Collisional production of sterile neutrinos via secret interactions and cosmological implications

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Secret interactions among sterile neutrinos have been recently proposed as an escape route to reconcile eV sterile neutrino hints from short-baseline anomalies with cosmological observations. In particular models with coupling $g_X \gtrsim 10^{-2}$ and gauge boson mediators X with $M_X \lesssim 10$ MeV lead to large matter potential suppressing the sterile neutrino production before the neutrino decoupling. With this choice of parameter ranges, big-bang nucleosynthesis is left unchanged and gives no bound on the model. However, we show that at lower temperatures when active-sterile oscillations are no longer matter suppressed, sterile neutrinos are still in a collisional regime, due to their secret self-interactions. The interplay between vacuum oscillations and collisions leads to a scattering-induced decoherent production of sterile neutrinos with a fast rate. This process is responsible for a flavor equilibration among the different neutrino species. We explore the effect of this large sterile neutrino population on cosmological observables. We find that a signature of strong secret interactions would be a reduction of the effective number of neutrinos $N_{\rm eff}$ at matter radiation equality down to 2.7. Moreover, for $M_X \gtrsim g_X$ MeV sterile neutrinos would be free-streaming before becoming nonrelativistic and they would affect the large-scale structure power spectrum. As a consequence, for this range of parameters we find a tension of an eV mass sterile state with cosmological neutrino mass bounds.

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I. INTRODUCTION

In recent years there has been a renewed attention on eV sterile neutrinos, suggested by anomalies found in the short-baseline neutrino experiments (see [1] for a recent review). In this context, it has been realized that these light sterile states would be efficiently produced in the early Universe by oscillations with active neutrinos, leading to a conflict with different cosmological observations [2,3]. This problem motivated the investigation of different mechanisms to suppress the sterile neutrino thermalization (see, e.g., [2,4–6]). In [7,8] has been recently proposed a novel suppression mechanism based on the introduction of secret interactions among sterile neutrinos, mediated by a massive gauge boson X, with $M_X \ll M_W$ (see [9] for the case of sterile neutrinos interacting with a light pseudoscalar). Indeed, these secret interactions would generate a large matter term in the sterile neutrino sector, which lowers the effective neutrino in-medium mixing angle.

However, as the matter potential declines, a resonance is eventually encountered. Sterile neutrinos would be produced by a combination of (damped) Mikheyev-Smirnov-Wolfenstein (MSW)-like resonant flavor conversions among active and sterile neutrinos [10] and the nonresonant processes associated with the secret collisional effects [11]. It particular, assuming for secret interactions a coupling constant $g_X \gtrsim 10^{-2}$ and masses of the mediator $M_X \gtrsim \mathcal{O}(10)$ MeV, the sterile neutrino production would occur at $T \gtrsim 0.1$ MeV, with nontrivial consequences on big-bang nucleosynthesis (BBN) [7]. In this context, in [12] it has been shown that BBN observations would lead to severe constraints on the parameter space of the model, reducing the range which satisfies cosmological bounds.

In [8] much lighter bosons were considered. The advantage of this choice seems twofold. On one hand the matter potential would be so strong to inhibit any sterile neutrino production during BBN, thus evading the corresponding constraints. Moreover, if the new interaction mediator X couples not only to sterile neutrinos but also to dark matter particles, for such small masses it might also possibly relieve some of the small-scale structure problems of the cold dark matter scenario [8,13,14] (nonstandard interactions were also introduced to alleviate these problems in [15–20] and in the references therein). Secret interactions among sterile neutrinos, mediated by very light (or even massless) pseudoscalars, and their connection with dark matter was explored in [9].

The aim of this paper is to show that also for these small masses a large sterile neutrino production is unavoidable at $T \ll 0.1$ MeV, when the matter potential becomes smaller than the vacuum oscillation term. Indeed, the small vacuum oscillations act as a *seed* for a *scattering-induced decoherent production* of sterile neutrinos, associated with the secret self-interactions in the sterile sector. This process is very rapid and leads to a quick flavor equilibration among the active and the sterile neutrino species after the active neutrino decoupling. We will explore the consequences of this large sterile neutrino abundance on cosmological observables, namely the effective number of neutrinos $N_{\rm eff}$ at matter radiation equality and recombination and the sterile neutrino mass bounds.

The paper is organized as follows. In Sec. II we present an overview of the flavor evolution for the active-sterile neutrino system in the presence of secret interactions. In Sec. III we discuss the scattering-induced decoherent production of sterile neutrinos, associated with the secret interactions in the postdecoupling epoch. In Sec. IV we discuss the observable signatures of this scenario for very fast sterile-sterile scattering processes. We find that active-sterile neutrino flavor conversions lead to a reduction of $N_{\rm eff}$ down to 2.7. Furthermore, for $M_X \gtrsim g_X$ MeV sterile neutrinos would be free-streaming before they become a nonrelativistic species. Thus, their large number density would affect the large-scale structure power spectrum. In this case, we compare the sterile neutrino abundance with the most recent cosmological mass bounds finding a tension for an eV mass range sterile state. Finally, in Sec. V we summarize our results and conclude.

II. FLAVOR EVOLUTION

We consider a $3+1$ neutrino mixing scenario, involving the three active families and a sterile species. As usual, we describe the neutrino (antineutrino) system in terms of 4×4 density matrices $\rho = \rho(p)$, for the different neutrino momenta p. The evolution equation of the density matrix ρ is ruled by the kinetic equations [4,21,22]

$$\mathrm{i}\frac{d\rho}{dt} = [\Omega, \rho] + C[\rho]. \tag{1}$$

Assuming no neutrino asymmetry, the dynamics of antineutrinos is identical to the one of neutrinos. The evolution in terms of the comoving observer proper time t can be easily recast in function of the neutrino temperature T_ν (see [4] for a detailed treatment). The first term on the right-hand side of Eq. (1) describes the flavor oscillations Hamiltonian, given by

$$\Omega = \frac{\mathsf{M}^2}{2p} + \sqrt{2}G_{\rm F}\left[-\frac{8p}{3}\left(\frac{\mathsf{E}_\ell}{M_{\rm W}^2} + \frac{\mathsf{E}_\nu}{M_{\rm Z}^2}\right)\right] + \sqrt{2}G_{\rm X}\left[-\frac{8p\mathsf{E}_s}{3M_{\rm X}^2}\right], \tag{2}$$

where $\mathsf{M}^2 = \mathcal{U}^\dagger \mathcal{M}^2 \mathcal{U}$ is the neutrino mass matrix, written in terms of the solar $\Delta m^2_{\rm sol}$, the atmospheric $\Delta m^2_{\rm atm}$ [23] and sterile $\Delta m^2_{\rm st}$ [24,25] mass-squared differences. We have $\Delta m^2_{\rm sol} \ll \Delta m^2_{\rm atm} \ll \Delta m^2_{\rm st} \sim \mathcal{O}(1)$ eV^2. Here $\mathcal{U}$ is the 4×4 active-sterile mixing matrix, parametrized as in [4], where the values of the different mixing angles are given by the global fits of the active [23] and of the sterile neutrino mixings [24,25], respectively. The terms proportional to the Fermi constant G_F in Eq. (2) are the standard matter effects in active neutrino oscillations. In particular, the two contributions E_ℓ and E_ν are the energy density of $e^\pm$ pairs, and ν and $\bar{\nu}$, respectively. Finally, the term proportional to G_X in Eq. (2) represents the new matter secret potential where E_s is the energy density of ν_s and $\bar{\nu}_s$.

Flavor evolution generally occurs at $T_\nu \ll M_X$ (we comment below for the case $T_\nu > M_X$), so that one can reduce it to a contact interaction, with an effective strength[1]

$$G_X = \frac{\sqrt{2}}{8}\frac{g_X^2}{M_X^2}. \tag{3}$$

The numerical factor $\sqrt{2}/8$ has been included in order to have in the X sector the same relation which holds among the Fermi constant G_F, the $SU(2)_L$ coupling constant g and the W mass in the Standard Model. When this matter term is of the order of the vacuum oscillation frequency, associated with $\Delta m^2_{\rm st}$, a MSW resonance among active and sterile neutrinos occurs [10]. In the following we will neglect the standard matter effects in neutrino oscillations, proportional to G_F, since we are in the limit $G_F \ll G_X$.

The last term in the right-hand side of Eq. (1) represents the collisional term. Since we will work at $T_\nu \ll 1$ MeV, the standard collisional term $\propto G_F^2$ in the active sector [26] can be neglected and we only consider the collisional effect in the sterile sector, associated with the secret self-interactions $\nu_s\nu_s \to \nu_s\nu_s$, which reads [27]

$$C[\rho] = -i\frac{\Gamma_X}{2}[\mathsf{S}_X, [\rho, \mathsf{S}_X]], \tag{4}$$

where

$$\Gamma_X \simeq G_X^2 T_\nu^5 \frac{p}{\langle p\rangle}\frac{n_s}{n_a} \tag{5}$$

is the scattering rate [26],[2] with $\langle p\rangle \simeq 3.15 T_\nu$ the average momentum for a thermal Fermi-Dirac distribution, and n_s and n_a the sterile and the active neutrino abundance, respectively. In the flavor basis, $\mathsf{S}_X = \mathrm{diag}(0,0,0,1)$ is the matrix with the numerical coefficients for the scattering process. Notice that in the range we consider for G_X the

[1]See [8] for an explicit calculation of the neutrino potential associated with the secret interactions.

[2]Note a typo in the scattering rate associated with secret interactions in [8], where it was written as $\propto G_F^2$ instead of G_X^2.

elastic scattering terms which redistribute momenta are much larger than the Hubble parameter, so we expect any initial sterile distribution will rapidly approach the standard Fermi-Dirac shape.

The interplay between flavor oscillations and collisions becomes more transparent considering the equations of motion [Eqs. (1) and (2)] in the case of mixing of only one active flavor with the sterile neutrinos. In this case one may write $\rho = \frac{1}{2}(1 + \mathbf{P}\cdot\boldsymbol{\sigma})$, $\Omega = \frac{1}{2}(\omega_0 + \Omega\cdot\boldsymbol{\sigma})$ and $\mathrm{S}_X = \frac{1}{2}(s_0 + \mathbf{S}_X\cdot\boldsymbol{\sigma})$, where σ are the Pauli matrices. One recovers the known evolution of the polarization vector $\mathbf{P}$, expressed by Stodolsky's formula [2,28]

$$\frac{d\mathbf{P}}{dt} = \Omega \times \mathbf{P} - D\mathbf{P}_T. \tag{6}$$

The first term at the right-hand side represents the precession of the polarization vector $\mathbf{P}$ around Ω. The effect of the collisions is to destroy the coherence of the flavor evolution, leading to a shrinking of the length of $\mathbf{P}$. More precisely, the component of the polarization vector "transverse" to the flavor basis $\mathbf{P}_T$ is damped with a rate $D = (1/2)\Gamma_X|\mathbf{S}_X|^2$ (we remind the reader that $\mathbf{P}_T$ represents the off-diagonal elements of ρ). As shown in [28,29], the combination of precession and damping can lead to different behaviors depending on the relative strength of the two effects.

In particular, when the typical oscillation rate ($t_{\rm osc}^{-1}$), collision rate ($t_{\rm coll}^{-1}$) and the expansion rate of the Universe (H) obey the hierarchy $t_{\rm osc}^{-1} \gg t_{\rm coll}^{-1} \gg H$ the flavor dynamics can be described as follows. Active neutrinos ν_α start to oscillate into sterile states ν_s. The average probability to find a sterile neutrinos in a scattering time scale (Γ_X^{-1}) is given by $\langle P(\nu_\alpha \to \nu_s)\rangle_{\rm coll}$. In each collision, the momentum of the ν_s component is changed, while the ν_α component remains unaffected. Therefore, after a collision the two flavors are no longer in the same momentum state, and then they can no longer oscillate. Conversely, they start to evolve independently. However, the remaining active neutrinos can develop a new coherent ν_s component which is made incoherent in the next collision, and so on. The process continues till one reaches a flavor equilibrium with equal number of ν_α and ν_s (corresponding to $|\mathbf{P}| = 0$, i.e. a completely mixed ensemble). Starting with a pure active flavor state, corresponding to $\rho = \mathrm{diag}(1, 0)$, the final density matrix would be $\rho = \mathrm{diag}(1/2, 1/2)$. The averaged relaxation rate to reach this chemical equilibrium is [11,28]

$$\Gamma_t \simeq \langle P(\nu_\alpha \to \nu_s)\rangle_{\rm coll}\Gamma_X. \tag{7}$$

Notice that this rate is nonzero as soon as an initial sterile neutrino density is produced when the matter term becomes of the order of the vacuum oscillation frequency. This initial sterile abundance is again proportional to the conversion probability. Thus Γ_t is, at first, proportional to the square of $\langle P(\nu_\alpha \to \nu_s)\rangle_{\rm coll}$.

In the following we will show that the conditions to trigger this dynamics are always fulfilled at $T_\nu \ll 1$ MeV, leading to a copious sterile neutrino production. The key observation is that, although the refractive potential is smaller than the oscillation rate and the collisional cross sections are even smaller, yet they exceed the Hubble rate and lead to scattering-induced decoherent production of sterile neutrinos.

III. STERILE NEUTRINO PRODUCTION BY SCATTERING-INDUCED DECOHERENCE

In Fig. 1 we show the behavior of the different neutrino refractive and collisional rates normalized to the Hubble rate $H(T_\gamma)$, versus photon temperature $T_\gamma = 1.4T_\nu$ (see [4] for details). For the sake of illustration, we show the quantity of Eqs. (2)–(5) averaged over thermal Fermi-Dirac distributions. Results are shown for $g_X = 10^{-1}$. The left panel is for $G_X = 10^8 G_F$, or $M_X = 1.2$ MeV, while the

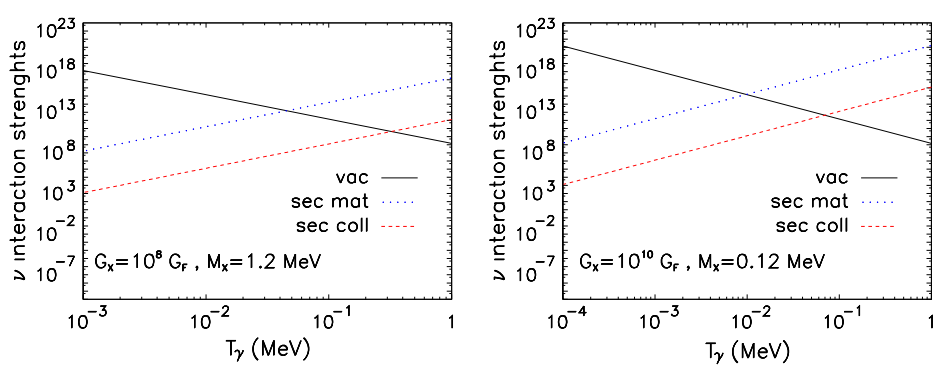

FIG. 1 (color online). Neutrino refractive and collisional rates (normalized in terms of the Hubble rate) versus photon temperature T_γ for $g_X = 0.1$. The left panel corresponds to $G_X = 10^8 G_F$ and $M_X = 1.2$ MeV, while the right panel to $G_X = 10^{10} G_F$ and $M_X = 0.12$ MeV. The curves are the active-sterile vacuum term (solid curve), the secret matter potential for $\rho_{ss} = 0.06$ (dotted curve), and the scattering-induced decoherent production rate Γ_t associated with G_X^2, assuming a sterile neutrino abundance from vacuum oscillations (dashed curve).

right panel corresponds to $G_X = 10^{10} G_F$, i.e. $M_X = 0.12$ MeV. We show the active-sterile vacuum term (solid curve) and the secret matter potential (dotted curve) assuming $\rho_{ss} = n_s/n_a = 0.06$, corresponding to the initial sterile neutrino abundance induced by vacuum oscillations (see later). We note that for $T_\nu > M_X$ the real form of the matter potential would deviate from the contact structure of Eq. (3) used in the figure (see [8]). In particular, for $T_\nu \simeq M_X$ the potential would vanish, leading to a possible production of ν_s when this condition is fulfilled. However, since the duration of this phase is expected to be shorter than the inverse of the sterile neutrino production rate Γ_t^{-1}, for simplicity we neglect this possible (small) extra contribution of sterile neutrinos. In the left panel a resonance would take place at $T_\gamma \simeq 5 \times 10^{-2}$ MeV, while in the right panel at $T_\gamma \simeq 1 \times 10^{-2}$ MeV. This resonance excites sterile states.

In principle one should perform numerical simulations in a $(3+1)$ scheme in order to calculate the resonant sterile neutrino abundance and the further flavor evolution. However, in the presence of the very large matter potential and collisional term, induced by the secret interactions, these would be computationally demanding. Moreover, our main argument is not related to the details of the corresponding dynamics. Therefore, for simplicity we assume that the resonance is completely nonadiabatic, so we have to take into account only the vacuum production of sterile neutrinos at lower temperatures when the matter term becomes smaller than the vacuum oscillation term, associated with $\Delta m^2_{\rm st}$. This is a very conservative assumption. However, it allows us to easily compute the flavor evolution and is enough to show the role of the damping term.

The active-sterile vacuum oscillation probability, averaged over a collision time scale, is given by [30]

$$\langle P(\nu_\alpha \to \nu_s)\rangle_{\rm coll} \simeq \frac{1}{2}\sin^2 2\theta_{\alpha s}. \tag{8}$$

Taking as the representative mixing angle $\sin^2 2\theta_{es} \simeq 0.12$ [24], one would expect a sterile neutrino abundance, $n_s \simeq 0.06 n_a$. This seems a negligible contribution but is enough to generate a large scattering rate proportional to G_X^2; see Eq. (7). This is shown in Fig. 1 as dashed curves. As one can see, at $T_\gamma \lesssim 10^{-2}$ MeV, $\Gamma_t \gg H(T_\gamma)$. Therefore, the scattering-induced decoherent production will lead to a quick flavor equilibrium. Starting with a density matrix having

$$(\rho_{ee}, \rho_{\mu\mu}, \rho_{\tau\tau}, \rho_{ss})_{\rm initial} = (1, 1, 1, 0), \tag{9}$$

one thus quickly reaches

$$(\rho_{ee}, \rho_{\mu\mu}, \rho_{\tau\tau}, \rho_{ss})_{\rm final} = (3/4, 3/4, 3/4, 3/4) \tag{10}$$

for all the parameter space associated with eV sterile neutrino anomalies. Notice that this result does not depend on the particular value of G_X but only on the fact that the condition of strong damping is realized. The same equilibrium value would remain valid for example, for smaller masses M_X by different order of magnitudes, at least until we can treat the collisional term as a four-point effective interactions, i.e. for $T_\nu > M_X$. As discussed before, the rate of sterile neutrino rethermalization is extremely fast, so that the process is instantaneous. Indeed, from Eqs. (5) and (7) one can estimate $\Gamma_t \gtrsim 10^{-18}$ MeV for the cases we have shown. This means that it is practically impossible to numerically follow the rise of the sterile neutrino production. However, it can be interesting to appreciate this dynamics in a case where the process is slower. We address the interested reader to the lower panel of Fig. 3 in [12], where it is shown the evolution of the diagonal elements of the density matrix ρ in a $(2+1)$ scheme, for a scenario with $g_X = 10^{-2}$ and $G_X = 10^3 G_F$. In this case the sterile neutrino production starts at $T \lesssim 1$ MeV and the final value would be $\rho = 2/3$, as expected with only three oscillating neutrino families.

Furthermore, we mention that the final equilibrium value does not depend on the exact values of the active-sterile neutrino mixing angles. Indeed, we explicitly checked that also assuming that the sterile species mixes only with an active one, all the flavors will participate to the equilibrium, due to the presence of the active mixing angles that connect the different species, and indirectly link them also to the sterile species.

Secret interactions mediated by a light (or even massless) pseudoscalar ($M_X \ll T$), with a Lagrangian $\mathcal{L} \sim g_s X \bar{\nu}\gamma_5 \nu$, were considered in [9]. It was found that couplings $g_s \sim 10^{-5}$ were sufficient to block thermalization prior to neutrino decoupling (and make dark matter sufficiently self-interacting). However, as long as $g_s \gtrsim 10^{-6}$ the $\nu_s - X$ plasma would be strongly interacting till sterile neutrinos become nonrelativistic. Therefore, the mechanism of ν_s production and flavor equilibration would apply also to this case. However, in the massless case, there would be also the production of a bath of X via the $\bar{\nu}_s \nu_s \to XX$ process (with a rate $\Gamma_X \sim g_s^4 T$). Therefore one would expect a chemical equilibrium over the five species, with an abundance $\rho = 3/5$ for each of them. As mentioned in [9], a plasma of strongly interacting ν_s and X can have interesting cosmological signatures, since the $\nu_s - X$ bath would behave as a single massless component with no anisotropic stress.

IV. COSMOLOGICAL SIGNATURES

A. Effective neutrino species $N_{\rm eff}$

The initial snapshot of active and sterile neutrino distribution soon after steriles are excited via oscillation is a shared gray-body distribution, a Fermi-Dirac function weighted by a factor 3/4 for each species. In the absence of any interaction, these distribution would remain frozen but for the effect of momentum redshift. In the model we

are considering, after their production sterile neutrinos are fastly rescattering among themselves via secret interactions of order G_X^2. They are, therefore, collisional. In fact, a gray-body distribution is not a solution of the collisional Boltzmann equation, so these scattering processes will push the sterile distribution towards a Fermi-Dirac shape, with the constraint that the total neutrino number density is kept constant. Furthermore, as soon as sterile neutrinos change their distribution this has a feedback on active neutrino distribution too, which are still efficiently oscillating into sterile states. This implies that all neutrino species in the presence of sterile-sterile scattering will adjust quite efficiently their distribution to a thermal equilibrium distribution. The constant number density (or entropy) constraint implies that their eventual temperature is reduced by a factor $(3/4)^{1/3}$ with respect to the initial active neutrino temperature $T_\nu = (4/11)^{1/3} T_\gamma$. Indeed, we have

$$n_\nu = 2\int \frac{d^3p}{(2\pi)^3}\frac{3}{4}\frac{1}{\exp(p/T_\nu)+1} = 2\int \frac{d^3p}{(2\pi)^3}\frac{1}{\exp[p/(T_\nu(3/4)^{1/3})]+1}. \quad (11)$$

There are, in fact, no pair production processes from the electromagnetic plasma which can refill active neutrino densities, since in the scenario we are considering, sterile neutrinos are excited well below the active neutrino decoupling phase.

From these considerations we see that the total energy density stored in active and sterile neutrinos is reduced. Indeed, before sterile neutrinos are excited, and after the $e^+ - e^-$ annihilation phase, active neutrino energy density is given by

$$\epsilon_{\nu,\mathrm{in}} = 3\times\frac{7}{8}\left(\frac{4}{11}\right)^{4/3}\epsilon_\gamma, \quad (12)$$

where the photon energy density is $\epsilon_\gamma = (\pi^2/15)T_\gamma^4$ and the effective number of neutrino species is $N_{\mathrm{eff}} \simeq 3$, neglecting the small effect due to partial neutrino heating of order $\Delta N_{\mathrm{eff}} = 0.046$ [31]. After sterile states are produced via oscillations and kinetic equilibrium is reached via secret interactions, we have four species which share a common temperature $T_\nu = (4/11)^{1/3}(3/4)^{1/3}T_\gamma$. Therefore, till all neutrinos are fully relativistic

$$\epsilon_{\nu,\mathrm{fin}} = 4\times\left(\frac{3}{4}\right)^{4/3}\times\frac{7}{8}\left(\frac{4}{11}\right)^{4/3}\epsilon_\gamma, \quad (13)$$

and correspondingly the value of N_{eff} decreases to

$$N_{\mathrm{eff}} \sim 4\times\left(\frac{3}{4}\right)^{4/3} \sim 2.7. \quad (14)$$

This value is only slightly reduced at the matter radiation equality, i.e. for $T_\gamma \sim 0.7$ eV, since at this energy scale only the high energy tail of sterile neutrino distribution counts as radiation. If we weight this contribution with the ratio of corresponding pressure over the pressure of a purely relativistic gas, as in [30] and assume relativistic active states at this epoch, we find with $m_{\mathrm{st}} \sim \sqrt{\Delta m_{st}^2} \simeq 1$ eV

$$N_{\mathrm{eff}} \sim 3\left(\frac{3}{4}\right)^{1/3}\left(\frac{3}{4}+\frac{1}{4}\frac{P}{P_0}\right) \sim 2.66, \quad (15)$$

where the first and second terms in parentheses are the active and sterile contribution, respectively, and

$$P = 2\int \frac{d^3p}{(2\pi)^3}\frac{p^2}{3\sqrt{p^2+m_{\mathrm{st}}{}^2}}\frac{1}{\exp[p/(T_\nu(3/4)^{1/3})]+1}, \quad (16)$$

$$P_0 = 2\int \frac{d^3p}{(2\pi)^3}\frac{p}{3}\frac{1}{\exp[p/(T_\nu(3/4)^{1/3})]+1}. \quad (17)$$

A value of the effective number of neutrino smaller than the expected standard result would be a signature in favor of strongly interacting sterile states mixed with active neutrinos, though there might be different models which can account for this result (e.g., low-reheating scenarios [32]). Presently, the most precise determination of N_{eff} is from the Planck experiment. In the standard ΛCDM model but allowing for a free number of relativistic species the result is $N_{\mathrm{eff}} = 3.30 \pm 0.27$ (68% C.L.) [33], which is compatible with Eq. (15) at about 2σ. A new data release by the Planck Collaboration is expected soon, including polarization data. This might put further constraint on N_{eff}. In the future even more tight bounds are foreseen to come by next-generation experiment such as Euclid [34], which should reach a sensitivity of order $\Delta N_{\mathrm{eff}} < 0.1$.

B. Cosmological mass bounds

The sterile neutrino production due to the scattering-induced decoherent effects is expected to affect the cosmic microwave background (CMB) and large-scale structures (LSSs), which are both sensitive to neutrino mass scale in the 10^{-1}–1 eV range. One of the main effects of a massive neutrino is due their free-streaming till the epoch when they become nonrelativistic, which suppresses the growth of perturbations on small scales. However, if sterile states scatter via secret interactions, the free-streaming regime is delayed until the scattering rate becomes smaller than the Hubble parameter. This means that if G_X is large enough so that this condition holds at the nonrelativistic transition, sterile neutrinos would never have a free-streaming phase

but always diffuse.[3] The smaller value of G_X for which this happens can be obtained from the condition that the scattering rate equals the value of H at a temperature $3.15T_\nu \sim \langle p \rangle \sim \sqrt{\Delta m^2_{\rm st}}$:

$$G_X^2 T_\nu^5 \sim H(T_\gamma). \tag{18}$$

Using the standard expression of the Hubble rate in the $T_\gamma \sim$ eV range and the fact that, as we have seen, the sterile temperature is given by $T_\nu = (4/11)^{1/3}(3/4)^{1/3}T_\gamma$, this gives $G_X \sim 10^{10} G_F$, which corresponds to $M_X \simeq 10^{-1}$ MeV for $g_X \simeq 10^{-1}$. Therefore, the mass bound discussed below only applies *as long as the coupling G_X is smaller than this value*.

Assuming that the active neutrinos are much lighter than the sterile species, one can define an effective sterile neutrino mass [35]

$$m_{\rm st}^{\rm eff} = \rho_{ss}\sqrt{\Delta m^2_{\rm st}} = \frac{3}{4}\sqrt{\Delta m^2_{\rm st}}. \tag{19}$$

The latest analysis in [24] of the sterile neutrino anomalies gives a best fit $\Delta m^2_{\rm st} = 1.6\ {\rm eV}^2$ with a 2σ range

$$1.08\ {\rm eV}^2 < \Delta m^2_{\rm st} < 1.99\ {\rm eV}^2. \tag{20}$$

Using Eq. (19) the lower value in the 2σ range gives $m_{\rm st}^{\rm eff} \simeq 0.78$ eV. This value has to be compared with the cosmological mass bounds. We also comment that the global analysis of sterile neutrino anomalies presented in [25] finds a discrepancy between appearance and disappearance sterile neutrino data. As a consequence only a small region around $\Delta m^2_{\rm st} \simeq 0.9\ {\rm eV}^2$ would be compatible with all data. This would correspond to $m_{\rm st}^{\rm eff} \simeq 0.7$ eV.

Cosmological bounds on sterile neutrino mass and abundance in the early Universe are rather sensitive to the data set used in the analysis, notably CMB data from Planck [33] and the recent but controversial BICEP2 experiment [36] (see also [37]), LSSs, H_0 measurements as well as lensing and cluster data (CFHTLenS + PSZ). In particular, the Planck Collaboration combines Planck with WMAP polarization data, baryon acoustic oscillation and high multipole CMB data. In this case the bound obtained is $m_{\rm st}^{\rm eff} < 0.42$ eV at 95% C.L. [33], which is in strong disagreement with the sterile neutrino abundance produced by the secret interactions. Moreover a possible nonzero sterile neutrino mass has been claimed in order to relieve the discrepancy between the CMB measurements and other observations, like current expansion rate H_0, the galaxy shear power spectrum and counts of galaxy clusters, providing a value $m_{\rm st}^{\rm eff} \simeq 0.7$ eV at 2σ [38–41] (see also [42]) or an upper bound of $m_{\rm st}^{\rm eff} < 0.6$ eV [43]. Stronger bounds have also be quoted in [41,43].

[3]We are very pleased to thank Basudeb Dasgupta for pointing us to this possibility.

In general, from the results presented here, one would conclude that the minimum $m_{\rm st}^{\rm eff} \simeq 0.78$ eV obtained from the secret collisional production and compatible with the $\Delta m^2_{\rm st}$ range would be in tension (at least at 2σ level) with the bounds on sterile neutrino mass from cosmology. A possible way out to this result is to consider extremely high couplings, $G_X \geq 10^{10} G_F$, since in this case sterile-sterile scatterings are in equilibrium till the eV scale, when they become nonrelativistic. As we mentioned, they would never experience a free-streaming regime and cosmological mass bounds do not apply.

To close, we remark that in deriving our constraint we have been conservative, since we have assumed that sterile neutrinos are produced only by vacuum oscillations at $T \ll 1$ MeV. However, it has been argued in [8,13] that there could be another colder primordial population of ν_s generated at $T \gg$ GeV by the decoupling of the $U(1)_X$ sector from standard particles. This additional contribution would increase the tension between the sterile neutrino abundance and the cosmological mass bound.

V. CONCLUSIONS

Secret interactions among sterile neutrinos mediated by a light gauge boson X have been recently proposed as an intriguing possibility to suppress the thermalization of eV sterile neutrinos in the early Universe. In particular, interactions mediated by a gauge boson with $M_X \leq 10$ MeV would suppress the sterile neutrino productions for $T \gtrsim 0.1$ eV and seemed therefore safe from cosmological constraints related to big-bang nucleosynthesis [12].

In the present work we have shown that when the matter potential produced by the sterile interactions becomes smaller than the vacuum oscillation frequency, sterile neutrinos are copiously produced by the scattering-induced decoherent effects in the sterile neutrino sector. This process would lead to a quick flavor equilibration, with a sterile neutrino abundance largely independent on the specific values of G_X and M_X. A possible complete rethermalization of sterile neutrinos and its impact on mass bound was already advocated in [13]. Indeed, we have shown that due to the large damping effects this is *always* the case in secret interactions among sterile neutrinos with low M_X masses.

We investigated the cosmological consequences of this huge sterile neutrino population. We find that a signature of secret interactions would be a reduction of the effective number of neutrinos $N_{\rm eff}$ down to 2.7. If this value is compatible with the 2σ range given by the Planck experiment [33], the future experiment Euclid [34], with a sensitivity to $\Delta N_{\rm eff} < 0.1$ may probe this small deviation with respect to the standard expectation. Moreover, for $M_X \gtrsim g_X$ MeV, sterile neutrinos would be free-streaming at the matter-radiation equality epoch. Then, the large sterile neutrino production would be in tension with the most recent cosmological mass bounds on sterile neutrinos.

We note that for the parameters $M_X \simeq g_X$ MeV, where secret interactions would play also an interesting role in relation to dark matter and small-scale structures [8], sterile neutrinos would be at the border between free-streaming and collisional regime at the neutrino decoupling. In this case a dedicated investigations is necessary to assess if mass bounds apply also in this case or can be evaded.

Finally, we mention that recently it also has been speculated that the background of sterile neutrinos produced by the collisions associated with the secret interactions would also modify the optical depth of ultrahigh energy recently observed by IceCube [44]. Therefore, future high-energy neutrino observations would be an interesting additional channel to probe this scenario.

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[1] K. N. Abazajian *et al.*, Light sterile neutrinos: A white paper, arXiv:1204.5379.

[2] S. Hannestad, I. Tamborra, and T. Tram, Thermalisation of light sterile neutrinos in the early universe, J. Cosmol. Astropart. Phys. 07 (2012) 025.

[3] J. Hamann, S. Hannestad, G. G. Raffelt, and Y. Y. Y. Wong, Sterile neutrinos with eV masses in cosmology: How disfavoured exactly?, J. Cosmol. Astropart. Phys. 09 (2011) 034.

[4] A. Mirizzi, N. Saviano, G. Miele, and P. D. Serpico, Light sterile neutrino production in the early universe with dynamical neutrino asymmetries, Phys. Rev. D **86**, 053009 (2012).

[5] N. Saviano, A. Mirizzi, O. Pisanti, P. D. Serpico, G. Mangano, and G. Miele, Multi-momentum and multi-flavour active-sterile neutrino oscillations in the early universe: Role of neutrino asymmetries and effects on nucleosynthesis, Phys. Rev. D **87**, 073006 (2013).

[6] C. M. Ho and R. J. Scherrer, Sterile neutrinos and light dark matter save each other, Phys. Rev. D **87**, 065016 (2013).

[7] S. Hannestad, R. S. Hansen, and T. Tram, How Secret Interactions Can Reconcile Sterile Neutrinos with Cosmology, Phys. Rev. Lett. **112**, 031802 (2014).

[8] B. Dasgupta and J. Kopp, A Ménage à Trois of eV-Scale Sterile Neutrinos, Cosmology, and Structure Formation, Phys. Rev. Lett. **112**, 031803 (2014).

[9] M. Archidiacono, S. Hannestad, R. S. Hansen, and T. Tram, Cosmology with self-interacting sterile neutrinos and dark matter—A pseudoscalar model, arXiv:1404.5915.

[10] L. Wolfenstein, Neutrino oscillations in matter Phys. Rev. D **17**, 2369 (1978); S. P. Mikheyev and A. Yu. Smirnov, Resonance enhancement of oscillations in matter and solar neutrino spectroscopy, Yad. Fiz. **42**, 1441 (1985) [Sov. J. Nucl. Phys. **42**, 913 (1985)].

[11] K. Kainulainen, Light singlet neutrinos and the primordial nucleosynthesis, Phys. Lett. B **244**, 191 (1990).

[12] N. Saviano, O. Pisanti, G. Mangano, and A. Mirizzi, Unveiling secret interactions among sterile neutrinos with big-bang nucleosynthesis, Phys. Rev. D **90**, 113009 (2014).

[13] T. Bringmann, J. Hasenkamp, and J. Kersten, Tight bonds between sterile neutrinos and dark matter, J. Cosmol. Astropart. Phys. 07 (2014) 042.

[14] P. Ko and Y. Tang, $\nu\Lambda$MDM: A model for sterile neutrino and dark matter reconciles cosmological and neutrino oscillation data after BICEP2, Phys. Lett. B **739**, 62 (2014).

[15] C. Boehm, P. Fayet, and R. Schaeffer, Constraining dark matter candidates from structure formation, Phys. Lett. B **518**, 8 (2001).

[16] P. Serra, F. Zalamea, A. Cooray, G. Mangano, and A. Melchiorri, Constraints on neutrino–dark matter interactions from cosmic microwave background and large scale structure data, Phys. Rev. D **81**, 043507 (2010).

[17] L. G. van den Aarssen, T. Bringmann, and C. Pfrommer, Is Dark Matter with Long-Range Interactions a Solution to All Small-Scale Problems of ΛCDM Cosmology?, Phys. Rev. Lett. **109**, 231301 (2012).

[18] C. Boehm, J. A. Schewtschenko, R. J. Wilkinson, C. M. Baugh, and S. Pascoli, Using the Milky Way satellites to study interactions between cold dark matter and radiation, Mon. Not. R. Astron. Soc. **445**, L31 (2014).

[19] C. Boehm and P. Fayet, Scalar dark matter candidates, Nucl. Phys. **B683**, 219 (2004).

[20] X. Chu and B. Dasgupta, Dark Radiation Alleviates Problems with Dark Matter Halos, Phys. Rev. Lett. **113**, 161301 (2014).

[21] B. H. J. McKellar and M. J. Thomson, Oscillating doublet neutrinos in the early Universe, Phys. Rev. D **49**, 2710 (1994).

[22] G. Sigl and G. Raffelt, General kinetic description of relativistic mixed neutrinos, Nucl. Phys. **B406**, 423 (1993).

[23] F. Capozzi, G. L. Fogli, E. Lisi, A. Marrone, D. Montanino, and A. Palazzo, Status of three-neutrino oscillation parameters, circa 2013, Phys. Rev. D **89**, 093018 (2014).

[24] C. Giunti, M. Laveder, Y. F. Li, and H. W. Long, Pragmatic view of short-baseline neutrino oscillations, Phys. Rev. D **88**, 073008 (2013).

[25] J. Kopp, P. A. N. Machado, M. Maltoni, and T. Schwetz, Sterile neutrino oscillations: The global picture, J. High Energy Phys. 05 (2013) 050.

[26] N. F. Bell, R. R. Volkas, and Y. Y. Y. Wong, Relic neutrino asymmetry evolution from first principles, Phys. Rev. D **59**, 113001 (1999).

[27] Y.-Z. Chu and M. Cirelli, Sterile neutrinos, lepton asymmetries, primordial elements: How much of each?, Phys. Rev. D **74**, 085015 (2006).

[28] L. Stodolsky, On the treatment of neutrino oscillations in a thermal environment, Phys. Rev. D **36**, 2273 (1987).

[29] K. Enqvist, K. Kainulainen, and M. J. Thomson, Stringent cosmological bounds on inert neutrino mixing, Nucl. Phys. **B373**, 498 (1992).

[30] T. D. Jacques, L. M. Krauss, and C. Lunardini, Additional light sterile neutrinos and cosmology, Phys. Rev. D **87**, 083515 (2013); **88**, 109901(E) (2013).

[31] G. Mangano, G. Miele, S. Pastor, T. Pinto, O. Pisanti, and P. D. Serpico, Relic neutrino decoupling including flavor oscillations, Nucl. Phys. **B729**, 221 (2005).

[32] K. Ichikawa, M. Kawasaki, and F. Takahashi, The oscillation effects on thermalization of the neutrinos in the Universe with low reheating temperature, Phys. Rev. D **72**, 043522 (2005).

[33] P. A. R. Ade *et al.* (Planck Collaboration), Planck 2013 results. XVI. Cosmological parameters, Astron. Astrophys. **571**, A16 (2014).

[34] R. Laureijs *et al.* (EUCLID Collaboration), Euclid definition study report, arXiv:1110.3193.

[35] A. Mirizzi, G. Mangano, N. Saviano, E. Borriello, C. Giunti, G. Miele, and O. Pisanti, The strongest bounds on active-sterile neutrino mixing after Planck data, Phys. Lett. B **726**, 8 (2013).

[36] P. A. R. Ade *et al.* (BICEP2 Collaboration), Detection of B-Mode Polarization at Degree Angular Scales by BICEP2, Phys. Rev. Lett. **112**, 241101 (2014).

[37] R. Adam *et al.* (Planck Collaboration), Planck intermediate results. XXX. The angular power spectrum of polarized dust emission at intermediate and high Galactic latitudes, arXiv:1409.5738 [Astron. Astrophys. (to be published)].

[38] J. Hamann and J. Hasenkamp, A new life for sterile neutrinos: Resolving inconsistencies using hot dark matter, J. Cosmol. Astropart. Phys. 10 (2013) 044.

[39] M. Wyman, D. H. Rudd, R. A. Vanderveld, and W. Hu, Neutrinos Help Reconcile Planck Measurements with the Local Universe, Phys. Rev. Lett. **112**, 051302 (2014).

[40] R. A. Battye and A. Moss, Evidence for Massive Neutrinos from Cosmic Microwave Background and Lensing Observations, Phys. Rev. Lett. **112**, 051303 (2014).

[41] M. Archidiacono, N. Fornengo, S. Gariazzo, C. Giunti, S. Hannestad, and M. Laveder, Light sterile neutrinos after BICEP-2, J. Cosmol. Astropart. Phys. 06 (2014) 031.

[42] R. A. Battye, T. Charnock, and A. Moss, Tension between the power spectrum of density perturbations measured on large and small scales, arXiv:1409.2769.

[43] B. Leistedt, H. V. Peiris, and L. Verde, No New Cosmological Concordance with Massive Sterile Neutrinos, Phys. Rev. Lett. **113**, 041301 (2014).

[44] J. F. Cherry, A. Friedland, and I. M. Shoemaker, Neutrino portal dark matter: From dwarf galaxies to IceCube, arXiv:1411.1071.