1 INTRODUCTION

In this introductory chapter, we first give a survey of literature comprising of current status of neutrino physics, theoretical and experimental sectors. Then a brief review of the Standard Model (SM) of Particle Physics along with its limitations for which we have to go for beyond Standard Model (BSM) frameworks. We also discuss the different BSM frameworks which can address the issues of the SM, mainly the neutrino mass and mixing, the different seesaw scenarios in brief which can explain the tiny neutrino mass. We mainly focus on the appealing BSM framework, left-right symmetric model and discuss the different phenomenology that could be addressed in its framework, mainly the neutrinoless double beta decay, lepton flavor violation, matter-antimatter asymmetry of the universe which we discuss in different sections. Finally, we discuss in brief the role of flavor symmetry in particle physics giving special emphasis on the discrete group Z_n which we have implemented in this thesis.

Eminent theoretical physicist of all times, Professor Stephen Hawking quoted "Look up at the stars and not down at your feet. Try to make sense of what you see, and wonder about what makes the universe exist. Be curious".

Indeed our inquisitive mind knows no bounds. Several questions like what the world is made up of, how the universe begin, what holds the world together arises at times in the curious minds of all ages since millennia. Particle physics can deal with such quests and can produce a better understanding of the fundamental physical laws, offering a plethora of tough challenges. Today the ultimate questions about the closure of the fundamental dynamical laws and the origin of the observed universe begin to seem accessible. The observed universe is understood to be made up of elementary particles, the first of its kind being the electron discovered by J.J Thomson in the year 1896. Experimentalists started probing the atoms and in 1911, Rutherford found that it consists of the protons and the neutrons. It is again known to most that the latter two are composed of even elementary quarks and gluons. Another fascinating particles which although are very tiny but plays a significant and exceptional role in building up the universe and unravel some of its best-kept secrets are the neutrinos. Owing to its indomitable importance in providing information about the basic structure of the universe, about the mysteries of life and most importantly our very existence has made neutrino physics a hotbed of research for the intrepid band of scientists across the world. Neutrinos are by far the most elusive of all the known occupants of the subatomic realm. In the Standard Model (SM) of particle physics, they belong to the family of leptons (same as that of an electron). But unlike their subatomic cousins, they do not carry an electric charge and are so pathologically shy that they hardly interact with any other particle. About a hundred trillions of neutrinos produced in the nuclear furnace at the Sun's core pass through our body every second without causing any harm and without leaving any trace of their presence. Due to their reluctance to mingle with any particle, they are very difficult to encase for which experimentalists are bound to build large detectors to study these "most wanted" of all the cosmic messengers. Neutrino physics has successively brought the prestigious Nobel prizes the most recent of which was awarded to eminent scientists Takaaki Kajita and Arthur B. McDonald for the discovery of a very significant property of this particle, that they can morph into another type while travelling through space, which is known by the term "neutrino oscillation" in the year 2015. This milestone discovery has pointed out that these ubiquitous particles must have mass and they mix during propagation. Whereas the very successful SM presumed that neutrinos are massless and comes in three different flavors $(\nu_e, \nu_\mu, \nu_\tau)$ and cannot change their form. So the discovery of the massive nature of neutrino and their chameleon-like behavior to change flavors has exposed a severe loophole in the models' elegant structure thereby pointing out its incompleteness. This lead to several theoretical model building beyond the SM and sophisticated instruments are set up across the globe for a detailed investigation of the properties of this prolific yet so elusive particle!

1.1 A brief historical background of Neutrinos

Over the past few decades, neutrinos have drawn the attention of some of the most brilliant minds in the history of physics. These mysterious phantoms are invented for the very first time in the form of scientific witchcraft to get rid of the energy crisis in nuclear physics long before its presence was ever detected by experiments. James Chadwick in the year 1914 [1] showed that beta particle in radioactive decay shows a continuous spectrum contrary to what was expected. When scientists could not account for these missing energy during beta decay, one theorist, the sharp-witted physicist Wolfgang Pauli (Wolfi) found it necessary to invent a new particle to compensate the energy, in the year 1930 [2]. Nurturing some doubt, Pauli even confessed to German astronomer Walter Baade that," I have done a terrible thing, I have postulated a particle that cannot be detected. That is something no theorists should ever do." In 1934, Fermi entitled them as "neutrinos" (meaning "little neutral one" in Italian) [3]. Being electrically neutral, neutrinos cannot be detected via electromagnetic means nor do they experience any strong force. Their chance of interaction via weak force is also very small making it almost near to impossible for theorists and experimentalists to detect these evasive particles. But then with the advent of nuclear power during the second world war, the intense source of radioactive nuclei became available. This not only led to the deadliest of nuclear weapons but also produced an immense source of neutrinos. This caught the attention of eminent socialist Bruno Pontecorvo who first proposed the use of large tanks of carbon tetrachloride (in around 1946), containing chlorine atoms to detect these ghostly particles. Because a neutrino when reacts with a chlorine atom, an argon atom would be produced (inverse beta decay), which is radioactive and its subsequent decay would thus hint that its predecessor was stuck by a neutrino. Despite the pioneering theoretical efforts of Pauli, Fermi and Pontecorvo, very little was known about the existence of these ubiquitous particles, until in 1956, Frederick Reines and Clyde Cowan successfully trapped the (anti)neutrinos for the very first time in a nuclear reactor situated at Savannah river site in south Carolina [4]. Reines bagged a share of the 1995 physics Nobel prize for the detection of the neutrinos with Cowan. It was after that, physical chemist, Ray Davis drew his attention towards the neutrinos produced inside the core of the Sun (solar neutrinos) and in detecting them unlike the ones produced by man-made reactors. Davis set up his experiment at the Homestake gold mine at Brookhaven national laboratory. It was in 1968 that Davis reported the detection of the solar neutrinos but he claimed to detect only one-third of what was predicted by the solar model which led to a great discrepancy between the theory and experiment. It was named as "the solar neutrino problem" [5]. The same problem was again detected by another solar neutrino experiment at Kamioka mine in Japan, also named as Kamiokande around 1989, thereby confirming it. It was Pontecorvo's insight that neutrinos have more than one flavor (i.e., apart form the electron neutrino) which could provide a hint to this problem of solar neutrino deficit. It was much earlier, that is in 1962, Leon Lederman, Melvin Schwartz and Jack Steinberger confirmed the existence of the muon neutrino [6]. The third variety of neutrino, the tau neutrino was observed in the year 2000 at the Fermilab [7]. Thus three "flavors" of the elusive neutrinos came into being. Pontecorvo first made the realization that these ghostly particles might change its flavor (or oscillate) during propagation, which is possible only if they have a tiny yet finite mass. He proposed that two-third of the solar (electron) neutrinos produced in the Sun's core could morph into another type while travelling through space. This proposal was allowed quantum mechanically. Any particle can behave as a wave with a particular wavelength that depends on the mass and speed of the particle. Mathematically each flavor of neutrino has an associated wave function, with different mass and wavelengths. A neutrino is a superposition of all the three flavors, which while travelling through space moves at different rates. Thus at different points in space, the flavor depends on the intrinsic degree of mixing between the different flavors. However, the most successful SM [8, 9, 10] of particle physics, formulated in 1970 would not permit this to happen as the model describes this particle as massless traveling at the speed of light. But theory and experiment should go at the same pace and finally in the late 1990s the phenomenon of neutrino oscillation was first experimentally verified in the Super-Kamiokande (Super-K) detector. Super-K could detect not only the solar neutrinos but also the atmospheric ones (those produced by the cosmic rays hitting the earth's atmosphere) and both of these types produce different experimental signatures thereby became easy to distinguish. It is then apart from the solar neutrino problem, another deficit of the atmospheric neutrinos was observed [11, 12]. Thus they concluded that neutrino oscillation might result in the morphing of the neutrino types. The Sudbery neutrino observatory (SNO) built inside a nickel mine in south Ontario focused on solving the issue of the solar neutrino problem once and for all. Then in the year 2002, the SNO team lead by Arthur B. McDonald, the then professor at Princeton University confirmed the change of form of the solar (electron) neutrinos into muon or tau neutrinos during their journey [13, 14, 15]. The observed flux of the neutrinos matched the theoretical predictions, confirming neutrino oscillation. For this remarkable discovery, the eminent scientists Arthur B. McDonald (for SNO) and Takaaki Kajita (for Super-K) jointly bagged the 2015 Nobel prize in physics.

This is the first compelling evidence of the massive nature of neutrinos and an urgent need to extend the otherwise successful SM and explore for some "new physics".

1.2 Current status of neutrinos

Leading German American physicist, Albert Einstein rightly quoted, "A theory can be proved by experiment, but no path leads from the experiment to the birth of a theory".

Experiments do provide dramatic confirmation of interesting theoretical predictions in the history of Science. Neutrino physics is no such exception. The most important is the experimental proof of Pontecorvo's prediction of neutrino oscillation and the corresponding realization of non-zero neutrino mass, contrary to the expectations of the SM. The remarkable discovery of neutrino oscillation also brought about several new theories and experiments, focusing on the precise measurements of how the mixing between different flavors occurs, denoted by parameters called the "mixing angles". For details about the neutrino mass and mixing, please see [16, 17] The matrix defining the mixing between the flavor eigenstates of two neutrinos and the mass eigenstates was first proposed by Maki, Nakagawa, Sakata and Pontecorvo (later extended to three neutrinos and parameterized in the so- called PMNS matrix). From Super-K and SNO, two of the mixing angles, namely the atmospheric mixing angle (θ_{23}) and the solar mixing angle (θ_{12}) could be determined but not the reactor mixing angle (θ_{13}). Having determined these mixing angles could help the researchers in evaluating the mass differences among the different neutrino types. The current data in neutrino oscillation can be well fitted in terms of two squared mass differences named as the solar mass splitting, Δm^2_{21} and the atmospheric mass splitting, Δm^2_{31}. Matter effects in the Sun suggests $\Delta m^2_{21} > 0$. But Δm^2_{31} is measured via oscillations in the vacuum which depends on the absolute value, its sign (whether + or -) is unknown at present. Thus there is a hierarchy in the neutrino mass states, known as normal and inverted ordering depending on the sign of the atmospheric mass splitting. The current oscillation experiments again can be classified depending on their sensitivities to neutrino parameters. The solar neutrino experiments sensitive to measurements of Δm^2_{31}, θ_{12} are Super-K [18], Borexino [19, 20, 21], Sage[22], SNO. The atmospheric and astrophysical neutrino experiments which are sensitive to the parameters, Δm^2_{23}, θ_{23} are Super-K [23], Icecube [24], ANTARES [25] etc. The reactor-based experiments are DayaBay [26], RENO [27], Double-Chooz[28], sensitive to Δm^2_{31}, θ_{12}. There are again the accelerator- based experiments like T2K [29], MINOS [30], NOVA [31], OPERA [32], ICARUS [33, 34, 35], Microboone [36] etc. measuring the parameters, Δm^2_{23}, θ_{23}, δ_{CP} and even the sterile neutrino. From the recent of these experiments, the non-zero reactor mixing angle θ_{13} has also been discovered. Its rather large value allows for more experiments to determine the neutrino mass hierarchy and the possible exploration of CP violating effects in the neutrino oscillation. Large underground detectors have been proposed to investigate both mass hierarchy and the phase of the possible CP violation. With the increase in development in the sector, the parameter space for the neutrino parameters has been reduced to even narrower ranges of values with increasing confidence levels. For a recent global fit of neutrino oscillation data, the reader can refer to [37]. We see that despite the good precision that neutrino experiments have reached in recent years, many of the neutrino parameters are yet to determined experimentally. With increasing data samples from running experiments, the maximal mixing of θ_{13} and the lower octant of θ_{23} are excluded now. They favour a normal hierarchy of neutrino mass or else there will be an exclusion of CP violation in neutrino sector. Notwithstanding, these would affect numerous other theories and experiments like neutrinoless double beta decay. At present, although the current preferred value of δ_{CP} for both normal and inverted mass orderings lies close to 3/2, the precise value of the CP violating phase in the leptonic sector remains unknown. Besides, several important issues in the neutrino sector are yet to be addressed. Amongst them notable are, the mass hierarchy problem, the intrinsic nature of the neutrinos, whether they are Dirac or Majorana particle (identical particle and antiparticle), the absolute mass of the neutrinos, the existence of CP violation in the leptonic sector, can sterile neutrino be a dark matter candidate etc. Besides the SM is considered an insufficient theory, owing to the fact that it fails to address some other vital questions like, the origin of the tiny neutrino mass, matter-antimatter asymmetry of the universe (BAU), dark matter (DM), lepton number violation (LNV), lepton flavor violation (LFV) and various other cosmological problems [38, 39, 40, 41, 42]. Future astrophysical, laboratory, accelerator and reactor probes are hopeful to address all these open questions that may further reinforce the physics beyond SM. The Higgs field in the SM which is responsible for generating masses to all known particles do not have coupling to neutrinos due to the absence of the right-handed (RH) neutrinos. One can generate a light Majorana mass term for light neutrinos in the SM through the dimension five Weinberg operator of type $(LLHH)/\Lambda$ with the introduction of an unknown cutoff scale [43]. Several beyond Standard Model (BSM) frameworks have been proposed which can provide a dynamical origin of such operators in a renormalizable theory. This is typically achieved in the context of seesaw models where a hierarchy or seesaw between the electroweak scale and the scale of newly introduced fields decide the smallness of neutrino masses. Popular seesaw models can be categorized as type I seesaw [44, 45, 46, 47], type II seesaw [48, 49, 50, 51, 52] , type III seesaw [53, 54], inverse seesaw [55, 56] etc. Other BSM frameworks include, SUSY, extra dimensions, left-right symmetric model (LRSM) etc. with some larger particle contents. The seesaw mechanism being the simplest way to understand the smallness of neutrino masses in BSM. Nevertheless, LRSM is widely used and is an appealing theory where the left and right chiralities are treated in equal footing at high energy scales. Herein, the seesaw mechanisms arise naturally. The minimal LRSM is based on the gauge group $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ and is widely studied in literature [57, 58, 59, 60, 61, 62, 63, 64, 65, 66]. We will basically consider this beautiful BSM model in this thesis.

Symmetry plays a very significant role in particle physics. It is utmost important to understand the underlying symmetry in order to understand the origin of neutrino mass and the leptonic mixing. Neutrino mass and mixing matrix have different forms based upon some flavor symmetries. Symmetries can relate the free parameters of the model thereby making them vanish and making the model even more predictive. The $\mu - \tau$ symmetric neutrino mass matrix giving zero θ_{13} is one of such scenario where the discrete flavor symmetries relates the terms in the neutrino mass matrix. Neutrino oscillation data before discovery of non-zero θ_{13} agrees perfectly with the $\mu - \tau$ symmetric neutrino mass matrix. Amongst the different realizations of $\mu - \tau$ symmetric neutrino mixing pattern, notable are tribimaximal mixing (TBM) [67, 68, 69, 70, 71], bimaximal mixing (BM) [72, 73, 74], hexagonal mixing (HM)[75], golden ratio mixing (GRM)[76, 77, 78, 79] matrices. Amongst them, the most popular one which is consistent with neutrino oscillation data is the Tribimaximal mixing (TBM) structure as proposed by Harison, Perkins and Scott [67]. The resulting mass matrix in the basis of a diagonal charged lepton mass matrix is 2-3 symmetric and magic. The reactor mixing angle (θ_{13}) vanishes in TBM because of the bimaximal character of the third mass eigenstate ν_3. However (θ_{13}) has been measured to be non-zero by experiments like T2K, Daya Bay, RENO and DOUBLE CHOOZ which demands a correction to the TBM form which may be a correction or some perturbation to this type. Henceforth, owing to the current scenerio of neutrino oscillation parameters several new models have been theorized and studied by the scientific communities. Symmetries plays an important role to impose texture zeroes in the mass matrix. Simplest case is one can consider the charged lepton mass matrix to be diagonal and then consider all the possible texture zeroes in the symmetric Majorana mass matrix. Certain one zero and two zero textures in neutrino mass matrix are only consistent with the neutrino data. Another important phenomena that have gained attention in neutrino physics are the processes, neutrinoless double beta decay (NDBD) (for a review, please see[80]) and charged lepton flavor violation (CLFV)[81]. The exact mechanism of LFV being unknown, its study is of large interest as it is linked to neutrino mass generation, CP violation and new physics BSM. The LFV effects from new particles at the TeV scale are naturally generated in many models and therefore considered to be a prominent signature for new physics. The observation of NDBD could also throw light on the absolute scale of neutrino mass and in explaining the matter-antimatter asymmetry of the universe. The different NDBD experiments are KamLAND-Zen [82], NEMO-3 [83, 84], GERDA [85, 86], EXO-200[87], CUORE [88], MAJORANA [89] which provides bounds on the effective neutrino mass parameter, m_{ee}. Apart from neutrino oscillation experiments, the neutrino sector is constrained by the data from cosmology as well. Together with BAO (Baryon Acoustic Oscillation) data, PLANCK mission provided in 2018, the most stringent cosmological constraint on the sum of the active neutrino masses as < 0.12 eV [90]. Furthermore, the evidence of neutrino mass also brought about the curiosity in the minds of the cosmologists if they could account for the mysterious dark matter, whose presence is only seen through its gravitational influence on galaxy clusters. But then from several experiments, it is found that neutrino mass is way too tiny to account for dark matter. Again, some exotic theories proposed for another fourth variety of a neutrino termed as the sterile neutrino, which would never interact with matter but maybe unveiled by indirect means. This type of neutrino has gained attention because it turns out to be massive enough, it may account for dark matter. However, the hunt for a viable dark matter candidate is still on. Thus we see that enormous progress has been made in neutrino physics in the last few decades. Nevertheless, what has been perceived so far is very less in comparison to what awaits to be accomplished in the coming days. The long history of surprise may continue.

1.3 The Standard Model (SM)

Formulated in the early 1970s by Weinberg and Salam [8, 9, 10], the standard electroweak model of particle physics described both the weak and electromagnetic interactions. Along with the theory of strong interactions in the form of quantum chromodynamics or QCD, it provides the "Standard Model (SM)" of all the particle interactions except gravity. Thus SM can be defined as a gauge theory that gives a unified description of all the known particles and their interactions. The gauge group for SM is $SU(3)_C \times SU(2)_L \times U(1)_Y$ representing the strong, weak and electromagnetic interactions. The subscripts represent color charge, left-handed (LH) and hypercharge respectively. It encompasses eighteen fundamental parameters, most of which are associated with the masses of the quarks and leptons, the gauge bosons and the Higgs. Numerous experiments have verified the predictions of SM with exquisite precision. The fabled Large Hadron Collider (LHC) at CERN, the most powerful atom-smasher ever was built to nail down the final missing piece of the model. The LHC successfully confirmed the existence of the Higgs boson, a particle hypothesized to be responsible for endowing other elementary particles with mass. The particle contents of the SM along with their transformation properties under the SM gauge group is shown in tabular form in table 1.1. ϕ in table 1.1 represents the scalar field doublet. The subscripts, L and R represents the left and right chirality respectively and the superscript in ϕ denotes its electric charge. The numbers in the brackets denote the transformation properties under the group $SU(3)_c \times SU(2)_L \times U(1)_Y$ which are the dimension of the field under $SU(3)_c, SU(2)_L$ and the hypercharge respectively.

Multiplets	Particle generation	$SU(3)_c \times SU(2)_L \times U(1)_Y$
l_L	$\begin{bmatrix} \nu_e \\ e^- \end{bmatrix}_L, \begin{bmatrix} \nu_\mu \\ \mu^- \end{bmatrix}_L, \begin{bmatrix} \nu_\tau \\ \tau^- \end{bmatrix}_\tau$	(1, 2, -1)
E_R	$e_R^-, \quad \mu_R^-, \quad \tau_R^-$	(1,1,-2)
Q_L	$\begin{bmatrix} u \\ d \end{bmatrix}_L, \begin{bmatrix} c \\ s \end{bmatrix}_L, \begin{bmatrix} t \\ b \end{bmatrix}_L$	(3, 2, $\frac{1}{3}$)
U_R	$u_R, \quad c_R, \quad t_R$	(3, 1, $\frac{4}{3}$)
D_R	$d_R, \quad s_R, \quad b_R$	(3, 1, $\frac{-2}{3}$)
$\Phi = \begin{bmatrix} \Phi^+ \\ \Phi^0 \end{bmatrix}$		(1, 2, 1)

Table 1.1: All the particle contents of the Standard Model.

1.3.1 The electroweak sector

The electroweak model is also known as the Glashow-Weinberg-Salam model describes all the electromagnetic and weak interactions represented by the gauge group, $SU(2)_L \times U(1)_Y$. The gauge bosons mediating these interactions are the massive W^+, W^-, Z^0 bosons and the massless photon A_μ. The basic principles of the SM are a)local gauge symmetry and b) Spontaneous symmetry breaking and Higgs mechanism.

The study of gauge theory deals with the Lagrangian density (Lagrangian) which encodes basically all the informations about the interactions along with the dynamics of the fields. The local gauge transformations which keep the SM Lagrangian invariant under $SU(2)_L \times U(1)_Y$ are,

$$\overline{\Psi}'_L = e^{(ig\frac{\tau}{2}\theta(x)+ig'\frac{Y}{2}\Theta(x))}\overline{\Psi}_L, \overline{\Psi}'_R = e^{ig'\frac{Y}{2}\Theta(x)}\overline{\Psi}_R. \tag{1.1}$$

By insisting on a local gauge transformation, we are required to replace the ordinary derivative by a covariant derivative given by,

$$D_\mu = \partial_\mu + i\frac{g'}{2}YB_\mu + i\frac{g}{2}\tau_a W^a_\mu \tag{1.2}$$

The symmetry groups are gauged by the introduction of the gauge fields, W_μ^a and B_μ corresponding to the symmetry groups SU(2) and U(1) respectively. The generators of the groups are τ_a and Y respectively, with $a = 1, 2, 3$ for the three different generations of the leptons. The coupling constants of electromagnetic and weak interactions are represented by, $g^{'}$ and g. The gauge field dynamics is given by the gauge part of the Lagrangian called the Yang-Mill's Lagrangian,

$$\mathcal{L}_{\mathrm{YM}} = -\frac{1}{4}\mathrm{W}^{\mathrm{a}}_{\mu\nu}\mathrm{W}^{\mathrm{a}\mu\nu} - \frac{1}{4}\mathrm{B}_{\mu\nu}\mathrm{B}^{\mu\nu}, \tag{1.3}$$

where, $B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu$ and $W^{a\mu\nu} = \partial_\mu W^a_\nu - \partial_\nu W^a_\mu - g f^{abc} W^b_\mu W^c_\nu$, f^{abc} being the structure constants representing the components of the antisymmetric levi-civita tensor for the SU(2) group.

The fermionic part of the Lagrangian is given by,

$$\mathcal{L}_{\mathrm{f}} = \overline{\mathrm{L}}\gamma^\mu(\mathrm{i}\partial_\mu - \mathrm{g}^{'}\frac{\mathrm{Y}}{2}\mathrm{B}_\mu - \mathrm{g}\frac{\tau}{2}\mathrm{W}_\mu)\mathrm{L} + \overline{\mathrm{R}}\gamma^\mu(\mathrm{i}\partial_\mu - \mathrm{g}^{'}\frac{\mathrm{Y}}{2}\mathrm{B}_\mu)\mathrm{R} \tag{1.4}$$

The gauge-invariant Lagrangian for the Higgs field is given by,

$$\mathcal{L}_\phi = |(i\partial_\mu - g\frac{\tau}{2}.W_\mu - g^{'}\frac{Y}{2}B_\mu)\phi|^2 - V(\phi) \tag{1.5}$$

The Yukawa Lagrangian for the Quark and lepton masses is given by,

$$\mathcal{L} = -Y_d[\overline{Q_L}\phi d_R] - Y_u[\overline{Q_L}\tilde{\phi} u_R] - Y_l[\overline{l_L}\phi l_R] + h.c, \tag{1.6}$$

where Y are the Yukawa couplings. Thus we arrive at the total SM Lagrangian by requiring a $SU(2)_L \times U(1)_Y$ invariant form given by,

$$\mathcal{L}_{SM} = \mathcal{L}_{YM} + \mathcal{L}_f + \mathcal{L}_\phi + \mathcal{L}_{yuk} \tag{1.7}$$

From the Yukawa Lagrangian, we see that due to the absence of the RH neutrinos, the mass term for the neutrino does not arise. However, the scalar potential and the Higgs field will provide an elegant solution for the mass problem of the fermions and the gauge bosons through spontaneous symmetry breaking (SSB) and the Higgs mechanism which we will discuss in the next section.

1.3.2 Spontaneous symmetry breaking and Higgs mechanism

In the non-abelian $SU(2)_L \times U(1)_Y$ electroweak theory, one needs to generate masses for the three gauge bosons $W^{\pm}$ and Z bosons but the mediator of the electromagnetic interaction remains massless and QED must remain an exact symmetry so that the electric charge is conserved. A symmetry is said to be spontaneously broken if the Lagrangian is invariant under some symmetry whereas the ground state does not possess the same symmetry as its Lagrangian. The term "spontaneous" comes from the fact that the system tends towards its ground state spontaneously. To break the symmetry spontaneously, a complex $SU(2)_L$ doublet of scalar fields is introduced given by,

$$\phi = \begin{bmatrix} \phi^+ \\ \phi^0 \end{bmatrix} = \frac{1}{\sqrt{2}} \begin{bmatrix} \phi_1 + i\phi_2 \\ \phi_3 + i\phi_4. \end{bmatrix}. \tag{1.8}$$

The Lagrangian describing the interaction and propagation of the scalars is given by,

$$\mathcal{L}_\phi = (D^\mu \phi)^\dagger (D_\mu \phi) - V(\phi) \tag{1.9}$$

where the scalar potential is given by,

$$V(\phi) = \mu^2 \phi^\dagger \phi + \lambda(\phi^\dagger \phi)^2, \phi^\dagger \phi = \frac{1}{2}\left[\phi_1^2 + \phi_2^2 + \phi_3^2 + \phi_4^2\right] \tag{1.10}$$

where $\lambda > 0$. We require the Lagrangian to remain invariant under the symmetry operation which replaces ϕ by $-\phi$. It suffices to keep the first two allowed terms in the general expression of V in powers of ϕ.

For $\mu^2 > 0$, no spontaneous symmetry breaking occurs and the ground state corresponds to $\phi = 0$. The Electroweak symmetry is spontaneously broken only when $\mu^2 < 0$. In this case, the fundamental state corresponds to a circle of degenerate fundamental states. The potential has two minima which satisfies,$\frac{\partial V}{\partial \phi} = 0$ which yields, $\phi = \sqrt{\frac{-\mu^2}{\lambda}}$. Gauge invariance allows us to choose the state of minimum energy,

$$< 0|\phi|0 >= \phi_0 = \frac{1}{\sqrt{2}} \begin{bmatrix} 0 \\ v \end{bmatrix}, v = \sqrt{\frac{-\mu^2}{\lambda}}, [\phi_1 = \phi_2 = \phi_4 = 0, \phi_3 = v] \tag{1.11}$$

The VEV v will break the gauge symmetry $SU(2)_L \times U(1)_Y$ into $U(1)_{em}$. To develop the theory around the minimum, we can parameterize the fluctuations around ϕ_0 in terms of four fields θ_1, θ_2, θ_3 and h(x) as,

$$< \phi(x) >= \begin{bmatrix} \theta_1 + i\theta_2 \\ \frac{1}{\sqrt{2}}(v + h(x)) - i\theta_3 \end{bmatrix} = e^{\frac{i\theta_a \tau_a}{v}} \begin{bmatrix} 0 \\ \frac{1}{\sqrt{2}}(v + h(x)) \end{bmatrix} \quad (1.12)$$

where, $h(x)$ represents the physical Higgs field. Then considering an $SU(2)_L$ gauge transformation on this field,

$$\phi(x) \to e^{-\frac{i\theta_a \tau_a}{v}} \phi(x) = \begin{bmatrix} 0 \\ \frac{1}{\sqrt{2}}(v + h(x)) \end{bmatrix}. \quad (1.13)$$

The three fields θ_1, θ_2, θ_3 represents the three Goldstone bosons that will give masses to the three weak gauge boson fields, $W^a_\mu(x)$, a=1, 2, 3. Spontaneous symmetry breaking (SSB) will lead to a massive boson from the field h(x). To determine the masses of the gauge bosons, it is sufficient to substitute ϕ_0 into the gauge-invariant Lagrangian,

$$\begin{aligned} \mathcal{L} = (\partial_\mu \phi + ig\frac{1}{2}\tau.W^a_\mu \phi + ig^{'}\frac{1}{2}Y.B_\mu \phi)^\dagger (\partial^\mu \phi + ig\frac{1}{2}\tau.W^{a\mu}\phi + ig^{'}\frac{1}{2}Y.B^\mu \phi) - V(\phi) \\ -\frac{1}{4}W_{\mu\nu}.W^{\mu\nu} - \frac{1}{4}B_{\mu\nu}.B^{\mu\nu} \end{aligned} \quad (1.14)$$

where the last term represents the kinetic energy term for the gauge fields. The relevant term in the Lagrangian providing the mass for the gauge bosons is

$$|(ig\frac{1}{2}\tau_i.W^a_\mu + ig^{'}\frac{1}{2}Y.B_\mu)\phi_0|^2. \quad (1.15)$$

Substituting the value of ϕ_0 wand ith some further simplification and then comparing with the mass terms for the charged gauge bosons, $\frac{1}{2}M^2_W W^+_\mu W^{-\mu}$ and $\frac{1}{2}M^2_W Z_\mu Z^\mu$, we can get the masses for the gauge bosons. Where,

$$W^\pm = \frac{W^1 \mp iW^2}{\sqrt{2}}, Z_\mu = \frac{gW^3_\mu - g^{'}B^\mu}{\sqrt{g^2 + g^{'2}}}, A_\mu = \frac{g^{'}W^3_\mu + gB^\mu}{\sqrt{g^2 + g^{'2}}}. \quad (1.16)$$

The masses are thus obtained as $M_W = \frac{vg}{2}$ and $M_Z = \frac{v\sqrt{g^2+g^{'2}}}{2}$. The SSB has produced a precise prediction of the gauge boson masses relating them to the VEV of the scalar field. The measured masses of these gauge bosons are, $M_Z = 91.1875 \pm 0.0021$ GeV and $M_W = 80.399 \pm 0.023$ GeV [91].

The Lagrangian breaks spontaneously the $SU(2)_L$ and $U(1)_Y$ gauge symmetry but ϕ_0 is chosen such that Q(ϕ_0) = 0), i.e., it remains neutral, thus $U(1)_{em}$ remains unbroken. Thus the generator of the group, photon remains massless. Whereas, the three goldstone bosons have been absorbed by the three gauge bosons to form their longitudinal components and get mass. This process is known as the Higgs mechanism.

1.3.3 Fermionic weak interactions and fermion masses

The interaction between the gauge fields and the fermions are completely determined by local gauge invariance. The electroweak charged current Lagrangian is given by,

$$\mathcal{L}_{CC} = \frac{g}{\sqrt{2}} \left(\overline{\nu_L} \gamma^\mu W_\mu^- l_L + \overline{l_L} \gamma^\mu W_\mu^+ \nu_L \right). \tag{1.17}$$

The neutral current Lagrangian consisting of the neutral gauge bosons is given by,

$$\begin{aligned} \mathcal{L}_{NC} = -\frac{1}{2}\overline{l_L} \left[(gc_W - g^{'} s_W)\gamma^\mu Z_\mu + (gs_W + g^{'} c_W)\gamma^\mu A_\mu \right] l_L + \\ \frac{1}{2}\overline{\nu_L}(gc_W + g^{'} s_W)\gamma^\mu Z_\mu \nu_L + g^{'}\overline{L_R}(c_W \gamma^\mu A_\mu - s_W \gamma^\mu Z_\mu) l_R. \end{aligned} \tag{1.18}$$

In the above equations, s_W, c_W represents sin θ_W and cos θ_W, where, θ_W is the Weinberg angle [92].

In a similar manner like the gauge bosons, the fermions (except the neutrinos) also attain their mass after SSB of the electroweak theory. The same Higgs doublet which generates masses for the gauge bosons is also sufficient to give masses to the leptons and the quarks. As we see in the original SM Lagrangian, a Dirac mass term (-m$\overline{\Psi}\Psi$) was forbidden by the gauge invariance. However electroweak symmetry breaking(EWSM) allows for Dirac mass terms to appear via Yukawa interactions. The Yukawa Lagrangian is given by,

$$\mathcal{L} = -Y_d[\overline{Q_L}\phi d_R] - Y_u[\overline{Q_L}\tilde{\phi} u_R] - Y_l[\overline{l_L}\phi l_R] + h.c, \tag{1.19}$$

where, $\phi = \frac{1}{\sqrt{2}} \begin{bmatrix} 0 \\ v + h \end{bmatrix}$, $\tilde{\phi} = i\tau_2 \phi^\dagger$ and Y_d, Y_u and Y_l are the Yukawa couplings which are generally complex square matrices. Thus after substitution of ϕ, the required charged lepton and quark masses are generated and given by,

$$M_u = \frac{Y_u v}{\sqrt{2}}, M_d = \frac{Y_d v}{\sqrt{2}}, M_l = \frac{Y_l v}{\sqrt{2}}. \tag{1.20}$$

Thus we have the mass matrices as complex. For arbitrary complex square matrix, we need unitary matrices to diagonalize the mass matrix to get the propagating eigenstates called the mass eigenstates. We can multiply it from the left and right with distinct unitary matrices. The 3×3 unitary matrices, V_L, V_R are the rotation matrices such that,

$$M_d(diag) = V_L^{d\dagger} M_d V_R^d, M_u(diag) = V_L^{u\dagger} M_u V_R^u, M_l(diag) = V_L^{l\dagger} M_l V_R^u \tag{1.21}$$

The rotations of the quarks will affect the quark charged current in a way that the quark gauge interactions are not diagonal. The corresponding Lagrangian is thus given by,

$$\mathcal{L} = -\frac{g}{\sqrt{2}}\overline{u'_L}V_L^{u\dagger}\gamma^\mu V_L^d d'_L W_\mu + h.c \tag{1.22}$$

where, $V_L^{u\dagger}V_L^d = V_{CKM}$ is the mixing matrix known as Cabibbo-Kobayashi-Maskawa (CKM) matrix [93, 94]. There are different (but equivalent) representations of the CKM matrix in literature. The Particle data group advocates the use of the following one as the standard CKM parametrization:

$$V = \begin{bmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -c_{23}s_{12} - s_{23}s_{13}c_{12}e^{i\delta} & -c_{23}c_{12} - s_{23}s_{13}s_{12}e^{i\delta} & s_{23}c_{13} \\ s_{23}s_{12} - c_{23}s_{13}c_{12}e^{i\delta} & -s_{23}c_{12} - c_{23}s_{13}s_{12}e^{i\delta} & c_{23}c_{13} \end{bmatrix} \tag{1.23}$$

The parameterized form of CKM matrix consists of three mixing angles and the Dirac CP violating phase, δ. The abbreviations used are $c_{ij}= \cos\theta_{ij}$, $s_{ij}=\sin\theta_{ij}$, (i,j = 1, 2, 3) for the three generations of quarks [95]. However, we see that due to the absence of the RH neutrinos, the Dirac mass term for the neutrinos is forbidden and hence neutrinos remain massless in the Standard Model.

1.4 Limitations of the Standard Model

In spite of its enormous success in explaining almost all experimental results, the SM of particle physics is found to be an incomplete theory owing to the various issues which cannot be addressed within its framework. Notable amongst them are listed below,

1.4.1 Theoretical

- It cannot incorporate gravitational interactions. We know that gravitational interactions are present in nature no matter how feeble it is in comparison to the other fundamental interactions which makes the model incomplete.
- SM could not address whether the gauge couplings unify at high energies, like GUT (SO(10)) as they do not unify in the SM. The SM is thus often referred to as an effective low-energy theory of the corresponding high-energy theory.

- There are a large number of free parameters in the model, namely the gauge coupling constants, scalar potential parameters, charged lepton masses, quark masses, mixing angles of the CKM matrix, CP violating phase and the strong CP violating parameter. The values of these parameters are not predicted by the model but to be determined by experiments. Thus, we cannot consider it to be a complete theory with so many inadequacies.

- The huge difference in the strength of fundamental forces is one aspect of the so-called "hierarchy problem" [96]. It also refers to the wide range in mass for the elementary particles. We know that the electron is about 200 times lighter than the muon and 3500 times lighter than the tau. Same thing for the quarks: the top quark is 75000 times heavier than the up quark. This wide spectrum of masses among the building blocks of matter could not be explained by the SM.

  Within the SM, the mass of the Higgs gets some very large quantum corrections self-interactions, gauge loops, and fermion loops (especially the top quark). These loops are quadratically divergent and go like $\int \frac{d^4k}{k^2-m^2} \sim \Lambda^2$, for some unknown cut off scale Λ. For large Λ (of the order of Planck mass), these corrections are much larger than the actual mass of the Higgs, which is termed as the Hierarchy problem. This means that the bare mass parameter of the Higgs in the SM must be fine-tuned in such a way that almost completely cancels the quantum corrections. This level of fine-tuning is deemed unnatural by many theorists. Among the beyond SM theories, supersymmetry can address this problem by naturally giving a solution to this problem by stabilizing the ratio, $\frac{\Lambda_{EW}}{M_{Planck}}$.

- There has been no experimental indication of the CP (charge, parity) symmetry in the strong interactions. But, theoretically, the QCD Lagrangian suggests otherwise. It goes as,

$$\mathcal{L}_{QCD} = \frac{\theta_{QCD} g_3^2}{32\pi^2} G^a_{\mu\nu} \tilde{G}^{a\mu\nu}, \tag{1.24}$$

  which is P and T violating leading to CPT invariance and CP violation. In the above equation, θ_{QCD} is a parameter, g_3 is the coupling constant of QCD and G represents the QCD field strength tensor. The strong CP problem is the puzzling question of why quantum chromodynamics (QCD) does not seem to break CP-symmetry. As there is no known reason for it to be conserved in QCD specifically, this is a "fine tuning" which is considered very unnatural and is known as the strong CP problem.

- Another problem is the problem of mass, the origin of particle masses. If they are due to a Higgs boson, why are the masses so small.
- Flavor: why are there so many different types of quarks and leptons and why do their weak interactions mix in a peculiar way observed?

1.4.2 Observational

- SM does not seem to have any explanation of the dark sector. Whereas about 27% (five times of the ordinary baryonic matter) and 68% of the matter in our universe consists of dark matter and dark energy from observation of the cosmic microwave background (CMB) as measured by WMAP [39, 40, 100, 101]. Unlike the other baryonic particles of the SM, they can interact only via gravity. These dark sectors which comprise a dominant contribution of the universe is an evidence of the incompleteness of the SM.
- In the SM, there are three generations of the neutrinos, also called three flavors: electron neutrinos, muon neutrinos and tau neutrinos like the other matter particles in the SM. But some experiments like Liquid Scintillator Neutrino Detector (LSND) [102, 103, 104] and the MiniBooNE [105] experiment have shown hints for a new type of neutrino that does not fit neatly into this simple picture. It has been termed to be "sterile" meaning it likely would not interact directly with any of the SM particles. It might, however, be a form of dark matter.
- Contrary to the expectation of the SM, the milestone discovery of neutrino oscillation has given compelling evidence that these tiny ubiquitous particles are massive and they mix during propagation. But as we have seen neutrinos are massless in the SM, which is a very important setback of the otherwise successful theory. Some experimental results have even suggested that there might be a fourth type of neutrino called a sterile neutrino that is yet to be discovered.
- The concept of matter-antimatter asymmetry (also known as baryon asymmetry) of the observable universe could not be addressed in SM scenario. Whereas compelling

evidence of this asymmetry is observed from various sources like the Big Bang nucleosynthesis (BBN) and Cosmic microwave background radiation (CMBR). A good amount of CP violation along with the violation of baryon number must occur. It can only be understood in theories where the universe evolves far out of equilibrium during its early age. Whereas in SM baryon number is a good symmetry and thus no question of baryon number violation arises within its framework.

- The intrinsic nature of the neutrinos, whether they are Dirac particles or Majorana particles (particle and antiparticle are identical) could not be explained by the SM.

Whatsoever, apart from all the failures that the SM withstands, it has always been appreciated by physicists for being a remarkable beginning to understanding and possibly unifying all of physics.

However, the thirst to know the unknowns always prevails. Keeping a note of the above-listed drawbacks of the SM, we reach a scenario where we need to address these issues by extending the SM. There are several interesting BSM frameworks which can explain one or more of these issues, like the Super Symmetry (SUSY), Extra Dimensions, Seesaw mechanisms, Left-right symmetric model, GUT etc.

1.5 Beyond Standard Model (BSM)

1.5.1 Neutrino Mass

To get the answers to the queries the SM could not address, various BSM frameworks have been proposed off late. Again the existence of neutrino mass is very firmly established from neutrino oscillations, though the absolute scale of the mass is not perceived yet. In this section, we are focusing on this very important issue, how the mass of the neutrino could be generated in an extension of the SM. To begin with, we very briefly discuss the phenomenon which brought into light the massive nature of the neutrinos.

1.5.1.1 Neutrino Oscillation

It is a well-established fact that the neutrinos and antineutrinos which takes part in the standard charged current (CC) and neutral current (NC) weak interaction are of three different types or flavors, ν_e and $\bar{\nu_e}$, ν_μ and $\bar{\nu_\mu}$, ν_τ and $\bar{\nu_\tau}$. The morphing between these three different flavors during its propagation is termed as "neutrino oscillation".

The data of the neutrino oscillation experiments as observed in solar, atmospheric and long baseline reactor and accelerator experiments like K2K, MINOS, Double Chooz are well fitted in the framework of three neutrino scheme

$$\nu_{lL}(x) = \sum_{i=1}^{3} U_{li}\nu_{iL}(x) \quad ; \quad (l = e, \mu, \tau) \tag{1.25}$$

Here, $\nu_i(x)$ is the field of neutrinos with mass m_i. The LH flavor field $\nu_{lL}(x)$ enters into the standard leptonic charged current as

$$j_\alpha^{cc}(x) = 2 \sum_{l=e,\mu,\tau} \bar{\nu}_{lL}(x)\gamma_\alpha l_L(x) \tag{1.26}$$

and determines the notion of a LH flavor ν_L which is produced in charged current weak processes together with a lepton, $l^\dagger$. The flavor ν_l is described by the mixed state

$$|\nu_l >= \sum_{i}^{3} U_{li}^*|\nu_i > \tag{1.27}$$

where, $|\nu_i >$ is the neutrino state with mass, m_i and a definite momentum. The flavor eigenstate and the mass eigenstate can be co-related by 3×3 rotation matrix,

$$\begin{pmatrix} \nu_e \\ \nu_\mu \\ \nu_\tau \end{pmatrix} = (U_{\text{PMNS}}) \begin{pmatrix} \nu_1 \\ \nu_2 \\ \nu_3 \end{pmatrix}, \quad U_{\text{PMNS}} = \begin{pmatrix} U_{e1} & U_{e2} & U_{e3} \\ U_{\mu 1} & U_{\mu 2} & U_{\mu 3} \\ U_{\tau 1} & U_{\tau 2} & U_{\tau 3}. \end{pmatrix}$$

U_{PMNS} is the 3×3 unitary mixing matrix first established by Pontecorvo, Maki, Nakagawa, Sakata similar to the CKM mixing matrix of the Quark sector and is given by,

$$U_{PMNS} = \begin{bmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -c_{23}s_{12} - s_{23}s_{13}c_{12}e^{i\delta} & -c_{23}c_{12} - s_{23}s_{13}s_{12}e^{i\delta} & s_{23}c_{13} \\ s_{23}s_{12} - c_{23}s_{13}c_{12}e^{i\delta} & -s_{23}c_{12} - c_{23}s_{13}s_{12}e^{i\delta} & c_{23}c_{13} \end{bmatrix} U_{Maj} \tag{1.28}$$

The abbreviations used are just like the CKM mixing matrix described before. The mixing matrix consists of the lepton mixing angles and the phases. The phase matrix, U_{Maj} is $diag(1, e^{i\alpha}, e^{i\beta})$ contains the Majorana phases α and β which would appear if neutrino is a Majorana particle. Further, we can write $U_{PMNS} = U_l^{\dagger} U_{\nu}$ where U_l and U_{ν} are the diagonalizing matrix of the charged lepton and neutrino respectively. The probability of $\nu_e \to \nu_{\mu}$ oscillation in time t is

$$P(t)_{\nu_e \to \nu_{\mu}} = sin^2\theta cos^2\theta [e^{-iE_2 t} - e^{-iE_1 t}]^2 = sin^2 2\theta .2 sin^2 (\frac{t.\Delta m^2_{21}}{4E}), \tag{1.29}$$

From the expression of probability, we see that it depends upon the mixing angle θ, the mass squared difference, time, t which can be expressed in terms of the propagation distance L as $t = \frac{L}{c}$ and the neutrino energy, E. Thus, it is quite obvious that for oscillation to occur, the neutrino mass should definitely be non-zero and there should be a definite mixing between the different flavors, defined by the mixing angle. So far, we know some of the neutrino oscillation parameters to some level of accuracy, for the recent best fit values of the parameters the readers can refer to [37]. However, the absolute scale of the neutrino mass is not yet known. Besides there arises a hierarchy in the neutrino mass also called the ordering. $\Delta m^2_{21} > 0$ is known from experiments. But the sign of Δm^2_{31} is still an unknown in neutrino physics sector. Thus there is a hierarchy/ordering in the neutrino mass states, known as normal and inverted ordering depending on the sign of the atmospheric mass splitting.

- Normal Hierarchy (NH) which corresponds to $m_1 < m_2 \ll m_3$.
- Inverted Hierarchy (IH) which corresponds to $m_3 \ll m_1 < m_2$.

1.5.1.2 Dirac and Majorana mass term

Masses and mixing are characterized by a mass term which as we have seen consists of both the LH and RH components of the fields. Due to the absence of the RH neutrino, we could not get massive neutrino in the SM. However, if we add RH neutrinos in the SM, we will get a different scenario. In the case of a neutrino which is neutral, unlike its other subatomic cousins, two distinct classes of neutrino mass terms are allowed in the Lagrangian of the electroweak interactions, these are called Dirac and Majorana mass terms. The main difference between Dirac and Majorana lies in the fact that in one lepton number is conserved

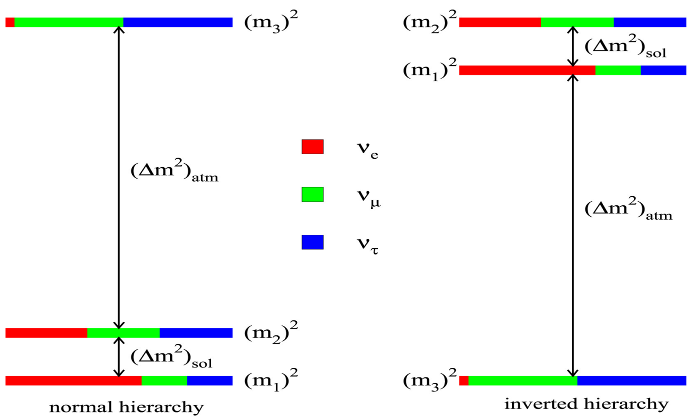

Figure 1.1: Mass splittings and the two possible neutrino mass hierarchies (normal and inverted). The colours in the figure represent the fraction of each flavor in the mass states, (Taken from [106]).

while in the other it is violated by two units. The Lagrangian corresponding to the Dirac neutrino mass term is,

$$\mathcal{L}^D = -\sum_{ij} \overline{\nu_{iL}} M_D \nu_{jR} + h.c \tag{1.30}$$

where M_D is a complex 3×3 matrix which can be diagonalized by the unitary matrix U and V as,

$$Diag(M_D) = V^\dagger M_D U \tag{1.31}$$

Thus, the mass term becomes,

$$\mathcal{L}^D = -\sum_{i=1}^{3} M_i \overline{\nu_{Li}} \nu_{Ri} \tag{1.32}$$

where, ν_i are the mass eigenstates with mass m_i. In the absence of a RH neutrino, we can write the mass term when we consider $\nu_L^C = C\overline{\nu_L}^T$, C being the charge conjugation matrix. Thus we can write the mass term as,

$$\mathcal{L}^L = -\frac{1}{2} \sum_{ij} \overline{\nu_{iL}} M_L \nu_{jL}^c + h.c. \tag{1.33}$$

The factor $\frac{1}{2}$ is introduced to avoid double counting of the independent fields. M_L is a complex symmetric mass matrix which can be diagonalized by the unitary matrix, U, which consist of the three mixing angles and three phases, such that,

$$\mathrm{M_L} = \mathrm{UM_L(Diag)U^\dagger}. \tag{1.34}$$

Thus, the mass term can be written as,

$$\mathcal{L}^L = -\frac{1}{2}\sum_{i=1}^{3} M_i \overline{\nu_i}\nu_i; \quad \nu_i^C = C\overline{\nu_i}^T \tag{1.35}$$

The mass term above is violating lepton number conservation by two units and is termed as the Majorana particle, which cannot distinguish between particle and anti particle. Similarly, we can write a RH Majorana mass term as,

$$\mathcal{L}^R = -\frac{1}{2}\sum_{i=1}^{3} M_i \overline{\nu_i}\nu_i \tag{1.36}$$

However, the RH neutrino field does not appear in the SM Lagrangian and it has to be extended to get such a term.

1.5.1.3 Seesaw Mechanism

A light Majorana mass term for light neutrinos in the SM through the dimension five Weinberg operator [43] of type $\frac{(LL\phi\phi)}{\Lambda}$ with the introduction of an unknown cutoff scale Λ. Several beyond BSM frameworks have been proposed henceforth which can provide a dynamical origin of such operators in a renormalisable theory. This is typically achieved in the context of the so-called seesaw models where a hierarchy or seesaw between the electroweak scale and the scale of newly introduced fields decide the smallness of neutrino masses. The simplest of all the BSM mechanisms to explain the origin of the tiny neutrino mass is the seesaw mechanism wherein different new heavy scalar particles are added to the SM like the heavy RH neutrinos (N_R), SU(2) triplet, Δ, fermion triplets, $\sum$, fermion singlet etc. Depending upon the different inclusions, the seesaw mechanism is thus categorized into type I SS, type II SS, type III SS, Inverse SS, radiative SS mechanisms etc. Below, we briefly describe some of the mentioned SS mechanisms.

Type I Seesaw: The type I seesaw mechanism [44, 45, 46, 47] is the simplest of all the seesaw mechanisms to realize the dimension five operator to explain the tiny neutrino mass within a renormalizable framework. In this case, RH neutrinos are added to generate the small neutrino masses. We can thus write a Majorana mass term as

$$-\mathcal{L}^I = Y^\nu \overline{\nu_R}\tilde{\phi}^\dagger L + \frac{1}{2}M_R\overline{\nu_R}\nu_R^C + h.c \tag{1.37}$$

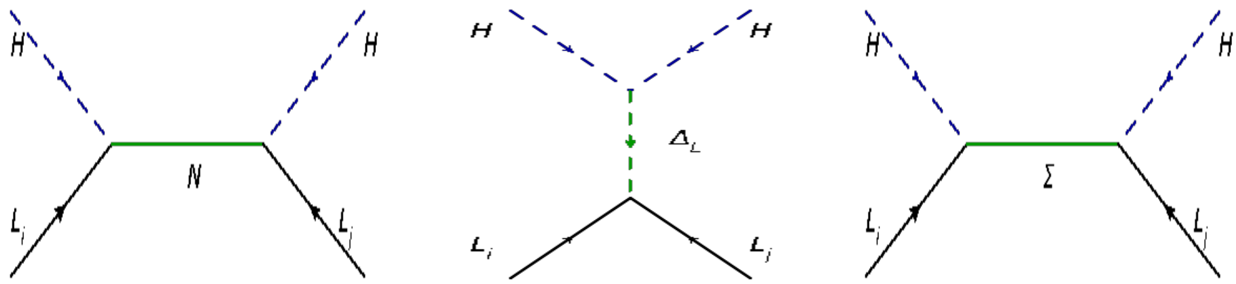

Figure 1.2: The schematic diagrams for type I, type II and type III seesaw are shown in the figure.

M_R is a symmetric RH Majorana mass matrix whereas the Yukawa matrix Y^ν is non-symmetric and non-hermitian. After EWSB, the usual Dirac mass terms arises,

$$\mathcal{L}_{Dirac} = -\overline{\nu_R} M_D \nu_L - \overline{\nu_L} {M_D}^\dagger \nu_R \tag{1.38}$$

which give the Dirac mass to the neutrinos as, $Y^\nu v$, where v is the VEV of the Higgs. The typical scale of M_R is much higher than M_D, as ν_R are gauge singlets, their masses are decoupled from the electroweak scale allowing them to be of the order of the cut off scale of the low energy theory and are unprotected by gauge symmetry.

The final Lagrangian of the mass term is obtained as,

$$\mathcal{L}^\nu{}_m = \frac{1}{2} \begin{pmatrix} \overline{\nu_L}^C & \overline{\nu_R} \end{pmatrix} \begin{pmatrix} 0 & M_D^T \\ M_D & M_R \end{pmatrix} \begin{pmatrix} \nu_L \\ \nu_R \end{pmatrix} \tag{1.39}$$

To make the full matrix of the above equation diagonal, we can perform a unitary transformation by R such that

$$R^T \begin{pmatrix} 0 & M_D^T \\ M_D & M_R \end{pmatrix} R = \begin{pmatrix} M_{light} & 0 \\ 0 & M_{heavy} \end{pmatrix} \tag{1.40}$$

We finally get the masses from type I seesaw as, $M_{light} = M_\nu = M_D^T M_R^{-1} M_D$ and $M_{heavy} = M_R$, from which it is evident why it is termed as seesaw. Heavier the M_R, lighter will be the M_ν.

Type II Seesaw: In the type II seesaw mechanism [48, 49, 50, 51, 52] , the SM is extended by including an additional scalar triplet Δ, the matrix representation of which is given by,

$$\Delta = \frac{1}{\sqrt{2}}\sum_i \sigma^i \Delta_i = \begin{pmatrix} \frac{\Delta^+}{\sqrt{2}} & \Delta^{++} \\ \Delta^0 & -\frac{\Delta^+}{\sqrt{2}} \end{pmatrix}, \tag{1.41}$$

where, σ_i represents the Pauli matrices and the three complex scalars are, $\Delta^0 = \frac{\delta_1+i\delta_2}{\sqrt{2}}$, $\Delta^+ = \delta_3$, $\Delta^{++} = \frac{\delta_1-i\delta_2}{\sqrt{2}}$. The corresponding Lagrangian describing the neutrino mass in type II seesaw is given by,

$$-\mathcal{L}^{II} = Y_\Delta L^T Ci\sigma_2 \Delta L + M_\Delta^2 Tr\left[\Delta^+\Delta\right] + \frac{1}{2}\left(\lambda_\Delta M_\Delta \tilde{\phi}^\dagger \Delta^+ \phi\right) + h.c. \tag{1.42}$$

In the above equation, M_Δ represents the mass of the Higgs triplet with Y_Δ as its Yukawa coupling. When the neutral component of the Higgs doublet, ϕ acquires a VEV by EWSB, it induces a tadpole term for Δ as given by the last term of equation (1.42) thereby generating a VEV for Δ as,

$$< \Delta >= v_\Delta = \frac{\lambda_\Delta v^2}{M_\Delta}. \tag{1.43}$$

Thus neutrino mass is generated via the type II seesaw mechanism and is given by,

$$M_\nu = \frac{Y_\Delta v_\Delta}{\sqrt{2}}. \tag{1.44}$$

Type III Seesaw: In type III seesaw [53, 54], the additional scalars added are three hyperchargeless triplets, $\sum$ with the following $SU(2)_L$ representation

$$\sum = \frac{1}{\sqrt{2}}\sum_i \sigma^i \Delta_i = \begin{pmatrix} \frac{\Sigma^0}{\sqrt{2}} & \Sigma^+ \\ \Sigma^- & \frac{-\Sigma^0}{\sqrt{2}} \end{pmatrix}. \tag{1.45}$$

Σ has a definition in terms of three components, η_1, η_2, η_3 such that $\Sigma^0 = \eta_3$, $\Sigma^\pm = \frac{\eta_1 \pm \eta_2}{\sqrt{2}}$ The relevant Lagrangian describing the interaction is given by,

$$-\mathcal{L}^{III} = Y_\Sigma \tilde{\phi}^\dagger \Sigma^a L + \frac{1}{2} M_\Sigma Tr\left[\Sigma^a \Sigma^b\right] + h.c \tag{1.46}$$

The neutrino mass generated in this case is similar to the type I seesaw, here the neutral component of Σ is similar to the RH neutrino in the type I seesaw. The difference is that RH neutrino in type I seesaw are singlets whereas here they are triplets. The neutrino mass matrix is thus given by,

$$M_\nu = M_D M_\Sigma^{-1} M_D^T \tag{1.47}$$

M_D in the above expression like before is $\frac{Y_\Sigma v}{\sqrt{2}}$. The value of M_Σ can reach the cutoff scale of the low energy theory as it is unprotected by any kind of symmetry. The type III seesaw mass term is a lepton number violating as the co-existence of M_Σ and Y_Σ does not provide any leptonic charge to Σ.

Inverse Seesaw: Likewise, in inverse seesaw (ISS) [55, 56], SM is extended by one or more generations of RH neutrinos, ν_R and additional fermionic singlets, S. The 9×9 neutrino mass matrix from the ISS mechanism is given by,

$$M_\nu = \begin{pmatrix} 0 & M_D^T & 0 \\ M_D & 0 & M_R^T \\ 0 & M_R & \mu \end{pmatrix}. \tag{1.48}$$

Considering $\mu \ll M_D \ll M_R$ and after block diagonalization of the above 9×9 matrix, we get the effective light neutrino mass matrix as,

$$M_\nu = M_D^T (M_R^T)^{-1} \mu M_R^{-1} M_D. \tag{1.49}$$

It generally happens that in ISS mechanism, neutrino mass is obtained by double suppression of M_R, unlike the other seesaw mechanisms. Besides the seesaw mechanisms there are other appealing beyond BSM frameworks, where the origin of the light neutrino mass can be elegantly explained. We will consider one such framework in this thesis, the model description of which we will discuss in the next section.

1.5.2 Left-Right Symmetric Model (LRSM)

As we have already seen how the neutrino mass could be generated in an extension of the SM by the addition of the non- renormalizable dimension five Weinberg operator. The primary goal is always to give the neutrino a small mass in the context of a model with no renormalization group needed and closely resembles the SM which could be at low energies accessible at present-day experiments. Our choice is to focus on a so-called left-right symmetric model (LRSM) [57, 58, 59, 60, 61, 62, 63, 64, 65] which appeals to be one of the most straightforward and natural extensions of the SM. The most attractive feature of the LRSM lies in its relation between the high and low mass scales of the theory. The main motivation behind the origin of the LRSM can be listed as follows

- Their gauge group is a very simple extension of the SM gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$ SM and is given by $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, where the subscripts represent color, LH, RH and the difference between baryon and lepton number. Lepton number violation and heavy particles can naturally appear after the left-right symmetry is spontaneously broken to electroweak symmetry.

- Parity is an explicit symmetry of the theory until spontaneous symmetry breaking takes place. Spontaneous breaking of P (and also CP) is possible but it requires a RH ν degrees of freedom which naturally appears in the model because the symmetry of the theory ensures the presence of the RH counterparts of every LH particle.

- The seesaw mechanisms could be naturally realized in the context of LRSM. The additional neutrino multiplets give rise to Majorana masses and allow for scenarios such that the neutrino masses are naturally light.

The LRSM was first introduced around 1974 by Pati and Salam. Again, Rabindra N. Mohapatra and Goran Senjanovic were also some pioneers of this very elegant theory. Initially proposed in the context of GUTs [60], these extensions are very much attractive. LRSM naturally arises from SO(10) based GUTs [107, 108]. There are two ways of introducing the symmetry: as generalized parity (P) or as generalized charged conjugation (C). In the case of fermions, they coincide with the usual parity and charge conjugation. Whereas, for gauge boson, they are chosen in such a way as to keep the gauge interactions invariant:

$$P : W_L, q_L, l_L \leftrightarrow W_R, q_R, l_R, C : W_L, q_L, l_L \leftrightarrow -W_R^\dagger, q_R^C, l_R^C, \tag{1.50}$$

It also leads to the equality of the gauge couplings, g_L and g_R. One of the important features of the model is the gauged U(1) B-L symmetry (difference in the baryon and lepton number) which is conserved until SSB occurs. The conservation is because there are no couplings between quarks and leptons in the theory due to which L and B will be conserved separately. The fact that the B-L is anomaly free, is sometimes used to signify why one should gauge the symmetry. In the SM, due to the absence of the RH neutrino, ν_R, B-L symmetry is only partially anomaly free, as $Tr[B-L]Q_a^2 = 0$ but $Tr[B-L]^3 \neq 0$ and cannot be gauged. But on the inclusion of ν_R to the SM, which happens in LRSM, $Tr[B-L]^3 = 0$ which implies that it can be gauged and thus appears in the gauge group of LRSM. Another important reason is that B-L automatically appears as a gauged symmetry in SO(10) grand unification. The presence of the RH ν further modifies the Gell-Mann-Nishijima like formula for the electric charges to be,

$$Q = T_{3L} + T_{3R} + \frac{B-L}{2} \tag{1.51}$$

Where T_{3L} and T_{3R} represents the isospin third components under $SU(2)_L$ and $SU(2)_R$.

1.5.2.1 Particle contents of LRSM and Electroweak Symmetry breaking

We are mainly focused on the simplest LRSM also known as the minimal left-right symmetric model (MLRSM). As the term suggests LRSM enforces us to have a symmetric particle content under $G_{LRSM} = SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. Imposing a discrete LRSM ensures the necessity of the scalar sector to have either a RH doublet or a triplet in order to break the $SU(2)_R$ symmetry and the LH counterparts should naturally be present to obey the symmetry. However, the LH scalar alone is not sufficient to generate the masses which lead to the necessity of a scalar bidoublet. Although it is seen that in most of the cases, the LH field is considered to be negligible with its VEV almost equal to zero. But, in our case, we will consider all the particle contents which are summarized in a tabular form in table 1.2.

Fields	Representations	Charge under G_{LRSM}
$Q_{L,R}$	$\begin{bmatrix} u' \\ d' \end{bmatrix}_{L,R}$	(3, 2, 1, 1/3), (3,1,2,1/3)
$L_{L,R}$	$\begin{bmatrix} \nu_l \\ l \end{bmatrix}_{L,R}$	(1, 2, 1, -1), (1, 1, 2, -1)
$\Delta_{L,R}$	$\begin{bmatrix} \delta_{\frac{L,R}{\sqrt{2}}}^{+} & \delta_{L,R}^{++} \\ \delta_{L,R}^{0} & -\delta_{\frac{L,R}{\sqrt{2}}}^{+}. \end{bmatrix}$	(1, 3, 1, 2), (1, 1, 3, 2)
ϕ	$\begin{bmatrix} \phi_1^0 & \phi_1^+ \\ \phi_2^- & \phi_2^0 \end{bmatrix} \equiv \left(\phi_1, \widetilde{\phi_2}\right)$	(1,2,2,0)

Table 1.2: All the particle contents of the MLRSM along with the charge assignments.

The superscripts in the representation denote the electric charge and the numbers in the brackets represent the respective quantum numbers under the gauge group $SU(3)_c \times SU(2)_L \times$

$SU(2)_R \times U(1)_{B-L}$. The fermionic gauge Lagrangian is given by,

$$\begin{aligned}\mathcal{L}_f = \overline{\psi}_L i\gamma^\mu \left(\partial_\mu + ig_L \frac{\tau_a}{2}.W^a_{L\mu} + ig^{'} \frac{B-L}{2} B_\mu\right)\psi_L + \\ \overline{\psi}_R i\gamma^\mu \left(\partial_\mu + ig_R \frac{\tau_a}{2}.W^a_{R\mu} + ig^{'} \frac{B-L}{2} B_\mu\right)\psi_R\end{aligned} \tag{1.52}$$

In the above equation, $g_{L,R}$ and $g^{'}$ represents the coupling constants of the $SU(2)_{L,R}$ and $U(1)_{B-L}$ gauge groups respectively. Like the SM, $(W_{L,R})_\mu{}^a$ and B_μ represents the gauge fields, which are in fact the generators of the groups. τ_a are the Pauli spin matrices and $a = 1, 2, 3$. Again the gauge interaction Lagrangian can be written as,

$$\mathcal{L}_g = -\frac{1}{4} W^{\mu\nu}_{Li} W_{Li\mu\nu} - \frac{1}{4} W^{\mu\nu}_{Ri} W_{Ri\mu\nu} - \frac{1}{4} B^{\mu\nu} B_{\mu\nu}. \tag{1.53}$$

For the fermions to attain mass, a Yukawa Lagrangian is necessary which couples to the bidoublet, ϕ. The corresponding Yukawa Lagrangian is given by,

$$-\mathcal{L}_{Dirac} = \overline{l_{iL}} \left(Y^l_{ij}\phi + \tilde{Y}^l_{ij}\tilde{\phi}\right) l_{jR} + \overline{Q_{iL}} \left(Y^q_{ij}\phi + \tilde{Y}^q_{ij}\tilde{\phi}\right) Q_{jR} + h.c. \tag{1.54}$$

$\tilde{\phi}$ in the above Yukawa Lagrangian has the same transformations as ϕ and is given by $\tilde{\phi} = \tau_2 \phi^* \tau_2$. Only incorporating ϕ into the model could not break the LRSM gauge group down to SM gauge group, so we have the scalar triplets. The Yukawa Lagrangian for the triplets which we will see in the next subsection that gives Majorana masses to the neutrinos is given by,

$$-\mathcal{L}_{Majorana} = f_{L,ij} \Psi_{L,i}{}^T C i\sigma_2 \Delta_L \Psi_{L,j} + f_{R,ij} \Psi_{R,i}{}^T C i\sigma_2 \Delta_R \Psi_{R,j} + h.c. \tag{1.55}$$

where, $f_{L,R}$ are the coupling matrices which are considered to be equal for a discrete left-right symmetry. Again, the Kinetic part of the scalar Lagrangian in LRSM can be written as,

$$\mathcal{L}_{scalar} = Tr\left[(D_\mu\phi)^\dagger (D^\mu\phi) + (D_\mu\Delta_L)^\dagger (D^\mu\Delta_L) + (D_\mu\Delta_R)^\dagger (D^\mu\Delta_R)\right]. \tag{1.56}$$

The covariant derivatives in the above expression are,

$$D_\mu\phi = \partial_\mu\phi - ig_L W^a_{L\mu}\frac{\tau_a}{2}\phi - ig_R \frac{\tau_a}{2} W^a_{R\mu}, \tag{1.57}$$

$$D_\mu\Delta_R = \partial_\mu\Delta_R - ig_R\left[\frac{\tau_a}{2}W^a_{R\mu}, \Delta_R\right] - ig^{'} B_\mu \Delta_R, \tag{1.58}$$

$$D_\mu\Delta_L = \partial_\mu\Delta_L - ig_L\left[\frac{\tau_a}{2}W^a_{L\mu}, \Delta_L\right] - ig^{'} B_\mu \Delta_L, \tag{1.59}$$

where the expression $[X, Y] = XY - YX$. The total Lagrangian for LRSM is thus the sum of all, that is

$$\mathcal{L}_{total} = \mathcal{L}_g + \mathcal{L}_f + \mathcal{L}_{Dirac} + \mathcal{L}_{Majorana} + \mathcal{L}_{scalar} + \mathcal{L}_V, \tag{1.60}$$

where, $\mathcal{L}_V$ is the Lagrangian corresponding to the Higgs potential, the details of it has been explained beautifully in [48].

1.5.2.2 Spontaneous Symmetry Breaking in LRSM

The scalar potential in LRSM is a combination of interaction terms consisting the potential and hence can be written as,

$$V_{\text{scalar}} = V_\Phi + V_{\Delta_{L,R}} + V_{\Phi,\Delta_{L,R}} \tag{1.61}$$

where the individual potentials are given by the following equations,

$$\begin{aligned} V_\Phi = -\mu_1^2 \text{Tr}[\Phi^\dagger \Phi] - \mu_2^2 \text{Tr}[\Phi^\dagger \tilde{\Phi} + \tilde{\Phi}^\dagger \Phi] + \lambda_1 \Big(\text{Tr}[\Phi^\dagger \Phi]\Big)^2 + \lambda_2 \Big\{ \Big(\text{Tr}[\Phi^\dagger \tilde{\Phi}]\Big)^2 + \\ \Big(\text{Tr}[\tilde{\Phi}^\dagger \Phi]\Big)^2 \Big\} + \lambda_3 \text{Tr}[\Phi^\dagger \tilde{\Phi}] \text{Tr}[\tilde{\Phi}^\dagger \Phi] + \lambda_4 \text{Tr}[\Phi^\dagger \Phi] \text{Tr}[\Phi^\dagger \tilde{\Phi} + \tilde{\Phi}^\dagger \Phi], \end{aligned} \tag{1.62}$$

$$\begin{aligned} V_{\Delta_{L,R}} = -\mu_3^2 \text{Tr}[\Delta_L^\dagger \Delta_L + \Delta_R^\dagger \Delta_R] + \rho_1 \Big\{ \Big(\text{Tr}[\Delta_L^\dagger \Delta_L]\Big)^2 + \Big(\text{Tr}[\Delta_R^\dagger \Delta_R]\Big)^2 \Big\} + \\ \rho_2 \Big\{ \text{Tr}[\Delta_L \Delta_L] \text{Tr}[\Delta_L^\dagger \Delta_L^\dagger] + \text{Tr}[\Delta_R \Delta_R] \text{Tr}[\Delta_R^\dagger \Delta_R^\dagger] \Big\} + \rho_3 \text{Tr}[\Delta_L^\dagger \Delta_L] \text{Tr}[\Delta_R^\dagger \Delta_R] + \\ \rho_4 \Big\{ \text{Tr}[\Delta_L \Delta_L] \text{Tr}[\Delta_R^\dagger \Delta_R^\dagger] + \text{Tr}[\Delta_L^\dagger \Delta_L^\dagger] \text{Tr}[\Delta_R \Delta_R] + \alpha_1 \text{Tr}[\Phi^\dagger \Phi] \text{Tr}[\Delta_L^\dagger \Delta_L + \Delta_R^\dagger \Delta_R], \end{aligned} \tag{1.63}$$

$$\begin{aligned} V_{\Phi,\Delta_{L,R}} = \alpha_1 \text{Tr}[\Phi^\dagger \Phi] \text{Tr}[\Delta_L^\dagger \Delta_L + \Delta_R^\dagger \Delta_R] + \Big\{ \alpha_2 \Big(\text{Tr}[\Phi^\dagger \tilde{\Phi}] \text{Tr}[\Delta_L^\dagger \Delta_L] + \\ \text{Tr}[\tilde{\Phi}^\dagger \Phi] \text{Tr}[\Delta_R^\dagger \Delta_R] \Big) + \text{h.c.} \Big\} + \alpha_3 \text{Tr}[\Phi \Phi^\dagger \Delta_L \Delta_L^\dagger + \Phi^\dagger \Phi \Delta_R \Delta_R^\dagger] + \beta_1 \text{Tr}[\Phi^\dagger \Delta_L^\dagger \Phi \Delta_R + \\ \Delta_R^\dagger \Phi^\dagger \Delta_L \Phi] + \beta_2 \text{Tr}[\Phi^\dagger \Delta_L^\dagger \tilde{\Phi} \Delta_R + \Delta_R^\dagger \tilde{\Phi}^\dagger \Delta_L \Phi] + \beta_3 \text{Tr}[\tilde{\Phi}^\dagger \Delta_L^\dagger \Phi \Delta_R + \Delta_R^\dagger \Phi^\dagger \Delta_L \tilde{\Phi}]. \end{aligned} \tag{1.64}$$

In the above equations for the scalar potential, we have introduced the scalar mass parameters μ_i and quartic scalar interaction strengths λ_i, ρ_i, α_i and β_i which are dimensionless coefficients.

The breaking pattern of the LRSM gauge group is given by,

$$SU(2)_L \times SU(2)_R \times U(1)_{B-L} \xrightarrow{<\Delta_R>} SU(2)_L \times U(1)_Y. \xrightarrow{<\Delta_L><\phi>} U(1)_{\rm em} \tag{1.65}$$

After spontaneous symmetry breaking, the neutral components of the scalars acquires a VEV as

$$\langle \Delta_{L,R} \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & 0 \\ v_{L,R} & 0 \end{pmatrix}, \langle \Phi \rangle = \begin{pmatrix} k_1 & 0 \\ 0 & e^{i\theta} k_2 \end{pmatrix}. \tag{1.66}$$

The magnitudes of the VEVs follows the relation, $|v_L|^2 < |k_1^2 + k_2^2| < |v_R|^2$. One can show that in this case the smallness of $|v_L|^2$ follows from the Higgs potential (as shown in [109]).

1.5.2.3 Masses in LRSM, Fermion and Gauge boson mass

Gauge boson masses and mixing The leptonic charged current interaction in flavor basis is given by,

$$\mathcal{L}_{CC}^{lepton} = \frac{g}{\sqrt{2}} \left[\overline{l'} \gamma^\mu P_L \nu' W_{L_\mu}^- + \overline{l'} \gamma^\mu P_R \nu' W_{R_\mu}^- \right] + h.c, \tag{1.67}$$

where, the LR charged gauge boson mixing is described by,

$$\begin{bmatrix} W_L^\pm \\ W_R^\pm \end{bmatrix} = \begin{bmatrix} \cos\xi & \sin\xi e^{i\alpha} \\ -\sin\xi e^{-i\alpha} & \cos\xi \end{bmatrix} \begin{bmatrix} W_1^\pm \\ W_2^\pm \end{bmatrix}, \tag{1.68}$$

where $W_{1,2}^\pm$ are the mass eigenstates for the gauge bosons, ξ is the gauge boson mixing angle and is related to the VEVs as

$$\tan 2\xi = -\frac{2k_1 k_2}{v_R^2 - v_L^2}. \tag{1.69}$$

The VEVs k_1, k_2 satisfy the VEV of the SM namely, $k = v_{\rm SM} = \sqrt{k_1^2 + k_2^2} \approx 246$ GeV. Whereas the VEV v_L plays a very significant role in neutrino mass mechanism and is generated after the electroweak symmetry breaking due to the following induced VEV relation,

$$\langle \Delta_L \rangle = v_L = \frac{\gamma v_{\rm SM}^2}{v_R}. \tag{1.70}$$

Here, γ is a dimensionless parameter given by [109],

$$\gamma = \frac{\beta_1 k_1 k_2 + \beta_2 {k_1}^2 + \beta_3 {k_2}^2}{(2\rho_1 - \rho_3) k^2}. \tag{1.71}$$

Again, considering the hierarchy, $v_R >> k_1, k_2 >> v_L$. the masses are thus,

$$\begin{aligned} M_{W_L}^2 &= \frac{g^2 k^2}{4}\left(1 - \frac{2k_1^2 k_2^2}{k^2 v_R^2}\right), \qquad M_{W_R}^2 = \frac{g^2}{2}|v_R^2|, \\ M_Z^2 &= \frac{g^2 k^2}{4\cos^2\theta_W}\left(1 - \frac{k_1^2}{4\cos^4\theta_Y v_R^2}\right), \qquad M_{Z'}^2 = g^2 v_R^2 , \end{aligned} \tag{1.72}$$

with θ_W and θ_Y being the mixing angles.

Fermion masses The Quark mass terms in LRSM are,

$$M_u = \frac{1}{\sqrt{2}}(k_2 Y_q + k_1 \tilde{Y_q}), M_d = \frac{1}{\sqrt{2}}(k_1 Y_q + k_2 \tilde{Y_q}). \tag{1.73}$$

The Dirac mass terms for the leptons comes from the Yukawa Lagrangian, which for the charged leptons and neutrinos are given by,

$$M_l = \frac{1}{\sqrt{2}}(k_2 Y_l + k_1 \tilde{Y_l}), M_D = \frac{1}{\sqrt{2}}(k_1 Y_l + k_2 \tilde{Y_l}) \tag{1.74}$$

The 6×6 neutrino mass matrix in LRSM is given, in the (ν_L, ν_R) gauge eigenbasis as

$$M_\nu = \begin{pmatrix} \sqrt{2} f v_L & M_D \\ M_D^T & \sqrt{2} f v_L \end{pmatrix} = \begin{pmatrix} M_{LL} & M_D \\ M_D^T & M_{RR} \end{pmatrix} \tag{1.75}$$

Assuming $M_{LL} \ll M_D \ll M_R$, the light neutrino mass after symmetry breaking is generated within a type I+II seesaw as,

$$\mathrm{M}_\nu = {\mathrm{M}_\nu}^{\mathrm{I}} + {\mathrm{M}_\nu}^{\mathrm{II}} \tag{1.76}$$

$$M_\nu = M_{LL} - M_D {M_{RR}}^{-1} {M_D}^T = \sqrt{2} v_L f_L - \frac{v_{\mathrm{SM}}^2}{\sqrt{2} v_R} Y_l {f_R}^{-1} Y_l^T, \tag{1.77}$$

$$M_{LL} = \sqrt{2} v_L f_L, M_{RR} = \sqrt{2} v_R f_R, \tag{1.78}$$

Again, the neutrino mass matrix M_ν can be digonalized by a 6×6 unitary matrix as follows,

$$\mathcal{V}^T M_\nu \mathcal{V} = \begin{bmatrix} \widehat{M_\nu} & 0 \\ 0 & \widehat{M}_{RR} \end{bmatrix}, \tag{1.79}$$

where, $\mathcal{V}$ represents the diagonalizing matrix of the full neutrino mass matrix, M_ν, $\widehat{M_\nu} = Diag(m_1, m_2, m_3)$, with $\mathrm{m_i}$ being the light neutrino masses and $\widehat{M}_{RR} = Diag(M_1, M_2, M_3)$, with M_i being the heavy RH neutrino masses. The diagonalizing matrix is represented as,

$$\mathcal{V} = \begin{bmatrix} U & S \\ T & V \end{bmatrix} \approx \begin{bmatrix} 1 - \frac{1}{2}RR^\dagger & R \\ -R^\dagger & 1 - \frac{1}{2}R^\dagger R \end{bmatrix} \begin{bmatrix} V_\nu & 0 \\ 0 & V_R \end{bmatrix}, \tag{1.80}$$

where, R describes the left-right mixing and given by,

$$R = M_D M_{RR}^{-1} + \mathcal{O}(M_D^3 (M_{RR}^{-1})^3). \tag{1.81}$$

The matrices U, V, S and T are as follows,

$$U = \left[1 - \frac{1}{2} M_D M_{RR}^{-1} (M_D M_{RR}^{-1})^\dagger\right] V_\nu, \tag{1.82}$$

$$V = \left[1 - \frac{1}{2} (M_D M_{RR}^{-1})^\dagger M_D M_{RR}^{-1}\right] V_R, \tag{1.83}$$

$$S = M_D M_{RR}^{-1} v_R f_R, \tag{1.84}$$

$$T = -(M_D M_{RR}^{-1})^\dagger V_\nu. \tag{1.85}$$

The seesaw mechanisms could be realized in the context of LRSM in a very elegant way. However, there are various experimental constraints in the masses of the particles in the model which comes from different sources, like $K_L - K_S$ mass difference, neutrinoless double beta decay, non-leptonic kaon decays, muon decays, lepton flavor violation, astrophysical constraints arising from nucleosynthesis etc. We will consider those constraints in the phenomenological studies we addressed in this thesis.

1.5.3 Neutrinoless double beta decay

By far, we know that neutrinos are massive but the question of their intrinsic nature is still not answered. We are not aware of whether these particles are Dirac or Majorana in nature. By Majorana, we mean the particles which have identical particle and antiparticle, first hypothesized by Ettore Majorana in 1937 [110]. It is closely connected to the one if lepton number is or is not a symmetry of nature, because a Majorana mass term violates lepton number by two units. This issue cannot be addressed by neutrino oscillation experiments which are not sensitive to the Majorana parameters. Therefore, an alternate experiment has to be performed to determine the nature of neutrinos. Because lepton number violating processes generically has very small amplitudes owing to its suppression by the tiny neutrino masses makes them very difficult to observe experimentally. Presently, the most promising attempts to find lepton number violation are the experiments on neutrinoless double beta decay (NDBD). We start with what this process is all about. By definition, NDBD is a very slow second order radioactive process in which two neutrons inside a nucleus transform into two protons emitting two electrons. The discovery of this process could lead to two very important consequences, firstly that neutrinos are Majorana particles and secondly total lepton number is not conserved in nature, two findings which could have far-reaching implications in particle physics and cosmology. Besides, the observation of NDBD could also throw light on several other issues which could not be addressed by the SM, like the tiny neutrino mass by the presence of a heavy scale by the introduction of Majorana neutrinos and LNV along with the CP asymmetry could explain the matter-antimatter asymmetry of the universe.

However, after about 80 years of experimental effort, no compelling evidence for the existence of NDBD has been obtained, but a new generation of experiments that are already running or about to run swears to push forward the limit exploring the different mass ordering for the neutrino masses. To do that, the experiments are using masses of NDBD isotopes ranging from tens of kilograms to several hundred, and will need to improve the background rates achieved by previous experiments by, at least, an order of magnitude.

1.5.3.1 Neutrinoless double beta decay theory

We know that a first order beta decay is a process in which a nucleus with atomic number Z transforms into a nucleus with atomic number Z+1 with the simultaneous emission of an electron and an anti-neutrino given by the reaction,

$$N\left(A,Z\right) \rightarrow N\left(A,Z+1\right) + e^{-} + \overline{\nu_{e}}. \tag{1.86}$$

But, in some arrangements of the nuclei, such a process is energetically forbidden or suppressed by selection rules. Thus arises the concept of double beta decay, which was first discussed by M. Goeppert-Mayer in the year 1935 in the form of,

$$N(A,Z) \rightarrow N(A,Z+2) + e^- + e^- + +\overline{\nu_e} + +\overline{\nu_e}, \tag{1.87}$$

in which an initial nucleus (Z, A), with proton number Z and total nucleon number A, decays to (Z + 2, A) emitting two electrons and two antineutrinos in the process. It is a second order process (also written as $2\nu\beta\beta$) and can be seen as two simultaneous beta decays. This can only happen for isotopes containing even-even nuclei. In nature, 35 isotopes are known which show the specific ground state configuration, necessary for double beta decay (DBD). Although the first evidence for DBD came in 1950 through a geochemical observation for the decay of 130-Te [111], while the experimental observation of double beta decay [112] was first observed in the laboratory in 1987 for 82-Se. Such processes conserve electric charge and lepton number and are allowed in the electroweak SM.

However, many models for beyond the SM predicts lepton number violation. If indeed lepton number is broken in nature, another form of double beta decay is possible, the so-called neutrinoless double beta decay (denoted by $0\nu\beta\beta$), given by the transition,

$$N(A,Z) \rightarrow N(A,Z+2) + e^- + e^-. \tag{1.88}$$

The Feynman diagram for $0\nu\beta\beta$ is shown in figure 1.3. It is evident from the figure that no neutrinos are emitted in the process. This is possible only if $\nu_e = \overline{\nu_e}$, i.e., when neutrino is a Majorana particle. If the two outgoing electrons in $0\nu\beta\beta$ are considered to have typical momentum as p_1 and p_2 respectively, the matrix element of the decay process, mediated by the light neutrinos can be written as,

$$\mathcal{M}_{\mu\lambda} = \lambda^* g_{\mu\lambda} \sum_i \frac{U_{ei}^2 m_i}{\langle p^2 \rangle} \overline{e_L}(p_1) C \overline{e_L}^T(p_2), \tag{1.89}$$

where, λ is a phase, C is the charge conjugation matrix, $\langle p^2 \rangle$ is the typical momentum transfer, which is of the order of nuclear scale and is inversely proportional to square of nuclear radius 10^{-13} cm, thus $\langle p^2 \rangle \approx (100MeV)^2$. The amplitude for the light neutrino exchange can thus be written as,

$$A_\nu = \left(\frac{g}{\sqrt{2}}\right)^4 \frac{1}{M_W^4} \frac{{U_{ej}}^2 m_j}{\langle p^2 \rangle} \tag{1.90}$$

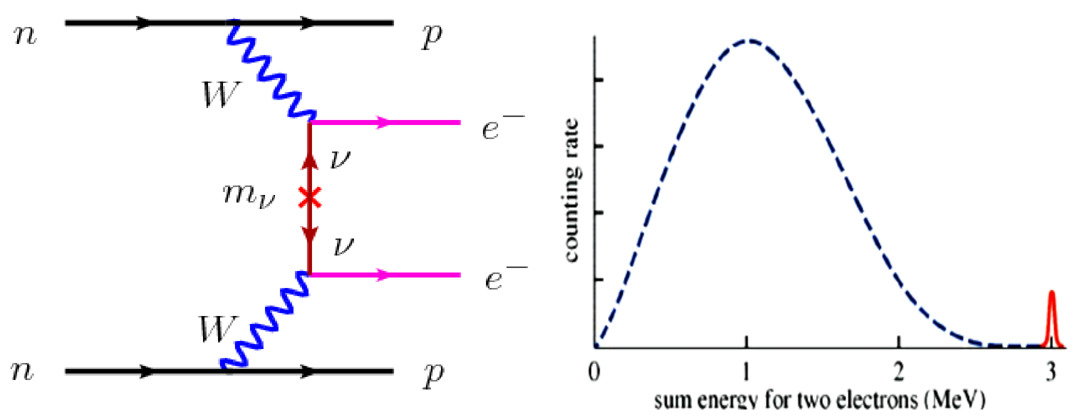

Figure 1.3: The standard diagram for $0\nu\beta\beta$ mediated by W bosons by exchange of light neutrinos(left) and the energy spectrum for $2\nu\beta\beta$(blue) and $0\nu\beta\beta$ (red)(right).

The general expression for the total decay width of $0\nu\beta\beta$, mediated by light neutrinos is given by,

$$\Gamma^{0\nu} = \frac{1}{T_{\frac{1}{2}}{}^{0\nu}} = G^{0\nu}(Q,Z)\left|M^{0\nu}\right|^2 \frac{|m_{\beta\beta}|^2}{m_e{}^2}, \tag{1.91}$$

where $T_{\frac{1}{2}}{}^{0\nu}$ is the half-life of the process, the terms $G_{0\nu}$, M_ν and m_e represents the phase space factor, the nuclear matrix element(NME) and the electron mass respectively. NME depends on the nucleus under consideration. The parameter $|m_{\beta\beta}|$ is the effective Majorana mass of the light neutrino which is a combination of the neutrino mass eigenstates and first row of the neutrino mixing matrix terms, given by,

$$m_{\beta\beta} = \sum_j U_{ej}{}^2 m_j \ , \ j = 1, 2, 3. \tag{1.92}$$

In the standard parameterization of the mixing matrix, $m_{\beta\beta}$ can be written as,

$$|m_{\beta\beta}| = m_1 c_{12}{}^2 c_{13}{}^2 + m_2 s_{12}{}^2 c_{13}{}^2 e^{2i\alpha} + m_3 s_{13}{}^2 e^{2i\beta}. \tag{1.93}$$

The Majorana phases α and β are the phase differences of U_{e_2} and U_{e_3} with respect to U_{e_1}.

The different NDBD experiments are KamLAND-Zen [82], NEMO-3 [83], GERDA [85], EXO-200 [87], CUORE [88], MAJORANA [89] etc. which provides good limits on the NDBD half-life. We see that just measuring the half-life does not allow to extract the effective Majorana neutrino mass $m_{\beta\beta}$. Besides, we need to have sufficient information about the so called NME [113, 114] and the phase space integrals. The calculation of NME is a nuclear physics problem and it is a many-body problem which uses approximations. Presently there are two approaches which are generally considered, the Nuclear Shell Model (NSM) [113] and the Quasi-Particle Random Phase Approximation (QPRA) [116, 117].

In recent years, experiments such as KATRIN hope to measure the mass of neutrinos in the next coming years. If a mass measurement is obtained, it would be a very impactful result. Furthermore, the current generation of experiments, such as MAJORANA, COURE and EXO explore to expedite whether or not neutrinos are indeed their own antiparticle, besides giving a stringent bound on the absolute scale of the neutrino mass which if found will be a landmark discovery in neutrino sector.

1.5.3.2 Neutrinoless double beta decay in LRSM

In the context of the BSM frameworks, where SM is extended with the inclusion of several new particles, many new contributions to NDBD amplitudes comes into the picture. In this thesis, we will mainly focus on the LRSM scenario, where due to the presence of new scalars and gauge bosons, various additional sources would give rise to contributions to NDBD process, which involves RH neutrinos and RH gauge bosons [118][119], Higgs triplets [120] as well as the mixed LH-RH contributions. There are several contributions to the neutrinoless double beta decay in addition to the standard contribution to NDBD via light Majorana neutrino exchange. There are many such studies in literature, where NDBD is studied in the framework of LRSM like [121, 122, 123, 124, 125, 126, 127, 128, 129, 130]. The different Feynman diagrams that contributes to the NDBD amplitude along with the standard light neutrino contribution that arises in the framework of LRSM are shown in the figures 1.4, 1.5, 1.6, 1.7. The detailed study of NDBD in LRSM will be presented in the subsequent chapters.

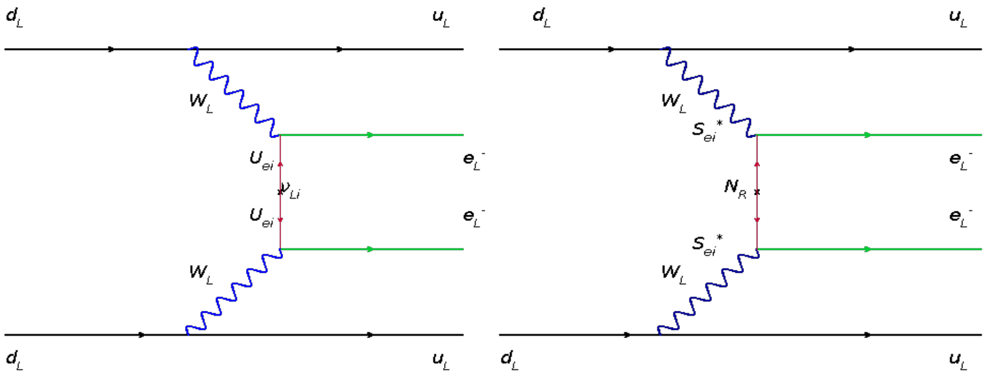

Figure 1.4: Feynman diagrams corresponding to neutrinoless double beta decay due to $\nu - W_L - W_L$, $N - W_L - W_L$contributions.

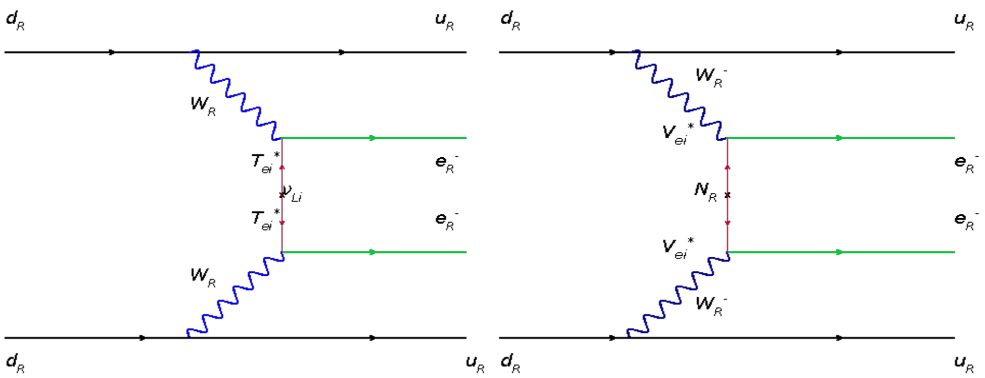

Figure 1.5: Feynman diagrams corresponding to neutrinoless double beta decay due to $\nu - W_R - W_R$, $N - W_R - W_R$ contributions.

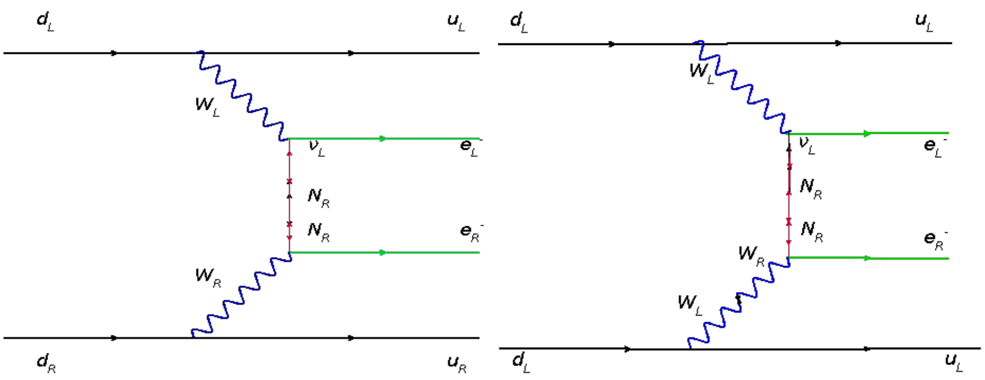

Figure 1.6: Feynman diagrams corresponding to neutrinoless double beta decay due to $N - W_L - W_R$ mediation with heavy-light neutrino exchange and $W_L - W_R$ mixing (λ and η contributions).

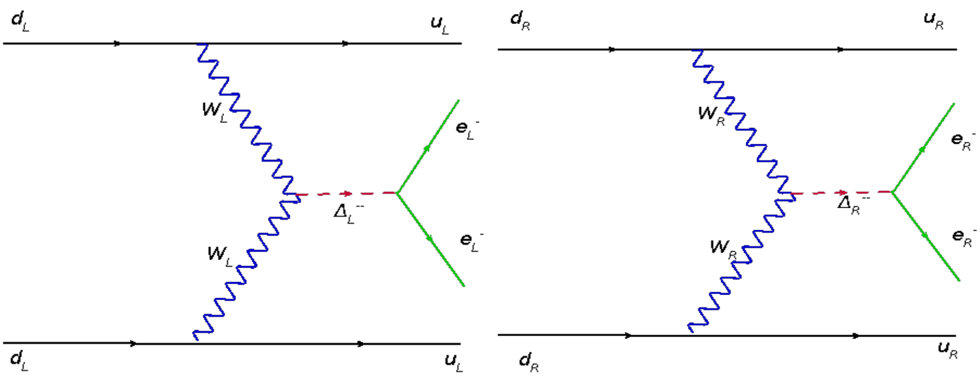

Figure 1.7: Feynman diagrams corresponding to neutrinoless double beta decay due to $\Delta_L - W_L$ $\Delta_R - W_R$ contributions.

1.5.4 Baryon Asymmetry of the Universe

Right after the big bang, a huge amount of energy was produced. In fact the universe then was so hot and dense and hence energetic for the spontaneous creation of particle and antiparticle pairs. It was expected that these particle-antiparticle pairs come together to leave a sea of radiation. But what is observed is otherwise. There must be a predominance of matter over antimatter produced for which the present observable universe is matter-dominated, the proof of our very existence. This matter-antimatter asymmetry is considered an important cosmological puzzle among the physics community. That no primordial antimatter is found in our observable universe has been experimentally determined. From the analysis of Wilkinson Microscopy Anisotropy Probe (WMAP), combined with the baryon acoustic oscillations data, it is found that the baryon to photon ratio of the number density has been measured to the unprecedented precision of less than 10% as [131]

$$\eta_B = \frac{n_B - n_{\overline{B}}}{n_\gamma} = (6.20 \pm 0.15) \times 10^{-10}, \tag{1.94}$$

where, n_B, $n_{\overline{B}}$, n_γ are the number densities of baryon, anti-baryon and photon respectively. This result is also consistent with the results obtained from big bang nucleosynthesis (BBN) [132, 133, 134, 135]. It is also called the baryon asymmetry of the universe (BAU). Many even suggest that neutrinos could play a significant role in this asymmetry. However, the most essential conditions for baryon asymmetry has been postulated long ago by Sakharov [136] which are,

1. Baryon number violation
2. Charge (C) and CP (Charge-Parity) violation
3. Departure from thermal equilibrium

Although SM contains the above mentioned three Sakharov ingredients, yet not in a sufficient amount to explain the observed asymmetry. Thus the explanation of the cosmological BAU requires new physics beyond the SM, either by the introduction of some new sources of CP violation and out of equilibrium conditions like the decay of some new particles or by modification of the electroweak phase transition. Various mechanisms has been proposed to explain the observed asymmetry, like GUT baryogenesis [137, 138], electroweak baryogenesis [139, 140], leptogenesis [141], Affleck-Dine mechanism [142, 143] etc. However, we have picked up two of these mechanisms in this thesis, which we will briefly discuss below.

1.5.4.1 Leptogenesis

Many appealing theoretical models have been proposed to explain this tiny value of η_B that quantifies the baryon asymmetry of the universe. One of the popular mechanism to generate the BAU is by leptogenesis. It is a mechanism where lepton asymmetry created before the electroweak phase transition gets converted to BAU via the B+L violating processes sphaleron processes, which converts any primordial L asymmetry or B-L asymmetry into a baryon asymmetry. The origin of this asymmetry may start from the leptons (including neutrino). A super-heavy counterpart of the light neutrino is assumed to be present in the early universe, the decay of whose producing more matter than antimatter could create the asymmetry. This heavy neutrino simultaneously plays an important role in explaining the tiny neutrino mass via the seesaw mechanism as has been discussed before. Such a realization of leptogenesis by the decay of the heavy neutrinos in out of equilibrium condition with Majorana masses considerably larger than the critical temperature $T_c \approx 100 - 200$ GeV transforming as singlets under the SM gauge group was first proposed by Fukugita and Yanagida [141]. Such a process satisfies all the three basic Sakharov's conditions for baryogenesis. The necessary CP asymmetry is provided by the interference between the tree-level and the one-loop decay diagrams. The departure from the thermal equilibrium occurs when the Yukawa interactions are passably slow and temperatures above the electroweak scale.

The lepton number violation in this regard comes from the decay of the heavy neutrinos into a lepton and a Higgs doublet, $N_i \rightarrow L + \phi^c$ and its respective CP conjugate process, $N_i \rightarrow L^c + \phi$ which can occur at both tree and one loop levels [145]. Hence, their CP-violating asymmetry ϵ_i which arises from the interference between the tree-level amplitude and its self-energy correction [144] is defined as,

$$\epsilon_i = \frac{\Gamma\left(N_i \rightarrow l + \phi^c\right) - \Gamma\left(N_i \rightarrow l^c + \phi\right)}{\Gamma\left(N_i \rightarrow l + \phi^c\right) + \Gamma\left(N_i \rightarrow l^c + \phi\right)}. \tag{1.95}$$

The scale of the masses of these RH neutrinos is model dependent and may have different range in different model. It is of the order of a few TeV in Left-Right Symmetric Model[39,40] or certain E6 [41] models up to 10^{16} GeV in typical Grand Unified Theories (GUTs) such as SO(10) [42,43] models.

1.5.4.2 Resonant Leptogenesis

In a TeV scale LRSM, what we are concerned with is another kind of leptogenesis mechanism known as resonant leptogenesis (RL) [146, 145, 147, 148, 149] in which self-energy effects of the heavy neutrino on the leptonic asymmetry becomes dominant and get resonantly enhanced up to order one. In this framework, the presence of the RH neutrinos (type I SS) and the scalar triplets (type II SS) suggests their decays which could give rise to the lepton asymmetry. That is, the net lepton symmetry generated will be a combination of the asymmetry created due to the two seesaw terms. However, in this thesis, we will basically consider the decay of the heavy RH neutrinos for generating lepton asymmetry, i.e., the contribution from the type I seesaw only. The decay of the scalar triplet Δ_L would not much affect our result as above TeV scale, the decay of RH neutrinos are in thermal equilibrium and hence they would wash out any kind of pre-existing lepton asymmetry and so we have ignored it. So the dominant contribution would come from the type I seesaw term, that is from the decay of the RH neutrino. For RL, a basic requirement is that a pair of heavy Majorana neutrinos must have a mass difference comparable to their decay widths (i.e., $M_i - M_j \approx \Gamma$). We can find several constructions of RL models appearing in the literature.

The decay rates of the heavy neutrino decay processes are governed by the Yukawa couplings, and is given by,

$$\Gamma_i = \left(Y_\nu^\dagger Y_\nu\right)_{ii} \frac{M_i}{8\pi}. \tag{1.96}$$

In case $M_i - M_j \approx \Gamma$, the CP asymmetry becomes very large (even of order 1). The CP violating asymmetry ϵ_i is thus given by,

$$\epsilon_i = \frac{Im\left[\left(Y_\nu^\dagger Y_\nu\right)_{ij}^2\right]}{\left(Y_\nu^\dagger Y_\nu\right)_{11}\left(Y_\nu^\dagger Y_\nu\right)_{22}} . \frac{\left(M_i^2 - M_j^2\right) M_i \Gamma_j}{\left(M_i^2 - M_j^2\right) + M_i^2 \Gamma_j^2}, \tag{1.97}$$

where,

$$\frac{Im\left[\left(Y_\nu^\dagger Y_\nu\right)_{ij}^2\right]}{\left(Y_\nu^\dagger Y_\nu\right)_{11}\left(Y_\nu^\dagger Y_\nu\right)_{22}} \approx 1. \tag{1.98}$$

The variables i, j run over 1 and 2, $i \neq j$. The CP violating asymmetries ϵ_1 and ϵ_2 can give rise to a net lepton number asymmetry, provided the expansion rate of the universe is larger than Γ_1 and Γ_2. This can further be partially converted into baryon asymmetry of the universe by B+L violating sphaleron [150] processes. The net baryon asymmetry is then calculated using [144],

$$\eta_B \approx -0.96 \times 10^{-2} \sum_i \left(k_i \epsilon_i\right), \tag{1.99}$$

k_1 and k_2 being the efficiency factors measuring the washout effects linked with the out of equillibrium decay of N_1 and N_2. We can define the parameters, $K_i \equiv \frac{\Gamma_i}{H}$ at temperature, $T = M_i$, $H \equiv \frac{1.66\sqrt{g_*}T^2}{M_{Planck}}$ is the Hubble's constant with $g_* \simeq 107$ (effective numbers of degrees of freedom available at temperature T) and $M_{Planck} \equiv 1.2 \times 10^{19}$ GeV is the Planck mass. For simplicity, the efficiency factors, k_i can be calculated using the formula [151],

$$k_1 \equiv k_2 \equiv \frac{1}{2}\left(\sum_i K_i\right)^{-1.2}, \tag{1.100}$$

which holds validity for two nearly degenerate heavy Majorana masses and $5 \leq K_i \leq 100$.

RL has several advantages, like their predictions for the BAU, are almost independent of the primordial L-number, B-number and heavy neutrino abundances. Because, in RL scenarios, the decay widths of the heavy Majorana neutrinos can be significantly larger than the Hubble expansion rate H of the Universe. As a result, the heavy Majorana neutrinos can rapidly come into thermal equilibrium from their decays, inverse decays and scatterings with the other SM particles even if there were no heavy Majorana neutrinos at high temperatures. Moreover, in this high-temperature regime, any pre-existing lepton asymmetry will rapidly be driven to zero, due to the L-violating inverse decays and scattering processes which are almost in thermal equilibrium. As the Universe cools down, a net lepton asymmetry can be created at temperatures just below the heavy neutrino mass as a consequence of the aforementioned CP-violating resonant enhancement that occurs in RL models. This L asymmetry will survive wash-out effects and finally get converted into the observed BAU.

1.5.5 Lepton Flavor Violation

Charged lepton flavor violation (CLFV) [81] is a clear signal of new physics; it directly addresses the physics of flavor and generations. The search for CLFV has continued from the early 1940s, when the muon was identified as a separate particle, until today. Certainly, in the LHC era, the motivations for continued searches are clear and have been covered in many reviews.

The discovery of neutrino oscillation has provided concrete evidence of the fact that neutrinos are massive as well as the violation of the lepton flavor during the propagation of the neutrinos. Lepton flavor is consequently a broken symmetry and the SM has to be adapted to incorporate massive neutrinos and thus we can also hope that LFV will be visible in the charged lepton sector. The exact mechanism of LFV being unknown, its study is of large interest as it is linked to neutrino mass generation, CP violation and new physics BSM. The LFV effects from new particles at the TeV scale are naturally generated in many models and therefore considered to be a prominent signature for new physics. In LRSM, where electroweak symmetry is broken dynamically, an experimentally accessible amount of LFV is predicted in a large region of parameter space. In a wide range of models for physics BSM, highest sensitivity in terms of BR is expected for $\mu \to 3e$ decay process. A new search for LFV decay $\mu \to 3e$ with an unprecedented sensitivity of $< 1.0 \times 10^{-16}$ as proposed would provide a unique opportunity for discoveries of physics BSM in the coming years.

The discovery of neutrino mass and neutrino oscillations guarantees that SM CLFV must occur through oscillations in loops. Such transitions are suppressed by$(\frac{\Delta m_\nu^2}{\Delta M_W^2})$. Thus giving potentially high sensitivity to LFV reactions in models BSM. For a review of CLFV in SM and beyond, please refer to [81]. Though the usual light neutrino contribution to CLFV is negligible, presence of heavy neutrinos in BSM frameworks can give rise to observable CLFV [152, 153, 155, 156, 157, 158, 159]. In LRSM, sizable CLFV occurs dominantly due to the contributions arising from the additional scalars and the heavy neutrinos. Within the LRSM, CLFV naturally occurs due to potentially large flavor violating couplings of the heavy RH neutrinos and Higgs scalars with charged leptons. Among the various processes that violate lepton flavor, the most relevant ones are the rare leptonic decay modes of the muon, notably, $(\mu \to e\gamma)$ and $(\mu \to 3\mathrm{e})$. The best upper limit for the branching ratio (BR) of these processes are provided by MEG collaboration [161] and SINDRUM experiment [162] which provide the corresponding upper limit as BR $(\mu \to e\gamma) < 4.2 \times 10^{-13}$ and BR $(\mu \to 3e) < 1.0 \times 10^{-12}$ respectively.

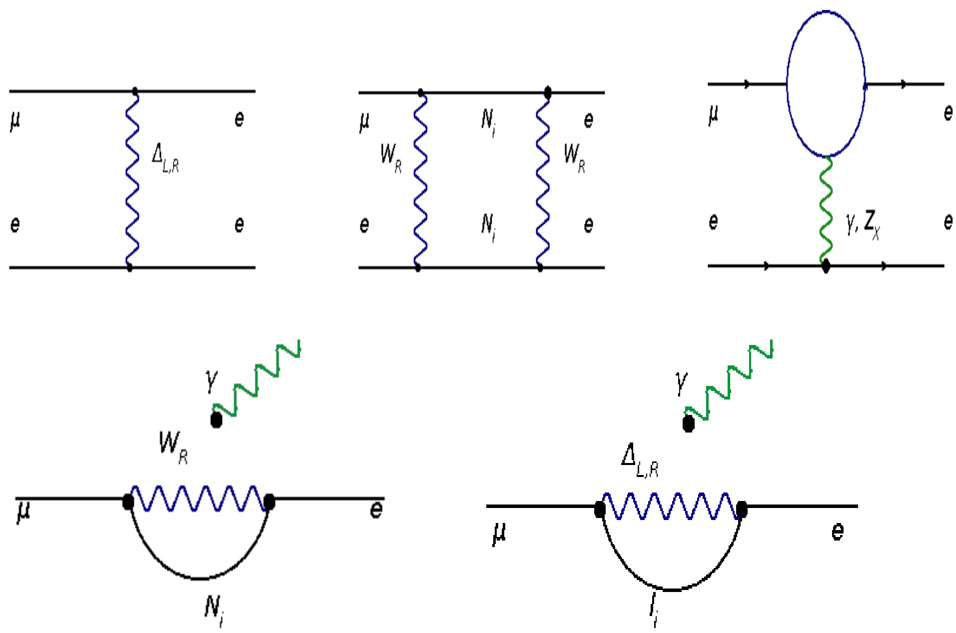

Figure 1.8: The Feynman diagrams corresponding to the decay modes, $(\mu \to 3\mathrm{e})$ (top) and $(\mu \to e\gamma)$ (bottom) in LRSM. The external photon line may be attached to any charged particle line. (Figures like [159])

Defining the decay rates (from reference [155])as,

$$\Gamma_\mu \equiv \Gamma\left(\mu^- \to \mathrm{e}^- \nu_\mu \overline{\nu_\mathrm{e}}\right), \Gamma^{\mathrm{Z}}_{\mathrm{capt}} \equiv \Gamma\left(\mu^- + \mathrm{A}\left(\mathrm{Z},\mathrm{N}\right) \to \nu_\mu + \mathrm{A}\left(\mathrm{Z}-1,\mathrm{N}+1\right)\right). \quad (1.101)$$

The relevant branching ratios (BR) for the processes are,

$$\mathrm{BR}_{\mu\to e\gamma} \equiv \frac{\Gamma\left(\mu^+ \to \mathrm{e}^+\gamma\right)}{\Gamma_\mu}, \tag{1.102}$$

$$\mathrm{BR}^{\mathrm{Z}}_{\mu\to \mathrm{e}} \equiv \frac{\Gamma\left(\mu^- + \mathrm{A}\left(\mathrm{N},\mathrm{Z}\right) \to \mathrm{e}^- + \mathrm{A}\left(\mathrm{N},\mathrm{Z}\right)\right)}{\Gamma^{\mathrm{Z}}_{\mathrm{capt}}}, \tag{1.103}$$

$$\mathrm{BR}_{\mu\to 3\mathrm{e}} \equiv \frac{\Gamma\left(\mu^+ \to \mathrm{e}^+\mathrm{e}^-\mathrm{e}^+\right)}{\Gamma_\nu}. \tag{1.104}$$

Adopting the notations of [160, 127] the branching ratio of the process $\mu \to 3e$ mediated by doubly charged scalars can be written as,

$$\mathrm{BR}\left(\mu \to 3\mathrm{e}\right) = \frac{1}{2}\left|\mathrm{h}'_{\mu\mathrm{e}}\mathrm{h}'^{\,*}_{\mathrm{ee}}\right|^2 \left(\frac{\mathrm{M}_{\mathrm{W_L}}{}^4}{\mathrm{M}^{++4}_{\Delta_\mathrm{L}}} + \frac{\mathrm{M}_{\mathrm{W_R}}{}^4}{\mathrm{M}^{++4}_{\Delta_\mathrm{R}}}\right), \tag{1.105}$$

where h'_{ij} describes the respective lepton-scalar couplings given by,

$$\mathrm{h}'_{\mathrm{ij}} = \sum_{\mathrm{n}=1}^{3} \mathrm{V}_{\mathrm{in}}\mathrm{V}_{\mathrm{jn}}\left(\frac{\mathrm{M_n}}{\mathrm{M}_{\mathrm{W_R}}}\right), \mathrm{i},\mathrm{j} = \mathrm{e},\mu,\tau. \tag{1.106}$$

with V being one of the lepton mixing matrices which will discuss and calculate in subsequent chapters.

The branching ratio for the CLFV process $\mu \to e\gamma$ is given by (as explained in [160]),

$$\mathrm{BR}\left(\mu \to \mathrm{e}\gamma\right) = \frac{3\alpha_{\mathrm{em}}}{2\Pi}\left(|\mathrm{G}^\gamma_\mathrm{L}|^2 + |\mathrm{G}^\gamma_\mathrm{R}|^2\right), \tag{1.107}$$

where, α_{em} is the fine structure constant defined as $\alpha_{em} = \frac{e^2}{4\pi}$, G^γ_L and G^γ_R are the form factors given by,

$$G^\gamma_L = \sum_{i=1}^{3}\left(S_{\mu i}{}^*S_{ei}G^\gamma_1(x_i) - V_{\mu i}S_{ei}\xi e^{i\zeta}G^\gamma_2(x_i)\frac{M_i}{m_\mu} + V_{\mu i}V_{ei}{}^*y_i\left[\frac{2}{3}\frac{M^2_{W_L}}{M^2_{{\Delta_L}^{++}}} + \frac{1}{12}\frac{M^2_{W_L}}{M^2_{{\Delta_L}^{+}}}\right]\right) \tag{1.108}$$

$$G^\gamma_R = \sum_{i=1}^{3}\left(V_{\mu i}V_{ei}{}^*|\xi|^2G^\gamma_1(x_i) - S_{\mu i}{}^*V_{ei}{}^*\xi e^{-i\zeta}G^\gamma_2(x_i)\frac{M_i}{m_\mu} + V_{\mu i}V_{ei}{}^*\left[\frac{M^2_{W_L}}{M^2_{W_R}}G^\gamma_1(y_i) + \frac{2y_i}{3}\frac{M^2_{W_L}}{M^2_{{\Delta_R}^{++}}}\right]\right). \tag{1.109}$$

The terms $x_i = \left(\frac{M_i}{M_{W_L}}\right)^2$ and $x_i = \left(\frac{M_i}{M_{W_R}}\right)^2$, $M_{\Delta_{L,R}}$ are the masses of the left and right scalar triplets, $\mathrm{M_i}(\mathrm{i} = 1, 2, 3)$ are the masses of the RH neutrinos. V is the mixing matrix of the RH neutrinos. ζ is the phase of the VEV k_2, whereas the left-right gauge boson mixing parameter, ξ is also very small. S being the light-heavy neutrino mixing. Again the loop functions $G^{\gamma}_{1,2}(a)$ are defined as,

$$G^{\gamma}_1(a) = -\frac{2a^3 + 5a^2 - a}{4(1-a)^3} - \frac{3a^3}{2(1-a)^4}\ln a \tag{1.110}$$

$$G^{\gamma}_2(a) = \frac{a^2 - 11a + 4}{1(1-a)^2} - \frac{3a^2}{(1-a)^3}\ln a \tag{1.111}$$

A new stringent upper bound on the decay rate of the process $\mu \to e\gamma$ was recently reported by the MEG collaboration. The BR ratio for this LFV process as given by MEG is $< 4.2 \times 10^{-13}$ at 90% CL [161]. While for the process $\mu \to 3e$ it is $< 1.0 \times 10^{-12}$ as obtained by the SINDRUM experiment[162].

1.6 Flavor symmetry in particle physics

Symmetries playsa very significant role in particle physics. The mathematical description of symmetries uses group theory, examples of which are SU(2) and SU(3). In particle physics, there are many examples of symmetries and their associated conservation laws which are given by the famous Noether's theorem. There are also cases where a symmetry is broken, and the mechanism has to be understood. The breaking of electroweak symmetry and the associated Higgs field has already been discussed before. Continuous (and local) symmetries such as Lorentz, Poincare and gauge symmetries are essential to understand several phenomena, which happen in particle physics like strong, weak and electromagnetic interactions among particles. Discrete symmetries such as Charge Conjugation (C), Parity (P) and Time Reversal (T) are also important. Furthermore, Abelian discrete symmetries, BSM Z_N, are also often imposed to control allowed couplings in model building for particle physics, in particular, model building BSM. Z_N is a cyclic group which is a finite subgroup of $SO(2)$ generated by $e^{\frac{2\pi}{n}}$ or we can say it describes a symmetry of a plane figure which remains invariant after a rotation of $\frac{2\pi}{n}$ degrees. In this thesis we will use this discrete abelian groups Z_8 and the simplest group Z_2 to restrict the appearance of certain elements of the neutrino mass matrix, also called the texture zero mass matrix. We have discussed it in chapter 5.

The basic motive behind the study of the gauge structure of the SM and its various extensions is understanding the fundamental workings of the universe. There are many free parameters in the SM and BSM frameworks in neutrino mass terms and their origin mainly comes from the flavor sector, i.e., the Yukawa couplings of quarks and leptons. The flavor symmetries has been introduced to govern these Yukawa couplings (the interactions among the leptons, quarks and Higgs boson) in the three generations although the origin of the generations is an unknown factor. With the discovery of neutrino masses and mixing, the study of the flavor symmetries has increased. Experiments of the neutrino oscillation measuring precisely the mixing angles and mass squared differences indicates a nearly tri-bimaximal mixing (TBM) for three flavors in the lepton sector. It is seen that the lepton sector is less hierarchical and has larger mixing angles in comparison to the quark sector which exhibits a strongly hierarchical mass spectrum and mixing angles which are comparatively small. So it is utmost important to find a natural model which leads to the observed mixing patterns of the quarks and leptons with sufficient accuracy. Models with non-abelian discrete flavor symmetries are mostly known to be able to describe the large mixing angles of the lepton sector to derive experimental values of quark/lepton masses and mixing angles. A very intensive discussion in non-abelian discrete flavor symmetries is seen in the review by Altarelli and Feruglio [163]. In general, for such model building, a discrete flavor group is broken to different subgroups in the charged lepton and neutrino sectors and the mismatch between the two subgroups allows one to predict the mixing matrix. The flavor symmetry may be a remnant of the higher dimensional space-time symmetry. It was shown how the flavor symmetry A4 (or S4) can arise if the three fermion generations are taken to live on the fixed points of a specific 2-dimensional orbifold.

The different BSM models address the challenge to give a reasonable description of the lepton sector. As previously stated, additional symmetries are added to the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$ and extend the field content. In the context of LRSM also, flavor symmetry has been addressed in [164], where the often used A_4 group has been implemented with some scalar fields (also called flavons) which are charged under the flavor symmetry but are neutral under the gauge symmetry. The general idea is to assign all fields to suitable irreducible representations of the overall symmetry and to construct invariants under it. The new structures, mainly arising from the new symmetry, are then used to explain the observed lepton mixing. Non-abelian discrete groups like S_4, A_4 are mostly used for model building, as they are the smallest group with a three dimensional (triplet) representation. The multiplication rules of the used irreducible representations are essential in this context. The reader is directed to [163] for further details.

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