Sterile neutrinos and their role in particle physics and cosmology

2014 Phys.-Usp. 57 503

(http://iopscience.iop.org/1063-7869/57/5/503)

Physics – Uspekhi **57** (5) 503 – 511 (2014)

CONFERENCES AND SYMPOSIA

PACS numbers: **12.60.–i**, 14.60.Pq, 95.30.Cq

100th ANNIVERSARY OF THE BIRTH OF B M PONTECORVO

Sterile neutrinos and their role in particle physics and cosmology

D S Gorbunov

DOI: 10.3367/UFNe.0184.201405i.0545

Contents

Abstract. Introducing new fermions into particles physics that are singlet under the gauge group of the Standard Model allows using the seesaw type-I mechanism to obtain small neutrino masses and explaining the phenomenon of neutrino oscillations. The mass scale of these fermions (sterile neutrinos) is a free parameter with values ranging widely from a subelectronvolt to 10^{15} GeV. The mass scale determines the direct search strategy for these particles, and depending on it, sterile neutrinos can influence the evolution of the Universe and can be responsible for the baryon asymmetry of the Universe and for the phenomenon of dark matter.

1. Introduction

Models involving sterile neutrinos constitute one of the most popular extensions of the Standard Model (SM) of elementary particle physics (see, e.g., Refs [1, 2]). These are new fermion degrees of freedom that have no charge under the SM gauge group $SU(3)_c \times SU(2)_W \times U(1)_Y$; hence the adjective *sterile*. These fermions are called neutrinos because they mix with the SM neutrinos (or *active* neutrinos), which provides the latter with mass and mixing between neutrinos of different sorts, or flavors (electron, muon, and tau neutrinos), resulting in neutrino oscillations.

The popularity of this line of research is due to the following reasons. It is a renormalizable generalization of the SM: the new constants are dimensionless and the theory, like the SM itself, can be valid at scales up to the gravitational one. This is a very ‘economical’ SM generalization: to explain the experimental results of neutrino oscillation observations, it suffices to introduce only two new fermions of the *Majorana type*, which adds four degrees of freedom to the SM. The values of model parameters can then be chosen so as to explain not only oscillations but also the baryon asymmetry of the Universe via the leptogenesis mechanism in the primordial plasma (mass terms of the Majorana type lead to lepton number violation). In such a generalization, only two of the three active neutrinos have mass, and to provide mass for all three SM neutrinos, three sterile neutrinos are also required. In this case, it turns out that for a certain region of the space of model parameters, the lightest of the three sterile neutrinos can be so long-lived that it can claim to be a dark matter particle.

Thus, introducing only three Majorana singlet fermions into the SM is also sufficient to resolve the three most important phenomenological SM problems (neutrino oscillations, dark matter, and the baryon asymmetry of the Universe) in the framework of a unique approach [3]. Finally, several anomalous results of oscillation experiments (the LSND (Liquid Scintillator Neutrino Detector) [4, 5], MiniBooNE (Mini Booster Neutrino Experiment) [6–9], gallium [10, 11], and reactor [12, 13] anomalies) can be explained in a natural manner as oscillations to sterile neutrinos with masses in the 1 eV range.

In general, most of the space of model parameters is inaccessible for direct experimental tests in the foreseeable future, which renders the model somewhat scholastic. The key parameter here is the sterile neutrino mass scale. Several assertions, depending on the sterile neutrino mass values, can be made depending on the expected phenomenological manifestations of the model:

1. The sterile neutrino mass scale has not been determined: masses up to 10^{15} GeV are permitted. The interaction of SM fields with sterile neutrinos, providing an explanation for neutrino oscillations, gives quantum corrections to the Higgs boson mass that are proportional to the sterile neutrino mass scale. The formal requirement that the corrections be small implies a restriction from above on the mass scale of the order of 10^7 GeV. Quantum corrections

---

D S Gorbunov Institute for Nuclear Research,
Russian Academy of Sciences,
prosp. 60-letiya Oktyabrya 7A, 117312 Moscow, Russian Federation
E-mail: gorby@ms2.inr.ac.ru

Received 23 March 2014
Uspekhi Fizicheskikh Nauk **184** (5) 545 – 554 (2014)
DOI: 10.3367/UFNr.0184.201405i.0545
Translated by G Pontecorvo; edited by A M Semikhatov

naturally cancel in supersymmetric generalizations of particle physics models that are being tested at the Large Hadron Collider.

2. Two sterile neutrinos are sufficient to explain not only the results of oscillation experiments but also the baryon asymmetry of the Universe via leptogenesis. Here, 'standard leptogenesis' [14] (the production of lepton asymmetry in the decays of heavy sterile neutrinos) occurs if the neutrino mass exceeds 10^9 GeV. For models with sterile neutrinos degenerate in mass, it is also possible to produce lepton asymmetry in the primordial plasma for lighter neutrinos with masses up to the TeV range [15] and even in the range from several hundred MeV to several tens of GeV [16]. Here, resonance oscillatory processes between active and sterile neutrinos in plasma are essential. In the case of three neutrinos with masses smaller than 100 GeV, such an asymmetry production mechanism does not require degeneracy [17].

3. The sterile neutrino accounting for dark matter is not the main fermion of the 'seesaw' mechanism providing active neutrinos with masses. The contribution of such a neutrino to the mass matrix of active neutrinos is small, such that one of them remains nearly massless. The mass parameter determining the rate of neutrinoless double β-decay and widely applied as an indicator of lepton number violation is also small. We note that the range of allowed dark matter sterile neutrino masses includes the 1–10 keV range; neutrinos with masses in this range form so-called *warm* dark matter. This is interesting from the standpoint of formation in the Universe of structures on the scale of dwarf galaxies, whose number, according to observations, is smaller than the number of such objects predicted by models involving *cold* dark matter [for example, models with weakly interacting massive particles (WIMPs)].

4. Sterile neutrinos with masses of the order of 1 eV can only make a small contribution to the total mass density of dark matter, at the same level as massive active neutrinos. In the early Universe, sterile neutrinos were part of the radiation plasma component (relativistic particles). A straightforward test of such models is only possible if the mixing with active neutrinos is not small. When the mixing is small, only cosmology can say whether such neutrinos exist. Accounting for uncertainties related to the absence of any understanding of the nature of dark matter (of the physics responsible for expansion of the Universe, involving acceleration), modern cosmological data allow the existence of a sterile neutrino with a mass of the order of several tenths of an electronvolt.

2. Mechanism leading to active neutrinos having mass

With the set of SM fields, massive neutrinos can be described by adding the following nonrenormalizable interaction to the Lagrangian:

$$\mathcal{L}^{(5)} = \frac{\beta_L}{4\Lambda}\, F_{\alpha\beta} \bar{L}_\alpha \tilde{H} H^\dagger L^{\mathrm{C}}_\beta + \text{h.c.}\,, \tag{1}$$

where L_α ($\alpha = \mathrm{e}\,, \mu\,, \tau$) are SM lepton doublets (left-handed chiral neutrinos and charged leptons form three $\mathrm{SU(2)_W}$ doublets); here and below, the superscript C indicates charge conjugation, $\tilde{H}_a = \epsilon_{ab} H^*_b$, $a, b = 1, 2$, H is a Higgs doublet (in the unitary gauge, $H^{\mathrm{T}} = (0, (v+h)/\sqrt{2})$ with $v = 246$ GeV, where h is the scalar Higgs boson observed in experiments at the Large Hadron Collider), ϵ_{ab} is the antisymmetric unit matrix, the parameter Λ with the dimension mass sets the scale of the new physics responsible for the appearance of nonremormalizable interaction (1), the dimensionless parameter β_L determines the strength of the new interaction, the numbers $F_{\alpha\beta}$ characterize the mixing between lepton generations, and h.c. stands for a Hermitian-conjugate term. As a result of the appearance of the nonzero vacuum expectation v of the Higgs field H, interaction (1) introduces mass terms in the active neutrino sector:

$$\mathcal{L}^{(5)}_{\nu\nu} = \frac{\beta_L v^2}{4\Lambda}\, \frac{F_{\alpha\beta}}{2}\, \bar{\nu}_\alpha \nu^{\mathrm{C}}_\beta + \text{h.c.} \tag{2}$$

Hence, using the square mass difference of 'atmospheric neutrinos' as the normalization constant, we obtain an estimate for the new physics scale:

$$\Lambda \sim \beta_L \left(\frac{3 \times 10^{-3}\,\mathrm{eV}^2}{\Delta m^2_{\mathrm{atm}}} \right)^{1/2} 3 \times 10^{14}\ \mathrm{GeV}\,. \tag{3}$$

At the energy scale Λ not exceeding 3×10^{14} GeV, interaction (1) must be substituted by a renormalizable theory similarly to how the four-fermion Fermi interaction is substituted by the electroweak theory at a scale of 100 GeV.

The physics underlying the neutrino scale Λ is not known. However, the fact that this scale is much lower than the gravitational one indicates that precisely the *new particle physics* is the most probable source of active neutrino masses. Neutrino oscillations are the only direct indication of the incompleteness of the SM. In all other respects, the scale is unknown: reducing the scale Λ and simultaneously decreasing the parameter β_L, we can still have relation (3) satisfied.

The introduction of sterile neutrinos N_I, $I = 1, 2, 3$, permits forming one of the possible renormalizable extensions of the SM, leading in the low-energy limit to the appearance of interaction (1). In this case, the Lagrangian acquires kinetic and mass terms for three Majorana neutrinos and the only possible renormalizable Yukawa-type coupling to SM fields [18]:

$$\mathcal{L}_N = \mathrm{i}\,\overline{N}_I \not{\partial} N_I - f_{\alpha I} \bar{L}_\alpha \tilde{H} N_I - \frac{M_{N_I}}{2}\, \overline{N}{}^{\mathrm{C}}_I N_I + \text{h.c.} \tag{4}$$

Here, $\not{\partial} = \partial_\mu \gamma^\mu$, where γ^μ are the Dirac matrices. Models with either three or two sterile neutrinos are appropriate. The sterile neutrino masses M_I and the dimensionless Yukawa coupling constants $f_{\alpha I}$ yield 18 new model parameters (11 in the case of two sterile neutrinos). These are more than necessary in order to express nine (respectively, seven) physical parameters that are the observables of the active neutrino sector: two square mass differences, one *CP*-violating phase, two (one) Majorana phases, and the mass of the lightest neutrino (equal to zero in the case of two sterile neutrinos). The remaining nine (four) free parameters can be used to resolve cosmological problems. Already in the case of two existing sterile neutrinos, an additional *CP*-violating phase appears, with a mixing angle with the active neutrinos and two Majorana masses, whose values can be chosen so as to ensure leptogenesis in the early Universe. In the case of three sterile neutrinos, explaining the phenomenon of dark matter also becomes possible.

Lagrangian (4) contains the Yukawa coupling to the SM Higgs boson, whose vacuum expectation gives the mass terms mixing sterile and active neutrinos:

$$\mathcal{V}_N = v\,\frac{f_{\alpha I}}{\sqrt{2}}\,\bar{\nu}_\alpha N_I + \frac{M_{N_I}}{2}\,\overline{N}_I^{\mathrm{C}} N_I + \mathrm{h.c.}$$
$$= (\bar{\nu}_1, \dots \overline{N}_1^{\mathrm{C}} \dots) \begin{pmatrix} 0 & v\,\dfrac{\hat{f}}{\sqrt{2}} \\ v\,\dfrac{\hat{f}^{\mathrm{T}}}{\sqrt{2}} & \hat{M}_N \end{pmatrix} (\nu_1, \dots N_1 \dots)^{\mathrm{T}} .$$

Here and below, hats indicate matrix quantities. In the SM, neutrinos are left-handed chiral, and therefore mixing separates the right-handed chiral component from the sterile neutrino, thus forming a Dirac mass term $\hat{M}_{\mathrm{D}} = v\hat{f}/\sqrt{2}$.

Introducing a certain hierarchy for the Majorana and Dirac components of the mass matrix, it is possible to explain the smallness of the active neutrino masses by what is known as the seesaw mechanism, in our case type I. Namely, for the matrix of Yukawa couplings $\hat{f}$ of the general form, assuming that

$$M_N \gg \hat{M}_{\mathrm{D}} = v\,\frac{\hat{f}}{\sqrt{2}}\,, \tag{5}$$

we perform an orthogonal transformation that eliminates mixing between the active and sterile sectors. As a result, mixing appears separately in the sterile and active sectors. The corresponding 3×3 matrices have the form

$$\hat{M}_N \simeq \mathrm{diag}\,(M_{N_1}, M_{N_2}, M_{N_3})\,,$$
$$\hat{M}^{\nu} = -\hat{M}_{\mathrm{D}}\,\frac{1}{\hat{M}_N}\,\hat{M}_{\mathrm{D}}^{\mathrm{T}} \propto f^2\,\frac{v^2}{M_N} \ll M_N\,.$$

Diagonalizing the matrix in the sterile sector leads to mass states that are very close to flavor states of sterile neutrinos, and we use the same notation N_I for them in what follows. The three states of the active sector are very close to the flavor states of active neutrinos ν_α. The mass terms in the sector of active neutrinos arise in the second order in the Yukawa couplings and are therefore doubly suppressed by the small ratio of the Dirac and Majorana masses introduced in (5). The hierarchy is enhanced, and precisely this is the seesaw mechanism.

From the standpoint of the phenomenology of the model, the mixing between the flavor state ν_α of active neutrinos (an eigenvector in the basis of weak gauge interactions) and the mass states of neutral fermions is of interest:

$$\nu_\alpha = U_{\alpha i}\nu_i + \theta_{\alpha I} N_I\,.$$

The matrix $\hat{U}$ defines mixing in the sector of active neutrinos, diagonalizing the mass matrix:

$$\hat{U}^{\mathrm{T}} \hat{M}^{\nu} \hat{U} = \mathrm{diag}\,(m_1, m_2, m_3)\,.$$

The mixing matrix $\hat{U}$ is called the Pontecorvo–Maki–Nakagawa–Sakata (PMNS) matrix by analogy with the Cabibbo–Kobayashi–Maskawa (CKM) matrix describing mixing between the upper and lower quarks. The mixing between active and sterile neutrinos describes by small matrix elements under the assumption of the seesaw mechanism:

$$\hat{\theta}_{\alpha I} = \hat{M}^{\mathrm{D}} \hat{M}_N^{-1} = \hat{f}\,\frac{v}{\hat{M}_N} \ll 1\,.$$

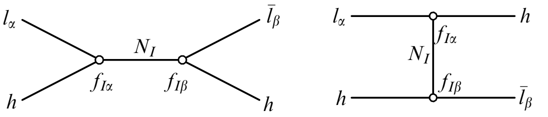

Figure 1. Diagrams for $lh \to lh$ scattering illustrating the appearance of nonrenormalizable interaction (1) in the case of an exchange by a virtual sterile neutrino N_I.

These elements determine the admixture of the heavy neutrino ($\approx$ sterile) mass state to the flavor state participating in weak interactions. The smaller the mixing is, the smaller the probability of sterile neutrinos participating in weak processes.

In the low-energy limit, when processes are considered with a momentum transfer smaller than the sterile neutrino masses, $|q^2| \ll M_N^2$, the model in (4) exactly yields nonrenormalizable interaction (1), with

$$\beta_L\,\frac{\hat{F}}{4\Lambda} = \hat{f}^{\mathrm{T}} \hat{M}_N^{-1} \hat{f}\,. \tag{6}$$

This can be illustrated by the Feynman diagrams for two-particle scattering presented in Fig. 1, where in the limit of small transferred momenta, we must write $1/M_{N_I}$ instead of the sterile neutrino propagator. In a similar limit, the electroweak model yields the Fermi theory of four-fermion interactions. In our case, the scale of new physics Λ in (1) sets the sterile neutrino masses M_N while the values of Yukawa couplings $\hat{f}$ determine the strength of the interaction, the parameter β_L.

Clearly, the *scale of sterile neutrino masses is not fixed*: the correct scale of active neutrino masses can be obtained with Yukawa couplings of the order of unity and sterile neutrino masses of the order of 10^5 GeV, and with smaller Yukawa couplings (for example, $\hat{f} \sim 10^{-6}$, as is the case with the Yukawa coupling of the electron in the SM) and, accordingly, lighter sterile neutrinos (100 GeV). Moreover, because the masses of active neutrinos depend on the quadratic form of Yukawa couplings $f_{\alpha I}$, the contributions of different terms can cancel, owing to which, in the case of not very heavy sterile neutrinos, the Yukawa couplings are not required to have very small values in order to ensure very small active neutrino masses. In the last case, mixing between the active and sterile components, $\hat{\theta}$, is not very small and interesting phenomenological consequences arise in the model. The allowed region in the parameter space of the model is schematically shown in Fig. 2.

3. Phenomenology and cosmological manifestations of sterile neutrinos

The phenomenology of sterile neutrinos and their cosmological manifestations depend essentially on their mass scale.

3.1 Sterile neutrino mass in the range $10^9 - 10^{15}$ GeV

The models with heavy sterile neutrinos are motivated by models of the Grand Unified Theory (GUT), in particular, those based on the gauge group SO(10). In the mass interval $\sim (10^9 - 10^{15})$ GeV, sterile neutrinos can be responsible for the baryon asymmetry of the Universe via the leptogenesis mechanism [14]. The idea is that the decays of nonrelativistic heavy sterile neutrinos in the early Universe could lead to the production of lepton asymmetry. It was then partially

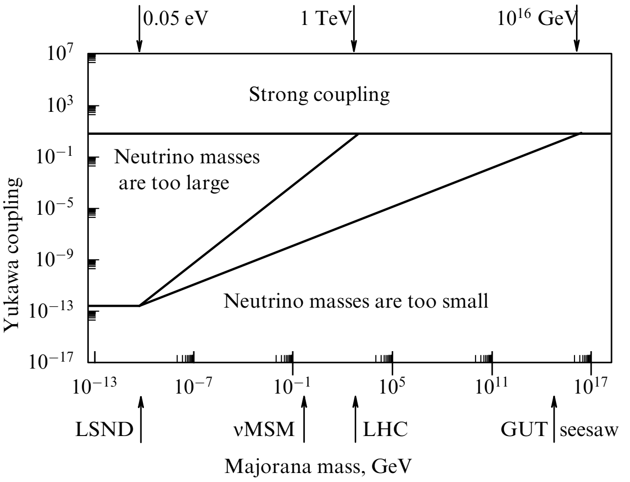

Figure 2. Allowed region in space of parameters of the seesaw mechanism of type I when condition (5) is satisfied; the Yukawa constants then become less than unity and give a correct scale for the active neutrino masses [19].

transformed into baryon asymmetry by so-called sphaleron processes [20], which were in equilibrium in the plasma up to the electroweak phase transition, under which the plasma temperature was ~ 100 GeV.

For the production of lepton asymmetry in the early Universe, three Sakharov conditions must have been satisfied simultaneously: lepton symmetry violation (guaranteed by the Majorana masses $M_{N_I} \neq 0$ of sterile neutrinos), charge C- and CP-symmetry violation (charge conjugation together with reflection of the spatial axes), and a deviation of these processes from equilibrium in the primordial plasma.

Lepton asymmetry production works as follows. In the early Universe, sterile neutrinos were produced *thermally* in processes of SM particle scattering owing to a nonzero mixing $\hat{\theta}$ or to other mechanisms, requiring a further modification of the model (for example, the introduction of a Yukawa sterile neutrino coupling to the inflaton field). When, as a result of expansion of the Universe, the plasma temperature decreases, the sterile neutrinos become nonrelativistic, they stop being produced in particle scattering in the plasma, but owing to the mixing $\hat{\theta}$, they decay into SM particles, which is a non-equilibrium process. The Yukawa couplings f take complex values in general, which ensures fulfilment of the second Sakharov condition. As a result, the sterile neutrino can decay with different probabilities into a lepton and a Higgs boson or via the conjugate channel (into an antilepton and a Higgs boson). The leading contribution to this *microscopic* asymmetry is due to interference of the tree amplitude and the single-loop contribution of the diagrams presented in Fig. 3. Here, we consider the N_1 neutrino decay, for which we can obtain the decay rate

$$\Gamma(N_1 \to lh) = \frac{M_1}{8\pi} \sum_{\alpha} \left| f_{1\alpha} + \frac{1}{8\pi} \sum_{\beta,I} g\left(\frac{M_1}{M_I}\right) f^*_{1\beta} f_{I\alpha} f_{I\beta} \right|^2 .$$

For the microscopic asymmetry value, this gives

$$\delta \equiv \frac{\Gamma(N_1 \to lh) - \Gamma(N_1 \to \bar{l}h)}{\Gamma(N_1 \to lh) + \Gamma(N_1 \to \bar{l}h)}$$

$$= \frac{1}{8\pi} \sum_{I=2,3} g\left(\frac{M_1}{M_I}\right) \frac{\mathrm{Im}\,\left(\sum_\alpha f_{1\alpha} f^*_{I\alpha}\right)^2}{\sum_\gamma |f_{1\gamma}|^2} .$$

The function $g(x)$ is defined by loop integrals and for the lightest sterile neutrino, $M_1 \ll M_{2,3}$, it gives

$$f\left(\frac{M_1}{M_I}\right) = -\frac{3}{2}\,\frac{M_1}{M_I} ,$$

$$\delta = -\frac{3\,M_1}{16\pi}\,\frac{1}{\sum_\gamma |f_{1\gamma}|^2} \sum_{\alpha\beta I} \mathrm{Im}\left[f_{1\alpha} f_{1\beta} \left(f^*_{I\alpha}\,\frac{1}{M_I}\,f^*_{I\beta} \right) \right] .$$

Taking relations (3) and (6) into account, we can obtain the following numerical estimate for the microscopic asymmetry value:

$$\delta \lesssim \frac{3\,M_1}{8\,\pi\,v^2}\,|\Delta m_{\mathrm{atm}}| \simeq 10^{-8}\,\frac{M_1}{10^8\ \mathrm{GeV}} . \qquad (7)$$

The resultant lepton asymmetry produced in the primordial plasma depends on whether the *relativistic* sterile neutrinos were in equilibrium and whether the decay processes of *nonrelativistic* sterile neutrinos were dominant (their production processes can be neglected). If the first condition is fulfilled, then satisfying the second condition requires that in the epoch under discussion, the expansion rate of the Universe, determined by the Hubble parameter, be higher than the decay rate of sterile neutrinos. Then the *macroscopic* lepton asymmetry produced is given by the microscopic asymmetry times the fraction of sterile neutrinos in the plasma, which gives

$$\Delta_L \sim \delta\,\frac{n_{N_1}(M_1)}{s(M_1)} \sim \frac{\delta}{g_*(M_1)} \sim 10^{-2}\,\delta ,$$

where $n_{N_1}(M_1)$ is the sterile neutrino concentration in the plasma, $s(M_1)$ is the entropy density in the plasma, and $g_*(M_1)$ is the effective number of plasma components.

With the estimated microscopic asymmetry value in (7), an explanation of the observed baryon asymmetry of the Universe is possible in the model where the mass of the lightest sterile neutrino exceeds 10^9 GeV. If the condition that the expansion rate be greater than the decay rate is not fulfilled, then the requirement of the sterile neutrino mass

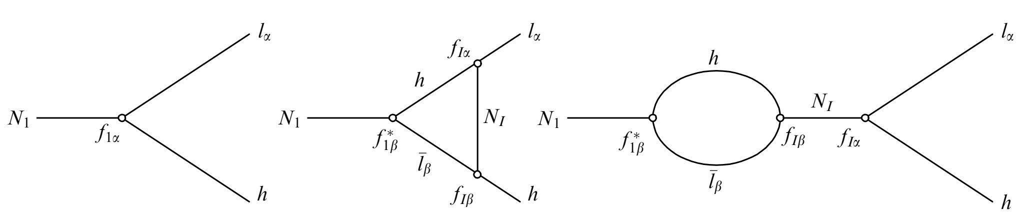

Figure 3. Feynman diagrams providing the leading contribution to the production of microscopic asymmetry in the decay of a heavy sterile neutrino.

becomes more stringent: successful leptogenesis is possible when $M_{N_1} > 10^{12}$ GeV.

Moreover, the resultant lepton asymmetry must not be subsequently 'washed away' in the plasma owing to lepton scattering on Higgs bosons, a process described by the diagrams in Fig. 1. The following estimate holds for the scattering cross section of this process at temperatures below the sterile neutrino mass scale:

$$\sigma_{\mathrm{lh}}^{\mathrm{tot}} \propto \sum_{\alpha\beta I} \left| \frac{f_{I\alpha} f_{I\beta}}{M_I} \right|^2 \propto \frac{\mathrm{tr}\,(\hat{M}^{\nu} \hat{M}^{\nu\dagger})}{v^4} \propto \frac{1}{v^4} \sum m_{\nu}^2 \,.$$

According to the results of numerical calculations, the scattering rate is lower than the expansion rate of the Universe for values of the active neutrino masses

$$m_{\nu} < 0.1-0.3 \text{ eV}.$$

Curiously, modern restrictions from above on the neutrino mass scale (following from cosmology) turn out to be consistent with this requirement, which can be interpreted in favor of the presented leptogenesis version and, consequently, of models with heavy sterile neutrinos. For the mechanism to be successful, at least two sterile neutrinos are necessary.

Regretfully, a straightforward experimental test of the seesaw mechanism by the production of sterile neutrinos with such large masses will not be possible in the foreseeable future. Moreover, the Yukawa coupling to sterile neutrinos gives quantum corrections to the mass of the SM Higgs boson of the order of $f M_N$. For the given range of neutrino masses, this is an unacceptably large value, which requires either a fine tuning of the 'tree' Lagrangian parameters at the scale of sterile neutrino masses or the introduction of a mechanism compensating these corrections (for example, supersymmetrization of the theory).

3.2 Sterile neutrino mass in the range 10^2-10^9 GeV

In the mass range $M_N \sim (10^2-10^9)$ GeV, the dynamics of sterile neutrinos in the early Universe can also lead to the production of the necessary amount of lepton asymmetry. This is possible in a spatial region of model parameters within which at least two sterile neutrinos are degenerate in mass: if the neutrino mass difference is of the order of their total width (inverse lifetime), resonance decay enhancement and equilibrium violation occur in the plasma [21]. As a result, the range of sterile neutrino masses allowing the problem of baryon asymmetry of the Universe to be resolved can be extended to TeV values [15]. At this lower boundary, direct sterile neutrino production is kinematically possible at the Large Hadron Collider, which permits performing direct experimental tests of the model.

We note that a large part of this region in the space of parameters is natural from the standpoint of mass stability of the Higgs boson with respect to quantum corrections due to coupling (4) to sterile neutrinos. The finite correction to the squared mass of the Higgs boson,

$$\delta m_{\mathrm{H}}^2 \propto f^2 M_N^2 \,,$$

should be smaller than the squared mass of the Higgs boson, $m_{\mathrm{H}}^2 \simeq (125\,\mathrm{GeV})^2$, which is fulfilled for masses $M_N \lesssim 10^7$ GeV.

We note that for Yukawa couplings of the general form, their values estimated on the basis of the seesaw mechanism (3) and (6) are quite small, the production rate of sterile neutrinos with masses of the order of 1 TeV is low, and a straightforward test of the model at the Large Hadron Collider is not possible. Testing the model is possible only for a mixing between active and sterile neutrinos that is not small, $\theta_{\alpha I}^2 \lesssim 10^{-3}$ [22].

3.3 Sterile neutrino mass in the range 1 keV – 10^2 GeV

For both phenomenology and cosmology, the range of small sterile neutrino masses (from ~ 1 keV to $\sim 10^2$ GeV) seems to be the most interesting one. In this case, the sterile neutrino masses lie in the same range as the masses of all SM particles, which seems to be natural. For the same reason, no question arises concerning quantum corrections to the Higgs boson mass. The production of sterile neutrinos is kinematically possible at the Large Hadron Collider and in other accelerator experiments.

In this mass range, it is also possible to produce the necessary amount of baryon asymmetry, for which the existence of two sterile neutrinos degenerate in mass is already sufficient. In the early Universe, before the electroweak phase transition, when the temperature exceeded 100 GeV, oscillations in the plasma occurred in the sector of neutral leptons, which for a certain region of the model parameter space led to a redistribution of the lepton charge between active and sterile neutrino flavors [23]. The lepton asymmetry produced in the active sector was transformed into baryon asymmetry by sphaleron transitions. Here, sterile neutrinos should not have reached equilibrium in the primordial plasma; otherwise, the asymmetry would have been 'washed away'. This requirement imposes an upper bound for the value of mixing between active and sterile neutrinos, which is responsible for the interaction of sterile neutrinos with plasma particles. Cosmology also imposes a bound from below for the mixing: the sterile neutrinos should have had time to decay before the onset of primordial nucleosynthesis (the production of light chemical elements in the early Universe); otherwise, energetic SM particles produced in decays (photons, charged leptons, and hadrons), by destroying the produced nuclei, would have affected the ultimate chemical content, which is mainly consistent with observations.

The allowed region in the parameter space of the model with two degenerate sterile neutrinos that explains the baryon asymmetry of the Universe is presented in Fig. 4 for two types of active neutrino mass hierarchies.

The existence of bounds for the mixing both from above and from below allows fully testing the model in experiments. Such a test seems realistic in the case of light mesons. Owing to mixing, sterile neutrinos with masses less than 5 GeV can be produced in weak decays of mesons and baryons and, as a result of mixing, can again decay into SM particles. In the case of cosmologically interesting mixing values, the partial widths of respective meson and baryon decays start at $\sim 5 \times 10^{-7}$ [25], and the sterile neutrino lifetime is presented in Fig. 5. Such values of partial widths and lifetimes determine the setup of an experiment for a direct search for sterile neutrinos: a beam of high-energy protons undergoes scattering on a fixed target in which mesons are produced, and sterile neutrinos are produced in the subsequent decays of the mesons. At a certain distance from the target, an 'empty' chamber is situated, and its volume is scanned in order to reveal pairs of charged particles, the result of sterile neutrino decays. The schematic layout of the experiment is shown in Fig. 6. Direct bounds for the parameters of the model shown

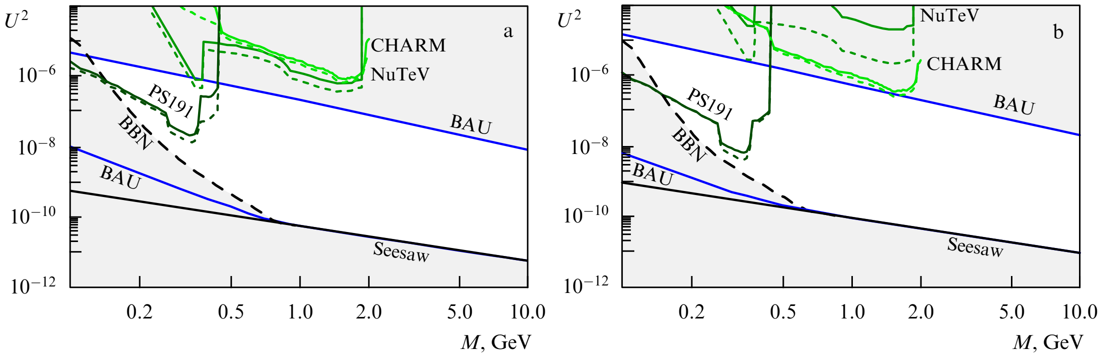

Figure 4. Allowed region in the space (degenerate sterile neutrino mass M, squared mixing strength of active and sterile neutrinos U^2) for (a) the direct and (b) the inverse active neutrino mass hierarchies [24]. The bound from above for the mixing strength ensures leptogenesis [upper curve, BAU (baryon asymmetry of Universe)], bounds from below ensure the primordial nucleosynthesis (BBN, Big Bang nucleosynthesis) and the seesaw mechanism of active neutrino mass formation due to mixing with sterile neutrinos (Seesaw curve). Moreover, presented are restrictions based on direct searches for sterile neutrinos in the experiments CHARM (CERN HAmburg Rome Moscow), NuTeV (Neutrino at the Tevatron), PS-191 (PS is a proton synchrotron): the solid and dashed curves corresponding to these experiments illustrate the dependence on the model parameters.

in Fig. 4 were obtained in precisely such experiments. Recently, a proposal was made for a new experiment of this type based on a 400 GeV proton beam to be provided by the SPS (Super Proton Synchrotron) at CERN [26–28]. Placing the decay volume 50 m long as close as possible to the target allows enhancing the sensitivity to the mixing angle by more than an order of magnitude. We here mean neutrinos with masses less than 2 GeV that are produced in the decays of charmed hadrons (mainly D- and D_s-mesons). To investigate the entire range of possible mixing angles, it would be necessary to increase the length of the decay volume to several kilometers and to make its width increase along the beam axis, at least in one transverse direction [26].

The introduction of a third sterile neutrino allows achieving the production of lepton asymmetry even without the masses of the two neutrinos being degenerate [17]. We note that the lower mixing boundary, both in this case and in the general case of the seesaw mechanism, lies not far from the boundary indicated in Fig. 4, and hence, in principle, it is also possible to comprehensively test the general approach to obtaining active neutrino masses via mixing with light sterile neutrinos [29].

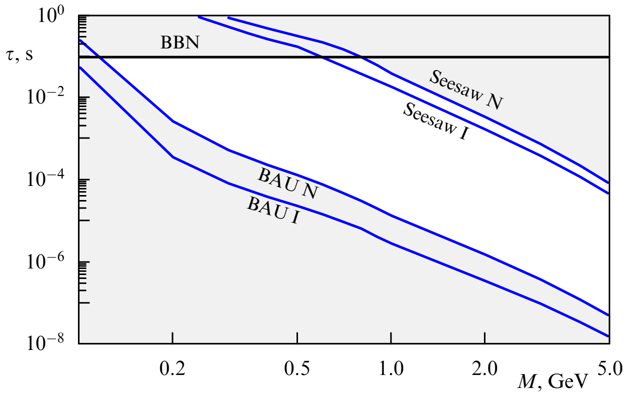

Figure 5. Allowed region in the space (degenerate sterile neutrino mass M, sterile neutrino lifetime τ) for models involving the direct (N) and inverse (I) active neutrino mass hierarchies [24].

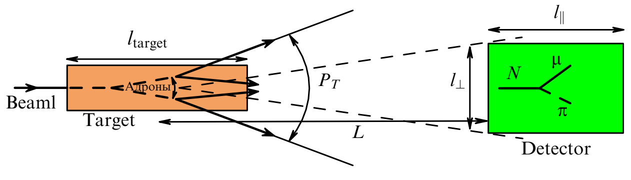

Figure 6. Schematic layout of the experiment for a direct search for sterile neutrinos with a mass less than several GeV [26].

In an alternative version of the values of model parameters, the third sterile neutrino can be used as a candidate for dark matter particles. Such an 'economic' modification of the SM, which allows *resolving the problems of neutrino oscillations, dark matter, and the baryon asymmetry of the Universe by introducing only three Majorana fermions*, has been called the νMSM (Neutrino Minimal Standard Model) in the literature [3].

A sterile neutrino aspiring to play the role of dark matter particles cannot take full part in the formation of active neutrino masses via the seesaw mechanism. The mixing that would then arise would lead to the sterile neutrinos rapidly decaying into three active neutrinos, while the lifetime of dark matter particles should exceed the age of the Universe (15 billion years). The lifetime

$$\tau_{N\to 3\nu} \sim \frac{1}{G_{\text{F}}^2 M_N^5 \theta_{\alpha N}^2} \sim \frac{1}{G_{\text{F}}^2 M_N^4 m_\nu} \sim 10^{11}\left(\frac{10\ \text{keV}}{M_N}\right)^4 \ \text{[years]},$$

where G_{F} is the Fermi constant. The region of small masses is then also forbidden owing to their overproduction in particle scatterings in the primordial plasma (Fig. 7). Moreover, the sterile neutrino undergoes radiative decay into a photon and an active neutrino, which occurs at the one-loop level of the perturbation theory. For the decay rate, we obtain [3]

$$\Gamma_{N\to\nu_\alpha\gamma} \simeq 5.5\times 10^{-22}\,\theta_{\alpha N}^2\left(\frac{M_1}{1\ \text{keV}}\right)^5 \text{s}^{-1}\,.$$

This process provides a test of the dark matter model: sterile neutrinos in galaxies and galaxy clusters decay, emitting a photon of the frequency $\omega_\gamma = M_N/2$. Because the velocities of dark matter particles in the galaxies are of the order of the velocities of stars, $10^{-3}-10^{-4}$ times the speed of light, a nearly monochromatic line is expected to be seen in the sky with such a relative width owing to the Doppler effect. The search for such lines based on the analysis of data obtained by orbital X-ray telescopes has not yet given positive results, but has

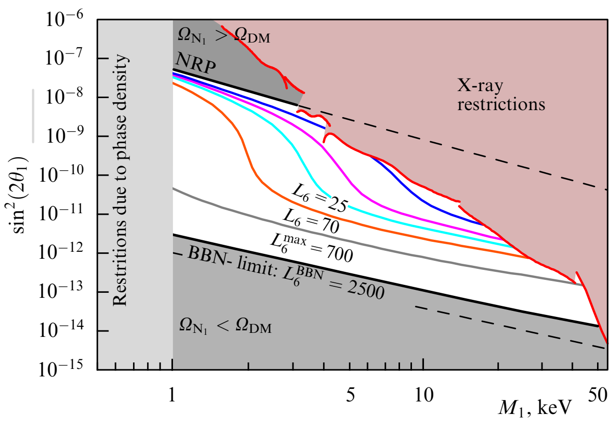

Figure 7. Region in the space of parameters (the lightest sterile neutrino mass M_1, the mixing angle θ_1) within which the sterile neutrino can resolve the dark matter problem [3]. For parameters on the curve denoted as NRP (nonresonant production), neutrino scattering and oscillations in the primordial plasma without lepton asymmetry lead to production of the number of neutrinos required to explain dark matter. The curves L_6 correspond to the value of lepton asymmetry in units of 10^{-6}. Also shown are bounds from direct searches for a monochromatic line in the X-ray spectra of galaxies, due to primordial nucleosynthesis (BBN), and from the phase density value of dark matter particles in galaxies [30].

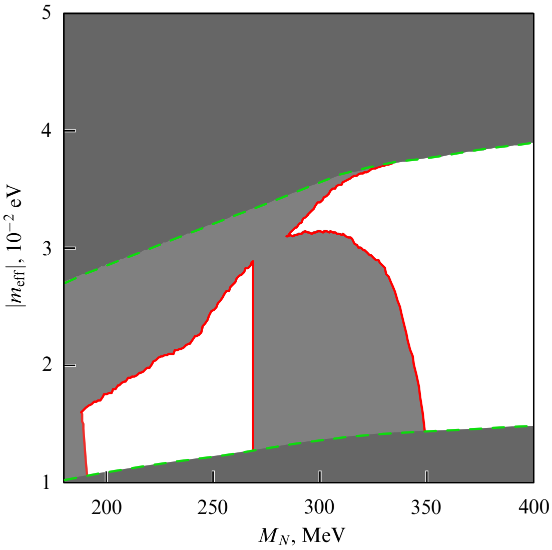

Figure 8. (Color online.) Allowed region in the parameter space (the mass of heavy degenerate sterile neutrinos M_N, the effective mass of double neutrinoless β-decay $m_{\rm eff}$) [36] for the inverse active neutrino mass hierarchy: cosmology limits the region by the green contour, direct searches give the red contours.

permitted ruling out the production mechanism of dark matter sterile neutrinos heavier than 5 keV in particle scatterings and neutrino oscillations in plasma (see Fig. 7). The key circumstance here is the direct dependence of the production rate (also including the number of dark matter particles) and of the decay rate (and, consequently, of the signal in the telescope) on the same mixing strength with active neutrinos. In the framework of the model of three sterile neutrinos, the possibility remains of producing the necessary number of neutrinos in the resonance oscillations in plasma if the lepton asymmetry essentially increases in the Universe after the electroweak phase transition, which is possible in the case of a very strong mass degeneracy of the two heavy sterile neutrinos, $\Delta M \lesssim 10^{-7}$ eV. The corresponding region of the parameter space is also shown in Fig. 7.

Alternative mechanisms for the production of dark matter sterile neutrinos can also be considered. In models with the inflationary stage in the early Universe, this can result in directly coupling the inflaton to sterile neutrinos ($\phi \bar{N}^C N$ for the inflaton that is a singlet under the SM group [31–33]; in the model of Higgs inflation [34], such a role is played by Yukawa coupling (4) [35]). The small contribution of dark matter neutrinos to the active neutrino masses remains common. The corresponding general prediction here is that one of the active neutrinos turns out to be very light; hence follows the small value of the effective mass $m_{\rm eff}$ in the double neutrinoless β-decay (Fig. 8). Nevertheless, for this region of sterile neutrinos, there is a fundamental possibility of directly testing a large part of the space of model parameters in accelerator and orbital experiments.

3.4 Sterile neutrino mass in the range 1 eV – 1 keV

From the standpoint of resolving cosmological problems, the mass range 1 eV–1 keV is of no interest: sterile neutrinos cannot help in producing baryon asymmetry and do not form the necessary amount of dark matter. They can nevertheless influence cosmological processes, because for successful functioning of the seesaw mechanism of active neutrino mass formation, the mixing between active and sterile sectors cannot be small. Such neutrinos thermalize in the primordial plasma, then decouple from it, but remain in the expanding Universe. As an additional ultra-relativistic component (radiation), they contribute to the total energy density of the Universe and can enhance the Universe expansion rate in the epoch of primordial nucleosynthesis and thus alter its predictions. Hence follows the bound for the admissible amount of sterile neutrinos, traditionally presented as a bound for the number of effective neutrino components $N_{\rm eff}$, which is equal to three for three active neutrinos or to four if a light sterile neutrino with a not very small mixing is added. As follows from the analysis of data (Fig. 9), only one neutrino with such properties is allowed in the model. Clearly, the seesaw mechanism would not function in this case: at least one more additional heavy sterile neutrino is required.

It is interesting that a number of anomalous results of the neutrino oscillation experiments LSND [4, 5], MiniBooNE [6–8], GALLEX (GALLium EXperiment) [11], and SAGE (Soviet–American Gallium Experiment) [10, 38] and measurements of reactor antineutrino fluxes [12, 13] can be

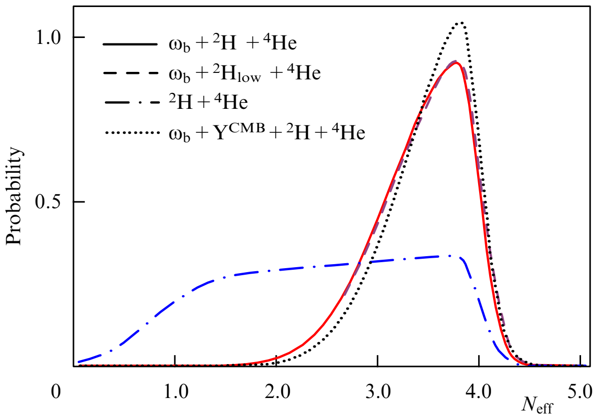

Figure 9. Region of most probable values of the effective number of neutrinos, $N_{\rm eff}$, from cosmological observational data [37].

explained within the model involving a light sterile neutrino. Although it seems highly improbable that the addition of only a single sterile neutrino would be sufficient to explain the entire set of anomalies, these results have roused interest in the mass region of the order of 1 eV. Cosmological observations play an important role here. Light massive neutrinos, including active ones, can influence the course of recombination (the formation of primordial hydrogen) in the early Universe as well as structure formation in the modern Universe. The first process occurs at plasma temperatures of the order of 0.3 eV; in an earlier epoch, particles with such masses contributed to the relativistic component of matter (radiation), and at a later epoch, to the nonrelativistic component (matter), thus altering the evolution of inhomogeneities in the baryon–photon plasma, and ultimately influencing the picture of CMB anisotropy. In the modern Universe, cosmological neutrinos with such masses have quite high velocities and freely travel through the halos of galaxies; however, they can be captured by gravitational potentials of galactic clusters and thus enhance the mass of clusters, affecting their evolution.

For SM neutrinos, the analysis of data from the Planck experiment, which measured the CMB anisotropy, together with data on the large-scale structure of the Universe (the abundance of galaxies and galactic clusters in the observable Universe) leads to a bound for the sum of the three SM neutrino masses at a 95% confidence level (CL) [39]:

$$\sum m_\nu < 0.23 \text{ eV} .$$

If a sole light sterile neutrino is added to the model, then the analysis of data yields the allowed range of cosmological parameters presented in Fig. 10. For certain sets of data the results of the analysis of cosmological observables do not contradict the existence of a light sterile neutrino and can even indicate that the version of the model with such a neutrino is preferable to the traditional SM version, but there are certain differences between the parameters that are consistent with cosmological data and those that are preferable for explaining the anomalies of oscillation experiments. We note that if the existence of these anomalies is confirmed or disproved by the results of direct searches (see the collection of proposals in [19] as an example concerning the gallium anomaly [41] and Ref. [42] for the reactor anomaly), then *promising cosmological experiments will permit us to investigate models involving small mixing angles, which are inaccessible for tests by oscillation experiments*. Thus, with the future catalogue of galaxies, Euclid, it will be possible to test the contribution of neutrinos to the radiative component with a one-percent accuracy (Fig. 11). Clearly, the contribution of such a CMB component to the total mass density of dark matter is small,

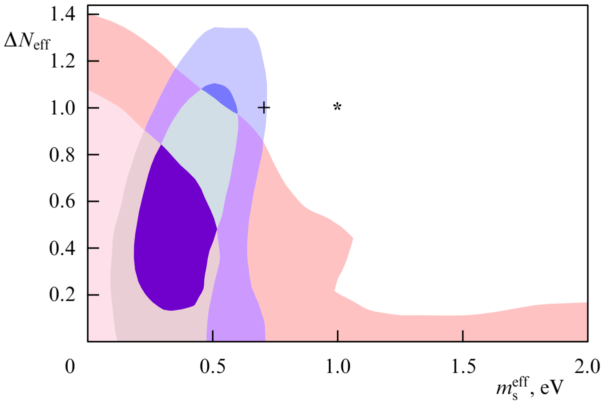

Figure 10. (Color online.) The region of most probable values in the space of parameters (the sterile neutrino mass $m_s^{\rm eff}$, the excess number of effective neutrino species over that expected in SM $\Delta N_{\rm eff} \equiv N_{\rm eff} - 3$) [40] from the analysis of the Planck experiment results (red and pink regions respectively correspond to 68 and 95% CL) and from the analysis of the whole set of cosmological data (the navy-blue and sky-blue regions respectively correspond to 68 and 95% CL), including data on the CMB anisotropy, observational data on galactic clusters via the Sunyaev–Zeldovich effect, results of measurements of the parameters of baryon acoustic (Sakharov) oscillations with the aid of analysis of the large-scale structure of the Universe and measurement of the Hubble parameter, determining the expansion rate of the Universe. The cross and small star respectively indicate the values in the parameter region that are preferable from the standpoint of resolving the problems of accelerator anomalies (LSND, MiniBooNE) and reactor anomalies (GALLEX, SAGE, reactor experiments).

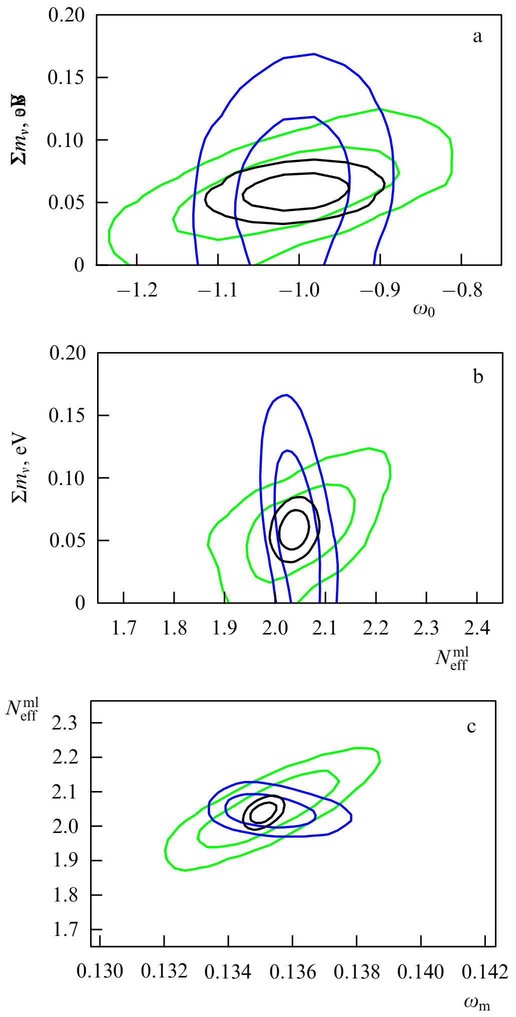

Figure 11. Estimate of the sensitivity of the data analysis for the future galaxy catalogue Euclid [43] to parameters of the neutrino sector: the sum of neutrino masses $\sum m_\nu$ and the number of relativistic degrees of freedom $N_{\rm eff}^{\rm ml}$ for the model involving the direct active neutrino mass hierarchy, when the difference of $N_{\rm eff}^{\rm ml}$ from two (to be precise, from 2.046) points to the existence of a new component. The quantity ω_0 parameterizes the dark matter equation of state (with $\omega_0 = 1$ for the cosmological constant); the value of $\omega_{\rm m}$ with high accuracy is equal to half the relative contribution of matter to the total energy density.

while for all massive neutrinos, we obtain

$$\Omega_\nu \approx \frac{\sum m_\nu}{46\,\mathrm{eV}} \ll \Omega_{\mathrm{DM}} \approx 0.25\,.$$

Hence, there can be only one sterile neutrino in the mass range 1 eV–1 keV, which is insufficient for explaining the results of oscillation experiments. Mixing angles that are not small can be tested in oscillation experiments. Here, the highest sensitivity to the mentioned mass range is exhibited by cosmological measurements.

4. Conclusion

In conclusion, we stress that the addition of sterile neutrinos in the SM is probably the most economical way, in the framework of a renormalizable generalization of the SM, to explain the three most acute phenomenological problems: neutrino oscillations, the baryon asymmetry of the Universe, and the dark matter phenomenon. To resolve the first two problems, it suffices to introduce only two sterile neutrinos. The production of asymmetry requires *CP*-symmetry violation in the neutrino sector; however, a corresponding parameter in the hidden sector is sufficient for this purpose: *the revelation of CP violation in the active neutrino sector or confirmation of its CP invariance will not confirm or disprove the explanation of baryon asymmetry via the leptogenesis model*. The third sterile neutrino can be stable for cosmological periods of time and form dark matter; however, its contribution to the masses of active neutrinos is small. An experimental indication of the existence of such dark matter would be given by a monochromatic line in the photon spectrum from galaxies and galactic clusters due to two-body radiative decays of sterile neutrinos. Cosmological observations (at least for the present) allow the existence of a sterile neutrino with a mass less than 1 eV; however, such a particle cannot simultaneously explain all the existing oscillation anomalies. There are good prospects for testing this model in the nearest future after the analysis of accumulated data (Planck) and as a result of implementation of next-generation experiments.

The addition of sterile neutrinos provides right-handed chiral components in the neutrino sector of the SM, in which both upper and lower quarks have left-handed and right-handed chiral components. Remarkably, part of the space of this model for neutrino masses less than several GeV can be tested in direct particle-physics experiments and orbital experiments.

References

1. Mohapatra R N, Pal P B *Massive Neutrinos in Physics and Astrophysics* (World Scientific Lecture Notes in Physics, Vol. 60) 2nd ed. (River Edge, N.J.: World Scientific, 1998)
2. Mohapatra R N, Pal P B *Massive Neutrinos in Physics and Astrophysics* (World Scientific Lecture Notes in Physics, Vol. 72) 3rd ed. (River Edge, N.J.: World Scientific, 2004)
3. Boyarsky A, Ruchayskiy O, Shaposhnikov M *Annu. Rev. Nucl. Part. Sci.* **59** 191 (2009); arXiv:0901.0011
4. Athanassopoulos C et al. (LSND Collab.) *Phys. Rev. Lett.* **75** 2650 (1995); nucl-ex/9504002
5. Aguilar A et al. (LSND Collab.) *Phys. Rev. D* **64** 112007 (2001); hep-ex/0104049
6. Aguilar-Arevalo A A et al. (MiniBooNE Collab.) *Phys. Rev. Lett.* **98** 231801 (2007); arXiv:0704.1500
7. Aguilar-Arevalo A A et al. (MiniBooNE Collab.) *Phys. Rev. Lett.* **102** 101802 (2009); arXiv:0812.2243
8. Aguilar-Arevalo A A et al. (MiniBooNE Collab.) *Phys. Rev. Lett.* **105** 181801 (2010); arXiv:1007.1150
9. Aguilar-Arevalo A A et al. (MiniBooNE Collab.) *Phys. Rev. Lett.* **110** 161801 (2013); arXiv:1207.4809; arXiv:1303.2588
10. Abdurashitov J N et al. *Phys. Rev. C* **73** 045805 (2006); nucl-ex/0512041
11. Kaether F et al. *Phys. Lett. B* **685** 47 (2010); arXiv:1001.2731
12. Mueller Th A et al. *Phys. Rev. C* **83** 054615 (2011); arXiv:1101.2663
13. Huber P *Phys. Rev. C* **84** 024617 (2011); *Phys. Rev. C* **85** 029901 (2012), Erratum; arXiv:1106.0687
14. Fukugita M, Yanagida T *Phys. Lett. B* **174** 45 (1986)
15. Pilaftsis A, Underwood T E J *Phys. Rev. D* **72** 113001 (2005); hep-ph/0506107
16. Asaka N, Shaposhnikov M *Phys. Lett. B* **620** 17 (2005); hep-ph/0505013
17. Drewes M, Garbrecht B *JHEP* (03) 096 (2013); arXiv:1206.5537
18. Minkowski P *Phys. Lett. B* **67** 421 (1977)
19. Abazajian K N et al., arXiv:1204.5379
20. Kuzmin V A, Rubakov V A, Shaposhnikov M E *Phys. Lett. B* **155** 36 (1985)
21. Pilaftsis A, Underwood T E J *Nucl. Phys. B* **692** 303 (2004); hep-ph/0309342
22. Dev P S B, Pilaftsis A, Yang U *Phys. Rev. Lett.* **112** 081801 (2014); arXiv:1308.2209
23. Akhmedov E Kh, Rubakov V A, Smirnov A Yu *Phys. Rev. Lett.* **81** 1359 (1998); hep-ph/9803255
24. Canetti L et al. *Phys. Rev. D* **87** 093006 (2013); arXiv:1208.4607
25. Gorbunov D, Shaposhnikov M *JHEP* (10) 015 (2007); *JHEP* (11) 101 (2013); arXiv:0705.1729
26. Gorbunov D, Shaposhnikov M, Proposal submitted to European Strategy Group (2012)
27. Gninenko S N, Gorbunov D S, Shaposhnikov M E *Adv. High Energy Phys.* **2012** 718259 (2012); arXiv:1301.5516
28. Bonivento W et al., arXiv:1310.1762
29. Gorbunov D, Panin A *Phys. Rev. D* **89** 017302 (2014); arXiv:1312.2887
30. Gorbunov D, Khmelnitsky A, Rubakov V *JCAP* **0810** 41 (2008); arXiv:0808.3910
31. Shaposhnikov M, Tkachev I *Phys. Lett. B* **639** 414 (2006); hep-ph/0604236
32. Bezrukov F, Gorbunov D *JHEP* (05) 010 (2010); arXiv:0912.0390
33. Bezrukov F, Gorbunov D *JHEP* (07) 140 (2013); arXiv:1303.4395
34. Bezrukov F, Shaposhnikov M *Phys. Lett. B* **659** 703 (2008); arXiv:0710.3755
35. Bezrukov F, Gorbunov D, Shaposhnikov M *JCAP* (10) 001 (2011); arXiv:1106.5019
36. Asaka T, Eijima S *Prog. Theor. Exp. Phys.* **2013** (11) 113B02 (2013); arXiv:1308.3550
37. Mangano G, Serpico P D *Phys. Lett. B* **701** 296 (2011); arXiv:1103.1261
38. Gavrin V N *Phys. Usp.* **54** 941 (2011); *Phys. Usp.* **181** 975 (2011)
39. Ade P A R et al. (Planck Collab.), arXiv:1303.5076
40. Hamann J, Hasenkamp J *JCAP* (10) 044 (2013); arXiv:1308.3255
41. Gavrin V N et al., arXiv:1006.2103
42. Danilov M (DANSS Collab.), arXiv:1311.2777
43. Basse T et al., arXiv:1304.2321