is due to energy loss caused by the unopposed escape of scalar particles produced from the stellar nucleus [60, 61]. On the other hand, fifth force experiments, that search for deviations to the Newtonian gravity set the strongest bounds for electron-scalar coupling with a very light mediator. The length scale of the gravitational experiments is related to the force mediator mass, which constraints can be directly obtained.
Experimental constraints at shorter lengths were reviewed and summarized by [39] with the results from [62]. Here, we depict the compendium of fifth force experiments within a single bound in Fig. 7. Additionally, we include the curves indicating the reference bound on \( g_e \) for a fixed \( g_e \)-value in the case of neutrinos propagating in supernovas and The Sun from [39].

VI. SUMMARY AND CONCLUSIONS

We have performed a robust analysis of the consequences of a possible large neutrino effective mass due to thermal corrections mediated by non-standard light scalar interactions among leptons in the context of the Early Universe. Such an effective mass is fed by CFS of propagating neutrinos through a thermal bath of neutrinos and electrons/positrons within the primordial plasma. At one-loop order, the effective neutrino mass is simply proportional to the respective scalar to neutrino/electron couplings, but inversely proportional to the square scalar mass. One can encode such dependencies in a couple of SNSI parameters, \( G_S \) and \( G_{\text{eff}} \). The effective neutrino mass also depends on the temperature of the corresponding thermal bath, through a monotonically increasing function, such that, as higher the temperatures the larger the effective mass contributions. Hence, even if none visible effects appear at small redshifts, possible changes on standard physics could arise as we look towards earlier times.

In the case where the neutrino effective mass gets comparable with neutrino temperature, their number and energy density drops significantly. However, the SNSI effect vanishes faster than the temperature drop and, in equilibrium, the standard neutrino density is recovered. Nevertheless, once neutrinos decouple from the primordial plasma its production gets largely suppressed, thus, their density at decoupling freezes out. This has an observable direct effect that is expressed as a smaller \( N_{\text{eff}} \) than expected.

BBN has shown to be sensitive to any non-standard physics that affects the expansion rate. We have exploited this feature and used BBN primordial nuclei outputs and observational data to set a constraint on the neutrino-scalar coupling. Our new bound on \( g_e \) is more restrictive that previously known bounds for the mass range \( 1.46 \times 10^{-15} \text{eV} \lesssim m_\nu \lesssim 2.54 \times 10^{-5} \text{eV} \).

Although our analysis is able to constraint the scalar-electron couplings, it does also involve scalar-neutrino coupling, and, thus, no straightforward bound to the former can be set without knowledge about the latter. Nevertheless, we have explored the parameter space assuming the saturation of our bound on the scalar-neutrino coupling to compare with other results from astrophysics and fifth force experimental limits.

Along with our analysis, we have assumed that the light scalar mediator would play no direct role in early cosmology, by looking upon the parameter range where it would stay out of equilibrium, and its production rate suppressed during the Early Universe. In the opposite scenario, neutrino NSI may have other consequences that can be further studied in cosmology. For instance, neutrino NSI may trigger active neutrino decays and annihilation into light bosons. Adding such effects to our analysis would probably amount to soften our bounds, since light scalars add to the relativistic degrees of freedom, rising \( N_{\text{eff}} \) and compensating the effect of thermal neutrino mass [39]. Furthermore, neutrino decay and annihilation during structure formation era could relax the bound on \( \Sigma m_\nu \) from LSS [33, 48], in contrast, larger bare neutrino masses impose more stringent constraints. Such analysis may be worthy of being pursued.

The physics of non-standard neutrino interactions is an active field of study due to its potential to solve current tensions in cosmology. The Early Universe can be used as a testing ground to study such interactions in environments unreachable by terrestrial or solar experiments. In this work, we used the indirect effect of neutrinos on the relic densities of light elements to impose bounds upon the possible interactions with a light scalar mediator. This bound is stronger than the previous bounds and contributes to a better understanding of the nature of neutrinos and its possible links to physics outside the standard model of particles.
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