Plan B for particle physics: finding long lived particles at CERN

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Chapter 1

Introduction: Physics beyond the Standard Model

1.1 Standard Model of particle physics

Our current understanding of elementary particles and interactions between them is summarized in the so called *Standard Model of elementary particle physics* (SM). This is a renormalizable quantum field theory, based on the gauge group $SU(3) \times SU(2) \times U(1)$ and includes spin-$\frac{1}{2}$ particles (matter), spin-1 particles (mediators of interaction) and a spin-0 Higgs boson, see Fig. 1.1. It contains three fermionic families (or *flavors*) of matter and one Higgs $SU(2)$ doublet (see e.g. [1, 2] for reviews). The Standard Model has been developed during most of the XXth century, and its predictions have been tested and confirmed by numerous experiments. The Large Hadron Collider's runs at 7 and 8 TeV culminated in the discovery of a Higgs boson-like particle with a mass of about 126 GeV – the last critical Standard Model component [3–7]. Thus, for the first time, we are in the situation when all the particles needed to explain the results of all previous accelerator experiments, have been found. At the same time, no significant deviations from the Standard Model were found in direct or in indirect searches for new physics (see e.g. the summary of the recent search results in [8–27] and most up-to-date information at [28–31]). For this particular value of the Higgs mass the Standard Model remains mathematically consistent and valid as an effective field theory up to a very high energy scale, possibly all the way to the scale of quantum gravity, the Planck scale [32–36].

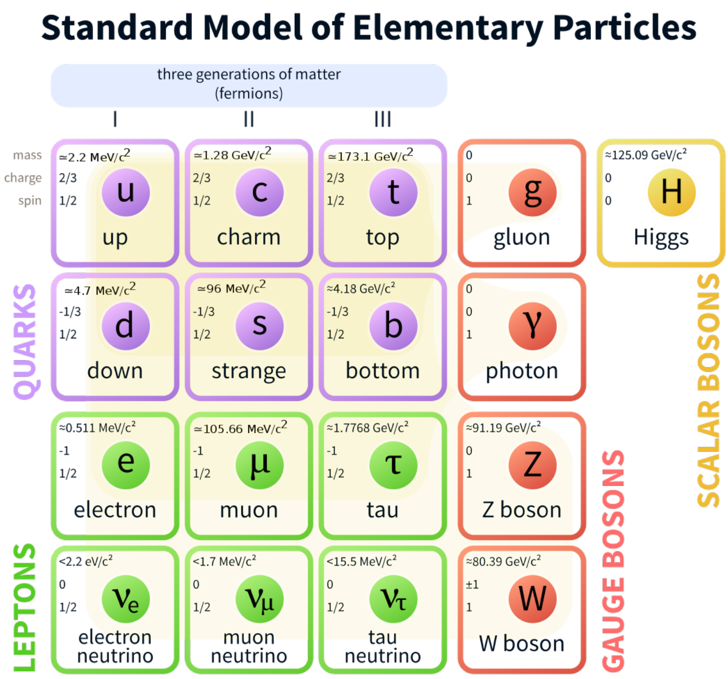

Figure 1.1: The Standard Model of particle physics. Matter particles (left 3 columns) are organised in three *generations* (or *flavors*). Gauge bosons are mediators of the strong (gluons), weak (W and Z bosons) and electromagnetic (photons) interactions. The Higgs mechanism is responsible for providing masses to fermions, and to W and Z bosons. The Higgs boson is the manifestation of this mechanism. Credits: Wikipedia

1.2 Beyond the Standard Model

However, it is clear that the SM is not a complete theory. It fails to explain a number of *observed phenomena* in particle physics, astrophysics and cosmology. These major unsolved challenges are commonly known as "beyond the Standard Model" problems:

Beyond-the-Standard-Model (=BSM) problems

- ▷ **Neutrino masses and oscillations**: what makes neutrinos disappear and then re-appear in a different form? Why do neutrinos have mass?
- ▷ **Baryon asymmetry of the Universe** (BAU): what mechanism has created the (tiny) matter-antimatter disbalance in the early Universe?
- ▷ **Dark Matter** (DM) : what is the nature of the most prevalent kind of matter in our Universe?
- ▷ **Initial conditions problem**: What is the origin of the initial state to which we trace back the evolution of the Universe? In particular, if the initial state was created during a stage of accelerated expansion (**cosmological inflation**), what was driving it?

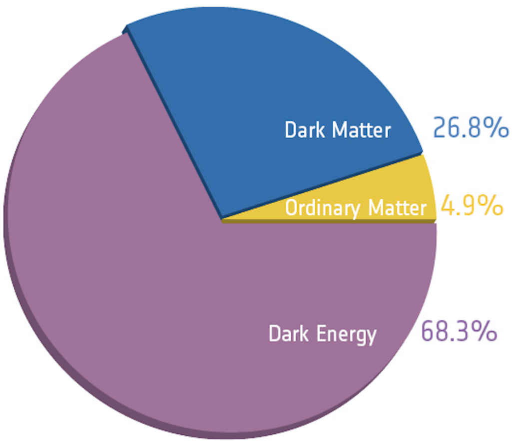

Figure 1.2: Composition of the Universe based on the precise measuremsnts of the Cosmic Microwave Background anisotropies by the Planck collaboration [37]. Credits: European Space Agency and Planck collaboration.

Some yet unknown particles or interactions would be needed to explain these puzzles and to answer these questions. But in that case, why haven't they yet been observed?

1.2.1 Dark matter

Using the laws of particle physics, combined with Einstein's gravity we can model the evolution of the Universe and see in much detail how its current state has emerged from very simple initial conditions. In the modern era of *precision cosmology* detailed predictions of this picture are confirmed with high accuracy using astronomical observations of various types [37]. Ironically, this success revealed one of the greatest mysteries of modern science: 95% of the total energy density of our Universe is composed of entities of unknown nature, see Fig. 1.2. In particular, we see that most of the matter in the Universe does not emit any light – **dark matter**. Indeed, numerous independent tracers of the gravitational potential (observations of the motion of stars in galaxies and galaxies in clusters; emissions from hot ionized gas in galaxy groups and clusters; 21 cm line in galaxies; both weak and strong gravitational lensing measurements) demonstrate that the dynamics of galaxies and galaxy clusters cannot be explained by the Newtonian potential created by visible matter only. Moreover, cosmological data (analysis of the cosmic microwave background anisotropies and of the statistics of galaxy number counts) show that the large scale structure of the Universe started to develop much before the decoupling of photons at the time of recombination of hydrogen and, therefore, much before ordinary matter could start clustering (for reviews see e.g. [38–40]). This body of evidence points at the existence of a new substance, distributed in objects of all scales and providing a contribution to the total energy density of the Universe at the level of about 25%. Various attempts to explain this phenomenon by the presence of macroscopic compact objects (such as, for example, old stars) or by modifications of the laws of gravity (or of dynamics) failed to provide a consistent description of all the above phenomena [41]. Therefore, a microscopic origin of the dark matter phenomenon (i.e., a new particle or particles) remains the most plausible hypothesis.

Neutrinos are the only electrically neutral and long-lived particles in the Standard Model. As the experiments show that neutrinos have mass, they could play the role of dark matter particles. Neutrinos are involved in weak interactions that keep these particles in the early Universe in thermal equilibrium down to temperatures of a few MeV. At lower temperatures, the interaction rate of weak reactions drops below the expansion rate of the Universe and neutrinos "freeze out" from the equilibrium. Therefore, a background of relic neutrinos was created just before primordial nucleosynthesis took off. As the interaction strength and, therefore, the decoupling temperature and concentration of these particles are known, their present day density is fully defined by the sum of the masses for all neutrino flavors. To constitute all of dark matter (DM), this mass should be about 11.5 eV (see e.g. [42]). Clearly, this mass is in conflict with the existing experimental bounds for neutrino mass: measurements of the electron spectrum of β-decay put the combination of neutrino masses below 2 eV [43], while from the cosmological data one can infer an upper bound of the sum of neutrino masses to be 0.58 eV at 95% confidence level [37]. The fact that SM neutrinos could not constitute 100% of DM follows also from the study of the phase space density of DM dominated objects that should not exceed the density of a degenerate Fermi gas: Fermionic particles could play the role of DM in dwarf galaxies only if their mass is above a few hundred eV (the so-called 'Tremaine-Gunn bound' [44], for review see [45] and references therein) and in galaxies, if their mass is tens of eV. Moreover, as the mass of neutrinos is much smaller than their decoupling temperature, they decouple relativistically and become non-relativistic only deeply in the matter-dominated epoch ("*hot dark matter*"). For such type of dark matter the history of structure formation would be very different and the Universe would look rather differently nowadays [46]. All these strong arguments prove convincingly that *the dominant fraction of dark matter* can not be made of the Standard Model neutrinos and therefore *the Standard Model of elementary particles does not contain a viable DM candidate.*

1.2.2 Matter-antimatter asymmetry of the Universe

One of the most important arguments for the existence of physics beyond the Standard Model is the baryon asymmetry of the Universe (BAU).

Historically, the first antimatter particle – positron – was observed in the cosmic rays [47]. Nevertheless, a wide range of observations: γ-ray spectra [48], measurements of the Cosmic Microwave Background(CMB) [49] and of the primordial abundances of light elements [50–52] indicate that there is only primordial (baryonic) matter in the Universe, see e.g. [53–55]. The best current determination of the baryon minus anti-baryon number density n_B, normalised to the entropy density s, is from the CMB (PLANCK [49]):

$$\frac{n_B}{s} = (8.59 \pm 0.13) \times 10^{-11}. \tag{1.2.1}$$

At the same time, all observed antimatter around us is consistent with being produced in the interaction of cosmic rays with the interstellar medium [56, 57]. So we clearly observe asymmetry between matter and antimatter.

It is unlikely that the Universe was born into an asymmetric state (in particular, because inflation – the period of exponentially fast acceleration – would dilute any primordial charge, including baryon charge, into an exponentially small quantity). One can think of a mechanism that would separate baryons from antibaryons in a baryon-symmetric Universe on scales of the order of the observable Universe today, but to invent such mechanism is at least as difficult as generating an asymmetry in the early Universe. Therefore, we will concentrate on the latter possibility.

In order to generate an asymmetry between baryons and anti-baryons in the early Universe, one needs to fulfill three conditions, refered to as "Sakharov conditions" [58]:

1. **A process that violates baryon number**

   Clearly, to evolve from a state with zero baryonic number $B = 0$, to a state with $B \neq 0$, requires non-conservation of the baryon number. It is simple to add a baryon-number-violating process to the Standard Model. However, such a process would lead to the decay of the proton (through a process such as $p \to \pi^0 + e^+$). The current limit on the proton lifetime [59], $\tau_p > 8.2 \times 10^{33}$ years for the process $p \to \pi_0 e^+$, puts severe constraints on such models.

   As it turns out, there is a non-perturbative process in the Standard Model that enables this violation at high temperature. Indeed, as it has been pointed out in [60], both baryon and lepton number symmetries are *anomalous* (as a consequence of chiral anomaly in the presence of $SU(2)$ gauge fields). At zero temperature with W-bosons being massive particles, the probability of such a process is exponentially suppressed. However, this rate is relatively fast before [61] and during the Electroweak Phase Transition [62].

   Thus, *the first Sakharov condition is fulfilled in the Standard Model*

2. C **and** CP **violation** (and violation of any other discrete symmetry that commutes with the Hamiltonian, but anti-commutes with the baryon number).

Particles and anti-particles must behave differently — otherwise, although there were processes that violate baryon number, particles and antiparticles could simultaneously use the B violation with equal rate and opposite total amount to make cancelling B and anti-B asymmetries. Since baryon number anticommutes with C and CP, the violation of both is required.

Charge symmetry C is broken in the Standard Model as a consequence of maximal parity breaking by the weak interactions [63]. Therefore, in any process that preserves CP, the charge symmetry would be broken. Moreover, CP violations have also been observed in the Standard Model, for example in the systems of kaons [64–66], B-mesons [67, 68], see [69] for review. However, the resulting CP-violation is too small to account for baryogenesis [70–72].

Thus, *although the necessary ingredients of the second Sakharov condition are present in the Standard Model, their values are, probably, too small to account for the observed baryon asymmetry.*

3. **Departure from thermal equilibrium**

   The above processes should be out of thermal equilibrium. Indeed, in thermal equilibrium, there are no asymmetries in unconserved quantum numbers, and B is not conserved by the first Sakharov condition. The only conserved global charges of the Standard Model are $B/3 - L_\alpha$ (where L_α is lepton number for generation $\alpha \in (e, \mu, \tau)$). If the corresponding chemical potentials are zero (as they should be in the absence of initial asymmetry), they will remain zero in thermal equilibrium. So generating the BAU is a dynamical process; phase transitions and the expansion of the Universe are sources of non-equilibrium.

   It has been shown that the non-equilibrium first order electroweak phase transition, necessary for baryogenesis in the Standard Model can only occur if the Higgs mass is smaller than $M_H \simeq 72$ GeV [73, 74] (see also [75]). Therefore, *the third Sakharov condition is not satisfied in the Standard Model given the mass of the Higgs particle. This suggests that generation of baryon-antibaryon asymmetry in the SM can not arise.*

Many extensions of the Standard Model are capable of incorporating the baryon asymmetry of the Universe (see e.g. [76–79]). In particular, BAU can be generated in classes of models with heavy neutral leptons. The idea of this scenario, called *leptogenesis* [80], has developed since 1980s (see reviews [81–85] and refs. therein). In particular, it was found that the Majorana mass scale of right-handed neutrinos can be as low as a TeV [86, 87] or even a GeV [88–90] (reaching all the way down to hundreds of MeV [91, 92]) and thus providing hope to probe the leptogenesis scenario directly at particle physics laboratories. Another possibility is to generate Baryon Asymmetry of the Universe using so-called hidden-sector particles, connected to the

SM through scalar portal. There are many other mechanisms for generating a baryon asymmetry of the Universe unrelated to the heavy neutral leptons or scalar portal (for reviews see, e.g. [77, 78, 93–97]). We will not discuss them here further.

1.2.3 Neutrino masses and oscillations

All massive fermions have left and right polarizations ("chiralities"). Neutrinos are the only exception to this rule. In the Standard Model neutrinos are massless and only left-chiral neutrino states participate in the weak interaction [98] – manifestation of the parity violation in the Standard Model [63]. However, starting from the famous "Solar neutrino problem" [99] numerous experiments have convincingly established that the SM neutrinos are massive and can experience transition between flavors (see [100] for the historical account or [101] for the most up-to-date status), unlike charged leptons – a quantum phenomenon, known as *neutrino oscillations*, first predicted in [102]. Neutrinos are produced and detected via weak processes; therefore, by definition, they are produced or detected as *flavor* (or *charge*) states ν_α (i.e. the states that couple to the e, μ and τ leptons, respectively). However, such states turn out to be a linear combination of the states ν_i, $(i = 1, 2, \dots)$ that obey the Klein-Gordon equation $(\Box + m_i^2)\nu_i = 0$ with the observed mass splittings [101]

$$\Delta m_\odot^2 \equiv \Delta m_{21}^2 = 7.55^{+0.20}_{-0.16} \cdot 10^{-5}\ \mathrm{eV}^2,\ \ \Delta m_{\rm atm}^2 \equiv \Delta m_{31}^2 = \begin{cases} 2.50^{+0.03}_{-0.03} \cdot 10^{-3}\ \mathrm{eV}^2 \ (\mathrm{NO}) \\ 2.42^{+0.03}_{-0.04} \cdot 10^{-3}\ \mathrm{eV}^2 \ (\mathrm{IO}) \end{cases}, \tag{1.2.2}$$

where $\Delta m_{ij}^2 \equiv |m_i^2 - m_j^2|$ and NO and IO mean normal and inverted neutrino mass ordering correspondingly. The transformation to this *propagation* or *mass basis* is given by the famous Pontecorvo-Maki-Nakagawa-Sakata matrix [103] that generically depends on three Euler angles, $(\theta_{12}, \theta_{13}, \theta_{23})$, a Dirac CP violating phase, δ, and two more Majorana phases (α_1, α_2), in case neutrinos are Majorana particles:

$$\begin{pmatrix} \nu_e \\ \nu_\mu \\ \nu_\tau \end{pmatrix} = U_{\rm PMNS}(\theta_{12}, \theta_{13}, \theta_{23}, \delta, \alpha_1, \alpha_2) \begin{pmatrix} \nu_1 \\ \nu_2 \\ \nu_3 \end{pmatrix}. \tag{1.2.3}$$

where

$$U_{\rm PMNS} = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -c_{23}s_{12} - s_{23}c_{12}s_{13}e^{i\delta} & c_{23}c_{12} - s_{23}s_{12}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{23}s_{12} - c_{23}c_{12}s_{13}e^{i\delta} & -s_{23}c_{12} - c_{23}s_{12}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix} \times \begin{pmatrix} 1 & 0 & 0 \\ & e^{i\alpha_1} & 0 \\ 0 & 0 & e^{i\alpha_2} \end{pmatrix}$$

with $s_{ij} = \sin\theta_{ij}$, $c_{ij} = \cos\theta_{ij}$.

At first sight the phenomenology of neutrino masses and mixing can be realised purely within the Standard Model. Indeed, one can write a Majorana mass term by making a Dirac spinor χ_α out of two (chiral) neutrino fields

$$\chi_\alpha = \frac{\nu_\alpha + \nu_\alpha^c}{\sqrt{2}},$$

where ν_α^c is a charged conjugated neutrino field ν_α. The mass term then is simply $m\bar{\chi}_\alpha\chi_\alpha$. Similarly, all the mass/mixing phenomenology can be described by the 3×3 matrix

$$\text{Neutrino mass term} = \sum_{\alpha,\beta} c_{\alpha\beta}\frac{v^2}{\Lambda}\bar{\chi}_\alpha\chi_\beta \tag{1.2.4}$$

Here v is the vacuum expectation value of the Higgs field and Λ is a parameter having the dimension of mass; $c_{\alpha\beta}$ is a dimensionless 3×3 matrix that encodes phenomenology of mixing and the hierarchy of neutrino masses. However, expression (1.2.4) is valid only in the broken (Higgs) phase of the Standard Model. Gauge invariant definition of the neutrino field $\nu_\alpha = (\tilde{H}\cdot L_\alpha)$ where L_α is the left lepton doublet (α is the flavor index $\alpha = \{e,\mu,\tau\}$); H is the Higgs doublet and $\tilde{H}_a = \epsilon_{ab}H_b$. As a result neutrino mass and mixing term can be written in the gauge invariant way only by introducing the so-called "Weinberg operator" [104]:

$$\Delta\mathcal{L}_{\text{osc}} = c_{\alpha\beta}\frac{(\bar{L}_\alpha^c\cdot\tilde{H})(\tilde{H}\cdot L_\beta)}{\Lambda} \tag{1.2.5}$$

where L_α^c is a charge-conjugation of the left lepton doublet, L_α. $\Delta\mathcal{L}_{\text{osc}}$ is a non-renormalizable dimension-5 operator. Generation of such operator clearly requires adding new particles to SM (for a review see e.g. [105]).

1.3 Portals to new physics

New particles that interact with SM particles may be directly responsible for some of the BSM phenomena or can serve as mediators (or "*portals*"), coupling to states in the "hidden sectors" and at the same time interacting with the Standard Model particles. Such portals can be renormalizable (mass dimension ≤ 4) or be realized as higher-dimensional operators suppressed by the dimensionful couplings Λ^{-n}, with Λ being the energy scale of the hidden sector. In the latter case the most promising are the dimension 5 operators that are not suppressed too much by the energy scale Λ.

In the Standard Model there can be only three *renormalizable* portals:

– the scalar portal that couples the gauge singlet scalar S to the $H^\dagger H$ term constructed of the Higgs doublet field H_a, $a=1,2$ with the portal Lagrangian

$$\mathcal{L}_{\text{scalar portal}} = \alpha_1 S H^\dagger H + \alpha S^2 H^\dagger H; \tag{1.3.1}$$

– the neutrino portal that couples the new gauge singlet fermion N to the $\epsilon_{ab}\bar{L}_a H_b$ term, where L_a is the SU(2) lepton doublet and ϵ_{ab} is absolutely antisymmetric tensor in 2 dimensions,

$$\mathcal{L}_{\text{neutrino portal}} = F_\ell(\epsilon_{ab}\bar{L}_{\ell,a}H_b)N, \tag{1.3.2}$$

with $\ell = e, \mu, \tau$;

– the vector portal that couples the field strength of a new U(1) field A'_μ to the U(1) hyperfield field strength B_μ,

$$\mathcal{L}_{\text{vector portal}} = \frac{\epsilon}{2} F'_{\mu\nu} F^{\mu\nu}, \tag{1.3.3}$$

where $F'_{\mu\nu} = \partial_\mu A'_\nu - \partial_\nu A'_\mu$ and $F_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu$;

Among the higher dimension portals it is worth to mention the very popular axion portal that is coupled to the Standard Model through interaction with gauge fields,

$$\mathcal{L}_{\text{axion portal}} = \frac{g_i}{\Lambda} a F_{i,\mu\nu} \tilde{F}^{i,\mu\nu}, \tag{1.3.4}$$

where a is a new pseudo Nambu-Goldstone boson, $F_{i,\mu\nu}$ is a field strength tensor of the SM gauge bosons, $\tilde{F}_{\mu\nu} = \frac{1}{2}\varepsilon_{\mu\nu\alpha\beta}F^{\alpha\beta}$ and g_i is a corresponding coupling constant.

These portals have different physical motivation and different phenomenology. To illustrate this, below we will concentrate on the scalar and neutrino portals.

1.4 Intensity and energy frontiers in searches for new physics

The development of the Standard Model has come to an end with the confirmation of one of its most important predictions – the discovery of the Higgs boson. The quest for new particles has not ended, however. The observed but unexplained phenomena in particle physics and cosmology that we discussed in previous section indicate that other particles exist in the Universe. It is possible that these particles evaded detection so far because they are too heavy to be created at accelerators. Alternatively, some of the hypothetical particles can be sufficiently light (lighter than the Higgs or W-boson), but interact very weakly with the Standard Model sector (we will use the term *feeble interaction* to distinguish this from the weak interactions of the Standard Model). In order to explore this latter possibility, the particle physics community is turning its attention to so-called *Intensity Frontier experiments*, rather than *Energy Frontier experiments* like at the LHC or Tevatron (Figure 1.3). Such experiments aim to create high-intensity particle beams and use large detectors to search for the rare interactions of feebly interacting hypothetical particles. Several Intensity Frontiers experiments have been proposed in recent years: DUNE [106], NA62 [107–109], SHiP [110, 111], etc.

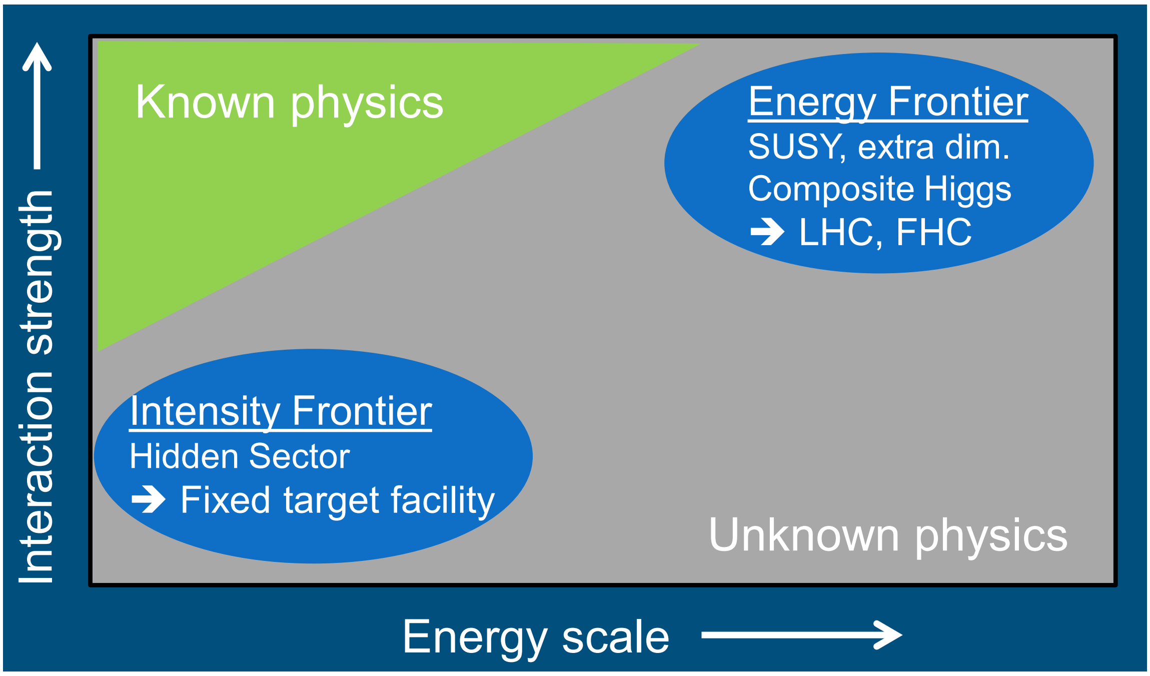

Figure 1.3: New physics that can be explored in intensity frontier experiments and its complimentarity with the energy frontier. Figure from [110].

Although LHC is the flagship of the Energy Frontier exploration, its high luminosity (especially in Run 3 and beyond) means that huge numbers of heavy, flavored mesons and vector bosons are created. This opens the possibility to supplement the High Luminosity phase of the LHC with companion experiments directed at the Intensity Frontier. Several such experiments have been proposed: CODEX-b [112], MATHUSLA [113, 114] and FASER [115, 116]. Therefore, it is important to compare LHC companion experiments to the proposed and specialized Intensity Frontier experiments. In Sec. 5 we will study and compare sensitivities of a specialized Intensity Frontier experiment (SHiP) and one of the most promising LHC companion experiments (MATHUSLA).

1.5 Neutrino portal

Heavy neutral leptons or sterile neutrinos N_I, $I = 1, \ldots, \mathcal{N}$ are singlets with respect to the SM gauge group and couple to the gauge-invariant combination $(\bar{L}^c_\alpha \cdot \tilde{H})$ (where L_α, $\alpha = 1, \ldots, 3$, are SM lepton doublets and $\tilde{H}_i = \varepsilon_{ij} H^*_j$ is the conjugated SM Higgs doublet) as follows

$$\mathcal{L}^c_{\text{neutrino portal}} = F_{\alpha I}(\bar{L}_\alpha \cdot \tilde{H}) N_I + h.c.\ , \tag{1.5.1}$$

with $F_{\alpha I}$ denoting dimensionless Yukawa couplings. The name "sterile neutrino" stems from the fact that the interaction (1.5.1) fixes the SM gauge charges of N_I

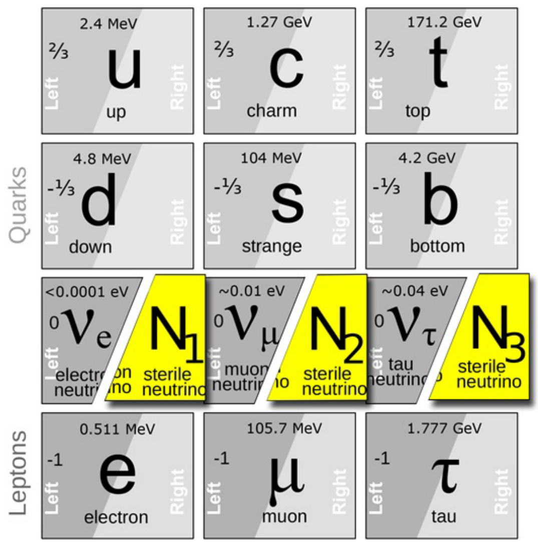

Figure 1.4: Three generations of Standard Model fermions with the three right-chiral sterile neutrinos N_1, N_2, N_3. Figure from [110].

to be zero. Also, as only the right-chiral part of the N field interacts with the active neutrino field the sterile neutrinos are often called *right-handed neutrinos*. Due to electroweak symmetry breaking, the SM Higgs field has a nonzero vacuum expectation value v and the interaction (1.5.1) provides heavy neutral leptons and SM (or *active*) neutrinos — with the mixing mass term ($v = 246\,\text{GeV}$)

$$M^D_{\alpha I} \equiv F_{\alpha I} v/\sqrt{2}\,.$$

The neutrino portal could be a connection between the SM and some hidden sector. But in the minimal realization of the hidden sector, the truly neutral nature of N allows one to introduce for it a Majorana mass term, consistent with the SM gauge invariance

$$\mathcal{L}_{\text{HNL}} = i\bar{N}_I \partial\!\!\!/ N_I + \left(M^D_{\alpha I}\bar{\nu}_\alpha N_I - \frac{M_{N,I}}{2}\bar{N}^c_I N_I + h.c. \right). \tag{1.5.2}$$

In this case the only parameters of the "hidden sector" are the masses of the sterile neutrinos $M_{N,I}$, $I = 1 \ldots \mathcal{N}$. If $\mathcal{N} = 3$ the number of sterile neutrinos coincides with the number of active neutrinos and every active neutrino gets its right-handed partner as all other fermions of the Standard Model. This case seems the most attractive possibility as the structure of the SM with three generations is restored, see Fig. 1.4.

The mass eigenstates of the active-plus-sterile sector are mixtures of ν and N, with small mixing angles and large splitting between mass scales of sterile and active neutrinos. The heavy mass eigenstates are "almost sterile neutrinos" while light mass eigenstates are "almost active neutrinos". In what follows we keep the same terminology for the mass states as for the gauge states. As a result of mixing, HNL couples to the SM fields in the same way as active neutrinos,

$$\mathcal{L}_{int} = \frac{g}{2\sqrt{2}} W^+_\mu \overline{N^c_I} \sum_\alpha U^*_{\alpha I} \gamma^\mu (1-\gamma_5) \ell^-_\alpha + \frac{g}{4\cos\theta_W} Z_\mu \overline{N^c_I} \sum_\alpha U^*_{\alpha I} \gamma^\mu (1-\gamma_5) \nu_\alpha + \text{h.c.}\ , \tag{1.5.3}$$

except the coupling is strongly suppressed by the small *mixing angles*

$$U_{\alpha I} = M^D_{\alpha I} M^{-1}_{N,I}. \tag{1.5.4}$$

Although, the neutrino portal could be a mediator to the large hidden sector, it is interesting to note that even the minimal realization discussed above is sufficient to solve the BSM problems. We will discuss this point in the next sections.

1.5.1 HNLs and neutrino masses, seesaw formula

If neutrinos are Majorana particles, the phenomenology of their oscillations can be described via the Weinberg operator (1.2.5). This operator of mass dimension 5 can be resolved in many different ways (for a review see e.g. [105]), but the simplest way is by the introducing a neutral fermion, i.e., sterile neutrino. In the Higgs phase, the combination of the neutrino portal interactions (1.5.2) lead to the mixing between sterile and active neutrinos. As a result *flavor eigenstates* do not coincide with the *mass eigenstates*. The latter are obtained by diagonalizing the matrix

$$\mathcal{M}_{\nu,N} = \left(\begin{array}{c|c} 0 & m_D \\ \hline m_D^T & M_I \end{array} \right), \tag{1.5.5}$$

where m_D is $3 \times \mathcal{N}$ Dirac mass matrix, $(m_D)_{\alpha I} = F_{\alpha I} v$, $v = \sqrt{2}\langle H \rangle$ and M_I is $\mathcal{N} \times \mathcal{N}$ matrix of Majorana masses. In the limit $m_D \ll M_I$, one can easily see that the Lagrangian (1.5.2), indeed, leads to the Weinberg operator (1.2.5) for neutrinos with

$$\frac{c_{\alpha\beta} v^2}{\Lambda} \equiv (\mathcal{M}_\nu)_{\alpha\beta} = -\sum_I (m_D)_{\alpha I} \frac{1}{M_I} (m_D)_{\beta I}. \tag{1.5.6}$$

The smallness of the Dirac mass term as compared to the Majorana masses M_I means that the active neutrino masses (3 eigenvalues of the matrix $(\mathcal{M}_\nu)_{\alpha\beta}$) become much smaller than the scale M_I and the electroweak scale. This mechanism is therefore known as the *seesaw mechanism* [117–120], see also [121] and refs. therein.[1]

[1]This mechanism is often called *Type-I seesaw mechanism* because there are other ways to "resolve" the Weinberg operator (see e.g. [105, 122]). For example, in the *Type-II seesaw mechanism* an extra SU(2) triplet scalar is introduced [123–126], in the *type-III seesaw mechanism* an extra fermion in the adjoint of SU(2) is added to the model [127].

Adding $\mathcal{N}$ new particles N_I to the Lagrangian $\mathcal{L}_{\rm SM}$ adds

$$N_{\rm parameters} = 7 \times \mathcal{N} - 3 \tag{1.5.7}$$

new parameters to the Lagrangian. These parameters can be chosen as follows: $\mathcal{N}$ real Majorana masses M_I *plus* $3 \times \mathcal{N}$ complex Yukawa couplings $F_{\alpha\, I}$ *minus* 3 phases absorbed in redefinitions of ν_e, ν_μ, ν_τ. The Pontecorvo–Maki–Nakagawa–Sakata matrix (1.2.3) plus three mass eigenstates m_1, m_2, m_3 of the active neutrino sector provide 9 parameters that can be determined experimentally. This shows that one needs $\mathcal{N} \geq 2$ to explain the neutrino oscillations by means of heavy neutral leptons.

To see how neutrino oscillation data limit the HNL parameters, let us first look at the *simplest (unrealistic) toy model with just one HNL*, N_1. In this case only one combination of neutrino flavors becomes massive, and its mass, m_ν, allows us to limit the sum of the Yukawa couplings, $|F_1|$:

$$|F_1|^2 \equiv \sum_\alpha |F_{\alpha 1}|^2 = \left(4.1 \times 10^{-8}\right)^2 \left(\frac{m_\nu}{m_{\rm atm}}\right)\left(\frac{M_1}{1\ \text{GeV}}\right) \qquad \text{(single HNL case)}, \tag{1.5.8}$$

where $m_{\rm atm} = \sqrt{\Delta m^2_{\rm atm}}$, see Eq. (1.2.2). We see that $|F_1|$ becomes smaller as M_1 becomes lighter. On the other hand, the mixing element of HNL in the weak interaction is given by

$$U^2 = 5.0 \times 10^{-11} \left(\frac{m_\nu}{m_{\rm atm}}\right)\left(\frac{1\ \text{GeV}}{M_1}\right) \qquad \text{(single HNL case)} \tag{1.5.9}$$

which is also very small and the dependence on the HNL mass, M_1 is opposite to Eq. (1.5.8).

In order to explain two mass differences $\Delta m^2_{\rm atm}$ and $\Delta m^2_\odot$ (see Eq. (1.2.2)), one needs $\mathcal{N} \geq 2$. As it turns out, in this case, much larger values of $|F|$ and U^2 are possible. Even in the simplest case of *two HNLs having the same mass* M_N, the mixing angle is expressed via one free parameter $X_\omega \geq 1$ (in the notations of [128]) that gives the following equation for $U^2 \equiv \sum_{\alpha,I} U^2_{\alpha I}$ [128–130]

$$U^2 = \frac{\sum_\nu m_\nu}{2M_N}(X_\omega^2 + X_\omega^{-2}) \simeq 5 \times 10^{-10} \kappa \left(\frac{1\ \text{GeV}}{M_N}\right)\left(\frac{X_\omega}{100}\right)^2, \tag{1.5.10}$$

where $\sum_\nu m_\nu = \kappa \times m_{\rm atm}$, i.e. $\kappa \simeq 1$ for normal hierarchy and $\kappa \simeq 2$ for inverted hierarchy. One sees that U^2 can be much larger that the naive estimate given by (1.5.9).

The Lagrangian (1.5.2) remains perturbative with Yukawa coupling $|F_{\alpha I}| \lesssim 1$ which corresponds to the scale $\Lambda \sim 10^{15}$ GeV in the Weinberg operator (1.5.6). Actually, the analysis shows that the theory stays perturbative up to $M_{N,I} \lesssim 10^{16}$ GeV [131]. In the opposite limit $M_{N,I} \to 0$ the neutrinos become massive Dirac fermions and the smallness of their masses is explained by small Yukawa coupling: $F_{\alpha I} \sim 10^{-11}$. The Lagrangian has in this case an exact, global, non-anomalous $U(1)_{B-L}$ symmetry.[2] This symmetry is broken when both the $F_{\alpha I}$ and the $M_{N,I}$ are nonvanishing, a fact that teaches us that any value of $M_{N,I}$ is technically natural [132], as defined by 'tHooft [133].

1.5.2 HNL and baryon asymmetry of the Universe. Leptogenesis

Leptogenesis [80] describes a class of mechanisms, which make use of the L violation that is present in Majorana neutrino mass models. Leptogenesis is defined here to include all scenarios which produce a lepton (anti-)asymmetry via CP-violating out-of-equilibrium processes, and rely on the equilibrium SM non-perturbative B+L violation, to partially transform the lepton deficit into a baryon excess. Leptogenesis therefore occurs before/at the electroweak phase transition. A small advantage is that there are no $\Delta B = 1$ interactions, so no concerns with proton decay. More importantly, a natural way to understand why neutrinos are much lighter than other SM fermions, is to suppose that their masses are Majorana, that is, L-violating. So in such extensions of the SM, the first Sakharov condition comes for free.

Baryon asymmetry of the Universe (BAU) can be generated through leptogenesis using HNLs of different mass scales, from sub-GeV to 10^{15} GeV (see [110] and references therein). In the case of GeV-scale HNLs, the masses of active neutrinos can also generate the Baryon Asymmetry of the Universe via HNL oscillations [88, 89]. Since HNLs at the GeV-scale possess only the Yukawa interaction at the unbroken phase in the Early Universe and its couplings are very suppressed as mentioned above, they can be out of equilibrium state. Then, the CP violation in the production and evolution of HNLs with oscillation effects generates the asymmetry of left-handed leptons which is partially converted into the baryon asymmetry of the Universe.

The successful baryogenesis requires an upper bound on the Yukawa couplings (to avoid fast washout of the baryon asymmetry due to the rapid scatterings of HNLs). This upper bound is more strong for $\mathcal{N} = 2$ case [134] and gets relaxed if more HNLs contribute to the neutrino masses.

1.5.3 HNL and dark matter

The only electrically neutral and long-lived particle in the Standard Model are neutrinos. To constitute all of the DM in the Universe the total mass of the neutrino of all flavors ($m_{\nu_e} + m_{\nu_\mu} + m_{\nu_\tau}$) should be about 11.5 eV (see e.g. [42]). But, if the DM particle is a fermion, the phase-space number density of DM in the faintest galaxies should not exceed the density of a degenerate Fermi gas. As the DM mass density is bounded from below observationally, this puts a lower bound on the mass

[2] This is the only global symmetry, exact at both classical and quantum levels, admitted by the Standard Model Lagrangian $\mathcal{L}_{\rm SM}$.

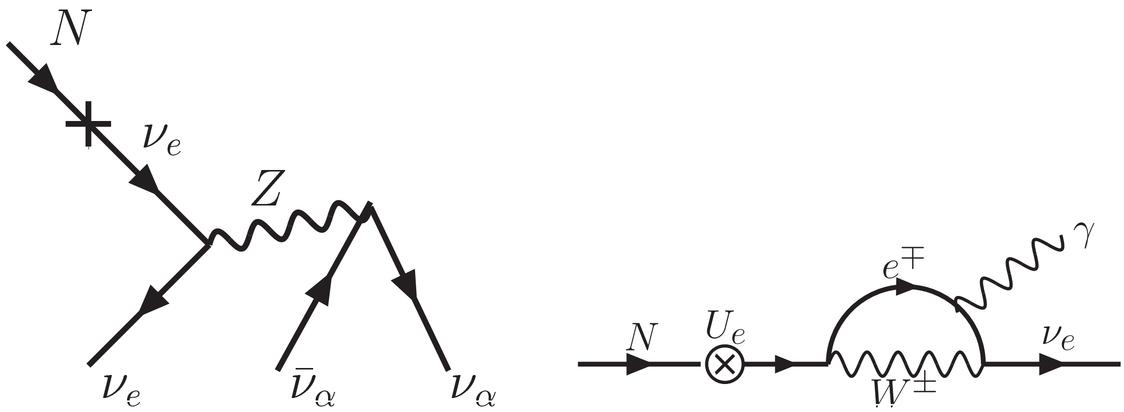

Figure 1.5: Decay channels of the sterile neutrino with the mass below twice the electron mass. Right panel shows radiative decay channel that allows to look for the signal of sterile neutrino dark matter in the spectra of dark matter dominated objects. Figure from [110].

of fermionic DM particles to be a few hundreds of eV (the so-called "Tremaine-Gunn bound" [44]). Therefore, cosmological and astrophysical requirements to neutrino dark matter contradict each other.

The *sterile neutrino* allows to resolve this contradiction. Its interaction with the Standard Model particles is similar to that of the active neutrino, but suppressed by the mixing angles $U_{\alpha,I}$ (see Eq. (1.5.3)). Therefore, the number density of the sterile neutrinos created in the early Universe can be much lower and account for the correct DM abundance with a much larger range of masses of the particle, easily satisfying the Tremaine-Gunn bound. The production through mixing with active neutrinos always contributes to the relic sterile neutrino abundance [135–138]. If a large lepton asymmetry or new particles and fields are present in the model then additional production mechanism are possible [139–143] (see [144–146] for review).

The sterile neutrino has a finite lifetime. It can decay to 3 active (anti)neutrinos and has also subdominant decay channel into a neutrino and a photon (Figure 1.5). To be a dark matter candidate its lifetime should be greater than the lifetime of the Universe. In view of the huge amount of dark matter particles in galaxies it will give rise to a strong signal that would be immediately observed, so the lifetime of sterile neutrino DM has to be at least $3 \cdot 10^6$ larger than the age of the Universe [147, 148].

Taking into account constraints from X-rays and cosmology, the allowed parameter space for the sterile neutrino DM candidate is shown in Fig. 1.6. The black line corresponds to the recently observed unidentified spectral line at the energy $E \sim 3.5$ keV in the stacked X-ray spectra of Andromeda galaxy, Perseus galaxy clusters, stacked galaxy clusters and the Galactic Center of the Milky Way [149–151], which is consistent with predictions for decaying dark matter with a mass $M_N \approx 7.1 \pm 0.1$ keV.

Importantly, the sterile neutrino DM candidate does not contribute to neutrino oscillations and cannot be searched in direct detection experiments [148].

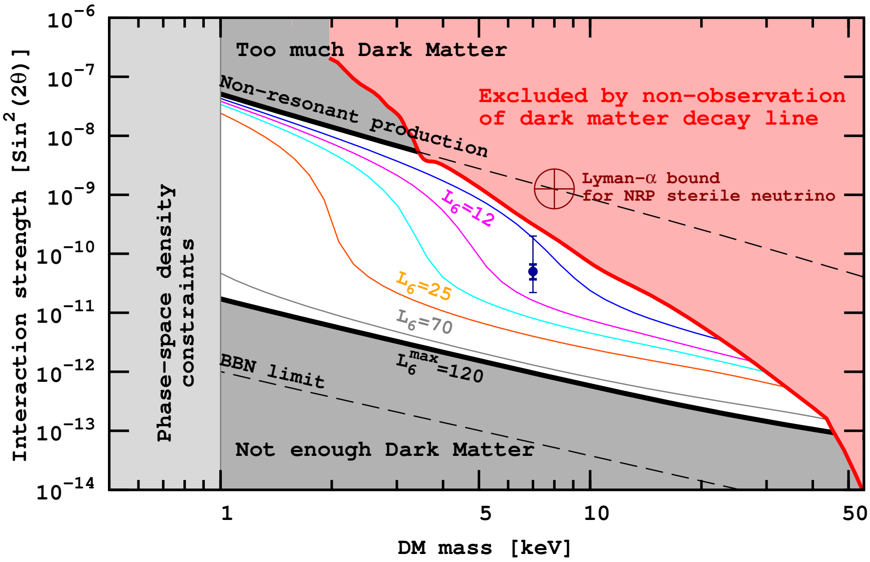

Figure 1.6: The parameter space of sterile neutrino dark matter produced via mixing with the active neutrinos (*unshaded region*). The two thick black lines bounding this region are production curves for non-resonant production [138] (*upper line*, "NRP") and for resonant production (RP) with the maximal lepton asymmetry, attainable in the νMSM (*lower line*, marked "$L_6^{max} = 120$") [90, 134, 140] (L_6 is defined as the ratio of the lepton density to the entropy density times 10^6). The thin coloured curves between these lines represent production curves for different values of lepton asymmetry. The red shaded upper right corner represents X-ray constraints [152–156] (rescaled by a factor of two to account for possible systematic uncertainties in the determination of DM content). The region below 1 keV is ruled out according to the phase-space density arguments [157] (see text for details). The point at ~ 7.1 keV corresponds to the unidentified spectral detected in stacked X-ray spectra of galaxies and galaxy clusters [149, 150]. Thick errorbars are $\pm 1\sigma$ limits on the flux as determine from data. Thin errorbars correspond to the uncertainty in the DM distribution. Figure from [110].

1.5.4 νMSM

Let us discuss here the framework of the so-called νMSM (neutrino Minimal Standard Model). This is a simple extension of the SM by introducing $\mathcal{N}$ right-handed neutrinos N_I ($I = 1, 2, \cdots, \mathcal{N}$) in order to explain the three observational phenomena which cannot be explained by the SM, *i.e.*, the non-zero masses of active neutrinos, the cosmic dark matter, and the baryon asymmetry of the universe. These right-handed neutrinos are introduced with Majorana masses M_I

$$|m_D|_{\alpha I} \ll M_I < \mathcal{O}(10^2)\ \text{GeV}\,, \tag{1.5.11}$$

where the Dirac mass is given by $|m_D|_{\alpha I} = |F_{\alpha I}|\langle H \rangle$. The first inequality, between Dirac and Majorana masses, is imposed by the seesaw mechanism. The scale of Majorana mass for the seesaw mechanism cannot be determined from active neutrino masses and can vary over a wide range (see Section 1.5.1). The possibility, discussed here, is to choose Majorana masses that are comparable to or smaller than the electroweak scale $\mathcal{O}(10^2)$ GeV, so that the masses of HNLs are comparable to, or smaller than, masses of quarks and charged leptons. Interestingly, even when HNLs are lighter than the electroweak scale, sufficient baryon asymmetry can be generated via oscillations, as was mentioned in Section 1.5.2. Even in the minimal option required for explaining two mass differences $\Delta m^2_{\rm atm}$ and $\Delta m^2_{\odot}$ (see Eq. (1.2.2)), say the two HNLs case, the sufficient baryon asymmetry can be generated as demonstrated in Ref. [89].

The two mass scales, $\Delta m^2_{\rm atm}$ and $\Delta m^2_{\odot}$, confirmed by various oscillation experiments, require that there must be at least two massive states of active neutrinos with different mass eigenvalues. This implies that the number of right-handed neutrinos must be equal or larger than two ($\mathcal{N} \geq 2$). Notice that the lightest active neutrino becomes exactly massless for the minimal choice $\mathcal{N} = 2$.

Our DM candidate is a HNL with $\mathcal{O}(10)$ keV mass (see the discussions in Sec. 1.5.3) We might expect that HNLs needed to explain $\Delta m^2_{\rm atm}$ and $\Delta m^2_{\odot}$ also may play the role of dark matter. However, that is impossible for the following reasons. First, it is shown [158] that HNLs that are responsible for the masses of active neutrinos must has sizable Yukawa couplings which would produce too much dark matter particles by the Dodelson-Widrow mechanism [135], so the present abundance would exceed the observed value. In addition, such HNLs cannot be dark matter since they give too much X-rays from their radiative decays [148]. Therefore, we must introduce right-handed neutrino(s) for dark matter, in addition to at least two right-handed neutrinos for active neutrino masses. In this case the number of right-handed neutrino must be $\mathcal{N} \geq 3$.

As a result, the minimal number of right-handed neutrinos explaining the neutrino masses, dark matter, and the baryon asymmetry at the same time equals $\mathcal{N} = 3$. In this case, the HNL N_1 plays a role of dark matter (see Section 1.5.3) and the heavier HNLs N_2 and N_3 are responsible for the seesaw mechanism and baryogenesis. The model with $\mathcal{N} = 3$ introduces 18 new parameters in addition to the parameters of the SM (see Eq. (1.5.7)), which are three Majorana masses M_I and 15 (physical) parameters in the neutrino Yukawa couplings $F_{\alpha I}$. The number of parameters associated with the heavier HNLs N_2 and N_3 is 11. Seven of these are parameters of the active neutrinos (two mass-squared-differences and three mixing angles of the active neutrinos and one Dirac-type phase and one Majorana-type phase in the PMNS matrix), and 4 are parameters of HNLs (their masses $M_{2,3} = M_N \pm \Delta M/2$ and one complex parameter). The residual 7 parameters are for dark matter N_1 (mass of N_1, three mixing elements $|U_{\alpha 1}|$ and three CP violating phases).

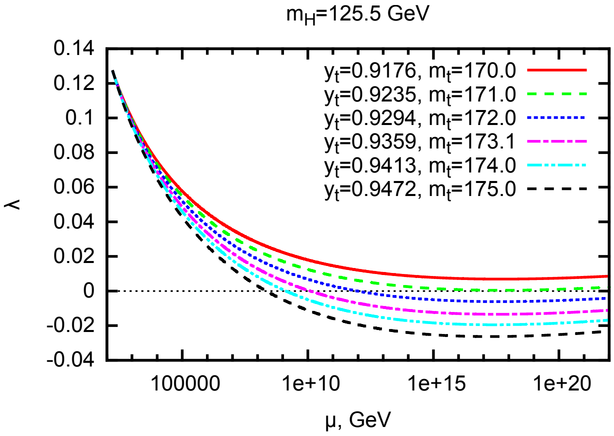

Figure 1.7: Renormalization group running of the Higgs coupling constant λ for the Higgs mass $M_h = 125.7$ GeV and several values of the top quark Yukawa $y_t(\mu = 173.2\text{GeV})$. Figure from [110].

One important consequence of this model is that the lightest active neutrino is lighter than $\mathcal{O}(10^{-5})$ eV [148, 158]. This comes from the fact that the dark matter HNL is allowed to give a tiny contribution to the seesaw. Therefore, the mass eigenvalues of heavier active neutrinos can be identified from Δm^2_{atm} and $\Delta m^2_{\odot}$ (the ordering of masses is still unknown).

It is important to note that the SM plus νMSM which introduces new feebly interacting particles below electroweak scale can be consistent up to the Planck scale. The *inflation* can be incorporated to νMSM (and in the Standard Model) through the non-minimal coupling of the SM Higgs field to the gravitational Ricci scalar R [159], $\frac{\xi h^2}{2}R$. The predictions of this inflation model are consistent with the recent Planck data [160].

The most minimal way to describe *the accelerated expansion of the Universe* at the present epoch in any theory, including the νMSM, is simply to add the cosmological constant Λ. The extremely small value of Λ remains without explanation (this is exactly the cosmological constant problem), but this "solution" fits all the cosmological data.

Finally, let us discuss the *vacuum stability* and the νMSM. The experimentally measured values of the Higgs mass and of the top quark's Yukawa coupling lead to quite a peculiar behaviour of the scalar self-coupling λ in the SM (and also in the νMSM, since HNLs couplings are small and can be safely neglected), see Fig. 1.7. This constant decreases with energy, reaches its minimum at energies close to the Planck scale, and then increases [161]. Depending on the values of the Higgs mass and top quark Yukawa coupling allowed by experiments, λ can cross zero at energies as small as 10^{10} GeV and remain negative around the Planck scale, or be positive at any energy, or just touch zero at an energy scale close to the Planck one [32–34, 162–

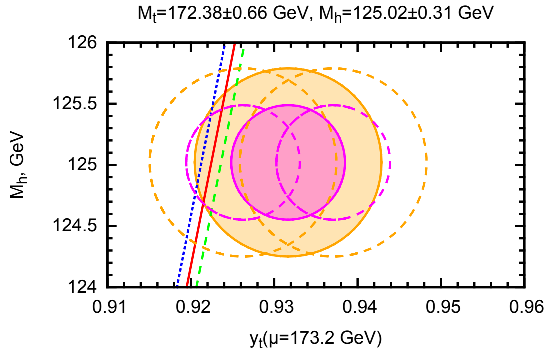

Figure 1.8: The figure shows the borderline between the regions of absolute stability and metastability of the SM vacuum on the plane of the Higgs boson mass and top quark Yukawa coupling in the $\overline{\text{MS}}$ scheme taken at $\mu = 173.2\,\text{GeV}$. The diagonal line stands for the critical value of the top Yukawa coupling y_t^{crit} as a function of the Higgs mass and the dashed lines account for the uncertainty associated to the error in the strong coupling constant α_s. The SM vacuum is absolutely stable to the left of these lines and metastable to the right. The filled ellipses correspond to experimental values of y_t extracted from the latest CMS determination [165] of the Monte-Carlo top quark mass $M_t = 172.38 \pm 0.10\,(\text{stat}) \pm 0.65\,(\text{syst})\,\text{GeV}$, if this is identified with the pole mass. The Higgs mass $M_h = 125.02 \pm 0.27\,(\text{stat}) \pm 0.15\,(\text{syst})\,\text{GeV}$ is taken from CMS measurements [166]. Dashed ellipses encode the shifts associated to the ambiguous relation between pole and Monte Carlo masses. See [36] and references therein for more discussion. Figure from [110].

164]. The behaviour of the Higgs self-coupling is closely related to the problem of stability of the SM vacuum: if λ is negative in some domain of energies, the effective potential of the scalar field without gravity develops a second, deeper minimum at the scalar field values of the order of Planck scale. In this case the SM vacuum becomes metastable. The situation is uncertain: the SM vacuum can be absolutely stable or metastable within experimental and theoretical error-bars, see Fig. 1.8.

If the SM vacuum is indeed metastable, there is a danger of transition from our vacuum to another, unwanted one, with Planck scale physics. Though the life-time of the SM vacuum exceeds the age of the Universe is by many orders of magnitude [167], it may happen that the Universe evolution during or after inflation could drive the system out of our vacuum. To prevent this, some kind of new physics should intervene to save the Universe from collapse in the Planck vacuum. In [35] has been demonstrated that the specific threshold effects in the SM (and in the νMSM) with non-minimal coupling to gravity at the energy scale M_P/ξ may lead to relaxation of the system in the SM metastable vacuum. No new heavy particles are needed for this to happen.

1.5.5 Summary

The motivation for the existence of a neutrino portal steams from both experiment and theory and points to the existence of heavy neutral leptons. These particles may play an essential role in cosmology, providing a dark matter candidate and producing baryon asymmetry of the Universe. In neutrino physics they provide the source of neutrino masses and mixings. If the masses of HNLs are smaller than ~ 5 GeV, it will be possible to search for them in decays of heavy mesons carrying strangeness, charm or beauty, created in high-intensity fixed-target experiments such as SHiP. Heavier HNL's can be searched for in collider experiments at the LHC and in future experimental facilities like FCC-ee [168]. Virtual HNLs lead to lepton flavor number violation that can potentially be seen in the processes like $\mu \to 3e, \mu \to e\gamma$ or $\tau \to 3\mu$. They also generically imply the lepton number violation that can manifest itself in neutrino-less double beta decays. It goes without saying that the discovery of such particles or indirect indication of their existence would revolutionise our understanding of particle physics and cosmology, whereas constraining their properties would help to elucidate various ideas on physics beyond of the Standard Model.

1.6 Scalar portal

The Higgs boson was recently found at the LHC at CERN [169–172] directly confirming the existence of the fundamental scalars in nature. Therefore, the idea of another fundamental scalar is viable and well-motivated. Such scalar particle appears in many extensions of the Standard Model. If the new particle has no SM charges it can be light and naturally obtain very suppressed couplings to other SM particles. This particle could be a portal between SM and hidden sector that is well motivated by such experimental observations as baryonic asymmetry of the Universe (see e.g. [76] for a review), dark matter [173–175] and the hierarchy problem [176–180].

If one considers a complex dark sector, many new avenues open up for new processes, such as dark baryogengesis [181], see Fig. 1.9. These processes appear naturally within the content of Asymmetric Dark Matter. The details of the baryogenesis depends on the properties of the dark sector and will not be discussed here.

Another interesting application of the additional scalar particle is to consider it to be an inflaton. In this case, the new scalar particle could be still light and, what is especially intriguing, its parameter space can be probed by high-intensity experiments [182, 183].

1.6.1 Scalar as a mediator between DM and the SM

Light scalars may certainly provide an interesting connection to the puzzle of dark matter (DM). It is quite conceivable that dark matter resides in the form of a SM singlet particle which is protected against decay by a (discrete) symmetry. In this case, a light scalar could mediate the interactions between DM and the SM.

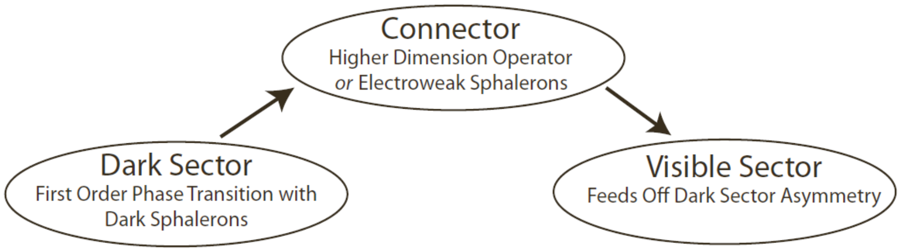

Figure 1.9: In a Hidden Valley with asymmetric dark matter, the dark baryon asymmetry created in a first-order phase transition in a Hidden Valley can be transferred to the Standard Model sector [181].

The Lagrangian for the minimal model reads

$$\mathcal{L} = \mathcal{L}_{\rm SM} - \theta \frac{m_f}{v}\, S\, \bar{f} f \;\; - \frac{1}{2}\kappa\, S \bar{\chi}\chi\,, \tag{1.6.1}$$

where χ denotes the dark matter particle which we assume to be a Majorana fermion, while S stands for the scalar mediator. The coupling of the scalar to SM fermions f arises from mixing with the SM Higgs particle, as we will discuss in Section 2.1.

In the hot early universe, the dark matter fermions are in thermal equilibrium due to their interactions with the SM bath. DM pairs can directly annihilate into SM particles via an intermediate scalar S. Alternatively, if kinematically allowed, they can annihilate into pairs of scalars, which subsequently decay to SM particles. As far as the experimental constraints on the scalar S are concerned, only its interactions with the SM are relevant. Indeed, the coupling θ is subject to strong bounds and, hence, processes involving this coupling must be suppressed. Therefore, we shall assume that the second class of processes dominates. The annihilation cross section for $\chi\chi \to SS$ can be estimated as [184]

$$\sigma\, v_{\rm rel} \simeq \sigma_1\, v_{\rm rel}^2 = \frac{\kappa^4 m_\chi}{24\,\pi}\sqrt{m_\chi^2 - m_S^2}\,\frac{9m_\chi^4 - 8m_\chi^2 m_S^2 + 2m_S^4}{(2m_\chi^2 - m_S^2)^4}\, v_{\rm rel}^2\,, \tag{1.6.2}$$

where $v_{\rm rel}$ denotes the relative velocity between two dark matter particles. Notice that this process is p-wave suppressed due to the CP properties of the initial and final state particles. Imposing that the relic density of χ matches the observed dark matter density $\Omega_\chi h^2 = 0.1199$ [49], we can determine κ. More specifically for p-wave suppressed annihilations and $m_\chi = 5$–$10\,\mathrm{GeV}$ the current value of the current DM abundance is achieved for $\sigma_1 \simeq 1.6\cdot 10^{-25}\mathrm{cm}^3/\mathrm{s}$ (see e.g. [185]).

Another consequence of the model (1.6.1) is an emergence of dark matter self-interaction, that could be relevant for cosmology. Current observations do not actually constrain the DM self-interaction cross section to be smaller than that of the strong interaction between nucleons (for a recent review, see [186]), which is many orders of magnitude less stringent than corresponding bounds on DM interacting with standard model particles [187, 188]. Self-interacting DM (SIDM) thus remains a fascinating option which, if confirmed observationally, would significantly reduce the number of possible DM candidates from particle physics. Such observations would, furthermore, offer a window into the particle properties of DM that may be impossible to access by other means – a fact which has created significant attention in recent years (see, e.g. Refs. [189–191]).

Observations of colliding *clusters* lead to the strongest currently existing constraints on SIDM, with $\sigma/m_\chi \lesssim 0.47\,\mathrm{cm}^2/\mathrm{g}$, and it has been argued that their small cores (if any) lead to even stronger bounds [190]. While not undisputed, this has triggered much phenomenological interest in velocity-dependent self-interactions in order to evade cluster bounds and at the same time allow for $\sigma/m_\chi \sim 1\,\mathrm{cm}^2/\mathrm{g}$ at (dwarf) galaxy scales, where the typical DM velocities are up to one order of magnitude smaller [192–195] (and more recently [190, 196, 197]). However, to obtain velocity dependence of the self-interaction one need to consider a sub-GeV scale mediator [186], which is a good motivation for search of light, feebly interacting particles.

However, the situation is not so clear because of large systematic uncertainties in observations. So one needs to deal with *ensembles* of many astrophysical objects, thereby reducing the systematic uncertainties related to individual objects. Such approach was introduced in [198]. As a result no velocity dependence was observed. However, the theoretical model used for description of SIDM halo shows significant deviations from simulations, so it should be also improved [199].

1.7 Summary

The Standard Model of particle physics is a highly successful theory that provides a consistent description of experimentally observed particles and their interactions. Its predictions have been tested and confirmed by numerous experiments. Nevertheless, SM cannot be a complete theory of nature because of the existence of phenomena that cannot be explained by it. To incorporate these phenomena new particles should be introduced. We do not observe new particles because they are either too heavy or light but superweakly interacting. The particles of the second type could be searched in intensity frontier experiments.

In this work we consider two specific models of feebly interacting particles, Higgs-like scalar and a heavy neutral lepton. We discuss their search by means of new generation intensity frontier experiments. In Chapter 2 and Chapter 3 we discuss the phenomenology for the scalar portal and HNLs correspondingly. In Chapter 4 we overview two specific proposals of such experiments, namely SHiP and MATHUSLA. In Chapter 5 we discuss the main features of their sensitivities analytically and present results of numerical simulations. Finally, in Chapter 6 we discuss a possibility to exploit the specialized neutrino detector iSHiP for the search of light dark matter and axion-like particles.