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A SEARCH FOR STERILE NEUTRINOS IN OSCILLATION PROCESSES

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Moscow 2021

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INTRODUCTION

The Standard Model of particle physics is currently the leading fundamental theory governing the laws of physics of our universe. It contains 6 leptons, 6 quarks, gauge $W^{\pm}$, Z^0, γ, gluon bosons, Higgs boson. This theory has been tested with very high precision. In the same time we know for sure that this theory isn't final theory of our world, still we have questions to it, for example, neutrino oscillations and existence of dark matter cannot be explained in terms of the Standard Model.

One of the extensions of the Standard Model that can serve as an explanation for these phenomena is the theory of existence of sterile neutrinos – neutral leptons that are supposed to interact only via gravitaty (since it have mass) but not via any of the other fundamental interactions so it means addition of new freedom degrees that are not charged by gauge group of Standard Model $\mathrm{SU}(3)_{\mathrm{c}} \times \mathrm{SU}(2)_{\mathrm{W}} \times \mathrm{U}(1)_{\mathrm{Y}}$. [1]

Today this theory is actively developed for several reasons:

- Sterile neutrinos generically appear in models that explain the smallness of neutrino masses. This is particularly true for the famous «seesaw mechanism» and its variants.[2][3]
- There is overwhelming evidence for the existence of new, extremely weakly interacting particles in the Universe that constitute about 80% of its total matter content. Given strong experimental constraints on this «dark matter», it very likely does not carry Standard Model gauge charges. The dark matter itself could be in the form of sterile neutrinos.[4]
- Finally, abnormal results of neutrino oscillation experiments such as a LSND[5], MiniBooNE[6] and discovery of reactor [7–10] and gallium [11–13] anomalies also suggest the existance of sterile neutrinos.

In this paper it's proposed to consider the results of oscillation experiments on the search for sterile neutrinos.

1. THEORETICAL MODELS WITH STERILE NEUTRINOS

1.1. NEUTRINO PORTAL

It is often argued that the observation of neutrino oscillations – and thus of non-zero neutrino mass – is the first evidence for physics beyond the SM of particle physics. Indeed, in the Lagrangian of the original SM, only left-handed neutrino fields appear. A single Weyl spinor was thus sufficient to describe each neutrino flavor, whereas for all other fermions, two Weyl spinors (corresponding to left-handed and right-handed polarizations) were required. Introducing neutrino masses suggests including new Weyl fermions.

$$\mathcal{L} \supset -y^{\alpha\beta}(i\sigma^2 H^*)L^\alpha N^\beta + h.c. \tag{1.1}$$

where H – Higgs doublet, L^α are the lepton doublets, $L^\alpha = (\nu^\alpha, e^\alpha)^T$, σ^2 is the second Pauli matrix, α and β are flavor indices, $\alpha = e, \mu, \tau$, $\beta = 1...n$ where $n \geq 2$ otherwise there would be two or more exactly massless neutrino states left, in conflict with the observed oscillation patterns. L^α and N^β should be interpreted as two-component spinors that transforms as $\left(\frac{1}{2}, 0\right)$ under Lorentz group. After the Higgs field acquires its vacuum expectation value v, the Yukawa couplings from eq. 1.1 yield mass terms of the form

$$\mathcal{L} \supset -M_D^{\alpha\beta}\nu^\alpha N^\beta + h.c. \tag{1.2}$$

where $M_D^{\alpha\beta} = \frac{y^{\alpha\beta}v}{\sqrt{2}}$ («D» means Dirac mass). Looking at the transormation properties of the field appearing in 1.1 under SM gauge symmetries, it straightforward to verify that $(i\sigma^2 H^*)$ and L^α both transform as doublets under $SU(2)_L$ and carry opposite hypercharges of $+\frac{1}{2}$ and $-\frac{1}{2}$. Thus the $(i\sigma^2 H^*)L^\alpha$

is a total singlet under the SM symmetries then N^β must be total singlets too. So N^β are called «sterile neutrinos» – «neutrinos» because they don't interact via electromagnetic or strong fields and «sterile» because they dont interact via weak fields, ν^α are «active neutrinos» that are well known for today. In fact, a coupling to the operator $(i\sigma^2 H^*)$ is the only direct renormalizable coupling that a SM singlet fermion can have with the SM. This type of operator is therefore also called the «neutrino portal».

1.2. SEESAW MECHANISM

The fact that the N fields in 1.1 and 1.2 are total SM singlets implies that they admit a Majorana mass term:

$$\mathcal{L}_{N_{mass}} = -\frac{1}{2} M_M^{\alpha\beta} N^\alpha N^\beta + h.c. \tag{1.3}$$

To get more clear the phenomenological consequences of the mass terms in 1.2 and 1.3 lets write ν and N as a vector $n = (\nu^e ... \nu^\tau, N^1 ... N^n)^T$ and write the mass term with matrix notation[14]:

$$\mathcal{L} \supset -\frac{1}{2} n^T M_n = -\frac{1}{2} n^T \begin{pmatrix} 0 & M_D \\ M_D^T & M_R \end{pmatrix} n + h.c. \tag{1.4}$$

where $M_D = (M_D^{\alpha\beta})$ is the general complex Dirac mass matrix from 1.2, $M_M = (M_M^{\alpha\beta})$ is the complex symmetric Majorana mass matrix from 1.3.

In the limit where the eigenvalues of M_M in 1.4 are significantly larger than those of M_D let's consider a seesaw mechanism with three neutrino mass eigenvalues of order $\frac{||M_D||^2}{||M_M||}$[1] [15]

[1] $||A||$ is a Euclidean matrix norm, for $A = (a_{ij})$ it's $||A|| = \sqrt{\sum_{ij} |a_{ij}|^2}$

We can set $||M_M|| \gg ||M_D||$ since M_M is not protected by any symmetry while M_D cannot be larger than electroweak scale (this would require Yukawa coupling $\gg 1$). If $||M_D|| \sim 100$ GeV and $||M_D|| \sim 10^{14}$ GeV:

$$\frac{||M_D||^2}{||M_M||} \sim \frac{10^4\,\mathrm{GeV}^2}{10^{14}\,\mathrm{GeV}} = 0.1\,\mathrm{eV} \tag{1.5}$$

According to this, we've got that the light neutrino mass eigenvalues are of order 0.1 eV so the seesaw mechanism gives an explanation for observed smallness of neutrino masses.

While the seesaw mechanism is one of the most widely discussed application of sterile neutrinos, the singlet states appearing in it are currently unobservable in practice. Even when their masses are within reach of current collider experiments, their extremely small mixing $\sim ||M_D||/||M_M||$ with active neutrinos prevents their efficient production.

Besides that sterile neutrino masses could be related to the generation of neutrino masses according to Inverse Seesaw Mechanism[16] and Extended Seesaw Mechanism model [17] if sterile neutrino masses below electroweak scale and their mixing angles $\gtrsim 10^{-2}$, but we will not consider these models in this paper.

2. NEUTRINO OSCILLATIONS WITH MORE THAN THREE FLAVORS

As it has been shown how sterile neutrino can appear in neutrino mass models now let's consider their phenomenology in experiments.

2.1. LAGRANGIAN FORMALISM

After extending of Standard Model by n sterile neutrinos, in general, we have the same formalism of neutrino oscillation as for three-flavor case. The weak interaction Lagrangian in the flavor basis is the same as in the Standard Model[18]:

$$\mathcal{L} = \sum_{\alpha=e,\mu,\tau} \left[\frac{g}{\sqrt{2}} (\overline{\nu_{\alpha,L}} \gamma^{\rho} e_{\alpha,L} W_{\rho}^{+} + h.c.) + \frac{g}{2cos\theta_w} \overline{\nu_{\alpha,L}} \gamma^{\rho} e_{\alpha,L} Z_{\rho} \right] \tag{2.1}$$

Here, g is the weak coupling constant and θ_W is the Weinberg angle. As seen from this Langragian, only left neutrinos interact with $W^{\pm}$ and Z^0 gauge bosons.

Transformation to the mass basis for (3+n) mass eigenstates:

$$\nu_{\alpha} = \sum_{j=1}^{3+n} U_{\alpha j} \nu_j \tag{2.2}$$

Here mixing matrix U is $(3+n) \times (3+n)$ matrix. An initial neutrino flavor state $|\nu_{\alpha}\rangle$ is created by acting on the vacuum $|0\rangle$ with operator $\nu_{\alpha}^{\dagger}$. So its decomposition to mass eigenstates $|\nu_j\rangle$ are:

$$|\nu_\alpha\rangle = \sum_j U^*_{\alpha j}|\nu_j\rangle \tag{2.3}$$

After time T and way's length the neutrino state $|\nu_\alpha(T,L)\rangle$:

$$|\nu_\alpha(T,L)\rangle = \sum_j e^{-iE_jT+ip_jL}U^*_{\alpha j}|\nu_j\rangle \tag{2.4}$$

Here E_j and p_j are the energy and momentum of j-th mass eigenstate.

If neutrino was detected in $\langle\nu_\beta|$ flavor state, the corresponding oscillation probability is:

$$P_{\alpha\beta} = |\langle\nu_\beta|\nu_\alpha(T,L)\rangle|^2 = \sum_{j,k} U^*_{\alpha j}U_{\beta j}U_{\alpha k}U^*_{\beta k}e^{-i(E_j-E_k)T+i(p_j-p_k)L} \tag{2.5}$$

Using relativistic energy-momentum realtion $p_j^2 = E_j^2 - m_j^2$ we can rewrite ths oscillation phase approximaterly as $-i(E_j - E_k)(T-L) - i\Delta m^2_{jk}L/(2E)$ where E is a suitably chosen average energy which should differ from E_j, E_k by no more that the neutrino wave packet width in momentum space. The first term here is neglible since its order $\Delta m^2_{jk}/E^2 \cdot \sigma_x$ (with σ_x the wave packet width in coordinate space). The second term is the standard oscillation phase. After that we have next expression for oscillation probability:

$$P_{\alpha\beta} = \sum_{j,k} U^*_{\alpha j}U_{\beta j}U_{\alpha k}U^*_{\beta k}e^{-i\Delta m^2_{jk}L/(2E)} \tag{2.6}$$

where j and k run from 1 to $(3+n)$

2.2. ELECTRON NEUTRINO APPEARANCE SEARCHES

In this chapter we will consider appearance of ν_e in $\nu_\mu \rightarrow \nu_e$ reaction. The results below are based on osclation probabilities calculated in the full 4-flavor model.

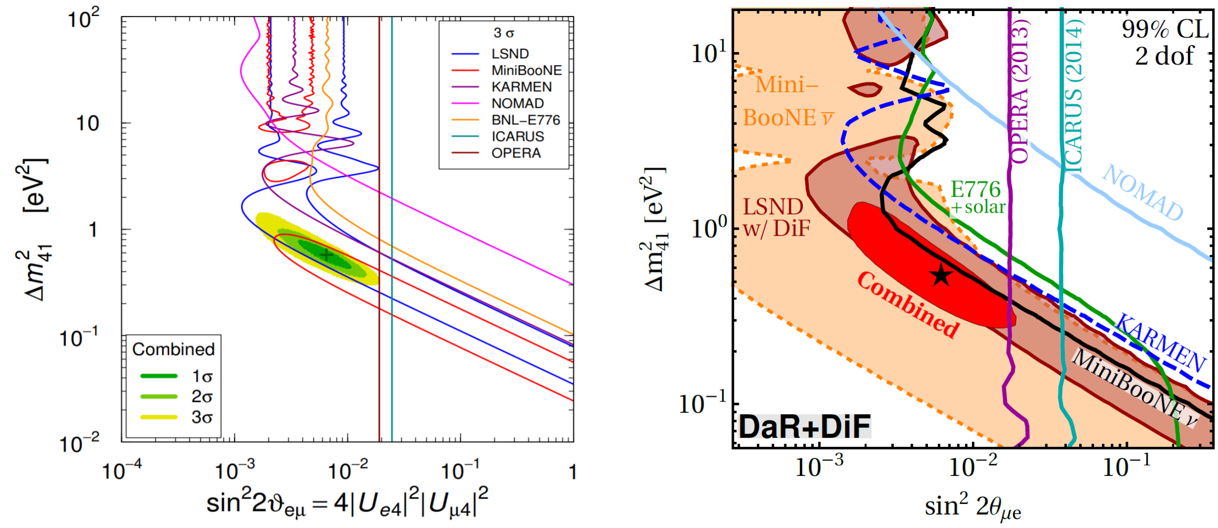
In the short-baseline limit($\Delta m_{21}^2 L/E \ll 1, \Delta m_{31}^2 L/E \ll 1$) the probability of $\nu_\mu \rightarrow \nu_e$ is:

$$P_{\mu e} \approx 4|U_{e4}|^2|U_{\mu 4}|^2 \sin^2\frac{\Delta m_{41}^2 L}{4E} \tag{2.7}$$

This is just the familiar two-flavor oscillation formula, with an effective mixing angle defined as

$$\sin^2 2\theta_{\mu e} = 4|U_{e4}|^2|U_{\mu 4}|^2 \tag{2.8}$$

We see that the effective mixing angle depends both on the mixing of sterile neutrinos with electron neutrinos U_{e4} and muon neutrinos $U_{\mu 4}$



(a) (b)

Figure 2.1 – Global constraints on $\nu_\mu \rightarrow \nu_e$ (and $\overline{\nu}_\mu \rightarrow \overline{\nu}_e$) oscillations at short baseline in the 3 + 1 framework [19][20]

In figure 2.2 the global constraints (as of spring 2018) on $\nu_\mu \rightarrow \nu_e$ (and

$\overline{\nu}_\mu \to \overline{\nu}_e$) oscillations at short baseline in the 3 + 1 framework, comparing the results from [19] (Nu-Fit/GLoBES) and [21] (Italy) are shown. Both group agree on the best fit region around effective mixing angle $\sin^2 2\theta_{\mu e} \sim 0.01$ and $\Delta m^2_{41} \lesssim 1\,\mathrm{eV}^2$. Here Δm^2_{41} is a measuring of the mass of the 4th neutrino and $\sin^2 2\theta_{\mu e} = 4|U_{e4}|^2|U_{\mu 4}|^2$ is measuring of the effective mixing which depends on the mixing between sterile neutrinos and electron neutrinos (U_{e4}) and on the mixing between sterile neutrinos and muon neutrinos ($U_{\mu 4}$).

The parameter region favored by MiniBooNE and LSND is in agreement with the null results from all other experiments, leading us to the conclusion that the global data on $\nu_\mu \to \nu_e$ oscillations, when viewed in isolation, is consistent and could point towards the existence of an eV-scale sterile neutrino.

2.3. ELECTRON NEUTRINO DISAPPEARANCE SEARCHES

Researchers on ν_e and $\overline{\nu}_e$ disappearance are driven by reactor experiments. In last years the following reactor experiments have been carried out or are still in progress: Double Chooz[9], Daya Bay[8], DANSS[22], Neutrino-4[10], STEREO[23], PROSPECT[24]. By comparing fluxes and spectra at different baselines, these experiments can search $\overline{\nu}_e$ disappearance due to oscillations into sterile neutrinos without having to rely on theoretical flux predictions.

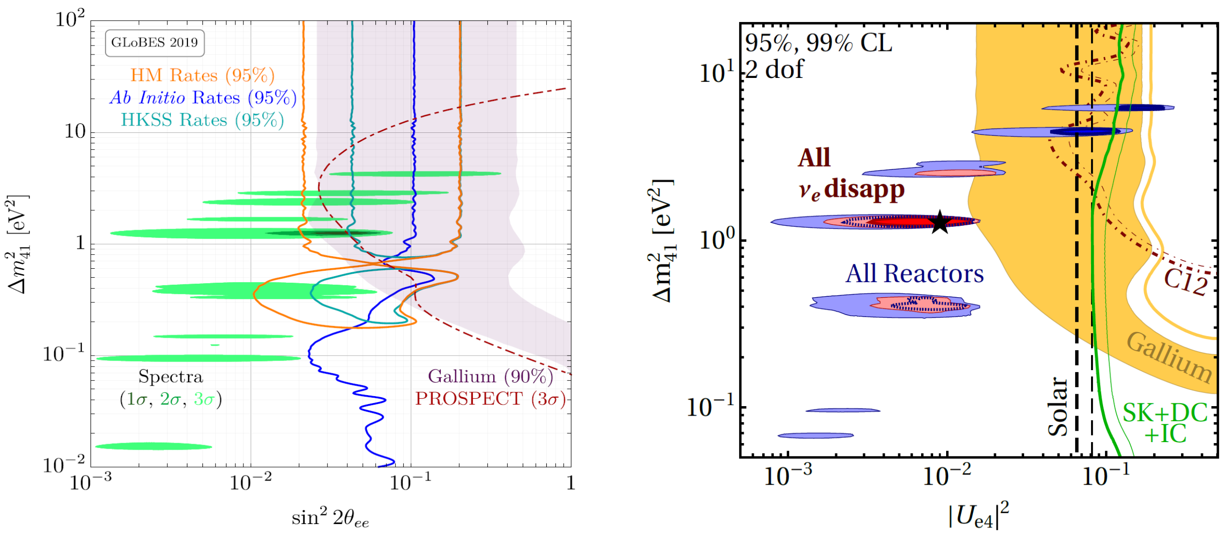
Consider again the short-baseline approximation as in 2.2. Then ν_e disappearance probability:

$$P_{ee} \approx 1 - 4|U_{e4}|^2(1 - |U_{e4}|^2)\sin^2\frac{\Delta m^2_{41} L}{4E} \tag{2.9}$$

And mixing angle θ_{ee}:

$$\sin^2 2\theta_{ee} = 4|U^2_{e4}|(1 - |U_{e4}|^2) \tag{2.10}$$

We see in panel (a) (from ref. [25]) that different theoretical flux predictions lead to significantly different results concerning the significance of the



(a) (b)

Figure 2.2 – Global constraints on short-baseline ν_μ and ν_e disappearance, shown as a function of the squared mass difference Δm^2_{41} and of the $\nu_e - \nu_s$ mixing. Panel (a) [25] displays the preferred parameter region based on comparing total measured reactor neutrino event rates to three different theoretical calculations (orange, blue and turquoise lines); it also displays as green filled contours the preferred regions based on reactor spectra alone (without considering total rate information), and in gray the preferred region corresponding to the “gallium anomaly”. The maroon dot-dashed curve shows the sensitivity of the PROSPECT experiment. Panel (b) [19] displays the fit to reactor data (blue ellipses) also to constraints from solar neutrinos (black dashed), atmospheric neutrinos (green), ν_e–^{12}C scattering (maroon dot-dashed), and to the best fit region based on a combination of all ν_e and $\overline{\nu}_e$ datasets (red ellipses, with the star showing the best fit point)

reactor anomaly: the Huber–Mueller calculation from refs. [26][27] lead to a significant apparent deficit of events, while the calculation from ref. [28] does not lead to a significant anomaly. This discrepancy between different theoretical calculations highlights the large and difficult-to-estimate systematic uncertainties in these calculation and prevents us from drawing firm conclusions from the rate anomaly.

2.4. MUON NEUTRINO DISAPPEARANCE SEARCHES

ν_μ disappearance case is similar to ν_e case and can be described by a simple two-flavor expression in the short-baseline limit. The ν_μ survival probability:

$$P_{\mu\mu} \approx 1 - 4|U_{\mu 4}|^2(1 - |U_{\mu 4}|^2)\sin^2\frac{\Delta m^2_{41}L}{4E} \tag{2.11}$$

And mixing angle:

$$\sin^2 2\theta_{\mu\mu} = 4|U^2_{\mu 4}|(1 - |U_{\mu 4}|^2) \tag{2.12}$$

Lets consider that $\Delta m^2_{21}L/E \ll 1, \Delta m^2_{31}L/E \ll 1$ so that the short baseline approximations from (2.7) (2.9) (2.11) are valid but simultaneously $\Delta m^2_{41}L/E \gg 1$ so the oscillating terms $\Delta m^2_{41}L/4E$ in these equations average to $\frac{1}{2}$ due to limited experimental energy resolution. In this limit (2.7) (2.9) (2.11) depend on two parameters: $|U_{e4}|^2$ and $|U_{\mu 4}|^2$. By observing all three oscillation channels (ν_e disappearance, ν_μ disappearance and $\nu_\mu \to \nu_e$ appearance), one can thus over-constrain the system, allowing for consistency tests. In practice, the above limits on Δm^2_{21}, Δm^2_{31}, Δm^2_{41} are simultaneously realized only in a small region of parameter space. However, as long as measurements at different energies are available (which is almost always the case), the conclusion that the system can be over-constrained remains valid.

It is illustrated in fig. 2.3 where authors of [19] compare the parameter region preferred by ν_e appearance experiments and ν_e disappearance experiments to the exclusion limits from ν_μ disappearance searches in the $|U_{\mu 4}|^2 - \Delta m^2_{41}$ plane. Evidently, there is stark tension in the global data set: the parameter region preferred by the short baseline anomalies is ruled out at high significance by the null results from ν_μ disappearance

This tension has been quantified in [19] using a parameter-goodness-of-fit (PG) test [29]. This test measures the statistical «penalty» one has to

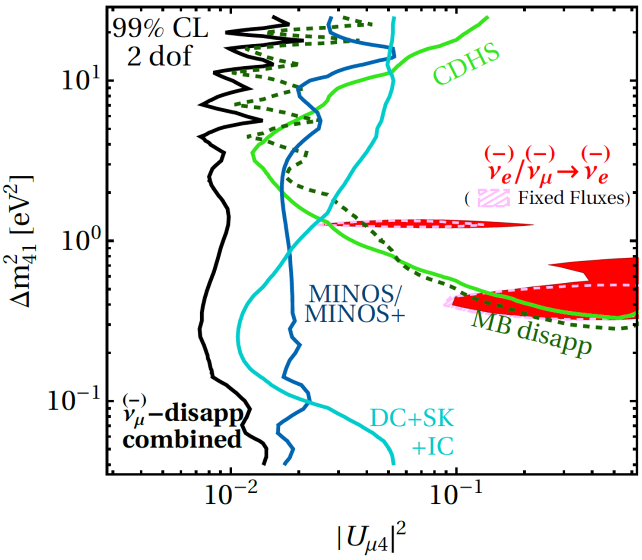


Figure 2.3 – Global constraints on short-baseline ν_μ and $\overline{\nu}_\mu$ disappearance in the $3+1$ scenario, shown as a function of the $\nu_s - \nu_\mu$ mixing $|U_{\mu 4}|^2$ and the squared mass difference Δm^2_{41}. Thick colored lines show exclusion limits from various searches, with the region to the left of the curves excluded. The red shaded region is preferred by the combination of $\nu_e/\overline{\nu_e}$ appearance and disappearance searches, where the anomalies are observed. There is clear tension between this data and the exclusion limits [19]

pay for combining two data sets. It does so by comparing the likelihood of the individual data sets at their respective best fit points to the likelihood of the combined data set at the global best fit point. If the global short-baseline neutrino oscillation data were fully consistent, the PG test should yield a large p-value for any portioning of the data into two independent subsets. The authors of [29] have in particular divided the data into appearance and disappearance data, and have found a small p-value of 3.71×10^{-7}. The p-value remains very small when any individual data set is removed from the fit. Only removing LSND has a significant impact, increasing the p-value to 1.6×10^{-3} which is still small.

The consistency of the fit also does not improve when more than one sterile neutrino is considered [30]. Even though the mixing matrix has many more parameters in this case, the only qualitatively new feature that appears in $3+2$ and $3+3$ models compared to the simple 3+1 scenario is the possibility of CP violation at short baseline. In the short-baseline limit, the mixing matrix in a 3+1 model reduces to a 2×2 matrix. which does not admit CP violation; in a 3+2 model, in contrast, CP violation is possible even at short baseline.

However, since there is no tension between neutrino and antineutrino data, adding CP violation does not improve the global fit. The main source of tension – namely the fact that explaining ν_e appearance data requires the sterile states to have sizable mixing with muon neutrino, in tension with ν_μ disappearance data – is not resolved by including extra sterile states.

It is therefore clear that vanilla 3+n scenarios are not sufficient to explain the entirety of the short-baseline anomalies. They remain a viable option to explain some of them, though.

CONCLUSION

In this abstract, simple theoretical models of sterile neutrinos, their searches in experiments on observing neutrino oscillations, and the results were briefly considered. The chapter 1 described how sterile neutrinos with a mass of eV can be introduced into theoretical models of the neutrino sector, in turn, the 2 chapter considered the search for neutrinos by observing the appearance of electron (anti-) neutrinos, as well as by observing the disappearance of electron and muon (anti -)neutrino.

There, in particular the tension between the simplest sterile neutrino scenario and non-anomalous data, in particular from disappearance searches, was highlighted. Given the strong and fairly consistent hints on the one side, and the exclusion of the minimal sterile neutrino models on the other side, it is clear that resolving the anomalies should be a top priority for neutrino physics in the coming years: either, a fundamental discovery with far-reaching consequences will be made, or we will learn important lessons about systematic effects in neutrino experiments, for instance about the subtleties of neutrino – nucleus interactions or about precision predictions for nuclear beta decay spectra. Such lessons would be indispensable for the successful operation of future facilities like DUNE and Hyper-Kamiokande. Fortunately, a large number of new experiments that are already running or are in advanced R&D stages will hopefully achieve this goal.

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