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 \( M_X \lesssim 10 \text{MeV} \) 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 \( N_{	ext{eff}} \) at matter radiation equality down to 2.7. Moreover, for \( M_X \gtrsim g_X \text{MeV} \) sterile neutrinos would be free-streaming before becoming non-relativistic and they would affect the large-scale structure power spectrum. As a consequence, for this range of parameters we find a tension of a eV mass sterile state with cosmological neutrino mass bounds.

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I. INTRODUCTION

In the recent years there has been a renewed attention on eV sterile neutrinos, suggested by anomalies found in the short-baseline neutrino experiments (see \cite{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 \cite{2,3}. This problem motivated the investigation of different mechanisms to suppress the sterile neutrino thermalization (see, e.g., \cite{4,5}). In \cite{6} it 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 \cite{7} 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 \cite{10}, and the non-resonant processes associated with the secret collisional effects \cite{11}. It particular, assuming for secret interactions a coupling constant \( g_X \gtrsim 10^{-2} \) and masses of the mediator \( M_X \gtrsim \mathcal{O}(10) \text{MeV} \), the sterile neutrino production would occur at \( T \gtrsim 0.1 \text{MeV} \), with non trivial consequences on big bang nucleosynthesis (BBN) \cite{7}. In this context, in \cite{8} 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 \cite{9} 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 \cite{10,11} (nonstandard interactions were also introduced to alleviate these problems in \cite{12} 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 \cite{9}.

The aim of this paper is to show that also for these small masses a large sterile neutrino production is unavoidable at \( T \ll 0.1 \text{MeV} \), when the matter potential becomes smaller than the vacuum oscillation term. Indeed, the small vacuum oscillations act as 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 $N_{\text{eff}}$ at matter radiation equality and recombination and the sterile neutrino mass bounds.

The paper is organized as follows. In Section II we present an overview of the flavor evolution for the active-sterile neutrino system in the presence of secret interactions. In Section III we discuss the scattering-induced decoherent production of sterile neutrinos, associated with the secret interactions in the post-decoupling epoch. In Section IV we discuss the observable signatures of this scenario for very fast sterile-sterile scattering processes. We find that active-sterile neutrino flavor conversions leads to a reduction of $N_{\text{eff}}$ down to 2.7. Furthermore, for $M_X \gtrsim g_X$ GeV sterile neutrinos would be free-streaming before they become a non relativistic 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 observables, respectively. In flavor basis, $S$ describes the Fermi constant $G_F$ instead of $G_f$. The last term in the right-hand side of Eq. (2) 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 (24) can be neglected and we only consider the collisional effect in the sterile sector, associated with the secret self-interactions $\nu_s \nu_s \rightarrow \nu_s \nu_s$, which reads (27)

$$C[\rho] = -i \frac{\Gamma_X}{2} \left[ S_X, [\rho, S_X] \right],$$

where

$$\Gamma_X \simeq G_X^2 t_\nu \left( \frac{p}{(p)} \right) \frac{n_s}{n_\nu},$$

is the scattering rate (26), with $\langle p \rangle \simeq 3.15 T_\nu$ the average-momentum for a thermal Fermi-Dirac distribution, and $n_s$ and $n_\nu$ the sterile and the active neutrino abundance, respectively. In flavor basis, $S_X = \text{diag}(0, 0, 0, 1)$ is the matrix with the numerical coefficients for the scattering process. Notice that, in the range we consider for $G_X$ the elastic scattering terms which redistribute momenta are much larger than Hubble parameter, so we expect any initial sterile distribution will rapidly approach the standard Fermi-Dirac shape.

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 $G_X^2$. 

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 $E_\nu$ and $E_\mu$ are the energy density of $e^\pm$ pairs, and $\nu$ and $\bar{\nu}$, respectively. Finally, the term proportional to $G_X$ in Eq. (3) represents the new matter secret potential where $E_\nu$ is the energy density of $\nu_s$ and $\bar{\nu}_s$.

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

$$G_X = \sqrt{\frac{2}{8}} \frac{g_X^2}{M_X^2}.$$ (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 $\Delta m^2_{\text{atm}}$, 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 (24) can be neglected and we only consider the collisional effect in the sterile sector, associated with the secret self-interactions $\nu_s \nu_s \rightarrow \nu_s \nu_s$, which reads (27)
The interplay between flavor oscillations and collisions becomes more transparent considering the equations of motion [Eq. (11–2)] in the case of mixing of only one active flavor with the sterile neutrinos. In this case one may write \( \rho = \frac{1}{2}(1 + \mathbf{P} \cdot \sigma) \), \( \Omega = \frac{1}{2}(\omega_0 + \Omega \cdot \sigma) \) and \( S_X = \frac{1}{2}(s_0 + S_X \cdot \sigma) \), where \( \sigma \) are the Pauli matrices. One recovers the known evolution of the polarization vector \( \mathbf{P} \), expressed by the Stodolsky’s formula [2, 28]

\[
\frac{d\mathbf{P}}{dt} = \Omega \times \mathbf{P} - D \mathbf{P}_T .
\]  

(6)

The first term at the right-hand-side represents the precession of the polarization vector \( \mathbf{P} \) around \( \Omega \). 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 \( D = (1/2)\Gamma X|S_X|^2 \) (we remind the reader that \( \mathbf{P}_T \) represents the off-diagonal elements of \( \rho \)). 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 \( t_{\text{osc}}^{-1} \), collision rate \( t_{\text{coll}}^{-1} \) and the expansion rate of the universe \( H \), obey the hierarchy \( t_{\text{osc}}^{-1} \gg t_{\text{coll}}^{-1} \gg H \) the flavor dynamics can be described as follows. Active neutrinos \( \nu_\alpha \) start to oscillate into sterile states \( \nu_s \). The average probability to find a sterile neutrinos in a scattering time scale \( (\Gamma_X^{-1}) \) is given by \( \langle P(\nu_\alpha \rightarrow \nu_s) \rangle_{\text{coll}} \). In each collision, the momentum of the \( \nu_s \) component is changed, while the \( \nu_\alpha \) 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 \( \nu_s \) 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 \( \nu_\alpha \) and \( \nu_s \) (corresponding to \( |\mathbf{P}| = 0 \), i.e. a completely mixed ensemble). Starting with a pure active flavor state, corresponding to \( \rho = \text{diag}(1,0) \), the final density matrix would be \( \rho = \text{diag}(1/2, 1/2) \). The averaged relaxation rate to reach this chemical equilibrium is \( t_{\text{osc}}^{-1} \)

\[
\Gamma_t \simeq \langle P(\nu_\alpha \rightarrow \nu_s) \rangle_{\text{coll}} \Gamma X .
\]  

(7)

Notice that this rate is non zero 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 \( \Gamma_t \) is, at first, proportional to the square of \( \langle P(\nu_\alpha \rightarrow \nu_s) \rangle_{\text{coll}} \).

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. STEERLE 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 \( H(T_\nu) \), versus photon temperature \( T_\gamma = 1.4T_\nu \) (see [2] for details). For the sake of illustration, we show the equation of \( \mathbf{P}_0 \) averaged over thermal Fermi-Dirac distributions. Results are shown for \( g_X = 10^{-1} \). Left panel is for \( G_X = 10^8 G_F \), or \( M_X = 1.2 \) MeV, while 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 \( \rho_{ss} = n_s/n_\alpha = 0.06 \), 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. (11) used in the Figure (see [8]). In particular, for \( T_\nu \simeq M_X \) the potential would vanish, leading to a possible production of \( \nu_s \) 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 \( \Gamma_2^{-1} \), 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 non-adiabatic, 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 \( \Delta m_{31}^2 \). 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 \( 10 \)

\[
\langle P(\nu_\alpha \rightarrow \nu_s) \rangle_{\text{coll}} \simeq \frac{1}{2} \sin^2 2\theta_{as} .
\]  

(8)

Taking as representative mixing angle \( \sin^2 2\theta_{as} \approx 0.12 \) [24] one would expect a sterile neutrino abundance, \( n_s \approx 0.06 n_\alpha \). 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 Figure [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

\[(\rho_{ee}, \rho_{\mu\mu}, \rho_{\tau\tau}, \rho_{ss})_{\text{initial}} = (1, 1, 1, 0)\] (9)

one thus, quickly reaches

\[(\rho_{ee}, \rho_{\mu\mu}, \rho_{\tau\tau}, \rho_{ss})_{\text{final}} = (3/4, 3/4, 3/4, 3/4)\] (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 re-thermalization is extremely fast, so that the process is instantaneous. Indeed, from Eq. (9) and (10) 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 \(\rho\) in a (2+1) scheme, for a scenario with \(g_X = 10^{-2}\) and \(G_X = 10^5 G_F\). 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) pseudo-scalar \((M_X \ll T)\), with a Lagrangian \(L \sim g_s X \bar{\nu} \gamma^5 \nu\), 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 non-relativistic. Therefore, the mechanism of \(\nu_s\) 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 \(\bar{\nu}_s \nu_s \rightarrow XX\) process (with a rate \(\Gamma_X \sim g_s^2 T\)). 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 \(\nu_s\) 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 \(N_{\text{eff}}\)

The initial snapshot of active and sterile neutrino distribution soon after sterilites are excited via oscillation is a shared grey body distribution, a Fermi-Dirac function weighted by a factor 3/4 for each species. In 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 $G_N^2$. They are therefore, collisional. In fact, a grey 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 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

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

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

(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 $e^+ - e^-$ annihilation phase, active neutrinos energy density is given by

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

where the photon energy density is $\epsilon_\gamma = (\pi^2/15)T_\gamma^4$ and the effective number of neutrino species is $N_{\text{eff}} \simeq 3$, neglecting the small effect due to partial neutrino heating of order $\Delta N_{\text{eff}} = 0.046$ [33]. 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

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

and correspondingly the value of $N_{\text{eff}}$ decreases to

$$N_{\text{eff}} \sim 4 \times \left(\frac{3}{4}\right)^{4/3} \sim 2.7 \; .$$

(14)

This value is only slightly reduced at the matter radiation equality, i.e. for $T_\nu \sim 0.7$ 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 $m_{\text{st}} \sim \sqrt{\Delta m_{\text{st}}} \simeq 1$ eV

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

(15)

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

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

$$P_0 = 2 \int \frac{d^3p}{(2\pi)^3} \frac{p}{3\sqrt{p^2 + m_{\text{st}}^2}} \frac{1}{\exp[p/(T_\nu(3/4)^{1/3})] + 1} \; .$$

(16)

(17)

A value of the effective number of neutrino smaller than the expected standard result would be a signature in favour 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 [31]). Presently, the most precise determination of $N_{\text{eff}}$ is from Planck experiment. In the standard $\Lambda$CDM model but allowing for a free number of relativistic species the result is $N_{\text{eff}} = 3.30 \pm 0.27$ (68 % C.L.) [33], which is compatible with [15] at about 2σ. A new data release by Planck collaboration is expected soon, including polarization data. This might put further constraint on $N_{\text{eff}}$. In the future even more tight bounds are foreseen to come by next generation experiment such as Euclid [32] which should reach a sensitivity of order $\Delta N_{\text{eff}} < 0.1$.

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 (LSS), which are both sensitive to neutrino mass scale in the $10^{-1}$ eV – 1 eV range. One of the main effect of a massive neutrino is due their free streaming till the epoch when they become non relativistic, which suppresses the growth of perturbations on small scales. However, if sterile states scatters via secret interactions, the free streaming regime is delayed until the scattering rate becomes smaller than the Hubble parameter. This means that if $G_N$ is large enough so that this condition holds at the non relativistic transition, sterile neutrinos would never have a free streaming phase, but always diffuse. 3

3 We are very pleased to thank Basudeb Dasgupta for pointing us this possibility.
The smaller value of $G_X$ for which this happens can be obtained from the condition that scattering rate equals the value of $H$ at a temperature $3.15 T_\nu \sim (p) \sim \sqrt{\Delta m^2_{\text{st}}}$

$$G^2 X T^5_\nu \sim H(T_\gamma).$$

(18)

Using standard expression of the Hubble rate in the $T_\gamma \sim$ eV range and the fact that, as we have seen, sterile temperature is given by $T_\gamma = (4/11)^{1/3} (3/4)^{1/3} T_0$ 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 $[34]\$

$$m^\text{eff}_{\text{st}} = \rho_{ss} \sqrt{\Delta m^2_{\text{st}}} = \frac{3}{4} \sqrt{\Delta m^2_{\text{st}}}. \quad (19)$$

Latest analysis in $[24]$ of the sterile neutrino anomalies gives a best-fit $\Delta m^2_{\text{st}} = 1.6$ eV$^2$ with a $2\sigma$ range

$$1.08 \text{ eV}^2 < \Delta m^2_{\text{st}} < 1.99 \text{ eV}^2. \quad (20)$$

Using Eq. (19) the lower value in the $2\sigma$ range gives $m^\text{eff}_{\text{st}} \simeq 0.78$ 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 $[22]$ finds a discrepancy between appearance and disappearance sterile neutrino data. As a consequence only a small region around $\Delta m^2_{\text{st}} \simeq 0.9$ eV$^2$ would be compatible with all data. This would correspond to $m^\text{eff}_{\text{st}} \simeq 0.7$ 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 $[35]$ and the recent but controversial BICEP2 experiment $[36]$ (see also $[37]$), LSS, $H_0$ 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 $m^\text{eff}_{\text{st}} < 0.42$ eV at 95% C.L. $[35]$, which is in strong disagreement with the sterile neutrino abundance produced by the secret interactions. Moreover a possible non-zero 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 $m^\text{eff}_{\text{st}} \simeq 0.7$ eV at $2\sigma$ $[35, 41]$ (see also $[12]$) or an upper bound of $m^\text{eff}_{\text{st}} < 0.6$ eV $[42]$. Stronger bounds have also been quoted in $[41, 43]$.

In general, from the results presented here, one would conclude that the minimum $m^\text{eff}_{\text{st}} \simeq 0.78$ eV obtained from the secret collisional production and compatible with the $\Delta m^2_{\text{st}}$ range would be in tension (at least at $2\sigma$ 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 non relativistic. 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 $\nu_s$ generated at $T \gg \text{GeV}$ 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 re-thermalization 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 $N_{\text{eff}}$ down to 2.7. If this value is compatible with the $2\sigma$ range given by the Planck experiment $[35]$, the future experiment Euclid $[32]$, with a sensitivity to $\Delta N_{\text{eff}} < 0.1$ 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 $M_X \approx g_X$ 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 neu-
trino 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 been speculated that the background of sterile neutrinos produced by the collisions associated with the secret interactions, would also modify the optical depth of ultra-high-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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