

# Sterile neutrinos and low reheating temperature

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**Abstract.** Cosmological and astrophysical constraints on the mixings of sterile neutrinos are commonly much more stringent than the ones coming from laboratory searches. Within the standard cosmological picture, this mixing needs to be very small to prevent the cosmological overabundance of sterile neutrinos. We present here a scenario, based on a low reheating temperature  $T_{RH} \ll 100$  MeV at the end of (the last episode of) inflation or entropy creation, and show that the production of sterile neutrinos becomes largely suppressed with respect to what happens in the standard framework. In this case, the dominant constraints come from laboratory measurements, which may render sterile neutrinos to be “visible” in future experiments.

## 1. Introduction

In inflationary models, the decay of the coherent oscillations of a scalar field represents the beginning of the radiation dominated era of the Universe. The subsequent thermalization of the decay products results in a thermal bath with the so called “reheating temperature”,  $T_{RH}$ . This temperature may have been as low as  $\sim 4$  MeV [1] once Cosmic Microwave Background (CMB) and Large Scale Structure (LSS) observations are taken into account, while Big Bang Nucleosynthesis (BBN)-only data impose a lower bound of  $\sim 2$  MeV on  $T_{RH}$  [2] if active neutrino oscillations are taken into account ( $\sim 0.7$  MeV [3] if they are not). Nonetheless, the standard computation of the relic densities relies on the assumptions that  $T_{RH}$  was large enough for the particles of interest to have reached thermal equilibrium and that the entropy of matter and radiation is conserved after they decouple. However, there are non-standard cosmological models in which these assumptions about the epoch of the Universe before BBN, an epoch from which we have no data, do not hold [4–6].

It is well known that the production rate of sterile neutrinos reaches its peak at a temperature  $T_{\max} \simeq 133 (m_s/\text{keV})^{1/3}$  MeV [7,8]. Hence, if the reheating temperature is smaller,  $T_{RH} < T_{\max}$ , the production of sterile neutrinos would be suppressed. In the case of models where there is a late episode of entropy production in which the Universe is reheated to a very low  $T_{RH}$ , the production of particles which are non-relativistic or decoupled at  $T \lesssim T_{RH}$  is inhibited [9–12]. The final number density of active neutrinos starts departing from the standard number for  $T_{RH} \lesssim 8$  MeV but stays within 10% of it for  $T_{RH} \gtrsim 5$  MeV [10,13].

In general, low  $T_{RH}$  cosmological scenarios are more complicated than the standard ones. Although different aspects of these models have been studied with interesting results, no consistent all-encompassing scenario exists yet. For instance, baryogenesis could be produced through the Affleck-Dine mechanism, in a model similar to that of Ref. [14] and dark matter could still consist of the lightest supersymmetric particle or other Weakly Interactive Massive

Particles (WIMPs), produced either thermally or non-thermally [15]. On the other hand, MeV scalars have been proposed as DM candidates [16, 17]: they would be in thermal equilibrium in the Early Universe with neutrinos, electrons, positrons and photons and would freeze-out at 100's keV– MeV temperatures.

If sterile neutrinos exist and have no extra Standard Model interactions, the dominant mechanism of production in the early Universe is through their mixing with active neutrinos [7]. Dodelson and Widrow [8] (see also Refs. [18, 19]) provided the first analytical calculation of the production of sterile neutrinos lighter than 100 keV in the early Universe, under the assumption (which we maintain here) of a negligible primordial lepton asymmetry.

Here we assume that the production of sterile neutrinos, through the conversion of active neutrinos produced in collisions, starts when the temperature of the universe is  $T_{\text{RH}} (< T_{\text{max}})$  and the active neutrinos are assumed to have the usual thermal equilibrium distribution. Thus, following Ref. [1], we restrict ourselves to reheating temperatures  $T_{\text{RH}} > 4$  MeV. This approach was shown [13] to give results which are correct within an order of magnitude [11]. In addition, by using different approximations, we are able to obtain an analytical result for the sterile neutrino abundance. In this way, we are able to write all our results in a simplified form which allows us to have a qualitative understanding of the problem.

In this talk we review the results obtained for two different mass regions for the heavy (mostly sterile) neutrinos, below [11] and above [12] 1 MeV, and show that the primordial abundance of sterile neutrinos does not necessarily impose their mixing with active neutrinos to be unaccessibly small. A very low reheating temperature scenario would suppress the sterile neutrino production, weakening the cosmological bounds and allowing the mixing with active ones to be as large as laboratory measurements constrain. These neutrinos could, therefore, be revealed in future experiments.

## 2. Sterile neutrino abundance

For simplicity, our analysis is based on the two-neutrino mixing approximation. Within this approximation, the vacuum mixing angle  $\sin\theta$  represents the amplitude of the heavy mass eigenstate, which for small  $\sin\theta$  is mostly sterile and whose mass we call  $m_s$ .

In order to obtain the distribution function of sterile neutrinos at a given temperature after the last episode of inflation, we start from the Boltzmann equation with the approximation that  $T \propto 1/a$ , with  $a$  the scale factor,

$$-HT \left( \frac{\partial f_s}{\partial T} \right)_{p/T} \simeq \Gamma_s(T) (f_s^{\text{eq}} - f_s), \quad (1)$$

where  $f_s^{\text{eq}} = (\exp(E/T) + 1)^{-1}$  is the Fermi-Dirac distribution that heavy neutrinos would have if they were in thermal equilibrium.

In this talk, we will consider two different mass regimes. For  $m_s < 1$  MeV, neutrinos cannot decay into  $e^+e^-$  pairs, whereas for  $m_s > 1$  MeV, this channel is open. In addition, above a few MeV, neutrinos are not fully relativistic at all times during their evolution before decoupling.

### 2.1. $m_s < 1$ MeV

For neutrino masses much smaller than the temperature of the plasma ( $m_s < 1$  MeV),  $\langle p \rangle \simeq \langle E \rangle$ , the averaged rate of sterile neutrino interactions is given by

$$\Gamma_s(T) \simeq \frac{1}{4} \sin^2 2\theta_m d_\alpha G_F^2 E T^4, \quad (2)$$

where  $\theta_m$  is the mixing angle in matter,  $G_F$  is the Fermi constant and  $d_\alpha = 1.13$  for sterile neutrino mixing with  $\nu_\alpha = \nu_e$  and  $d_\alpha = 0.79$  with  $\nu_\alpha = \nu_{\mu,\tau}$ . For  $T < 1.5$  GeV ( $m_s/\text{MeV})^{1/3}$

matter effects are negligible [20, 21] and hence  $\sin^2 2\theta_m \simeq \sin^2 2\theta$  is a very good approximation. Plugging Eq.(2) into Eq.(1), and solving for  $f_s$  in the limit  $f_s \ll f_s^{\text{eq}}$  (see Ref. [12] for the case when this approximation fails), the obtained distribution function results in a number density of massive (mostly sterile) neutrinos ( $m_s < 1$  MeV) given by [11]

$$n_s \simeq 10 d_\alpha \sin^2 2\theta \left( \frac{T_{\text{RH}}}{5 \text{ MeV}} \right)^3 n_\alpha, \quad (3)$$

where  $n_\alpha = 0.09 g T^3$  is the number density of a relativistic fermion with  $g$  degrees of freedom in thermal equilibrium. Notice that the number density of sterile neutrinos depends on both the active-sterile mixing angle and the reheating temperature. A low reheating temperature implies a small sterile neutrino number density. Note that the  $\nu_s$ -number density is independent of the mass of the sterile neutrinos, contrary to the standard result [8, 18, 19]. Thus, the mass density of non-relativistic sterile neutrinos depends linearly on the mass and on  $\sin^2 2\theta$ .

## 2.2. $m_s > 1$ MeV

For heavier sterile neutrinos, with  $m_s > 1$  MeV, and for the range of temperatures explored here, the heavy neutrino mass needs to be taken into account and the averaged production rate of sterile neutrinos  $\Gamma_s$  is given by [20, 21]

$$\Gamma_s(T) = \frac{1}{\tau_s} \left[ \frac{m_s}{E} + \frac{3 \times 2^7 T^3}{m_s^3} \left\{ \frac{3 \zeta(3)}{4} + \frac{7 \pi^4}{144} \left( \frac{ET}{m_s^2} + \frac{p^2 T}{3 E m_s^2} \right) \right\} \right], \quad (4)$$

where the first term is due to inverse decay and the other terms correspond to two-to-two particle processes. The function  $\tau_s$  in the denominator is the heavy neutrino lifetime. As in the case of  $m_s < 1$  MeV, Eq.(4) is valid when matter effects are not important.

For  $m_s < m_\pi \sim 140$  MeV, the heavy (mostly sterile) neutrino can decay into a light neutrino and two leptons  $\nu_s \rightarrow \nu_\alpha + l + \bar{l}$ , mainly  $\nu_\alpha \bar{\nu} \nu$  and  $\nu_\alpha e^+ e^-$ . If the active neutrino mixing with the sterile is  $\nu_\tau$  or  $\nu_\mu$ , the decay happens through neutral currents and the lifetime is  $\tau_s = \frac{1.0 \text{ sec}}{\sin^2 2\theta} \left( \frac{10 \text{ MeV}}{m_s} \right)^5$ . If instead  $\nu_s$  mixes mostly with  $\nu_e$ , the factor 1.0 sec should be replaced by 0.7 sec [20, 21], due to the presence of charged currents. However, we are not going to keep this distinction in the following. For larger masses, other decay modes open up, so for simplicity, in the following we will restrict ourselves to the range  $1 \text{ MeV} < m_s < 140 \text{ MeV}$ .

We solve analytically Eq.(1) for  $T_{\text{RH}} \leq m_s$ . In order to do so, we make several approximations. First, we assume that the actual distribution function of the heavy neutrinos is always much smaller than the equilibrium distribution,  $f_s \ll f_s^{\text{eq}}$ . Then, we approximate the Fermi-Dirac distribution by a Boltzmann distribution,  $f_s^{\text{eq}} \simeq e^{-E/T}$  and take neutrinos to be either purely non-relativistic, i.e.,  $E = m_s$  if  $p < m_s$ , or purely relativistic, i.e.,  $E = p$  if  $p > m_s$ . We integrate analytically Eq.(1) with  $p/T$  constant over temperatures  $T$  in the interval  $0 \leq T \leq T_{\text{RH}}$ . This procedure overestimates the final abundance, thus providing an upper bound on the actual abundance. We find that the number density  $n_s$  of heavy (mostly sterile) neutrinos for  $m_s < 140$  MeV is given by [12]

$$\begin{aligned} n_s(x_{\text{RH}}, m_s, T) &\simeq n_\alpha(T) \sin^2 2\theta \left( \frac{m_s}{\text{MeV}} \right)^3 2.1 \times 10^{-3} e^{-x_{\text{RH}}} \\ &\times \left[ \left( \frac{7}{3} + \frac{6 + 144 \zeta(3)}{\pi^4} \right) + \frac{2^3 \times 7}{3} \left\{ 1 + x_{\text{RH}} + \left( \frac{3}{2^3} + \frac{3^4 \zeta(3)}{7 \pi^4} \right) x_{\text{RH}}^2 \right\} \frac{1}{x_{\text{RH}}^3} \right. \\ &\left. + \left( 24 + 144 \zeta(3) + 12 x_{\text{RH}} + \frac{7}{2} x_{\text{RH}}^2 + \frac{1}{2} x_{\text{RH}}^3 \right) \frac{x_{\text{RH}}}{4 \pi^4} \right], \quad (5) \end{aligned}$$

where we have defined  $x_{\text{RH}} \equiv m_s/T_{\text{RH}}$ . Taking into account the subsequent decay of sterile neutrinos, the actual number density is  $n_s(T) \simeq n_s(x_{\text{RH}}, T) e^{-t/\tau_s}$ . In order to obtain analytical results for the cosmologically and astrophysically allowed regions in the  $(\sin^2 2\theta, m_s)$ -plane, we will use the instant-decay approximation for  $\tau_s$  smaller than the age of the Universe.

### 3. Laboratory Bounds

In laboratory searches, no positive evidence of heavy (mostly sterile) neutrinos has been found so far. Here, we review the most stringent bounds on the mixing angle with active neutrinos (for further details see a comprehensive discussion in Refs. [23,24]). They are shown in Fig. 1 as blue solid areas and in Figs. 2 and 3 as blue solid areas for  $\nu_e - \nu_s$  mixing and as red (cyan) dashed lines for  $\nu_\mu - \nu_s$  ( $\nu_\tau - \nu_s$ ) mixing.

Let us consider first sterile neutrinos mixing with  $\nu_e$ . For masses up to  $m_s \simeq 10$  MeV, an important bound is provided by searches of kinks in the electron spectrum of  $\beta$ -decays. For higher masses, very robust bounds can be set by looking for additional peaks in the spectrum of electrons in leptonic decays of pions and kaons.

In neutrino-oscillation, fixed-target and collider experiments, if sterile neutrinos mix with active ones, a beam of  $\nu_s$  would be produced and would subsequently decay into visible particles. Assuming only charged current and neutral current interactions for  $\nu_s$ , the absence of  $\nu_s$ -decay signatures in past and present experiments allows one to put limits on the mixing term which controls the intensity of the  $\nu_s$  beam and the decay time. For masses below 1 MeV, reactor- $\bar{\nu}_e$  experiments CHOOZ [25] and Bugey [26] set the most stringent constraints on the mixing angle. For higher masses (up to 140 MeV), a reanalysis of the Borexino Counting Test Facility, and Bugey and PS191 data yield the most constraining limits for this type of experiments [27,28].

Finally, if sterile neutrinos are Majorana particles, they would contribute to the mediation of neutrinoless double-beta decay. The limit on the half-life time of this process can be translated into a bound on the mixing with  $\nu_e$ ,  $\sin^2 \theta$ , which scales as  $m_s$  for  $m_s \lesssim 30$  MeV and as  $m_s^{-1}$  for  $m_s \gtrsim 400$  MeV, with  $\sin^2 2\theta < 5 \times 10^{-8}$  at  $m_s \sim 100$  MeV. Possible cancellations among contributions due to different mass eigenstates would weaken the inferred bound. Barring this possibility, this limit would exclude the region to the right of the dashed blue line shown in the left panel of Fig. 1 and the hatched blue area in Figs. 2 and 3.

Let us now consider the case of mixing with  $\nu_\mu$ . As for the case of sterile neutrinos mixing with  $\nu_e$ , peak searches provide very robust and stringent bounds on sterile neutrinos (for a detailed review see Fig. 1 of Ref. [23]). Motivated by the KARMEN anomaly, which could be explained by the existence of a heavy (mostly sterile) neutrino with mass 33.9 MeV, very sensitive searches have been performed for this neutrino mass [29]. Similarly, sterile neutrinos with heavier masses can be probed in kaon decays [30].

By searching for the production of  $\nu_s$  in pion and kaon decays and their subsequent decay, a reanalysis [23] of fixed-target data led to a stringent mass-dependent bound on  $\sin^2 2\theta$ . Similar bounds can be set on the product of the mixing angle with  $\nu_\mu$  and with  $\nu_e$  from the analysis of the same decay-searches data [23].

As in the case of  $\nu_e - \nu_s$  mixing, accelerator- $\nu_\mu$  disappearance experiments [31] set also bounds for the  $\nu_\mu - \nu_s$  case, in the low mass range considered here.

Let us consider finally the case of heavy (mostly sterile) neutrinos mixed with  $\nu_\tau$ . The only limits on these sterile neutrinos come from searches for  $\nu_s$  decays. The most stringent bound is obtained by the reanalysis of data from the CHARM experiment [32] (see also the analysis of data from the CLEO experiment [33]), in which  $\nu_s$  could be produced in  $D$  and  $\tau$  decays. However, in the mass range we are considering, this upper bound is rather weak.

#### 4. Cosmological and Astrophysical Bounds

In what follows we obtain six different types of cosmological and astrophysical bounds on the parameter space of active-sterile neutrino mixing,  $(\sin^2 2\theta, m_s)$ . Each of these limits is valid for a certain range of values of the heavy neutrino lifetime and mass. Whereas the four of them have to do with the photons which are produced in the decay, the other two have to do with a bound on the sterile neutrino abundance at a particular epoch. Finally, we will also comment on bounds from supernova (SN) observations.

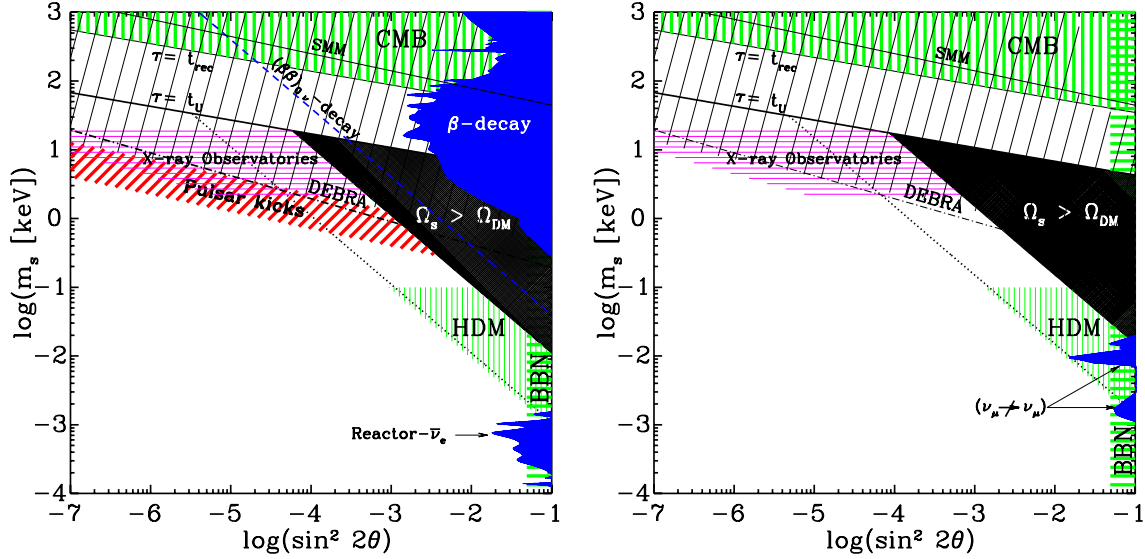
In the first place, we obtain approximate bounds by using LSS arguments. Secondly, we consider the Diffuse Extragalactic Background RAdiation (DEBRA) spectrum and set bounds on the basis of not finding any unexpected result from heavy (mostly sterile) neutrino decay. Then, we obtain bounds using the non-observation of a distortion of the CMB spectrum caused by the radiation from decay. Then, we calculate the limits based on the data from the primordial light element abundances and how neutrino decay would affect them. We also use the upper bound from BBN on the extra number of relativistic degrees of freedom at that epoch to set a limit on active-sterile neutrino mixing. Finally, we comment on the limits from the non-observation of photon lines in the different X-ray observatories. In our analysis, for  $\tau_s < t_U$  with  $t_U \simeq 14$  Gyr the present age of the Universe, we use the instant-decay approximation, i.e., all decays happen at  $t = \tau_s$ . This will allow us to obtain all the results analytically, while giving rise to accurate enough calculations for our purposes. We refer the reader to Refs. [11] and [12] for details.

##### 4.1. $m_s < 1$ MeV

The results for  $m_s < 1$  MeV and for  $T_{RH} = 5$  MeV are depicted in Fig. 1 [11]. For this case, the dominant decay mode of the mostly-sterile  $\nu$  is into three neutrinos. Assuming neutrinos are Majorana particles, a lifetime  $\tau_s$  equal to the age of the universe is indicated in Fig. 1 with the full thick line (for Dirac neutrinos,  $\tau_s$  is larger by a factor of 2). Equivalent lines corresponding to shorter or longer lifetimes can be easily obtained bearing in mind that the lifetime is proportional to  $(\sin^2 2\theta m_s^5)^{-1}$ . In the mass range in which sterile neutrinos can be part of the hot dark matter ( $\tau_s > t_U$ ), we can apply the bounds on the sum of the contributions of active and sterile neutrinos to the dark matter density. Combining this limit with an estimate of the light neutrino masses we obtain an upper limit on  $m_s$ . If the neutrino mass spectrum is normal hierarchical, oscillation data impose the sum of the active neutrino masses to be about 0.05 eV. This provides the most conservative bound on  $m_s$ , which we plot for  $m_s \lesssim 100$  eV (vertically hatched region with thin lines labeled as “HDM” in Fig. 1).

The decay mode into a neutrino and a photon happens with a branching ratio  $0.8 \times 10^{-2}$  [34]. Photons decouple from the plasma filling the Universe at the recombination epoch,  $t_{rec} \simeq 1.3 \times 10^{13}$  sec (the line  $\tau_s = t_{rec}$  is also shown in Fig. 1). Thus, if the sterile neutrino decays happen after recombination,  $\tau_s > t_{rec}$ , the photons produced do not interact ever after and could leave an imprint in the DEBRA spectrum [35]. Such a signature has not been observed, and thus the photon flux must not be larger than the observed DEBRA energy flux. Hence, this imposes a bound (see Ref. [36] for a more detailed analysis). For decays with  $\tau_s > t_{rec}$ , the bound obtained for unclustered neutrinos would reject the region above the dot-dashed line labeled as “DEBRA” in Fig. 1.

On the other hand, the CMB radiation is emitted at recombination. Electromagnetic decay products produced sometime before recombination may distort the CMB spectrum [37,38]. Non-thermal photons produced before the thermalization time  $t_{th} \simeq 10^6$  sec are rapidly incorporated into the Planck spectrum. This happens through processes that change the number of photons, such as double Compton scattering ( $\gamma e \rightarrow \gamma \gamma e$ ). If non-thermal photons are produced after  $t_{th}$ , i.e., if  $t_{th} < \tau_s < t_{rec}$ , the CMB Planck spectrum would be distorted. Current data pose very stringent upper bounds on possible distortions of this spectrum. For the earliest part of this time



**Figure 1.** *Left panel:* Bounds for  $\nu_e \leftrightarrow \nu_s$  mixing in the  $(\sin^2 2\theta, m_s)$ -plane for  $T_{RH} = 5$  MeV. See text. *Right panel:* Same as left panel but for  $\nu_{\mu,\tau} \leftrightarrow \nu_s$ . For  $\nu_\tau \leftrightarrow \nu_s$  the darkest gray-blue excluded region does not apply.

interval, i.e., for  $t_{th} \simeq 10^6$  sec  $< \tau_s < 10^9$  sec, photon number preserving processes, like elastic Compton scattering, are still efficient. These processes thermalize the photons not into a Plank spectrum but into a Bose-Einstein spectrum with a non-zero chemical potential  $\mu$ . For later decays,  $10^9$  sec  $< \tau_s < t_{rec}$ , the photon number preserving processes can no longer establish a Bose-Einstein spectrum. The energy released in this case is not thermalized but simply heats the electrons. For both epochs, the COBE satellite has provided stringent limits on any fractional increase in the photon energy density [39], as for instance that due to the decay of the sterile neutrinos. Hence, they can be used to exclude all the vertically hatched region with thick lines labeled as “CMB” in Fig. 1.

On the other hand, BBN data provides an upper bound on any source of extra energy density present in the Universe during the BBN epoch,  $t_{startBBN} \simeq 0.1$  sec  $< \tau_s < t_{endBBN}$ , as well as on extra radiation present during that period. The bounds are complicated in detail, but it is safe to say that if the extra energy density due to the presence of sterile neutrinos is very small, BBN will not be affected in any way. The bounds on extra contributions to the energy density during BBN are customarily presented in terms of the equivalent extra number of relativistic active neutrino species,  $\Delta N_\nu$ . Thus, to be on the safe side, we simply require the  $\Delta N_\nu$  due to the presence of sterile neutrinos to be smaller than 1 during BBN. This excludes the green horizontally hatched region with thick lines labeled as “BBN” in Fig. 1.

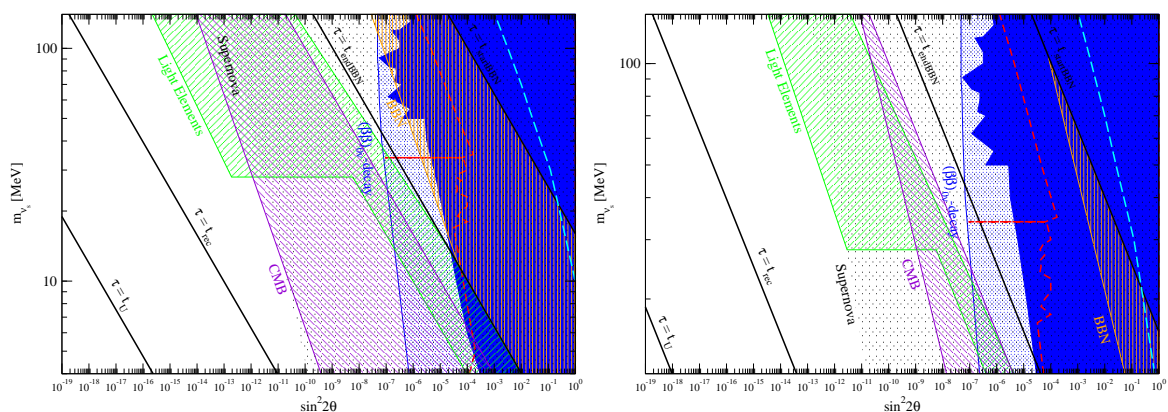
Some years ago, Ref. [40] proposed to observe clusters of galaxies with the Chandra and XMM-Newton observatories, in their high sensitivity range for X-ray photon detection of 1–10 keV. During the last years, a lot of work has been devoted to set limits by searching for  $\gamma$ -ray lines using different astrophysical observations [41]. This has allowed to exclude the horizontally hatched region with thin lines labeled as “X-ray Observatories” in Fig. 1. Here the density fraction of sterile neutrinos within clusters is assumed to coincide with the cosmological energy fraction ( $\Omega_s/\Omega_{DM}$ ).

Finally, through  $\nu_\alpha \leftrightarrow \nu_s$  oscillations, sterile neutrinos can be produced in SN cores and escape, carrying away a large amount of the released energy. The observations of  $\nu_e, \bar{\nu}_e$  from

SN1987A constrain the energy loss in  $\nu_s$  and yield a bound on the mixing angle. For  $m_s \lesssim 45$  keV,  $\nu_\alpha \leftrightarrow \nu_s$  oscillations are matter suppressed, whereas for  $m_s \gtrsim 45$  keV, the matter effects are negligible and neutrinos oscillate as in vacuum. These bounds exclude the diagonally hatched region with thin lines in Fig. 1. On the other hand, for small mixing angles sterile neutrinos are not trapped within the SN, thus they are emitted from the whole collapsing star, mostly from its core. The decay of these heavy neutrinos emitted in SN explosions would produce a flux of photons. The non-observation by the Solar Maximum Mission (SMM) of any  $\gamma$ -ray counts in excess of the background for a time interval of  $t_{\max} = 223.2$  sec after the arrival of the first  $\bar{\nu}_e$ 's from SN1987A [42], allows us to enlarge the disfavored regions of the parameter space (region above the line labeled as ‘‘SMM’’ in Fig. 1). Let us mention, however, that so much is not understood about the neutrino transport and flavor transformation in hot and dense nuclear matter, that conservatively the implicated  $m_s - \sin^2 2\theta$  region can only be considered disfavored but not excluded.

If the sterile neutrinos produced in non-resonant  $\nu_e \rightarrow \nu_s$  conversion, in fact, carry away a sizable fraction of the energy emitted in a SN explosion, asymmetric emission of  $\nu_s$  due to the presence of a strong magnetic field, could explain the very large velocities of pulsars [43] (see diagonally hatched region with thick lines labeled as ‘‘Pulsar kicks’’ in the left panel of Fig. 1).

#### 4.2. $m_s > 1$ MeV



**Figure 2.** *Left panel:* Bounds in the  $(\sin^2 2\theta, m_s)$ -plane of heavy (mostly sterile) neutrinos mixed with active ones for  $T_{\text{RH}} = m_s$ . The dark gray (blue) solid area and dark gray (blue) hatched area (for Majorana neutrinos) represent the experimentally excluded region for  $\nu_e - \nu_s$  mixing, the (red) short-dashed line and the (cyan) long-dashed line are the experimental upper bounds for  $\nu_\mu - \nu_s$  and  $\nu_\tau - \nu_s$  mixing, respectively. The cosmological and astrophysical bounds are indicated by the corresponding labels. Isolines for values of the heavy neutrino lifetime  $\tau_s$  equal to the start and end of the BBN epoch, the recombination time and the present age of the Universe are also shown. See text for details. *Right panel:* Same as left panel but for  $T_{\text{RH}} = m_s/3$ .

The results for sterile neutrino masses in the range of  $1 \text{ MeV} < m_s < 140 \text{ MeV}$  are shown in Figs. 2 and 3 [12]. For these masses, the decay branching ratio of the heavy (mostly sterile) neutrino into  $e^+e^-$  is about 10% (40%) for mixing with  $\nu_\mu$  or  $\nu_\tau$  ( $\nu_e$ ). Due to inverse Compton scattering on the CMB with a Thompson cross section, the interaction length of these electrons is about 1 kpc (see for example [44] or Fig. 5 of [45]) and the result of each of the interactions is a photon that shares a large portion of the incoming electron energy. Photons at these energies propagate for cosmological distances undisturbed by the CMB or infrared backgrounds and are



neutrinos are more restrictive than for Dirac neutrinos (see the hatched blue area labeled as “ $(\beta\beta)_{0\nu}$ -decay”). Values of the heavy neutrino lifetime  $\tau_s$  equal to the relevant epochs in the history of the Universe are also shown.

In Figs. 2 and the left panel of Fig. 3, we display the results for three different values of the ratio  $x_{\text{RH}} \equiv m_s/T_{\text{RH}}$ ,  $x_{\text{RH}} = 1, 3, 10$ , respectively. As expected, for  $T_{\text{RH}}$  increasingly smaller than the heavy (mostly sterile) neutrino mass, the cosmological bounds become less restrictive, and when  $T_{\text{RH}} \leq m_s/10$ , the cosmological bounds become completely irrelevant and only experimental data are able to restrict the parameter space. This result can also be seen in the right panel of Fig. 3, in which the reheating temperature is fixed to be  $T_{\text{RH}} = 5$  MeV. In this case all cosmological bounds become irrelevant for  $m_s \geq 30$  MeV.

## 5. Conclusions

It is commonly assumed that the cosmological bounds on the mixings of sterile and active neutrinos restrict the range of their allowed values much more than laboratory data. In fact, sterile neutrinos with parameters suitable to be found in the near future in different experiments would have mixings with active neutrinos too large to be allowed by the standard cosmological assumptions about the pre-BBN era in the Universe, an era about which we do not have any observational information. The standard assumptions are few but very powerful: it is usually assumed that the temperature reached in the radiation dominated epoch before BBN was very high, that the Universe was radiation dominated then and that the entropy of radiation and matter is conserved.

Here, we show that it is possible to evade most of the cosmological bounds by assuming that the temperature at the end of (the last episode of) inflation or entropy production, the so-called reheating temperature  $T_{\text{RH}}$ , is low enough. We concentrate on heavy (mostly sterile) neutrinos with masses below 140 MeV [11, 12]. For low  $T_{\text{RH}}$ , the production of sterile neutrinos is suppressed. We present the experimental bounds on sterile neutrino mixings, the disfavored regions from core-collapse SN data and the cosmological bounds imposed by LSS, DEBRA, the  $\gamma$ -line search, CMB, BBN and the abundance of light elements. We show our results for two different mass regimes, for  $m_s < 1$  MeV (in Fig. 1) [11] and for  $1 \text{ MeV} < m_s < 140$  MeV (in Figs. 2 and 3) [12]. We show that, unlike in standard cosmology, in low reheating temperature cosmologies it is possible to accommodate “visible” sterile neutrinos, i.e., sterile neutrinos which could soon be found in experiments. Hence, finding a particle, such as a “visible” sterile neutrino, whose existence would contradict the usual assumptions about the pre-BBN era and SN physics, would give us not only invaluable information for particle physics and astrophysics, but also an indication of enormous relevance in cosmology: it would tell us that the usual assumptions must be modified, as for instance in the way presented here.

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