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LEPTOGENESIS AS THE ORIGIN OF MATTER-ANTIMATTER ASYMMETRY IN EXTRA DIMENSIONAL AND SUPERSYMMETRIC MODELS

by

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Contents

Abstract

The Standard Model of particle physics is currently the best description of fundamental particles and their interactions. All particles save the Higgs boson have been observed in particle accelerator experiments over the years. Despite the predictive power the Standard Model there are many phenomena that the scenario does not predict or explain. Among the most prominent dilemmas is matter-antimatter asymmetry, and much effort has been made in formulating scenarios that accurately predict the correct amount of matter-antimatter asymmetry in the universe. One of the most appealing explanations is baryogenesis via leptogenesis which not only serves as a mechanism of producing excess matter over antimatter but can also explain why neutrinos have very small non-zero masses.

Interesting leptogenesis scenarios arise when other possible candidates of theories beyond the Standard Model are brought into the picture. In this thesis, we have studied leptogenesis in an extra dimensional framework and in a modified version of supersymmetric Standard Model. The first chapters of this thesis introduce the standard cosmological model, observations made on the photon to baryon ratio and necessary preconditions for successful baryogenesis. Baryogenesis via leptogenesis is then introduced and its connection to neutrino physics is illuminated. The final chapters concentrate on extra dimensional theories and supersymmetric models and their ability to accommodate leptogenesis. There, the results of our research are also presented.

List of publications

This PhD thesis is based on the research articles [1, 2, 3]:

I J. Maalampi, I. Vilja and H. Virtanen, *Thermal leptogenesis in a 5D split fermion scenario with bulk neutrinos*, Phys. Rev. D**82**, 013009 (2010).

II H. Kuismanen, J. Maalampi and I. Vilja, *Numerical study of leptogenesis in a 5D split fermion model with bulk neutrinos*, Phys. Rev. D**83**, 053005 (2011).

III H. Kuismanen, J. Pelto and I. Vilja, *Leptogenesis in B-L gauged SUSY with MSSM Higgs sector*, Phys. Rev. D**83**, 055001 (2011).

Chapter 1

Introduction

Our observable universe is described by the interplay between the large scales governed by general relativity and microscopic scales that are described by the Standard Model (SM) of particle physics. These two seemingly disparate realms came together 13.7 billion years ago in the Big Bang that gave birth to the universe. While to this date the fundamental theory of the very early universe remains a mystery also the current composition of the universe is a puzzle. Observations relying on the standard cosmological model show that 5 % of the energy density of the universe consists of baryonic matter while the rest of the energy density consists of dark matter and dark energy whose origin and composition are currently unknown. Baryonic matter is ordinary visible matter made up of nucleons but the SM does not explain its origin or its abundance.

According to the SM, equal amounts of particles and antiparticles should have been generated in the early universe. A particle and its antiparticle are equal in mass but have opposite quantum numbers, and upon collision they may annihilate each other producing mainly electromagnetic radiation. Thus, 13.7 billion years ago elementary particles and their antiparticles can be expected to have annihilated, resulting in a universe with only radiation. Yet, we are very much here and are made of baryonic matter, and on larger scales up to tens of Mpc galaxies are thought to consist of matter rather than antimatter [5]. One endeavor that is expected to settle the issue is the ongoing AMS-02 experiment that is designed to look for indications of large antimatter domains. The smoking gun pointing towards the existence of antimatter regions in the universe is the detection of antihelium or larger antinuclei [6].

To explain the matter-antimatter asymmetry of the universe we need to go beyond the SM. Extensions of the SM materialized in Grand Unified Theories (GUT) that were first proposed in the 1970s. They unify electromagnetic, weak and strong interactions and allow for baryogenesis. The hint towards theories beyond the SM followed in this thesis comes from neutrino physics and a mechanism of generating baryon asymmetry via lepton asymmetry. This is known as baryogenesis via leptogenesis [7]. The driving force behind leptogenesis is that somehow in the early universe a net lepton number was produced and this asymmetry is further converted to the baryon asymmetry we observe today.

The SM is widely considered to be an effective theory description of elementary particles and interactions so that at higher energy scales some fundamental theory takes over. This fundamental theory would explain the fermion masses and mixings, provide a solution to the strong CP problem, and on more aesthetic grounds possibly unify all four interactions and remove the need of fine tuning which comes about because of the wide gap between the electroweak and gravity scales. Leptogenesis on its own introduces extensions to the SM, but among the best known candidates for a fundamental theory are supersymmetric version of the SM [8] and superstrings [9]. The former doubles the particle content of the SM by relating each fermionic particle with a bosonic partner and vice versa. It modifies the SM in such a way that the three interactions, electromagnetic, weak and color interactions become equal in strength at an energy scale far beyond the realm of the SM. Superstring theories have supersymmetry built in them and they postulate that the fundamental building blocks of the universe are one-dimensional strings instead of point particles. These scenarios can incorporate all four interactions, namely a consistent description of gravity as well as the three interactions of the SM can be obtained. One peculiar feature of string theories is that they require several additional dimensions besides the four spacetime dimensions we are used to. One ramification of string theories is that these extra dimensions could be as large as the millimeter scale and thus observable to us.

The objective of this thesis is to find ways how leptogenesis can be accommodated with the other two possible extensions of the SM, namely supersymmetric and extra dimensional models. The research articles [1] and [2] both consider an extra dimensional model that has a mechanism of generating lepton asymmetry. While article [1] concentrates on the particle structure and inherent leptogenesis mechanism of the model, article [2] explores the parameter space which enables leptogenesis in the first place. Paper [3] contemplates a supersymmetric model that can generate lepton asymmetry through an intermediate step before giving rise to the conventional supersymmetric version of leptogenesis.

The structure of this thesis is as follows: Chapter 2 gives a general introduction to the currently known early stages of the universe with emphasis on the generation of matter-antimatter asymmetry. Chapter 3 introduces leptogenesis and also the connection between leptogenesis and neutrino masses. Chapter 4 presents extra dimensional theories and the research done in papers [1] and [2]. Supersymmetry and supersymmetric leptogenesis is considered in chapter 5 where also the research of paper [3] is discussed. Finally, chapter 6 summarizes the topics covered and presents some interesting prospects for leptogenesis.

Notation

Natural units $\hbar = c = k_B = 1$ are used throughout the thesis.

Chapter 2

The early universe and matter-antimatter asymmetry

The basic setup of cosmology is based on the homogeneous, isotropic and expanding universe while particles and their interactions are described by the SM. In this chapter, the evolutionary stages of the early universe with emphasis on the generation of net baryon number are outlined. In section 2.1 the standard cosmological model is presented, section 2.2 showcases the premises of generating net baryon number and section 2.3 reviews several mechanisms of baryogenesis.

2.1 Standard cosmological model

The cornerstones of the standard cosmological model are Einstein's theory of General Relativity and the SM that determines how the building blocks of matter were formed. Important clues about the dynamics of the universe came from observations made by Edwin Hubble in 1929 [13]. His observations suggest that the universe is expanding as galaxies seem to be receding away from us. This notion makes a compelling case for the fact that the universe has been hotter and denser in the past, which implies that space and time were born in a so called Big Bang.

Another essential ingredient originates from the discovery of Penzias and Wilson [14]. They observed the redshifted Cosmic Microwave Background (CMB) radiation which is an echo from the early universe when it was a few hundred thousand years old. CMB is a perfect fit to the blackbody spectrum with temperature $T = 2.725$ K and anisotropies $\sim 10^{-5}$ [15].

The final component completing the picture concerns the large scale structure of the universe. The Sloan Digital Sky Survey [16] and 2dF [17] are mappings of the large scale structure and they confirm the universe is isotropic, *i.e.* it looks the same from all directions, at scales larger than 100 Mpc.

2.1.1 Friedmann-Robertson-Walker universe

The Friedmann-Robertson-Walker (FRW) universe is homogeneous and isotropic, *i.e.* the energy density is uniform and the universe looks the same from all angles at large scales. It is also a description of an expanding universe.

The Robertson-Walker metric and Einstein equations

The Robertson-Walker (RW) metric is

$$ds^2 = -dt^2 + a(t)^2 \left(\frac{dr^2}{1 - Kr^2} + r^2 d\theta^2 + r^2 \sin^2\theta d\phi^2 \right). \tag{2.1}$$

Here, r, θ and ϕ are polar comoving coordinates while K measures the spatial curvature and $a(t)$ is the scale factor. The metric enters the Einstein equation

$$R_{\mu\nu} - \frac{1}{2} R g_{\mu\nu} = 8\pi G T_{\mu\nu}, \tag{2.2}$$

that determines the dynamics. Newton's gravitational constant is denoted by G. The left hand side of (2.2) depends on the metric while the right hand side describes the matter content. The energy-momentum tensor must contain the same symmetries as the metric, which gives the perfect fluid form for the energy-momentum

$$T_{\mu\nu} = (\rho + p) u_\mu u_\nu + p g_{\mu\nu}. \tag{2.3}$$

Here, ρ and p denote energy density and pressure, respectively, with $u^\mu = (1, 0, 0, 0)$ as the four velocity of the fluid in comoving coordinates.

The Friedmann equations

The Einstein equation (2.2) describes the evolution of the scale factor in terms of ρ and p of the cosmic fluid. By invoking (2.1), (2.2) and (2.3) we obtain for the 00-component

$$H^2 = \frac{8\pi G}{3}\rho - \frac{K}{a^2}, \tag{2.4}$$

where the Hubble parameter $H = \dot{a}/a$. The spatial component of the Einstein equation gives

$$\frac{\ddot{a}}{a} = -\frac{4\pi G}{3}(\rho + 3p). \tag{2.5}$$

Equations (2.4) and (2.5) are called the Friedmann equations.

The energy continuity equation can be derived from the Friedmann equations and it reads

$$\dot{\rho} = -3(\rho + p)\frac{\dot{a}}{a}. \tag{2.6}$$

Rewriting (2.4) as

$$\rho = \rho_c + \frac{3K}{8\pi G a^2} \tag{2.7}$$

introduces the critical density $\rho_c \equiv 3H^2/(8\pi G)$. The relation between the energy density ρ and ρ_c also determines the nature of the curvature

$$\begin{aligned} \rho &< \rho_c \Rightarrow K < 0 \\ \rho &= \rho_c \Rightarrow K = 0 \\ \rho &> \rho_c \Rightarrow K > 0. \end{aligned} \tag{2.8}$$

By defining the density parameter $\Omega(t)$ as

$$\Omega(t) \equiv \frac{\rho(t)}{\rho_c(t)} \tag{2.9}$$

we can rewrite (2.4) yet in another form

$$\Omega(t) = 1 + \frac{K}{H^2 a^2}. \tag{2.10}$$

In solving (2.6) we need to consider different alternatives for the equation of state $p(\rho)$. The cases that arise are

- Matter refers to nonrelativistic matter consisting of particles with velocities $v \ll 1$ and $p \ll \rho$. Thus, we can approximate $p \approx 0$, which yields $\rho \propto a^{-3}$.

- Radiation refers to ultrarelativistic particles whose rest masses are much smaller than the particle energy. For ultrarelativistic particles $p = \rho/3$ that upon insertion into (2.6) gives $\rho \propto a^{-4}$.

- Vacuum energy has constant energy density, $\rho = constant$, and (2.6) gives the equation of state $p = -\rho$.

Based on the scale factor dependences of the different energy density components we can deduce that the early universe was radiation-dominated. As the scale factor a grows with time, the other components (matter, curvature and radiation) begin to dominate.

2.1.2 Inflation

There are various problems in the Big Bang model (next subsection). First, the CMB anisotropies are very small in magnitude $\sim 10^{-5}$ [15] and they are correlated at all scales. However, at the time of last scattering, *i.e.* when the photons decouple from matter, there were regions that did not have a causal connection. This is known as the horizon problem [18]. Second, there is a spatial flatness problem in the universe [18]. To demonstrate this we have written (2.4) in the form

$$\Omega - 1 = \frac{K}{a^2 H^2}. \tag{2.11}$$

In the ΛCDM model, the present day density parameter is very close to one, $|\Omega(t_0) - 1| \lesssim 10^{-2}$. At the time of Big Bang nucleosynthesis (BBN, reviewed in the next subsection) $|\Omega(t_{\mathrm{BBN}}) - 1| \lesssim 10^{-17}$, which indicates strong fine tuning. Finally, the relic problem deals with theories beyond SM. Phase transitions in GUTs can produce magnetic monopoles, cosmic strings and domain walls. The energy density would thus receive additional contributions that are much higher than the observed energy density. What could remedy all these problems related to causal connection, spatial flatness, structure formation [19] and relics is inflation [20, 21], a period of accelerated expansion of the universe prior to the Hot Big Bang.

The definition of inflation is an era of accelerated expansion of the universe. The rate of increase of the scale factor is increasing which means $\ddot{a} > 0$ or equivalently

$$\frac{d}{dt}\frac{1}{aH} < 0. \tag{2.12}$$

This shows that the comoving Hubble length decreases with time bringing more regions into causal connection. As for the flatness problem, it can be seen from equation (2.11) that with inflation the density parameter $\Omega(t)$ is driven towards 1.

During inflation the energy density is dominated by a scalar field φ, the inflaton, that can have negative pressure, $p = w\rho,\ w < 0$. The so called slow roll approximation states that $\frac{1}{2}\dot{\varphi}^2 \ll V(\varphi)$ so it is in fact the potential $V(\varphi)$ term that dominates the energy density. Initially, φ is far from the minimum of $V(\varphi)$. Gradually φ is being pulled towards the minimum by the potential. The amount of inflation can be measured by

$$N(t) \equiv \ln \frac{a(t_{\mathrm{end}})}{a(t)} \tag{2.13}$$

which tells us how much the scale factor has grown by the end of inflation.

After inflation, the inflaton field begins to oscillate at the bottom of the potential $V(\varphi)$ and the energy stored in the inflaton potential $V(\varphi)$ is transferred to particles which become thermalized. This process is called reheating which is responsible for the particle content of the universe in the subsequent Hot Big Bang. The allowed range for the reheating temperature is $10^{-2}\mathrm{GeV} < T_R < 10^{16}\mathrm{GeV}$ where the lower limit is set by Big Bang nucleosynthesis and upper limit by GUTs [19].

Inflation combined with supersymmetric models places more restrictions to the reheating temperature. These models contain the supersymmetric partner of the graviton, the gravitino. The overproduction of gravitinos during reheating would destroy the predictions of BBN by the decay of gravitinos via the gravitational interaction after BBN. Thus, to avoid the production of unwanted relics in supersymmetric theories the reheating temperature T_R should satisfy $T_R \lesssim 10^7$ GeV [22].

2.1.3 Hot Big Bang

The early universe is radiation-dominated and particle reaction rates keep up with the expansion of the universe. Thus, different particle species remain in thermal equilibrium. As the universe expands, it becomes harder to maintain sufficient reaction rates. With falling temperature the photon energies decrease and particle-antiparticle generation becomes energetically impossible. One after another particles and antiparticles annihilate. In the following a brief description of each annihilation process and subsequent stable matter formation is given.

Thermal history

After electroweak (EW) phase transition at temperature $T \sim 100$ GeV the universe undergoes the following stages

- t quarks annihilate at temperature $T \simeq 100$ GeV
- Higgs boson and electroweak gauge bosons $W^{\pm}$ and Z annihilate at temperatures $T < 80$ GeV.
- At $T \sim 10$ GeV the b and c quarks undergo annihilation, followed by the heaviest lepton, τ.
- QCD phase transition takes place at $T \sim 150$ MeV.
- pions and muons annihilate at temperature $T \sim 100$ MeV.
- At temperature $T \sim 1$ MeV neutrinos decouple and electron and positrons annihilate, and an excess of electrons remains.

During QCD phase transition quark-gluon interactions become important and quarks form bound states with the gluons. Three-quark systems are baryons and the lightest of these are the proton and neutron. Quark-antiquark pairs are called mesons out of which the lightest are pions, $\pi^{\pm}$ and π^0.

Big Bang Nucleosynthesis

After the QCD phase transition nucleons and antinucleons annihilate each other and a small excess of nucleons remain. In the process of Big Bang

Nucleosynthesis (BBN) these nucleons combine to form nuclei of the lightest elements in practice up to lithium. Due to the neutrino decoupling weak interactions needed to keep the proton-neutron-proton conversion going become less efficient and consequently neutrons decay into protons. This starts to happen at $T \sim 1$ MeV. As the decay reactions take place gradually, a portion of neutrons can combine with protons to form deuterium. This process is however hindered by the presence of a large number of photons w.r.t. baryons, and high-energy photons constantly break nuclei. Thus, the binding energy 2.2 MeV for deuterium is easily overcome until the temperature drops to $T = 0.06$ MeV. After a sufficient amount of deuterium is produced heavier nuclei are formed in secondary reactions. The most stable of these nuclei is ^{4}He.

The baryon number of the universe resides in the newly formed nuclei and the observed baryon-photon ratio $\eta \equiv n_B/n_\gamma$ in the universe is [23]

$$5.1 \times 10^{-10} < \frac{n_B}{n_\gamma} < 6.5 \times 10^{-10}. \tag{2.14}$$

A plot of baryon-to-photon ratio versus the abundances of light elements is shown in Fig. 2.1 where observations of light element abundances are presented alongside predictions from nucleosyntehsis.

Recombination and photon decoupling

As the temperature falls further, the leftover electrons and nuclei form neutral atoms (recombination), which causes the density of free electrons to fall. Photons have less and less electrons to interact with and so the photon mean free path rapidly grows and exceeds the horizon distance. Photons decouple from matter and the universe becomes transparent at $t =$380 000 years. These decoupled photons form the CMB and their temperature today is $T_0 = 2.725$ K.

Baryon and lepton number

We have given brief descriptions of the various stages in the early universe. After nucleon and electron annihilation processes a small excess of matter particles has remained; this has undergone recombination to form light elements up to lithium. The SM cannot explain the observed excess of

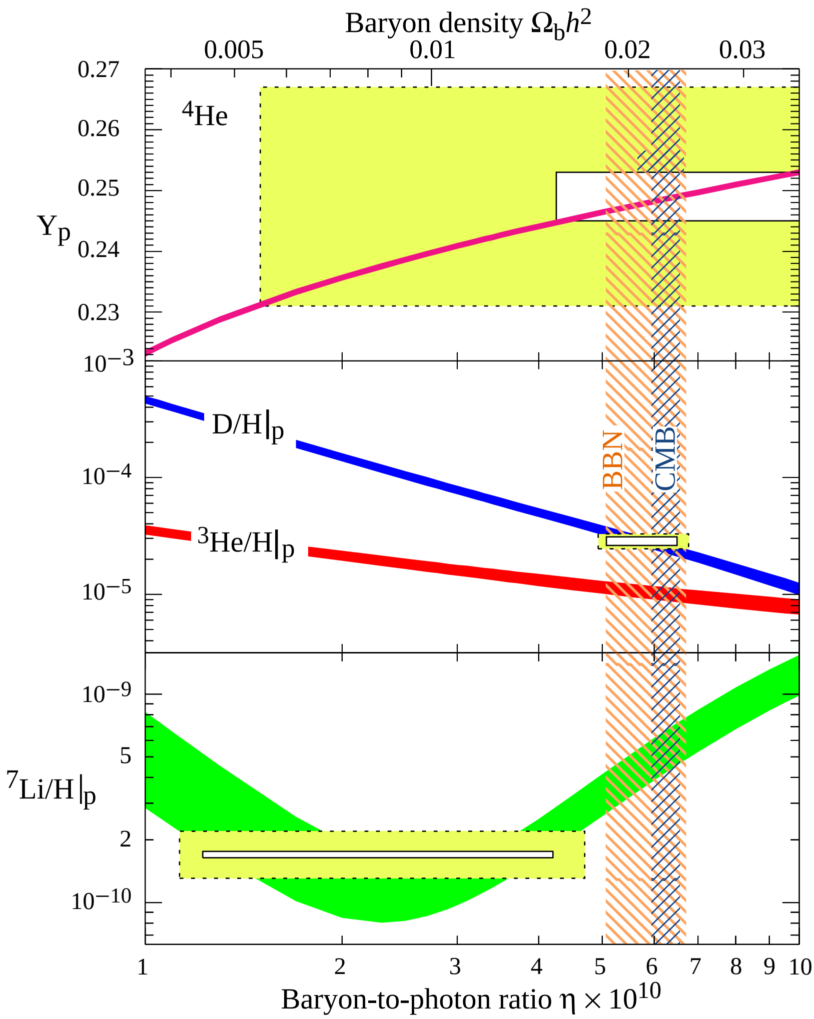

Figure 2.1: Picture from [23] shows the abundances of ^{4}He, D, ^{3}He and ^{7}Li as predicted by BBN. The bands show the 95 % CL range. Boxes indicate the observed light element abundances (smaller boxes: $\pm 2\sigma$ statistical errors; larger boxes: $\pm 2\sigma$ statistical and systematic errors). The narrow vertical band indicates the CMB measure of the cosmic baryon density, while the wider band indicates the BBN concordance range (both at 95 % CL).

baryons over antibaryons. If we assume no asymmetry between baryons and antibaryons at $T \sim \mathcal{O}(1 \text{ GeV})$, the annihilation $B + \bar{B} \to 2\gamma$ begins to dominate and [24]

$$\frac{n_B}{n_\gamma} = \frac{n_{\bar{B}}}{n_\gamma} \simeq 10^{-18}. \tag{2.15}$$

Yet, we observe the baryon-to-photon ratio (2.14) and structure in the universe in the form of *e.g.* stars and planets etc. Thus, the question is how should the SM be extended to include the mechanism of generating sufficient matter-antimatter asymmetry.

2.2 Sakharov conditions

Having gone through the crucial steps in the evolution of the early universe we are left with the problem matter-antimatter asymmetry. In 1967 Sakharov discovered three conditions that must be met in a model in order for it to produce net baryon number [25]:

- baryon number B violation
- C and CP violation
- departure from thermal equilibrium

Baryon number violating processes can generate net baryon number as opposed to baryon number conservation which would imply that the baryon number observed today dates back to baryon asymmetric initial conditions. Charge C and combined charge and parity CP violation ensure an asymmetry between matter and antimatter and that B violating processes proceed through unequal rates. In thermal equilibrium, inverse processes would wash out any pre-existing asymmetry based on the unitarity of the S matrix and CPT theorem. Thus, departure from thermal equilibrium is required in order to produce an asymmetry $n_B \neq n_{\bar{B}}$.

These crucial elements are actually found in the SM. CP violation is present in quark mixing depicted by the Kobayashi-Maskawa mechanism [26]. The Cabibbo-Kobayashi-Maskawa (CKM) quark mixing matrix contains one nonvanishing CP violating phase that is responsible for CP violation observed in K and B meson systems [23]. Departure from thermal equilibrium could be provided by the electroweak phase transition in the early universe. $B + L$ can also be violated in the SM and the mechanism is explained below.

2.2.1 $B + L$ violation in the Standard Model

Due to the chiral nature of the non-Abelian weak interactions B and L can be violated by the chiral anomaly [27] above the electroweak phase transition temperature [28]. The $SU(2)$ interacting fermions have global $U(1)$ symmetries [29, 30]

$$\psi_L^i(x) \to e^{i\beta}\psi_L^i(x) \tag{2.16}$$

and the related chiral currents are

$$j_\mu^B = \frac{1}{3}\sum_{generations} (\bar{q}_L\gamma_\mu q_L + \bar{u}_R\gamma_\mu u_R + \bar{d}_R\gamma_\mu d_R) \tag{2.17}$$

$$j_\mu^L = \sum_{generations} (\bar{\ell}_L\gamma_\mu \ell_L + \bar{e}_R\gamma_\mu e_R) \tag{2.18}$$

are conserved at tree level. In quantum theory the currents become anomalous and non-vanishing [27]

$$\partial^\mu j_\mu^B = \partial^\mu j_\mu^L = \frac{N_f}{32\pi^2}(-g^2 W_{\mu\nu}^I \tilde{W}^{I\mu\nu} + g'^2 B_{\mu\nu}\tilde{B}^{\mu\nu}). \tag{2.19}$$

The number of generations is denoted by N_f, and W_μ^I and B_μ are the $SU(2)_L$ and $U(1)_Y$ gauge fields with couplings g and g', respectively. The change in the topological charge of the gauge field is then [27]

$$B(t_f) - B(t_i) = \int_{t_i}^{t_f} dt \int d^3x \partial^\mu j_\mu^B = N_f(N_{cs}(t_f) - N_{cs}(t_i)), \tag{2.20}$$

where

$$N_{cs}(t) = \frac{g^3}{96\pi^2}\int d^3x \varepsilon_{ijk}\varepsilon^{IJK} W^{Ij} W^{Jj} W^{Kk}. \tag{2.21}$$

is the Chern-Simons number.

An infinite number of degenerate ground states exist and they differ in their Chern-Simons numbers, $\Delta N_{cs} = \pm 1, \pm 2, \ldots$. These points in field space are separated by a potential barrier, Fig. 2.2. At zero temperature, gauge field configurations that correspond to nonvanishing charges (2.20) are tunneling configurations, instantons [31]. Transitions between different vacua are proportional to the Chern-Simons number and the number of families (2.20) and these are related to changes in baryon and lepton number

$$\Delta B = \Delta L = N_f \Delta N_{cs}. \tag{2.22}$$

Since the number of families in the SM is 3, the smallest jumps between vacua are $\Delta B = \Delta L = \pm 3$, and $B + L$ is violated whereas $B - L$ is not. However, at zero temperature, the transition rate between vacua is very small [27]

$$\Gamma \sim e^{-S_{\text{int}}} \sim \mathcal{O}(10^{-165}) \tag{2.23}$$

and so $B + L$ violation does not occur in the SM in considerable amounts.

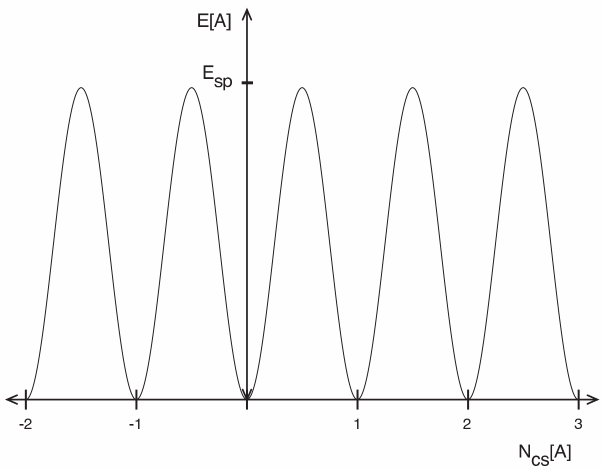

Figure 2.2: The picture is from [32]. The graph depicts the energy of the field configuration as a function of N_{cs}.

At nonzero temperature, the situation changes and transitions between vacua can occur by thermal fluctuations over the potential barrier rather than tunneling alone [28]. The transition rate is determined by the sphaleron configuration which is a saddle point of the field energy of the gauge-Higgs system [33], Fig. 2.2. The sphaleron energy is thus the height of the potential barrier between different vacua, Fig. 2.2 and it is given by

$$E_{\text{sph}}(T) \simeq \frac{8\pi}{g} v_F(T), \tag{2.24}$$

where $v_F(T)$ is the finite temperature vacuum expectation value (VEV) of the Higgs field. At temperatures

$$T_{\text{EW}} \sim 100 \text{ GeV} < T < T_{\text{sph}} \sim 10^{12} \text{ GeV} \tag{2.25}$$

the sphaleron processes are in thermal equilibirum and the suppression from the sphaleron rate is removed. The rate in the unbroken phase is given by [34]

$$\frac{\Gamma_{B+L}}{V} \sim \alpha^5 (\ln \alpha^{-1}) T^4, \tag{2.26}$$

where $\alpha = g^2/(4\pi)$ while a more detailed expression is given by [35]

$$\frac{\Gamma_{B+L}}{V} = (10.8 \pm 0.7) \left(\frac{gT}{m_D}\right)^2 \alpha^5 T^4 \left[\ln \left(\frac{m_D}{\gamma}\right) + 3.041 + \left(\frac{1}{\ln(1/g)}\right) \right] \tag{2.27}$$

with the Debye length $m_D \sim gT$ and hard gauge boson damping rate $\gamma \sim g^2 T \ln(1/g)$. Below the critical temperature T_c of electroweak phase transition the Higgs field acquires a VEV and the sphaleron rate per unit volume is [36]

$$\frac{\Gamma_{B+L}}{V} \simeq \kappa \frac{M_W^7}{(\alpha T)^3} e^{-\beta E_{sph}(T)}, \tag{2.28}$$

where $\beta = 1/T$, $M_W = g v_F(T)/2$ (temperature-dependent W boson mass) and κ =constant.

2.2.2 The dependence between baryon and lepton numbers

Chemical potentials can be assigned to the SM particles, quarks and leptons and Higgs fields, that exist in the weakly coupled plasma of the early universe. By imposing constraints based on the Yukawa, gauge and non-perturbative sphaleron interactions in thermal equilibrium, relations can be derived between the given chemical potentials [37]. The processes that introduce the constraints are the effective interaction induced by electroweak

$SU(2)$ instantons, $SU(3)$ quantum chromodynamics (QCD) instanton processes [38], vanishing of the total hypercharge, Yukawa and gauge interactions are in equilibrium [30, 32]. The results for the baryon and lepton numbers read

$$B = -\frac{4N_f}{3}\mu_l, \quad (2.29)$$

$$L = \frac{14N_f^2 + 9N_f}{6N_f + 3}\mu_l, \quad (2.30)$$

where N_f denotes the number of generations and μ_l is the chemical potential assigned to individual left-handed leptons. This gives the connection between B and L [39] in the pure minimal SM:

$$\begin{aligned} B &= c_s(B - L), \quad (2.31) \\ L &= (c_s - 1)(B - L), \end{aligned}$$

where $c_s = (8N_f + 4)/(22N_f + 13)$, and (2.31) hold for temperatures much larger than the Higgs VEV, $v_F \ll T$.

The relations in (2.31) suggest that $B-L$ violation is needed for producing excess B. Since sphalerons do not violate $B-L$, the net $B-L$ survives until present day and determines the amount of baryon excess according to (2.31).

2.3 Various mechanisms of baryogenesis

Here a variety of different baryogenesis mechanisms are briefly introduced. The main focus of this thesis is thermal leptogenesis which is the topic of the next chapters.

2.3.1 Electroweak baryogenesis

The electroweak phase transition could have provided the out-of-equilibrium condition given that the transition is first order. However, electroweak baryogenesis in the SM has been ruled out by observations that have set a limit to Higgs mass that is in conflict with electroweak baryogenesis scenarios. Namely, the lower bound for the Higgs mass is $m_H > 114.4$ GeV [23] and, on the other hand, electroweak baryogenesis requires $m_H < 45$ GeV [40].

Despite the fact that electroweak baryogenesis does not exist in the SM, several extensions of the SM do contain electroweak baryogenesis. These extensions are cold electroweak baryogenesis [41]-[45], the two-Higgs doublet model and Minimal Supersymmetric Standard Model [46].

2.3.2 GUT baryogenesis

GUTs can naturally accommodate the Sakharov conditions. Baryon number violation arises as quarks and leptons live in the same representation of a single group. Within a multiplet, the new gauge bosons carrying color and fractional electric charge can mediate transitions between quarks and leptons thus inducing B violation. CP violation can be generated by numerous complex phases. Finally, the average lifetimes of the heavy gauge bosons and or scalars are long compared to the age of the universe at the early epoch. Thus, the out-of-equilibrium condition is met.

Problems with GUT baryogenesis stem from possible relic production and difficulties in testing the models due to the high energy scale $M_{GUT} \sim 10^{16}$ GeV. Also, there must be sufficient $B-L$ produced prior to the onset of sphaleron transitions that violate $B+L$ but conserve $B-L$. Otherwise, any baryon excess will be erased by sphalerons. A candidate group that enables $B-L$ violation is $SO(10)$ that contains the gauged subgroup $U(1)_{B-L}$ [47, 48].

2.3.3 Affleck-Dine baryogenesis

In this scenario [49]-[52] scalar condensates with baryonic charge in the early universe can decay and produce net baryon number. The scalar field is a combination of flat directions of the scalar potential of some supersymmetric model and so the potential is virtually independent of this particular field. These scalar fields, *e.g.* scalar quarks, develop large VEVs during inflation and after inflation they can store baryon and lepton number. When the expansion rate of the universe becomes comparable to the scalar field masses, the flat directions start to oscillate around the minimum of the potential. Eventually, the scalar fields decay to SM particles that carry the B.

2.3.4 Baryogenesis via leptogenesis

Leptogenesis is based on models that produce net lepton number in the early universe and subsequently this lepton number is converted to net baryon number. Many of these models are based on heavy $SU(3)_c \times SU(2)_L \times U(1)_Y$ singlet neutrinos that are introduced to SM. These neutrinos have Yukawa couplings to left-handed SM neutrinos, which induces nonzero masses for SM neutrinos. The right-handed neutrinos can also decay to SM leptons and Higgs doublet in the out-of-equilibrium decays in the expanding early universe. The resulting net lepton number can then be converted to baryon number in electroweak sphaleron transitions. This is thermal leptogenesis and the subject of next chapters. Here, we give brief descriptions of non-thermal leptogenesis scenarios.

From inflaton decays

The obvious constraint between the reheating temperature and the right handed neutrino mass, $M_R \lesssim T_R$ can be lifted if the primordial right-handed neutrinos are produced non-thermally. One possible production channel is inflaton decay [53]. The inflaton can decay to the lightest right-handed neutrino, which subsequently decays into $H^\dagger + l_L$ and $H + l_L^\dagger$. The two channels have different rates, which generates CP violation that translates into lepton number that is given by [54]

$$\begin{aligned} \frac{n_L}{s} &\simeq -\frac{3}{2}\varepsilon\frac{T_R}{m_\Phi} \\ &\simeq 3\times 10^{-10}\Big(\frac{T_R}{10^6\mathrm{GeV}}\Big)\Big(\frac{M_1}{m_\Phi}\Big)\Big(\frac{m_3}{0.05\mathrm{eV}}\Big), \end{aligned} \tag{2.32}$$

where m_Φ is the inflaton mass and the CP violating phase $\delta_{\mathrm{eff}} = 1$ in

$$\varepsilon \simeq -2\times 10^{-6}\left(\frac{M_1}{10^{10}\mathrm{GeV}}\right)\left(\frac{m_3}{0.05\mathrm{eV}}\right)\delta_{\mathrm{eff}}. \tag{2.33}$$

Sphalerons convert the lepton number to baryon number which is given by [37]

$$\frac{n_B}{s} \simeq -\frac{8}{23}\frac{n_L}{s}. \tag{2.34}$$

The crucial point about generating lepton asymmetry via inflaton decay is that it does not require $T_R \sim M_{N_1}$ but instead that the inflaton mass has to exceed the masses of two right-handed neutrinos it decays into, $m_\Phi > 2M_{N_1}$. The lower reheating temperature $T_R < 10^7$ GeV satisfies the constraint on the gravitino abundance [55], and together with $m_\Phi > 2M_{N_1}$ (2.14) gives a constraint on the heaviest SM neutrino [30]

$$m_3 > 0.01 \text{ eV}. \tag{2.35}$$

This constraint agrees well with the results of atmospheric neutrino oscillation experiments, which are described in the next chapter.

Affleck-Dine leptogenesis

Affleck-Dine leptogenesis [51, 56, 30] can arise in supersymmetric theories, like Affleck-Dine baryogenesis mentioned earlier. The flat direction is $\phi \sim (H\ell_i)^{1/2}$, where H is the Higgs field and ℓ_i is the ith lepton doublet. Both H and ℓ_i represent the scalar components of the corresponding supermultiplet (discussed in chapter 5). Instead of baryon number, the scalar field ϕ can now store lepton number.

The scalar potential for the flat direction can be given, for example, by [50, 56]

$$V_{\text{SUSY}} = \frac{m_\nu^2}{4|\langle H\rangle|^4}|\phi|^6 \tag{2.36}$$

with diagonal neutrino mass matrix m_ν. The supersymmetry breaking potential is

$$\delta V = m_\phi|\phi|^2 + \frac{m_{\text{SUSY}} m_\nu}{8|\langle H\rangle|^2}(a_m\phi^4 + \text{h.c.}), \tag{2.37}$$

where $a_m \in \mathbb{C}$, $m_\phi \simeq m_{\text{SUSY}} \simeq 1$ TeV and $|a_m| \sim 1$. The amplitude of ψ decreases during inflation and starts to oscillate around the minimum of the potential when the Hubble parameter H_{exp} becomes comparable to m_ϕ. The initial value for the flat direction is $|\phi_0| \simeq \sqrt{m_\phi|\langle H\rangle|^2/m_\nu}$ which sets the initial value for lepton number generation.

The lepton number density obeys

$$\frac{\partial n_L}{\partial t} + 3H_{\text{exp}} n_L = \frac{m_{\text{SUSY}} m_\nu}{2|\langle H\rangle|^2}\text{Im}(a_m^*\phi^{*4}). \tag{2.38}$$

Lepton number generation commences when the AD field ϕ begins its coherent oscillations at a time $t_{\rm osc} \simeq H_{\rm osc} \simeq 1/m_\phi$ because $|\phi| \sim t^{-1}$. Thus,

$$\begin{aligned} n_L &\simeq \frac{m_{\rm SUSY} m_\nu}{2|\langle H\rangle|^2} \delta_{\rm eff} |a_m \phi_0^4| t_{\rm osc} \\ &\approx \delta_{\rm eff} m_\phi^2 \frac{|\langle H\rangle|^2}{2m_\nu}, \end{aligned} \tag{2.39}$$

where $\delta_{\rm eff} = \sin(4\arg\phi + \arg a_m)$ is the effective CP violating phase. The last line comes from the approximations $m_{\rm SUSY} \simeq m_\phi$, $|\phi_0| \simeq \sqrt{m_\phi |\langle H\rangle|^2/m_\nu}$ and $t_{\rm osc} \simeq 1/m_\phi$.

Once the inflaton decays and reheats the universe, the lepton asymmetry becomes

$$\frac{n_L}{s} \simeq \delta_{\rm eff} \frac{3T_R}{4M_{\rm Pl}} \frac{|\langle H\rangle|^2}{6m_\nu M_{\rm Pl}}. \tag{2.40}$$

Baryon asymmetry arises in sphaleron processes that convert n_L/s into

$$\frac{n_B}{s} \simeq \frac{1}{23} \frac{|\langle H\rangle|^2 T_R}{m_\nu M_{\rm Pl}^2}. \tag{2.41}$$

With $n_B/s \simeq 0.9 \times 10^{-10}$ and $T_R \simeq 10^6$ GeV, the lightest neutrino mass $m_\nu \simeq 10^{-9}$ eV. In light of neutrino oscillation experiments (described in the next chapter) this limit for the lightest neutrino mass is difficult to reconcile with observations.

Chapter 3

Thermal leptogenesis

Many models beyond SM predict superheavy particles that decay to SM particles directly creating an asymmetry between baryons and antibaryons. Another approach is to produce lepton number that can be converted to baryon number. The nonthermal scenarios have been briefly mentioned and the rest of this thesis focuses on thermal leptogenesis.

Thermal leptogenesis scenarios [7, 57] are based on the existence of heavy particles beyond SM that can decay to SM particles much like the GUT based scenarios and have Majorana masses. An exception to this is the Dirac leptogenesis scenario [58]. Much of the attraction around thermal leptogenesis lies in the fact that sufficient baryon number can be produced but also, on the same token, the small yet nonzero SM neutrino masses can be generated.

3.1 Massive neutrinos

3.1.1 Neutrinos in the Standard Model

The electroweak [10] theory of the SM is based on the group $SU(2)_L \times U(1)_Y$. The group $SU(2)_L$ is generated by the weak isospin operators $T^j = \sigma_j/2$, where σ_j are the Pauli matrices. In order to make $SU(2)_L$ a local symmetry of the Lagrangian gauge particles must be introduced. There are three of them A^1_μ, A^2_μ and A^3_μ and they only couple to isodoublet particles, *i.e.* the left-handed particles. The gauge field associated with $U(1)_Y$ is B_μ.

SM neutrinos belong to the fundamental representation of $SU(2)_L$ and they interact via the weak interaction only. Thus, they are found in the left-handed isodoublet alongside their electrically charged counterparts, e^-, μ^- and τ^-. There are no right-handed neutrinos in the SM. The right-handed charged leptons belong to the isosinglet representation of $SU(2)_L$ and hence do not couple to the $SU(2)_L$ gauge fields. The electroweak portion of the SM Lagrangian consisting of the kinetic terms and gauge interaction sector reads

$$\mathcal{L} = i\bar{\ell}_L \not{D}\, \ell_L + i\bar{e}_R \not{D}\, e_R - \frac{1}{4}F^{a\mu\nu}F^a_{\mu\nu} - \frac{1}{4}B^{\mu\nu}B_{\mu\nu}, \tag{3.1}$$

where e_R denotes the right-handed $SU(2)_L$ singlet charged leptons and l_L are the doublets. The covariant derivatives are

$$D_\mu \ell_L = (\partial_\mu - igA^a_\mu T^a - ig'\frac{Y}{2}B_\mu)\ell_L, \tag{3.2}$$

$$D_\mu e_R = (\partial_\mu - ig'\frac{Y}{2}B_\mu)e_R \tag{3.3}$$

and the gauge field tensors are

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu, \tag{3.4}$$

$$F^a_{\mu\nu} = \partial_\mu A^a - \partial_\nu A^a + gf^{abc}A^b_\mu A^c_\nu. \tag{3.5}$$

The three gauge fields of the local $SU(2)$ theory, A^1_μ, A^2_μ and A^3_μ couple only to left-handed SM leptons, and the gauge field associated with local $U(1)_Y$ couples to all leptons.

Mass terms for the fermions as well as gauge bosons are missing in (3.1). Bare mass terms such as $m_e(\bar{e}_L e_R + \bar{e}_R e_L)$ cannot be written because left-handed and right-handed fermions belong to different $SU(2)_L$ representations and have different $U(1)_Y$ charges. Thus, these terms would violate global symmetries. By introducing an $SU(2)$ doublet scalar field [11], the Higgs field, with hypercharge $Y = 1$ and that in the unitary gauge reads as

$$H(x) = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v + h(x) \end{pmatrix} \tag{3.6}$$

the fermions and bosons one can obtain masses in a gauge invariant fashion. The Higgs field H has a VEV $v/\sqrt{2}$ while $\langle h(x)\rangle = 0$, and this nonzero VEV generates the masses for all SM fermions and electroweak gauge bosons and thus breaks the symmetry group $SU(2)_L \times U(1)_Y$. The Lagrangian (3.1)

Table 3.1: The quantum numbers for SM leptons from electroweak theory.

	T	T^3	Y
e_{L_i}	1/2	-1/2	-1
ν_{L_i}	1/2	1/2	-1
e_{R_i}	0	0	-2

is now augmented with interactions between the leptons, gauge bosons and the Higgs field

$$\begin{aligned}\mathcal{L} &= i\bar{\ell}_L \not{D} \ell_L + i\bar{e}_R \not{D} e_R - f_i H^\dagger \bar{e}_{iR}\ell_{iL} + \text{h.c.} \\ &\quad -\frac{1}{4}F^{a\mu\nu}F^a_{\mu\nu} - \frac{1}{4}B^{\mu\nu}B_{\mu\nu} + (D_\mu H)^\dagger D_\mu H + \mu H^\dagger H - \lambda(H^\dagger H)^2\end{aligned} \tag{3.7}$$

where e_{iR} and ℓ_{iL} have the indices to show explicitly which generations and associated Yukawa couplings f_i are in question. The potential of the Higgs field is

$$V(H) = -\mu^2 H^\dagger H + \lambda(H^\dagger H)^2 \tag{3.8}$$

and the minimum is $|H| = v/\sqrt{2}$ with $v = \sqrt{\mu^2/\lambda}$.

Because the combination $T^3 + \frac{Y}{2}$ of the group generators annihilates the Higgs VEV $(0, v/\sqrt{2})^T$, the remaining symmetry is $U(1)_{em}$. The corresponding charge of the remnant symmetry $U(1)_{em}$ is the electric charge Q which is related to weak isospin T of $SU(2)_L$ and hypercharge Y of $U(1)_Y$ through

$$Q = T^3 + \frac{Y}{2}, \tag{3.9}$$

where T^3 is the third component of T. Table 3.1 summarizes the T, T^3 and Y quantum numbers for SM leptons.

SM leptons obtain their masses from the Yukawa couplings to the Higgs field and the gauge fields become massive through the kinetic term of the Higgs field shown in (3.7). The terms

$$\begin{aligned}\mathcal{L}_{G,M} &= \frac{1}{2}\begin{pmatrix}0 & v\end{pmatrix}\left(gA^a_\mu T^a + \frac{1}{2}g'B_\mu\right)\left(gA^{b\mu}T^b + \frac{1}{2}g'B^\mu\right)\begin{pmatrix}0 \\ v\end{pmatrix} \\ &= \frac{1}{2}\frac{v^2}{4}[g^2(A^1_\mu)^2 + g^2(A^2_\mu)^2 + (-gA^3_\mu + g'B_\mu)^2]\end{aligned} \tag{3.10}$$

generate gauge boson masses. Three of the bosons become massive

$$W^{\pm}_{\mu} = \frac{1}{\sqrt{2}}(A^1_\mu \mp i A^2_\mu), \qquad m_W = g\frac{v}{2}, \tag{3.11}$$

$$Z_\mu = \frac{1}{\sqrt{g^2+g'^2}}(gA^3_\mu - g'B_\mu), \qquad m_Z = \sqrt{g^2+g'^2}\,\frac{v}{2}, \tag{3.12}$$

and the fourth vector boson remains massless

$$A_\mu = \frac{1}{g^2+g'^2}(g'A^3_\mu + gB_\mu) \qquad m_A = 0. \tag{3.13}$$

This massless particle is the photon which is the gauge boson of the remnant electromagnetic symmetry group $U(1)_{em}$. This can be seen by looking at the covariant derivative (acting on the Higgs field H) written in terms of the mass eigenstate vector bosons. There the operator $T^3 + \frac{Y}{2}$ multiplies the massless A_μ. By defining the weak mixing angle θ_w

$$\begin{aligned} \cos\theta_w &= \frac{g}{\sqrt{g^2+g'^2}}, \\ \sin\theta_w &= \frac{g'}{\sqrt{g^2+g'^2}} \end{aligned} \tag{3.14}$$

the neutral vector bosons can be expressed as

$$\begin{aligned} Z_\mu &= \cos\theta_w A^3_\mu - \sin\theta_w B_\mu, \\ A_\mu &= \sin\theta_w A^3_\mu + \cos\theta_w B_\mu. \end{aligned} \tag{3.15}$$

Fermion masses are generated by the Yukawa coupling terms $f_i H^\dagger \bar{e}_{R_i} l_{L_i}$. The electron mass is then given by

$$m_e = \frac{1}{\sqrt{2}} f_e v \tag{3.16}$$

and other charged lepton masses are generated in a similar manner as a product of $v = 246$ GeV and the Yukawa coupling f_i.

The interaction terms between fermions and gauge bosons are found in the covariant derivatives of (3.7). By replacing the gauge fields A^a_μ and B_μ with the mass eigenstates A_μ, Z_μ and $W^\pm_\mu$ we get

$$gA^3_\mu T^3 + g'\frac{1}{2}YB_\mu = \frac{g}{\cos\theta w}[T^3 - \sin^2\theta_w Q]Z_\mu + g\sin\theta_w Q A_\mu, \tag{3.17}$$

which gives the neutral current interactions for the leptons

$$\mathcal{L}_{NC} = \frac{g}{\cos\theta_w}\bar{\psi}\gamma^\mu[T^3 - \sin^2\theta_w Q]\psi Z_\mu + g\sin\theta_w\bar{\psi}\gamma^\mu Q\psi A_\mu. \qquad (3.18)$$

Similarly, the interactions with the charged gauge bosons are obtained from the covariant derivative terms

$$gA^1_\mu T^1 + gA^2_\mu T^2 = \frac{g}{\sqrt{2}}[W^+_\mu(T^1 + iT^2) + W^-_\mu(T^1 - iT^2) \qquad (3.19)$$

and so

$$\begin{aligned} \mathcal{L}_{CC} &= i\frac{g}{\sqrt{2}}W^{+\mu}J^+_\mu + \text{h.c.}, \qquad (3.20) \\ J^+_\mu &= \bar{\nu}_L\gamma_\mu e_L \end{aligned}$$

is the charged current Lagrangian for the leptons.

At the time when SM was formulated (see *e.g.* [59]), there were no experimental evidence stating that neutrinos should have mass. This is why neutrinos in (3.7) lack a Yukawa term similar to $f_i H^\dagger \bar{e}_{iR} l_{iL}$ that generates mass when the Higgs field develops a VEV.

3.1.2 Heavy singlet neutrinos beyond SM and the seesaw mechanism

In the previous section we learned that SM neutrinos are massless. This contradicts neutrino oscillation experiments (described in the next section) that suggest neutrinos must have a mass. A way to circumvent the problem is to introduce a beyond SM Dirac mass term that is similar to the charged lepton mass terms. This would require a right-handed $SU(3)_C \times SU(2)_L \times U(1)_Y$ singlet neutrino that appears in the Yukawa interaction

$$\mathcal{L}_Y = -f_{ij}\bar{\ell}_{iL}\widetilde{H}N_{jR} + \text{h.c.}, \qquad (3.21)$$

where $\tilde{H} = i\sigma^2 H^*$ and which gives the neutrino mass terms

$$\mathcal{L}_{\text{mass}} = -f_{ij}\frac{v}{\sqrt{2}}\bar{\nu}_{iL}N_{jR} + \text{h.c.}. \qquad (3.22)$$

To explain the sub-eV masses, $m_\nu = 0.01$ eV, for neutrinos one would have to fine tune the couplings $f_{ij} \sim 10^{-13}$, which does not seem natural.

Small neutrino masses can be generated in an alternative scenario where the assumption that neutrinos are Dirac fermions is relaxed.

Since neutrinos are electrically neutral, they can have mass terms of the type $m_{ij}\overline{\nu_{iR}^c}\nu_{jL}$. However, this Majorana [60] mass term is not allowed for SM neutrinos because it is not gauge invariant. Thus, new gauge singlet neutrinos N_{iR} must be introduced with Majorana mass term

$$M_{ij}\overline{N_{iL}^c}N_{jR}. \tag{3.23}$$

Leptogenesis scenarios require that $B - L$ is broken, which could be the natural manifestation of the breaking some underlying symmetry. The singlet right-handed neutrinos N_R arise naturally in theories beyond the SM where there is an extra $U(1)$ gauge symmetry. The group $SO(10)$ is known to have $U(1)_{B-L}$ as a subgroup [47, 48] so the simplest extensions of the SM in this regard are considered to stem from $SO(10)$. To cancel the triangle anomaly $[U(1)_{B-L}]^3$ singlet right-handed neutrinos must be included [61, 62]. Now, the neutrino mass term for one SM neutrino flavor becomes

$$\begin{aligned} -\mathcal{L}_{mass} &= \frac{1}{2}\begin{pmatrix}\overline{\nu_L} & \overline{N_L^c}\end{pmatrix}\widetilde{\mathcal{M}}\begin{pmatrix}\nu_R^c \\ N_R\end{pmatrix} + \text{h.c.} \\ &= \frac{1}{2}\begin{pmatrix}\overline{\nu_L} & \overline{N_L^c}\end{pmatrix}\begin{pmatrix}0 & M_D \\ M_D^{\mathrm{T}} & M_R\end{pmatrix}\begin{pmatrix}\nu_R^c \\ N_R\end{pmatrix} + \text{h.c.} \end{aligned} \tag{3.24}$$

with real M_D and M_R for simplicity. By diagonalizing the mass matrix $\widetilde{\mathcal{M}}$ with U we compute $U^T\widetilde{\mathcal{M}}U$ and obtain the masses

$$\begin{aligned} m_1 &\simeq M_R, \\ m_2 &\simeq \frac{M_D^2}{M_R}. \end{aligned} \tag{3.25}$$

The light neutrino mass m_2 is suppressed by the mass of the singlet neutrino M_R if $M_D \ll M_R$, and the observed sub-eV neutrino masses would require that the heavy mass is $\sim 10^{15}$ GeV which is close to the GUT scale. This is the Type I seesaw mechanism [63, 64]. Type II [65]-[68] and III [69]-[71] seesaw mechanisms include $SU(2)$ triplet scalar and $SU(2)$ triplet fermions, respectively. The M_D blocks in the mass matrix are Dirac masses and they arise in the Yukawa interactions

$$\mathcal{L}_Y = -h_{ij}\bar{\ell}_{jL}\widetilde{H}N_{iR} + \text{h.c.}$$

when the the Higgs field develops a VEV as we learned before.

The existence of the Majorana mass term (3.23) implies two things. First, there is lepton number violation by two units $\Delta L = \pm 2$. This is not a problem since lepton number is not gauged within the SM. Second, if neutrinos are Majorana fermions, they are their own antiparticles, which reduces their number of degrees of freedom to half compared to Dirac fermions.

3.1.3 Neutrino oscillations and experiments

The motivation behind introducing Majorana neutrinos, neutrino masses (3.24) and other neutrino mass models is the observation that neutrinos oscillate from one flavor to another. The first hints toward flavor oscillations came from Davis's Homestake experiment based on a chlorine detector in 1968 that observed a deficit in the flux of solar electron neutrinos [73]. This was dubbed the solar neutrino puzzle. Several other experiments such as SAGE [74], GALLEX [75] and GNO [76] based on gallium detectors and Kamiokande [77] and Super-Kamiokande [78] using Cherenkov detectors tackled the solar neutrino puzzle. In 2001 the heavy water detector at Sudbury Neutrino Observatory (SNO) confirmed neutrino oscillations as the cause for the deficit in solar neutrinos [79]. The apparent deficit of ν_es exists because ν_es are converted to the other flavors ν_μ and ν_τ.

Another phenomenon related to neutrino oscillations is connected to atmospheric neutrinos produced in the atmosphere by cosmic ray interactions. The ratio of ν_μ to ν_e was observed to have a deficit at Super-Kamiokande [80] as ν_μ oscillate into other flavors. Other studies on atmospheric neutrinos have been conducted at long baseline experiments. These include the K2K [81] and MINOS [82] with neutrino path length extending hundreds of kilometers.

The determination of one of the neutrino mixing angles θ_{13} described in (3.28) has been the objective for some experiments. Antineutrino $\bar{\nu}_e$ disappearance experiments have been conducted at CHOOZ [83] and Palo Verde Nuclear Generating Station in Arizona [84]. Neither CHOOZ nor Palo Verde observed antineutrino disappearance. The Super-Kamiokande [85] and SNO [86] experiments have set limits to θ_{13} and [87] suggests θ_{13} is indeed nonzero.

Since neutrinos can obtain masses via the seesaw mechanism, (3.24) and

(3.25), the mass eigenstates are related to the flavor eigenstates through

$$\nu_{iL} = \sum_{\alpha=1}^{3} U^\dagger_{i\alpha} \nu_{\alpha L}, \tag{3.26}$$

where i and α denote the mass and flavor eigenstates, respectively. The fact that neutrino flavor states do not coincide with the mass eigenstates alters the charged current interaction (3.20) if we wish to write it entirely in terms of mass eigenstates

$$\mathcal{L}_{CC} = i\frac{g}{\sqrt{2}} \sum_{l,i} \bar{\ell}_L \not{W}^- U_{li} \nu_{iL} + \text{h.c.} \tag{3.27}$$

The Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix U is unitary and it can be parameterized as [72]

$$\begin{aligned} U \;=\; & \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix} \\ & \times \text{diag}(1, e^{i\alpha_{21}/2}, e^{i\alpha_{31}/2}), \end{aligned} \tag{3.28}$$

where $c_{ij} = \cos\theta_{ij}$, $s_{ij} = \sin\theta_{ij}$, the angles $\theta_{ij} = [0, \pi/2]$, $\delta = [0, \pi]$ is the Dirac CP violation phase and α_{21} and α_{31} are the Majorana CP violation phases. The Majorana phases are relevant if neutrinos are Majorana fermions, although these phases do not affect transition probabilities. The major difference between neutrino and quark mixing is that two Majorana neutrino generations are sufficient to induce CP violation whereas with quarks three generations are needed. More phases are left in the former case because they cannot be absorbed in Majorana fermion fields whereas with Dirac fermions more phases can be absorbed by redefining the fermion fields and thus more Dirac fermions are needed to have at least one phase left.

Neutrino experiments have determined limits for the mixing angles θ_{ij} and squared mass differences $\Delta m^2_{ij} \equiv m^2_i - m^2_j$. They have been found to be [88, 23]

$$\left.\begin{aligned} (\sin^2\theta_{23})_{\text{best fit}} &= 0.5, \quad 0.36 \le \sin^2\theta_{23} \le 0.67, \\ (|\Delta m^2_{31}|)_{\text{best fit}} &= 2.4 \times 10^{-3}\text{eV}^2, \\ 2.07 \times 10^{-3}\text{eV}^2 &\le |\Delta m^2_{31}| \le 2.75 \times 10^{-3}\text{eV}^2 \end{aligned}\right\} \text{atmospheric } \nu \tag{3.29}$$

and

$$\left.\begin{array}{rl} (\sin^2\theta_{12})_{\text{best fit}} & = 0.304, \quad 0.25 \le \sin^2\theta_{12} \le 0.37, \\ (|\Delta m^2_{21}|)_{\text{best fit}} & = 7.65 \times 10^{-5}\text{eV}^2, \\ 7.05 \times 10^{-5}\text{eV}^2 & \le \Delta m^2_{21} \le 8.34 \times 10^{-5}\text{eV}^2 \end{array}\right\}\text{solar } \nu. \quad (3.30)$$

CHOOZ [83] gives an upper limit to the remaining mixing angle

$$\sin^2 2\theta_{13} < 0 - 0.15 \text{ at } 90\% \text{ CL} \quad (3.31)$$

with $|\Delta m^2_{31}| \simeq 2.4 \times 10^{-3}$ eV^2 while a combined analysis of global data yields [88]

$$\sin^2\theta_{13} < 0.035\ (0.056) \qquad \text{at } 90\%\ (99.73\%) \text{ CL.} \quad (3.32)$$

Existing data does not restrict the sign of $\Delta m_{31(2)}$, and the two sign alternatives imply different hierarchies for the neutrino mass spectrum. With normal ordering $m_1 < m_2 < m_3$ the mass differences $\Delta m^2_{31} > 0$ and $\Delta m^2_{21} > 0$ whereas with inverted ordering $m_3 < m_1 < m_2$ the atmospheric mass difference is negative $\Delta m^2_{32} < 0$.

3.2 CP violation and leptogenesis

Models including right-handed singlet neutrinos can accommodate CP violation and lepton number violation

$$\mathcal{L} = -h_{ij}\bar{\ell}_{jL}\widetilde{H}N_{iR} + \frac{1}{2}M_{ij}\overline{N_{iR}}{N^c}_{jL} + \text{h.c.} \quad (3.33)$$

The Yukawa coupling matrix h_{ij} contains 9 CP violating phases, three of which can be absorbed into the wave function of ℓ. The interaction $-h_{ij}\bar{\ell}_{jL}\widetilde{H}N_{iR}$ allows a heavy right-handed neutrino N_i to decay to leptons and antileptons, see Fig. 3.1. The out-of-equilibrium decays at temperatures $T < M$, where M denotes the heavy Majorana mass scale, can generate net lepton number which is subsequently converted to baryon number in the sphaleron processes [7].

With hierarchical right-handed neutrinos $M_1 \ll M_2, M_3$ one can assume that the lightest neutrino N_1 produces the lepton asymmetry and that any asymmetry generated by the heavier neutrinos N_2 and N_3 has been washed

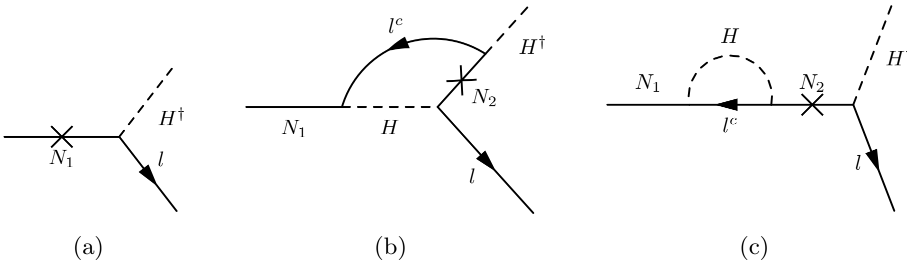

Figure 3.1: The relevant tree level and loop diagrams depicting the decay of the heavy neutrinos N_1 to SM leptons and Higgs doublet.

out. The CP violation parameter emerging from the interference between the tree level and one loop diagrams of Fig. 3.1(b) is given by [7], [57], [89]-[92]

$$\varepsilon_1 = \frac{\sum_\alpha [\Gamma(N_1 \to \ell_\alpha H) - \Gamma(N_1 \to \bar{\ell}_\alpha H^\dagger)]}{\sum_\alpha [\Gamma(N_1 \to \ell_\alpha H) + \Gamma(N_1 \to \bar{\ell}_\alpha H^\dagger)]} \tag{3.34}$$

$$\simeq \frac{1}{8\pi}\frac{1}{(hh^\dagger)_{11}} \sum_{i=2,3} \mathrm{Im}\Big[(hh^\dagger)^2_{1i}\Big]\left[f\left(\frac{M_i^2}{M_1^2}\right) + g\left(\frac{M_i^2}{M_1^2}\right)\right], \tag{3.35}$$

where

$$f(x) = \sqrt{x}\left[1 - (1+x)\ln\left(\frac{1+x}{x}\right)\right], \tag{3.36}$$

$$g(x) = \frac{\sqrt{x}}{1-x}. \tag{3.37}$$

The functions $f(M_i^2/M_1^2)$ and $g(M_i^2/M_1^2)$ arise from the vertex and mixing diagrams, respectively. For hierarchical right-handed neutrinos, $M_1 \ll M_2, M_3$, ε reduces to

$$\varepsilon_1 \simeq -\frac{3}{16\pi}\frac{1}{(hh^\dagger)_{11}} \sum_{i=2,3} \mathrm{Im}\big[(hh^\dagger)^2_{1i}\big]\frac{M_1}{M_i}. \tag{3.38}$$

The decays of the right-handed neutrinos have to be out-of-equilibrium because otherwise the asymmetry generated in (3.34) would be washed out. This means the Hubble rate has to exceed the decay width $\Gamma_{D_1} = (hh^\dagger)_{11}M_1/(8\pi)$ of the decaying neutrino N_1

$$r \equiv \frac{\Gamma_{D_1}}{H|_{T=M_1}} = \frac{M_{\mathrm{Pl}}}{1.7 \times 32\pi\sqrt{g_*}}\frac{(hh^\dagger)_{11}}{M_1} < 1. \tag{3.39}$$

with the number of degrees of freedom is denoted by g_* and for SM $g_* \simeq 106.75$, and the Planck mass $M_{\rm Pl} = 1.22 \times 10^{19}$ GeV. The effective neutrino mass is defined as

$$\tilde{m}_1 = (hh^\dagger)_{11} \frac{v^2}{M_1}$$

and the equilibrium mass by

$$m_* = 8\pi \frac{v^2}{M_1^2} H\big|_{T=M_1} \simeq 1.1 \times 10^{-3}\text{eV}.$$

with $v = 174$ GeV. The condition (3.39) can be expressed in an equivalent form

$$\tilde{m}_1 < m_* \simeq 1.1 \times 10^{-3}\text{eV}. \tag{3.40}$$

The generated lepton number is given by

$$Y_L = \frac{n_L - \bar{n}_L}{s} = \kappa \frac{\varepsilon}{g_*}, \tag{3.41}$$

where κ parametrizes the amount of washout that is generated by inverse decays and scattering processes. The generated baryon number is related to the lepton number through

$$Y_{B-L} \simeq -Y_L = -\kappa \frac{\varepsilon_1}{g_*}. \tag{3.42}$$

Thus, a relation between net baryon number and CP violation due to heavy neutrino decays has been established.

3.2.1 Boltzmann equations

In addition to the lepton number producing processes in Fig. 3.1, there are processes that strive to wash out any generated lepton number. These processes consist of [57], [91]

1. inverse decays $\ell H \to N_1$, $\bar{\ell} H^\dagger \to N_1$

2. off-shell $\Delta L = 1$ scatterings

$$\begin{array}{lll} N_1 \ell \leftrightarrow t\bar{q}, & N_1 \bar{\ell} \leftrightarrow t\bar{q} & \text{s channel} \\ N_1 t \leftrightarrow \bar{\ell} q, & N_1 \bar{t} \leftrightarrow \ell \bar{q} & \text{t channel} \end{array}$$

3. off-shell $\Delta L = 2$ scatterings

$$\ell H \leftrightarrow \bar{\ell} H^\dagger, \qquad \ell\ell \leftrightarrow H^\dagger H^\dagger, \qquad \bar{\ell}\bar{\ell} \leftrightarrow HH.$$

The Boltzmann equations relevant for leptogenesis are [93]-[95], [30]

$$\frac{dN_{N_1}}{dz} = -(D+S)(N_{N_1} - N_{N_1}^{\mathrm{eq}}), \tag{3.43}$$

$$\frac{dN_{B-L}}{dz} = -\varepsilon_1 D(N_{N_1} - N_{N_1}^{\mathrm{eq}}) - W N_{B-L}, \tag{3.44}$$

where $z = M_1/T$. These equations hold when all SM states that emerge in scattering and decay processes are in thermal equilibrium. The relativistic equilibrium N_1 number density is $N_{N_1}^{\mathrm{eq}}(z \ll 1) = 3/4$.

The N_1 abundance is modified by decays, inverse decays and $\Delta L = 1$ scatterings that tend to drive N_{N_1} towards its equilibrium value $N_{N_1}^{\mathrm{eq}}$. The term $D = \Gamma_D/(Hz)$ accounts for decays and inverse decays and $S = \Gamma_S/(Hz)$ denotes the contributions made by $\Delta L = 1$ scatterings. The decay term D also acts as the source for $B-L$ asymmetry. The total washout term $W = \Gamma_W/(Hz)$ receives contributions from all processes and it competes with the decay source term.

In a simplified picture only decays and inverse decays are considered and this gives the solution for N_{B-L} [96]

$$N_{B-L}(z) = N_{B-L}^{\mathrm{i}} e^{-\int_{z_{\mathrm{i}}}^{z} dz' W_{ID}(z')} - \frac{3}{4}\varepsilon_1 \kappa(z; \tilde{m}_1), \tag{3.45}$$

where W_{ID} denotes the contribution of inverse decays to the washout, and the efficiency factor κ is given by

$$\kappa(z) = \frac{4}{3}\int_{z_{\mathrm{i}}}^{z} dz' D(N_{N_1} - N_{N_1}^{\mathrm{eq}}) e^{\int_{z'}^{z} dz'' W_{ID}(z'')}. \tag{3.46}$$

The decay parameter K is defined as

$$K = \frac{\Gamma_D(z=\infty)}{H(z=1)} = \frac{\tilde{m}_1}{m_*} \tag{3.47}$$

and the contribution of inverse decays to the washout is

$$W_{ID}(z) = \frac{1}{2} D(z) \frac{N_{N_1}^{\mathrm{eq}}(z)}{N_l^{\mathrm{eq}}}. \tag{3.48}$$

Cases are distinguished as whether they are in the weak $K \ll 1$ or strong $K \gg 1$ washout regimes. In the regime of weak washout the efficiency factor is

$$\kappa_f(K) \simeq \frac{9\pi^2}{64} K^2, \tag{3.49}$$

where the initial abundance of N_1 is assumed to vanish, $N_{N_1}(z_\mathrm{i}) \equiv N^\mathrm{i}_{N_1} \simeq 0$. In the region of strong washout the integral for the efficiency factor becomes

$$\kappa(z) = \frac{2}{K} \int_{z_\mathrm{eq}}^{z} dz' \frac{1}{z'} W_{ID}(z') e^{-\int_{z'}^{z} dz'' W_{ID}(z'')}, \tag{3.50}$$

where the integrand has a maximum at $z = z_B$ and $N_{N_1}(z_\mathrm{eq}) = N^\mathrm{eq}_{N_1}(z_\mathrm{eq})$. The integral (3.50) receives its main contribution around z_B which is given by

$$z_B(K) \simeq 1 + \frac{1}{2} \ln\Big(1 + \frac{\pi K^2}{1024}\Big[\ln\Big(\frac{3125\pi K^2}{1024}\Big)\Big]^5\Big). \tag{3.51}$$

The efficiency factor from (3.50) is approximated by

$$\kappa_\mathrm{f}(K) \simeq \frac{2}{z_B(K) K}\big(1 - e^{-\frac{1}{2} z_B(K) K}\big). \tag{3.52}$$

By including $\Delta L = 1$ and $\Delta L = 2$ scatterings as well as thermal corrections [97] and scattering processes with gauge bosons [98, 97] one can solve the Boltzmann equations for the full theory. The efficiency factor receives corrections as compared to the cases shown in equations (3.49) and (3.52), and in the strong washout region

$$\kappa_\mathrm{f} = (2 \pm 1) \times 10^{-2} \Big(\frac{0.01\mathrm{eV}}{\tilde{m}_1}\Big)^{1.1 \pm 0.1}. \tag{3.53}$$

In some cases the decays of the other heavy neutrinos N_2 and N_3 can become significant [99]-[102] even if the heavy neutrino mass spectrum is hierarchical $M_1 \ll M_{2,3}$. These are known as flavor effects and they arise when the charged lepton Yukawa couplings are large enough. This influences the choice of basis in writing the Boltzmann equations: when the charged lepton couplings are large one should use the flavor basis whereas with weak coupling one should use a basis where two states are orthogonal to the lepton state produced in the heavy neutrino decay.

3.3 Quasidegenerate heavy neutrinos and resonant leptogenesis

The previous section dealt with hierarchical heavy neutrinos $M_1 \ll M_{2,3}$. In this case the contribution to CP violation comes from the decays of the lightest heavy neutrino N_1 and the amplitude is computed from the tree level and vertex loop diagrams of Figs. 3.1(a) and 3.1(b).

A different scenario arises when the heavy neutrinos are nearly mass-degenerate, *e.g.* $|M_1 - M_2| \ll M_{1,2}$ [103]-[107]. Both heavy states decay to SM leptons and the Higgs doublet and so the total CP violation parameter can be defined either as [103]

$$\begin{aligned} \varepsilon &= \varepsilon_{N_1} + \varepsilon_{N_2}, \\ \varepsilon_{N_i} &= \frac{\sum_\alpha [\Gamma(N_i \to \ell_\alpha H) - \Gamma(N_i \to \bar{\ell}_\alpha \bar{H})]}{\sum_\alpha [\Gamma(N_i \to \ell_\alpha H) + \Gamma(N_i \to \bar{\ell}_\alpha \bar{H})]} \end{aligned} \tag{3.54}$$

or as [105]

$$\varepsilon = \frac{\sum_\alpha [\Gamma(N_1 \to \ell_\alpha H) - \Gamma(N_1 \to \bar{\ell}_\alpha \bar{H}) + \Gamma(N_2 \to \ell_\alpha H) - \Gamma(N_2 \to \bar{\ell}_\alpha \bar{H})]}{\sum_\alpha [\Gamma(N_1 \to \ell_\alpha H) + \Gamma(N_1 \to \bar{\ell}_\alpha \bar{H}) + \Gamma(N_2 \to \ell_\alpha H) + \Gamma(N_2 \to \bar{\ell}_\alpha \bar{H})]}. \tag{3.55}$$

The latter definition (3.55) measures the net lepton number produced in any decay process irrespective of the source whereas the former (3.54) adds the separate contributions from N_1 and N_2 together.

With quasidegenerate heavy neutrinos the contribution given by the mixing diagram Fig. 3.1(c) can be significantly enhanced and the CP violation parameter can become as high as $\varepsilon \lesssim 1/2$ [105, 106]. The mixing diagrams have the N_2 running in the propagator in N_1 decays and vice versa. The heavy neutrino propagator includes all possible one-particle irreducible (1PI) diagrams and they are obtained by resumming an infinite series of heavy neutrino self-energy graphs. The resummed amplitudes for lepton and antilepton production via N_1 decay, Figs. 3.1(a) and 3.1(c) with

$N_j = N_2$, are [105]

$$\begin{aligned}
\mathcal{T}_{N_1} &= h_{l1}\bar{u}_l P_R u_{N_1} - i h_{l2}\bar{u}_l P_R[\not{p} - M_2 + i\Sigma_{22}^{\rm abs}(\not{p})]^{-1}\Sigma_{21}^{\rm abs}(\not{p})u_{N_1} \\
&= \bar{u}_l P_R u_{N_1}\left[h_{l1} - i h_{l2}\frac{M_1^2(1+iA_{22})A_{21}^* + M_1M_2A_{21}}{M_1^2(1+iA_{22})^2 - M_2^2}\right], \qquad (3.56) \\
\bar{\mathcal{T}}_{N_1} &= h_{l1}^*\bar{u}_l P_L u_{N_1} - i h_{l2}^* P_L[\not{p} - M_2 + i\bar{\Sigma}_{22}^{\rm abs}(\not{p})]^{-1}\Sigma_{21}^{\rm abs}(\not{p})u_{N_1} \\
&= \bar{u}_l P_L u_{N_1}\left[h_{l1}^* - i h_{l2}^*\frac{M_1^2(1+iA_{22})A_{21} + M_1M_2A_{21}^*}{M_1(1+iA_{22})^2 - M_2^2}\right]. \qquad (3.57)
\end{aligned}$$

The absorptive parts of the mixing loops are given by

$$\begin{aligned}
\Sigma_{ij}^{\rm abs}(\not{p}) &= A_{ij}(p^2)\not{p}P_L + A_{ij}^*(p^2)\not{p}P_R, \\
\bar{\Sigma}_{ij}^{\rm abs}(\not{p}) &= A_{ij}(p^2)\not{p}P_R + A_{ij}^*(p^2)\not{p}P_L \qquad (3.58)
\end{aligned}$$

with

$$A_{ij} = \frac{h_{l'i}h_{l'j}^*}{16\pi} \qquad (3.59)$$

in the limit of vanishing Higgs mass $M_H \to 0$. The CP asymmetry due to the decays of N_1 Fig. 3.1(c) with quasidegenerate heavy neutrinos $N_{1,2}$ becomes [105, 106, 29]

$$\varepsilon_{N_1} = -\frac{M_1}{M_2}\frac{\Gamma_2}{M_2}\frac{M_2^2\Delta M_{21}^2}{(\Delta M_{21}^2)^2 + M_1^2\Gamma_2^2}\frac{{\rm Im}[(h^\dagger h)_{12}^2]}{(h^\dagger h)_{11}(h^\dagger h)_{22}}. \qquad (3.60)$$

The resonance condition is

$$|M_1 - M_2| \sim \frac{\Gamma_{1,2}}{2}. \qquad (3.61)$$

In the high temperature regime the decay rates of heavy neutrinos can exceed the Hubble rate and the heavy neutrinos can quickly thermalize with SM particles. Also, washout effects are strong and they erase any preexisting lepton asymmetry. The advantage of resonant leptogenesis is that large amounts of lepton number can be produced as the universe cools to temperatures close to the heavy neutrino mass. This lepton asymmetry survives washout effects and is converted to baryon asymmetry by sphalerons. Leptogenesis could be possible with TeV-scale heavy neutrino masses.

Our research articles [1]-[3] consider models where CP violation is mainly generated via the mixing diagram in Fig. 3.1(c) and the definition of the

CP violation parameter follows (3.55). The decaying heavy particles form a quasidegenerate spectrum where the mass splittings are much smaller than the masses of the decaying states $|\Delta M_{ij}| \ll M_{i,j}$. Extra dimensions with SM singlet neutrinos [1, 2] and $B-L$ gauged supersymmetry [3] produce the desired heavy particle spectra, and these scenarios are discussed in the next two chapters.

Chapter 4

Leptogenesis in models with flat extra dimensions

Unifying the fundamental interactions with extra dimensons is an idea that dates back to 1914 when Gunnar Nordström proposed a 5-dimensional theory that ties electromagnetism and scalar gravity together [108]. In 1921 Theodor Kaluza realized that a 5-dimensional generalization of Einstein's theory of general relativity can describe both gravity and electromagnetism [109]. Oskar Klein considered gauge invariance and physical implications of compactification of the extra dimension [110]. Since the late 1970's superstring theories have been considered as a way to incorporate the SM of particle physics and gravity. These theories introduce extra dimensions on top of the familiar four spacetime dimensions. The work of Horava and Witten [111, 112] suggested that some of the extra dimensions could be larger than the Planck length $\ell_{\rm Pl} = 1/M_{\rm Pl}$. In late 1990's it was discovered that large extra dimensions in the millimeter scale could exist and they could be accessed in current particle colliders.

Section 4.1 reviews large extra dimensions and their observation [113, 23]. The topic of section 4.2 is how particles and especially neutrinos can exist in extra dimensions. Finally, section 4.3 presents our research [1, 2].

4.1 Large extra dimensions

The basic idea is that we live in a 3-brane hypersurface which is embedded in a higher dimensional world, the bulk. The fundamental higher dimensional Planck scale can be in the TeV scale and this translates to the strong hierarchies we observe in three spatial dimensions. According to Arkani-Hamed, Dimopoulos and Dvali, this can happen if the size of the extra dimension is large, even at the millimeter scale [114]-[116]. An alternative to producing large hierarchies is given in [117, 118] where the extra dimension is small and the metric of the 4-dimensional spacetime depends on the coordinates of the extra dimension.

4.1.1 The hierarchy problem

The large separation between the electroweak scale of $m_W \sim 1$ TeV and Planck scale $M_{\mathrm{P}l} = 1.2 \times 10^{19}$ GeV is dubbed the hierarchy problem. It arises from the radiative corrections to the Higgs boson mass

$$\delta m_H^2 = \frac{1}{8\pi^2}(\lambda_H^2 - \lambda_F^2)\Lambda^2 + \text{logarithmic divergence} + \text{finite terms} \quad (4.1)$$

which are quadratically divergent and depend on the cutoff scale Λ which lies in the GUT scale $\sim 10^{16}$ GeV or Planck scale $\sim 10^{19}$ GeV. Fermionic Yukawa couplings are denoted by λ_F and self-couplings to H by λ_H.

In order to produce a TeV scale mass for the Higgs boson, one has to get rid of the quadratic divergence. In dimensional regularization $1/\varepsilon$ singularities appear and these can be absorbed into the counterterms. The use of dimensional regularization could be justified if the SM was the fundamental theory. However, a complete theory should contain gravity alongside the SM and so a cutoff Λ has to be introduced. Thus, we are compelled to cancel the quadratic divergence in (4.1), which requires that the counterterm must be highly fine tuned.

4.1.2 The Arkani-Hamed-Dimopoulos-Dvali model

By considering our perceivable 4-dimensional world as a hypersurface embedded in a higher dimensional bulk we can circumvent the hierarchy problem. This can be seen if we write Newton's gravitational potential between two masses m_1 and m_2 in $D = 4 + d$ dimensions:

$$V(r) = -\frac{m_1 m_2}{M^{d+2} r^{d+1}}, \; r \ll R, \qquad (4.2)$$
$$V(r) = -\frac{m_1 m_2}{M_{\rm Pl}^2 r}, \; r \gg R,$$

where R represents the size of the largest dimensions, r denotes the distance between the masses and d denotes the number of extra spatial dimensions. The 4-dimensional Planck scale $M_{\rm Pl}$ is thus related to the higher dimensional fundamental scale M through

$$M_{\rm Pl}^2 = M^{d+2} V_d, \qquad (4.3)$$

where $V_d = (2\pi R)^d$ is the volume of the extra space.

If we demand that the fundamental scale has to be the electroweak scale $M = m_{EW}$, we can deduce the size of the extra dimension. One extra spatial dimension would result in $R = 10^{13}$ cm, which implies deviations from Newtonian gravity over solar distance scales. With $d = 2$ the size of extra dimensions can lie in the millimeter scale. This is also the limit of current experiments testing gravity.

4.1.3 Observation of large extra dimensions

Collider signatures include missing energy that is transferred into the bulk in graviton emission processes whose rates are suppressed by $M_{\rm Pl}$, and corrections to standard cross sections from graviton exchange [119]-[128]. Colliders have searched for extra dimensional gravitons in the processes $e^+e^- \rightarrow \gamma + \not{E}$ and $e^+e^- \rightarrow Z + \not{E}$ at LEP, and $p\bar{p} \rightarrow$ jet $+ \not{E}_T$ and $p\bar{p} \rightarrow \gamma + \not{E}_T$ at the Tevatron [23].

If the impact parameter b is smaller than the Schwarzschild radius for a collision with center of mass energy $\sqrt{s} \gg M$, black hole formation becomes possible [129, 130, 23]. The newly emerged black hole emits thermal radiation with Hawking temperature $T_H = (d+1)/(4\pi R_S)$, where R_S is the Schwarzschild radius. Thus, detectors should be sensitive to this thermal radiation. An example of black hole evaporation time is $10^{-26} - 10^{-27}$ s when the fundamental gravity scale is 1 TeV.

The possibility that gravity feels extra dimensions has fuelled much interest in experiments looking for deviations from Newton's gravitational inverse square law [131]. These deviations are parametrized by the modified Newtonian potential

$$V(r) = -G_N \frac{m_1 m_2}{r}[1 + \alpha \exp(-r/\lambda)] \tag{4.4}$$

with the parameter regions presented in Fig. 4.1. It has been discov-

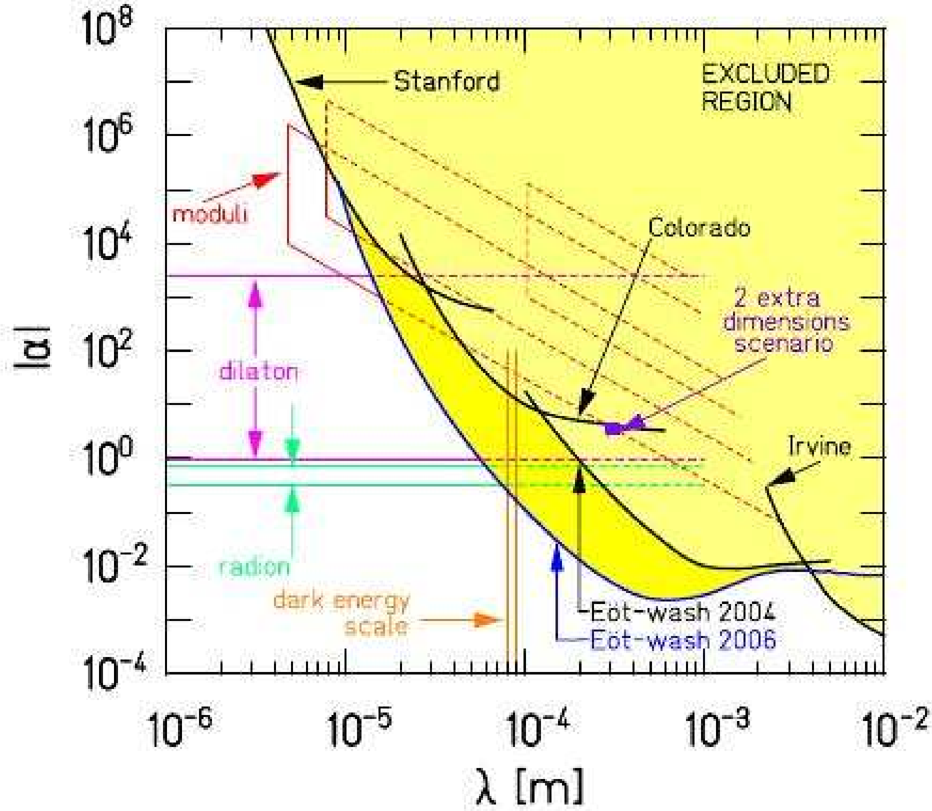

Figure 4.1: The figure taken from [132] shows experimental limits for λ and α from (4.4).

ered that astrophysical constraints prevent the observation of deviations to Newton's law. However, *e.g.* gauge bosons could mediate modifications to Newton's law and these modifications could produce a detectable signal in the coming experiments.

The fact that gravitons can exist in extra dimensions and excite KK modes places restrictions on extra dimensional gravity because these new degrees of freedom have to fulfill a number of astrophysical constraints [133]. One constraint comes from the supernova SN1987A where the KK gravitons cannot carry away more than half of the emitted energy [134]. The EGRET satellite is set out to measure the γ radiation that is produced the decays of KK gravitons from all supernovae in the universe [135]. Gravitons can also be trapped by supernova remnants and neutron stars. These gravitons decay into γs on occasion and limits on this radiation imply bounds for the fundamental scale and number of extra dimensions [133].

4.2 Brane-bulk models

Extra dimensions alter several features of high energy particle theories. Neutrino mixing, gauge coupling unification [136, 137] and gravity [119, 120] have to be modified. Different variations make some SM particles extra dimensional so that they can propagate in the bulk while the rest are restricted to the brane. Models where SM gauge bosons live in higher dimensions have been considered in [136, 137] and [138]-[144]. Here, we review [113] some generic features of brane-bulk models and later on we study how neutrinos can be accommodated to extra dimensional models.

Brane-bulk theories describe how SM particles residing on the brane interact with bulk particles that could be gravitons or SM singlets. The action describing the interactions of a bulk field $\phi(x, y)$ in a higher dimensional theory can be written as

$$\mathcal{S}_{\text{bulk}}[\phi] = \int d^4x d^d y \sqrt{|g_{(4+d)}|} \mathcal{L}(\phi(x, y)), \tag{4.5}$$

where x denotes the coordinates of 4-dimensional spacetime, y are the coordinates of the extra dimensions and $g_{(4+d)}$ is the metric of the $4 + d$-dimensional universe. Recall that we are restricted to flat extra dimensions and require that the metric is factorizable. Thus, the line element given by the metric becomes

$$ds^2 = g_{\mu\nu}(x) dx^\mu dx^\nu - \delta_{ab} dy^a dy^b. \tag{4.6}$$

The brane fields in turn form the action

$$\mathcal{S}_{\text{brane}}[\varphi] = \int d^4x d^d y \sqrt{|g_(4)|} \mathcal{L}(\varphi(x)) \delta^d(y - y_0), \tag{4.7}$$

where the d-dimensional δ function fixes the position of the brane to y_0 in the bulk and $g_{(4)}$ denotes the effective 4-dimensional metric. Finally, the brane-bulk action between a bulk field $\phi(x,y)$ and a fermionic field $\psi(x)$ on the brane is

$$\mathcal{S}_{\text{brane-bulk}} = \int d^4x d^dy \sqrt{|g_{(4)}|}\phi(x,y)\bar{\psi}(x)\psi(x)\delta^d(y-y_0). \tag{4.8}$$

To better understand the dynamics of the model one usually integrates over the extra dimensions. This procedure yields a 4-dimensional effective action. To demonstrate how this works we consider a 5-dimensional toy model, where the extra dimension is compactified on a circle of radius R and a scalar field $\phi(x,y)$ propagates in the bulk. The action is

$$\mathcal{S}[\phi] = \frac{1}{2}\int d^4x dy \left(\partial^A\phi\partial_A\phi - m^2\phi^2\right), \tag{4.9}$$

where the index $A = 0, ..., 4$. If the compactification of the extra dimension is reflected in the periodicity of the scalar field $\phi(y) = \phi(y+2\pi R)$, then the scalar field can be Fourier expanded as

$$\begin{aligned}\phi(x,y) \quad &= \quad \frac{1}{\sqrt{2\pi R}}\phi_0(x) + \sum_{n=1}^{\infty}\frac{1}{\sqrt{\pi R}}\left[\phi_n(x)\cos\left(\frac{ny}{R}\right)\right. \\ &\quad \left. +\hat{\phi}_n(x)\sin\left(\frac{ny}{R}\right)\right].\end{aligned} \tag{4.10}$$

Here $\phi_0(x)$ is referred to as the zero mode and it is independent of the extra dimension. The Fourier modes ϕ_n and $\hat{\phi}_n$ are the Kaluza-Klein (KK) modes. Introducing (4.10) into (4.9) and integrating over the extra dimension y yields

$$\begin{aligned}\mathcal{S}[\phi] \quad &= \quad \sum_{n=0}^{\infty}\frac{1}{2}\int d^4x\left(\partial^\mu\phi_n\partial_\mu\phi_n - m_n^2\phi_n^2\right) \\ &\quad +\sum_{n=1}^{\infty}\frac{1}{2}\int d^4x\left(\partial^\mu\hat{\phi}_n\partial_\mu\hat{\phi}_n - m_n^2\hat{\phi}_n^2\right)\end{aligned} \tag{4.11}$$

with $m_n^2 = m^2 + n^2/R^2$. Thus, in four dimensions the higher dimensional field appears as an infinite tower of fields with masses m_n.

4.2.1 Localization of fermions inside fat branes

A novel mechanism that can suppress proton decay induced by the variations of extra dimensional models was introduced in [145]-[147]. The model suggests that SM fermions are trapped inside a wall of size L at different points while SM gauge fields and the Higgs boson are free to propagate inside the wall. Couplings between fermions are suppressed because of the exponentially small overlaps of their wave functions. In particular, this suppression is needed in interactions between quarks and leptons that induce proton decay, and exponential suppression results if leptons and quarks sit at the opposite ends of the wall. The requirement that quarks and leptons live at the opposite ends of the wall comes from proton decay and not from the scenario itself.

We illustrate how fermion wave functions can be localized on the brane in a 5-dimensional model. Translational invariance in the fifth dimension is broken by a bulk scalar field Φ that develops a VEV $\langle\Phi\rangle$ that depends on the fifth dimension coordinate y. The VEV assumes the form of a domain wall centered at $y = 0$, and the zero mode of a bulk fermion is localized at the zero of $\langle\Phi\rangle(y)$. A single fermion case is described by the action

$$\mathcal{S} = \int d^4x dy \bar{\Psi}[i\Gamma^A\partial_A + \langle\Phi\rangle]\Psi, \tag{4.12}$$

where

$$\Gamma^\mu = \begin{pmatrix} 0 & \sigma^\mu \\ \bar{\sigma}^\mu & 0 \end{pmatrix}, \ \mu = 0, ..., 3, \qquad \Gamma^4 = i\begin{pmatrix} -\mathbb{I} & 0 \\ 0 & \mathbb{I} \end{pmatrix}. \tag{4.13}$$

The bulk fermion Ψ is a four component spinor with the chiral decompositions

$$\begin{aligned} \Psi_L(x,y) &= \sum_n f_n(y)\psi_{nL}(x), \\ \Psi_R(x,y) &= \sum_n g_n(y)\psi_{nR}(y), \end{aligned} \tag{4.14}$$

where $\psi_{nL,R}$ are four component spinors. By solving

$$(\partial_y + \langle\Phi\rangle)f_0 = 0, \qquad (-\partial_y + \langle\Phi\rangle)g_0 = 0 \tag{4.15}$$

we obtain the zero mode profiles

$$f_0(y) \sim \exp\Big[-\int_0^y ds\langle\Phi\rangle(s)\Big], \qquad g_0(y) \sim \exp\Big[\int_0^y ds\langle\Phi\rangle(s)\Big] \quad (4.16)$$

that are normalizable in finite space. For $\langle\Phi\rangle(y) = 2\mu^2 y$ the mode f_0 becomes centered at $y = 0$ and has the Gaussian form

$$f_0(y) = \frac{\mu^2}{(\pi/2)^{1/4}} \exp[-\mu^2 y^2] \quad (4.17)$$

and thus the left-handed massless fermion has been localized. If Ψ couples to $-\Phi$, the right-handed part is localized and the left-handed part is pushed to the other end of space.

4.2.2 Neutrinos in brane-bulk models

SM singlet neutrino can reside in the bulk while SM particles live on the brane [148]-[167]. The bulk neutrino can be expressed by two Weyl spinors $\Psi = (\psi_1, \psi_2)^T$. If the extra dimension is compactified on a circle with Z_2 orbifold, S^1/Z_2, then ψ_1 can be taken to be even under $y \to -y$ and ψ_2 is taken to be odd. If now the left-handed brane neutrino is restricted to the orbifold fixed point $y = 0$, then only the coupling to the even right-handed ψ_1 survives and the action is given by

$$\mathcal{S} \;=\; \int d^4x dy \Big[i\bar{\Psi}\Gamma^A\partial_A\Psi + \Big(\frac{h}{\sqrt{M}}\delta(y)\nu_L^\dagger H\psi_1 + \text{h.c.}\Big)\Big], \quad (4.18)$$

where $A = 0, ..., 4$.

The limit placed on neutrino couplings by the orbifold condition can be lifted if the brane is shifted. This mechanism exists in Type I string theory [169]. If the brane neutrinos can be located at a general coordinate y^* instead of orbifold fixed points and Fourier expansions are written for the higher dimensional spinors $\psi_{1,2}$

$$\begin{aligned} \psi_1(x,y) &= \frac{1}{\sqrt{2\pi R}} \sum_{n=0}^{\infty} \psi_1^{(n)}(x) \cos\Big(\frac{ny}{R}\Big) \\ \psi_2(x,y) &= \frac{1}{\sqrt{2\pi R}} \sum_{n=1}^{\infty} \psi_2^{(n)}(x) \sin\Big(\frac{ny}{R}\Big), \end{aligned} \quad (4.19)$$

then the coupling term $\bar{\nu}_L(\Psi + \Psi^c)\big|_{y^*}$+h.c. gives

$$m\nu_L^\dagger \left\{ \psi_1^{(0)} + \sqrt{2} \sum_{n=1}^{\infty} \left[\cos\left(\frac{ny^*}{R}\right)\psi_1^{(n)} + \sin\left(\frac{ny^*}{R}\right)\psi_2^{(n)c} \right] \right\} + \text{h.c.}, \quad (4.20)$$

where the upper index c refers to charge conjugation.

The effective 4-dimensional brane-bulk coupling in (4.18) is suppressed by the size of the extra dimension R. The brane neutrino couples to the kth mode with strength $h/\sqrt{2\pi RM}$ and in the general case of d extra dimensions this becomes $h/\sqrt{V_d M^d}$. If we recall the relation (4.3), then the effective coupling is hM/M_{Pl} which holds for any number of extra dimensions.

After electroweak symmetry breaking the Higgs field develops a VEV and this induces a mass for the brane neutrino. This mass is given by [148]-[150] (see also the review in [168])

$$m_\nu = \frac{hvM}{M_{\text{Pl}}}. \quad (4.21)$$

With $M \sim 10 - 100$ TeV the neutrino mass is $m_\nu \simeq (10^{-3} - 10^{-2})h$ eV and by setting the coupling $h \sim 10$ the neutrino mass m_ν attains values relevant in neutrion oscillations. The kth KK mode has mass $m_k = kR^{-1}$ and the mixing angle is $\sqrt{2}m_\nu R/k$.

The existence of an infinite tower of KK modes of singlet neutrinos can influence neutrino oscillations. The KK modes emerge with increaing mass and decreasing mixing. Analyses performed in [151, 155, 157, 160, 162] and [165]-[167] suggest the size of the extra dimension has to satisfy

$$\begin{aligned} R^{-1} &\geq 0.02 \text{ eV} \quad \text{hierarchical spectrum,} \\ R^{-1} &\geq 0.22 \text{ eV} \quad \text{inverted spectrum,} \\ R^{-1} &\geq 4.1 \text{ eV} \quad \text{degenerate spectrum.} \end{aligned} \quad (4.22)$$

4.3 Leptogenesis from bulk neutrinos in a split fermion scenario

The research articles [1] and [2] consider a model with bulk neutrinos and split fermions, which have been previously studied in [170]-[177], from the point of view of leptogenesis. The first version of this model was originally introduced in [178] and it had to be modified in order to produce CP violation. Leptogenesis in large flat extra dimensions has been considered before [179] with all the KK modes but without split fermions. Our treatment, on the other hand, concentrates on the effective model that includes the brane neutrinos and the bulk zero mode. In addition, a key feature of our model is the fat brane that accommodates split fermions, which allows for brane-bulk couplings away from orbifold fixed points.

4.3.1 The effective model

The model consists of a Dirac neutrino $\Psi = (\psi_+, \psi_-)^T$ propagating in the bulk and left-handed brane neutrinos $\nu_\alpha$with α denoting flavor. The split fermion idea introduced in chapter 4.1 also emerges as the brane neutrinos reside at different locations y_α in the extra dimensions and their wave functions are described by the Gaussian

$$\nu_\alpha(x,y) = \frac{1}{\sqrt{\sigma}} \exp\Big(-\frac{\pi}{2}\frac{(y-y_\alpha)^2}{\sigma^2}\Big)\nu_\alpha(x). \tag{4.23}$$

Brane-bulk interactions between the brane neutrinos, bulk neutrino and Higgs field give rise to the 4D mass terms [1]

$$\begin{aligned}\mathcal{S}_{\rm mass} &= \int d^4x \Big\{ \sum_{\alpha=1}^{n_f} \nu_\alpha^\dagger(x)\Big[m\psi_0^{(0)c} + \sum_{n>0}\big(m^\alpha_{n,+}\psi_+^{(n)c}(x) + m^\alpha_{n,-}\psi_-^{(n)}(x)\big)\Big] + \text{h.c.} \\ &\quad + \sum_{n>0}\frac{n}{R}\big(\psi_+^{(n)\dagger}(x)\psi_-^{(n)}(x) + \psi_-^{(n)\dagger}(x)\psi_+^{(n)}(x)\big)\Big\},\end{aligned} \tag{4.24}$$

where $\psi_+^{(n)}$ are now left-handed and even under the Z_2 symmetry $y \to -y$ and the right-handed $\psi_-^{(n)}$ are odd. In contrast to [178], the couplings $m^\alpha_{n,-}$ are now complex which is necessary for CP violation. The Yukawa interactions with the Higgs field are

$$\begin{aligned}\mathcal{S}_Y &= \int d^4x \sum_{\alpha=1}^{n_f} \nu_\alpha^\dagger(x)\Big[\frac{mh(x)}{v}\psi_+^{(0)c}(x) \\ &\quad + \sum_{n>0}\Big(\frac{m^\alpha_{n,+}h(x)}{v}\psi_+^{(n)c}(x) + \frac{m^\alpha_{n,-}h(x)}{v}\psi_-^{(n)}(x)\Big)\Big] + \text{h.c.}\end{aligned} \tag{4.25}$$

One of the objectives in [178, 1] was to determine the mass spectrum of the system when the higher KK modes decouple. What remains are the zero mode $\psi_+^{(0)}$ and the brane neutrinos ν_α. In order to obtain a viable mass spectrum from the effective mass matrix

$$\widetilde{\mathcal{M}}_L = \begin{pmatrix} m_{\alpha\beta} & m \\ m & 0 \end{pmatrix}, \tag{4.26}$$

$$m = M_* \sqrt{\frac{\sigma}{\pi R}}, \tag{4.27}$$

$$m_{\alpha\beta} = -\frac{M_*^2 \sigma}{2} \Bigg\{ e^{i\delta_\beta} \left[\mathrm{Erf}\left(\frac{\sqrt{\pi}}{2\sigma}(y_\alpha + y_\beta) \right) - \mathrm{Erf}\left(\frac{\sqrt{\pi}}{2\sigma}(y_\alpha - y_\beta) \right) \right] + e^{i\delta_\alpha} \left[\mathrm{Erf}\left(\frac{\sqrt{\pi}}{2\sigma}(y_\alpha + y_\beta) \right) + \mathrm{Erf}\left(\frac{\sqrt{\pi}}{2\sigma}(y_\alpha - y_\beta) \right) \right] \Bigg\} \tag{4.28}$$

with at least one light state χ_1 corresponding to the observed light neutrinos and two heavy unstable states, χ_2 and χ_3, the number of brane neutrinos has to be $n_f = 2$. Thus, one extra left-handed brane neutrino that is a SM singlet needs to be included.

By solving the eigenvalues of (4.26 in the weak coupling limit $mR \ll 1$ one obtains the masses for the light neutrino χ_1 and two heavy neutrinos $\chi_{2,3}$. The weak coupling implies to $M_*^2 \sigma \ll m$ which allows us to find approximate expressions for the heavy masses

$$m_2 \simeq \left| -\sqrt{2}m + \frac{1}{2}m_{12} + \frac{1}{4}(m_{11} + m_{22}) \right|, \tag{4.29}$$

$$m_3 \simeq \left| \sqrt{2}m + \frac{1}{2}m_{12} + \frac{1}{4}(m_{11} + m_{22}) \right| \tag{4.30}$$

and the light neutrino mass

$$m_1 \simeq \left| \frac{1}{2}(m_{11} + m_{22}) - m_{12} \right|. \tag{4.31}$$

Thus, the spectrum contains two quasidegenerate heavy neutrinos that can decay to the light neutrino and produce lepton asymmetry.

4.3.2 CP violation

The nonvanishing CP violation arises from the out-of-equilibrium decays of the two heavy mass eigenstates $\chi_{2,3}$ into the light state χ_1 and Higgs doublet, see Feynman diagrams in Fig. 4.2 for the decays of χ_2. The

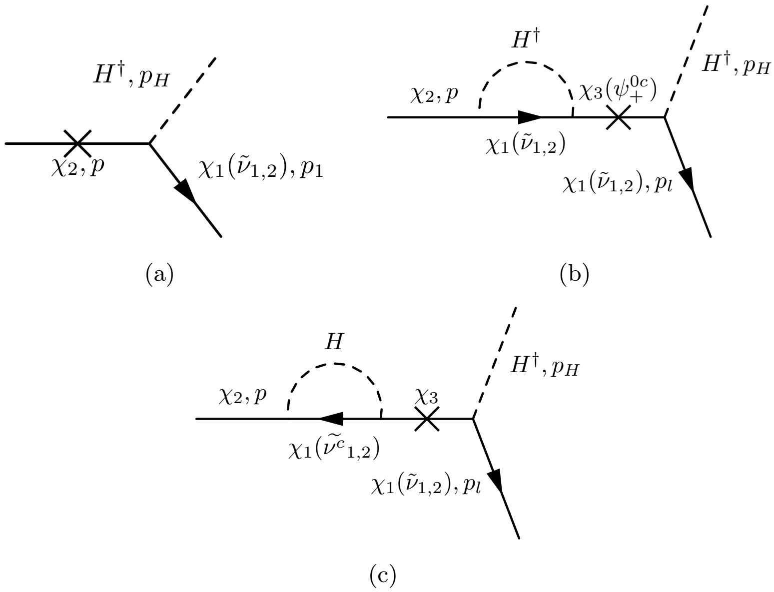

Figure 4.2: The relevant Feynman diagrams for the process $\chi_2 \to \chi_{1\mathrm{L}} H^\dagger$. The tree level diagram due to the decay of χ_2 to a neutrino and Higgs is in Fig. 4.2(a). Fig. 4.2(b) and 4.2(c) depict the mixing diagrams due to the the decay of χ_2 to a light neutrino and Higgs.

complex phases we have introduced in the brane-bulk couplings (4.24) and (4.25) are essential in producing the asymmetry.

The hierarchy in the neutrino mass spectrum reveals which diagrams in [2] are the chief sources of CP violation. Since we have a quasidegenerate spectrum of heavy neutrinos with the difference between the heavy neutrino masses being small, $|m_2 - m_3| \sim M_*^2\sigma \ll m_{2,3}$, the major contribution to the CP violation parameter is estimated to come from the diagrams in Fig. 4.2 similar to the tree level Fig. 3.1(a) and mixing diagrams Fig. 3.1(c) [103]-[106].

Our framework departs from earlier considerations [103]-[106] as the condition $A_{ij} = A_{ji}^*$ of [105] familiar from the $SO(10)$ motivated models is not fulfilled in our scenario [1]. Models in [103]-[106] have right-handed neutrinos with masses close to the GUT scale and thus the SM particles are massless whereas the framework depicted in (4.24) and (4.25) shows that the electroweak phase transition has taken place. Thus, the brane neutrinos have acquired mass terms through the Higgs VEV and the scale of the effective theory containing the zero mode and two brane neutrinos lies in the TeV region. The neutrino spectrum now includes a light neutrino mass that potentially corresponds to one of the mass eigenstates observed in neutrino experiments. Because all the neutrino masses are now included, the couplings are different as we have to pick up the correct interaction component at leading order from the light neutrino χ_1 running in the loop and from the propagating heavy states $\chi_{2,3}$.

With the above considerations we have obtained a lengthy expression for ε [1] where higher order Yukawa terms do not cancel as they do in [105]. This is indeed a result of the fact that the light neutrino is now part of the neutrino mass spectrum alongside the unstable heavy states.

4.3.3 Monte Carlo and parameter regions

To better understand the behavior of the CP violation parameter with varying parameters $\tilde{y}_{1,2},\ \delta_{1,2},\ R$ and the light neutrino mass scale $M_*^2\sigma$ we performed Monte Carlo analysis. The idea behind Monte Carlo analysis is the randomization of the various parameters in desired intervals and then imposing a number of constraints. In our case randomized variables were $\tilde{y}_{1,2},\ \delta_{1,2},\ R$ and $M_*^2\sigma$ [2]. The constraints come from the out-of-equilibrium condition of the decays, the washout parameter, the light neutrino mass and most importantly from the observed baryon asymmetry of the universe (2.14) [2].

The parameter scans predict that the radius of the extra dimension R has to be small for the model to produce a correct amount of lepton number and subsequent baryon number [2]. This brings the masses of the heavy states to the TeV scale, which is also required.

The higher order Yukawa terms $\sim m^4 A_{ij}^2$ need to be included in ε so that sufficient CP violation is generated. These terms correspond to higher order diagrams with more loops but these higher order diagrams would be suppressed by the higher powers of m/v which are small, $(m^2/v^2)m_{2,3} \ll m_{2,3}$.

4.3.4 Conclusions

We have studied CP violation and leptogenesis in an extra dimensional model with one flat extra dimension of size R where a bulk neutrino propagates while two neutrinos are restricted to the brane. Our setup constrains R to values that are a few orders of magnitude away from the Planck scale. With more extra dimensions the radii could be larger. On the other hand, adding more dimensions would change the hierarchy as the brane bulk coupling becomes suppressed by additional factors σ/R.

Direct comparisons to neutrino oscillation experiments could be drawn if the number of brane neutrinos is increased to four. This would produce three light neutrinos and two heavy ones that are unstable and decay to the lighter states. When it comes to sub-eV sterile neutrinos [180, 181], even more brane neutrinos have to be added to the framework so that comparisons could be drawn to 3+2 models.

As the light neutrino and unstable heavier neutrinos become intertwined in the mass spectrum, higher order term become important in the CP violation parameter. This scenario departs from earlier considerations where the unstable neutrinos have masses in the GUT scale while SM states remain massless because the electroweak phase transition has not occurred.

Chapter 5

Leptogenesis in supersymmetric models

Supersymmetry is another approach used to solve the hierarchy problem discussed earlier, and there are high expectations of discovering supersymmetric particles at the Large Hadron Collider (LHC). Supersymmetry doubles the particle content compared to the SM, giving each SM fermion a bosonic partner, a sfermion, and likewise each SM boson receives a fermionic partner. If the loop corrections to the Higgs boson mass are reconsidered, the bosonic and fermionic contributions cancel and no fine tuning is needed. Also, within supersymmetry all three SM gauge couplings become equal at an energy scale $\sim 10^{16}$ GeV thus making supersymmetry a viable GUT theory. As much attention is devoted to supersymmetric models and LHC is driving to observe supersymmetric particles, this chapter is devoted to supersymmetric models that can also accommodate leptogenesis.

The first section of this chapter gives a general idea of supersymmetry following the reviews [182, 8, 183] and supersymmetric leptogenesis. Soft leptogenesis is reviewed in the second section and the third introduces a new supersymmetric source of CP violation and lepton asymmetry and our research [3].

5.1 Unbroken supersymmetry and leptogenesis

Although up to eight supersymmetries are possible in four dimensions, the case $N = 1$ is probably the only one relevant for particle physics. The reason is that $N > 1$ theories are nonchiral and it is difficult to introduce supersymmetry breaking into them. Two irreducible representations are found in $N = 1$ supersymmetry, and these are the chiral and vector superfields. Chiral superfields $\Phi(\theta, \bar{\theta}, x)$ contain a Weyl spinor $\psi_\alpha(x)$, a supertpartner complex scalar field $\phi(x)$ and an auxiliary field $F(x)$ and also anticommuting Grassmann variables $\theta, \bar{\theta}$ are included in their expansions. Vector superfields $V(\theta, \bar{\theta}, x)$ contain a Weyl spinor $\lambda_\alpha(x)$, a superpartner vector field $A_\mu(x)$ and an auxiliary field $D(x)$.

5.1.1 Supersymmetric Lagrangian and scalar potential

To write a supersymmetric Lagrangian chiral superfields Φ_i transforming in different representations of a gauge group G are considered. A vector superfield V^a is introduced for every generator of the gauge group. The most general superspace Lagrangian is

$$\mathcal{L} = \sum_i \int d^4\theta \Phi_i^\dagger e^V \Phi_i + \sum_a \frac{1}{4g_a^2} \int d^2\theta W_\alpha^2 + \text{c.c.} + \int d^2\theta W(\Phi_i) + \text{c.c.} \tag{5.1}$$

where the chiral and vector superfields $V = V^a T^a$ in Wess-Zumino gauge are

$$\begin{aligned} \Phi(x, \theta, \bar{\theta}) &= \phi(x) + \sqrt{2}\theta\psi(x) + \theta^2 F + ..., \\ V(x, \theta, \bar{\theta}) &= -\theta\sigma^\mu\bar{\lambda}A_\mu + i\theta^2\bar{\theta}\bar{\lambda} - i\bar{\theta}^2\theta\lambda + \frac{1}{2}\theta^2\bar{\theta}^2 D, \end{aligned} \tag{5.2}$$

the chiral superfield

$$W_\alpha = -i\lambda_\alpha + \theta_\alpha D - \frac{i}{2}(\sigma^\mu\bar{\sigma}^\nu\theta)_\alpha F_{\mu\nu} + \theta^2\sigma^\mu_{\alpha\dot{\beta}}\partial_\mu\bar{\lambda}^{\dot{\beta}} \tag{5.3}$$

is the analog of gauge invariant field strength and $W(\Phi)$ is a holomorphic function of chiral superfields known as the superpotential. The superpotential can be at most cubic in the chiral fields to ensure renormalizability.

The scalar potential describes the interactions between the scalar fields that belong to the chiral supermultiplets. It is given by

$$V(\phi_i, \phi_i^*) = \sum_i |F_i|^2 + \frac{1}{2}\sum_a D_a D^a, \tag{5.4}$$

where the auxiliary fields F_i and D_a are given by the equations of motion

$$\begin{aligned} F_i &= -\frac{\partial W^*}{\partial \phi_i}, \\ D_a &= -g_a \phi^\dagger T^a \phi \end{aligned} \tag{5.5}$$

with $\phi = (\phi_1, ..., \phi_n)^T$.

5.1.2 Minimal Supersymmetric SM

The simplest supersymmetric version of the SM is known as the Minimal Supersymmetric SM (MSSM) and it is built on the same gauge group as the SM, $SU(3)_C \times SU(2)_L \times U(1)_Y$. All SM fermions are assigned to separate chiral supermultiplets according to whether they are $SU(2)_L \times U(1)_Y$ doublets or singlets. The fermions sit in the supermultiplets with their scalar partners, the sfermions. Gauge fields are in vector supermultiplets and they receive fermionic spin 1/2 partners, the gauginos. Two Higgs supermultiplets are also included.

Interactions among the chiral superfields are described by the superpotential

$$W = \bar{u}\mathbf{y_u}QH_u - \bar{d}\mathbf{y_d}QH_d - \bar{e}\mathbf{y_e}LH_d + \mu H_u H_d, \tag{5.6}$$

where Q and L stand for the chiral quark and lepton supermultiplets, respectively. The former contains left-handed quarks and their superpartners, the squarks, and similarly the latter includes $SU(2)_L \times U(1)_Y$ doublet left-handed leptons and scalar sleptons. The supermultiplets $\bar{u}$ and $\bar{d}$ contain $SU(2)_L \times U(1)_Y$ singlet right-handed quarks and corresponding squarks, which are now different from the squarks in the left-handed supermultiplets, and $\bar{e}$ stands for the supermultiplets that contain right-handed charged leptons and their superpartners that are again different from the sleptons in L. Two Higgs supermultiplets H_u with $Y = +1/2$ and H_d with $Y = -1/2$ are needed to cancel the electroweak gauge anomaly and to make the superpotential analytic. Both supermultiplets comprise fermionic partners of the scalar Higgses, the higgsinos.

More terms besides the ones in (5.6) can be included in the superpotential. The additional terms

$$\begin{aligned} W_{\Delta L=1} &= \frac{1}{2}\lambda^{ijk}L_iL_j\bar{e}_k + \lambda'^{ijk}L_iQ_j\bar{d}_k + \mu'^i L_iH_u, \\ W_{\Delta B=1} &= \frac{1}{2}\lambda''^{ijk}\bar{u}_i\bar{d}_j\bar{d}_k \end{aligned} \tag{5.7}$$

violate lepton and baryon numbers by one unit. What speaks against these additions is the fact that the proton is stable. Excluding baryon and lepton number violating terms can be done by invoking a new symmetry called R parity [184] or matter parity [185]-[188]. The corresponding matter parity quantum number for each particle is

$$P_M = (-1)^{3(B-L)}. \tag{5.8}$$

Quark and lepton supermultiplets have $P_M = -1$ while the Higgs supermultiplets, gauge bosons and gauginos are all even $P_M = +1$. Terms in the superpotential are allowed if the product of all $P_M = +1$. Moving from matter parity to R parity introduces the quantum number

$$P_R = (-1)^{3(B-L)+2s}, \tag{5.9}$$

where s is the spin of the particle. Particles within the same supermultiplet have different R parities, and all SM particles and Higgs bosons have $P_R = +1$ whereas squarks, sleptons, gauginos and higgsinos have $P_R = -1$.

The importance of R parity becomes obvious if MSSM conserves R parity. First, the lightest supersymmetric particle with $P_R = -1$ has to be stable. A neutral lightest supersymmetric particle would provide for a promising candidate for nonbaryonic dark matter [189, 190]. Each supersymmetric partner must decay to an odd number of light supersymmetric particles. Finally, collider experiments can produce processes where even numbers of sparticles can emerge at a time.

5.1.3 MSSM extended with right-handed neutrinos

A supersymmetric version of leptogenesis was first considered in [90] where singlet neutrino chiral superfields were added to the superpotential. The additional piece to the superpotential is

$$W = \frac{1}{2}M_iN_iN_i + h_{ij}\varepsilon_{\alpha\beta}L_i^\alpha H^\beta N_j. \tag{5.10}$$

Here the masses M_i provide lepton number violation and the complex Yukawa couplings are the source for CP violation as before. New contributions to CP violation now arise as the supersymmetric partners of leptons and the Higgs bosons can now run in the loops of Figs. 3.1(b) and 3.1(c). Also, the final states in diagrams in Fig. 3.1 can now consist of sleptons and higgsinos, the superpartners of leptons and Higgs bosons. There is also another source for the asymmetry as the superpartners of the singlet neutrinos can decay into leptons and sleptons.

5.2 Supersymmetry breaking and leptogenesis

The fact that no superpartners of the SM particles have been observed implies that supersymmetry is broken. If global supersymmetry is spontaneously broken in the vacuum state, then the total energy has the expectation value

$$\langle H \rangle = \frac{1}{4}\big(\|Q_1^\dagger|0\rangle\|^2 + \|Q_1|0\rangle\| + \|Q_2^\dagger|0\rangle\|^2 + \|Q_2|0\rangle\|^2\big) > 0, \qquad (5.11)$$

where Q_α and $Q_{\dot\alpha}^\dagger$ are supersymmetry generators. Supersymmetry is spontaneously broken if either the F-term F_i [191] or the D-term D^a [192, 193] is nonvanishing in the ground state. Spontaneous supersymmetry breaking predicts the existence of Nambu-Goldstone particle, which in this case is a fermionic Weyl fermion, the goldstino. The goldstino is the fermionic partner of the auxiliary field and these two reside in the same supermultiplet.

The source for supersymmetry breaking in the MSSM is believed to come from a hidden sector whose particles do not interact with the chiral fields of the MSSM or these sectors couple very weakly. The mediated supersymmetry breaking is thought to have two possible routes. The first setup considers gravity-mediated supersymmetry breaking where the supersymmetry breaking scale sits somewhere between the electroweak and Planck scales. The second approach is gauge-mediated supersymmetry breaking where new chiral supermultiplets have MSSM gauge interactions and couple to the supersymmetry breaking $\langle F \rangle$. In this case the supersymmetry breaking scale can be as low as the electroweak scale.

5.2.1 Soft supersymmetry breaking

Mechanisms that break supersymmetry at energy scales above the electroweak scale produce terms that explicitly break supersymmetry in the low energy effective theory. By ignoring the actual source of supersymmetry breaking, soft supersymmetry breaking terms can be written for the scalar fields of the theory and the couplings should have positive mass dimension which ensures the hierarchy between the electroweak and Planck scales. The general form of the soft supersymmetry breaking Lagrangian is

$$\begin{aligned}\mathcal{L}_{\rm soft} &= -\frac{1}{2}\Big(M_a\lambda^a\lambda^a + \frac{1}{3}a^{ijk}\phi_i\phi_j\phi_k + b^{ij}\phi_i\phi_j + t^i\phi_i\Big) + \text{h.c.} \\ &\quad -(m^2)^i_j\phi^{j*}\phi_i.\end{aligned} \tag{5.12}$$

The tadpole couplings t^i are reserved for singlet fields ϕ_i.

Applying the above priciples to the MSSM gives the supersymmetry preserving and supersymmetry breaking Lagrangians with the total being $\mathcal{L} = \mathcal{L}_{\rm SUSY} + \mathcal{L}_{\rm soft}$. Writing the soft terms for the MSSM gives

$$\begin{aligned}\mathcal{L}^{\rm MSSM}_{\rm soft} &= -\frac{1}{2}\big(M_3\tilde{g}\tilde{g} + M_2\widetilde{W}\widetilde{W} + M_1\tilde{B}\tilde{B} + \text{c.c}\big) \\ &\quad -\tilde{\bar{u}}\mathbf{a_u}\tilde{Q}H_u - \tilde{\bar{d}}\mathbf{a_d}\tilde{Q}H_d - \tilde{\bar{e}}\mathbf{a_e}\tilde{L}H_d + \text{c.c.}\big) \\ &\quad -\tilde{Q}^\dagger\mathbf{m}^2_Q\tilde{Q} - \tilde{L}^\dagger\mathbf{m}^2_L\tilde{L} - \tilde{\bar{u}}\mathbf{m}^2_{\bar{u}}\tilde{\bar{u}}^\dagger - \tilde{\bar{d}}\mathbf{m}^2_{\bar{d}}\tilde{\bar{d}}^\dagger - \tilde{\bar{e}}\mathbf{m}^2_{\bar{e}}\tilde{\bar{e}}^\dagger \\ &\quad -m^2_{H_u}|H_u|^2 - m^2_{H_d}|H_d|^2 - (bH_uH_d + \text{c.c.}),\end{aligned} \tag{5.13}$$

where tilde stands for the superpartner and $\tilde{g}$, $\widetilde{W}$ and $\tilde{B}$ are the gluino, wino and bino, respectively. The soft parameters are expected to be ~ 1 TeV so that the Higgs mass would not receive too large corrections that cause the hierarchy problem.

5.2.2 Soft leptogenesis

The inclusion soft supersymmetry breaking causes new effects to arise [194, 195]. The premise is the superpotential

$$W = h_{ij}N_iL_jH + \frac{1}{2}M_{ij}N_iN_j$$

introduced before and the soft supersymmetry breaking terms

$$\mathcal{L}_{\rm soft} = \tilde{m}_{ij}\tilde{N}^*_i\tilde{N}_j + \big(A_{ij}h_{ij}\tilde{N}_i\tilde{\ell}_jH + \frac{1}{2}B_{ij}M_{ij}\tilde{N}_i\tilde{N}_j + \text{h.c.}\big). \tag{5.14}$$

The new effect is that only one family of singlet neutrinos are needed. The superpartner of the singlet neutrino, corresponding to the sneutrino, is sufficient to produce the CP asymmetry.

The sneutrino mass eigenstates $\tilde{N}_+$ and $\tilde{N}_-$ are the decaying states of interest. Their masses are split by the soft parameters, and the squared masses are

$$M_\pm^2 = M^2 + \tilde{m}^2 \pm |BM|, \tag{5.15}$$

where M, $\tilde{m}$ and B arise from the matrices M_{ij}, $\tilde{m}_{ij}$ and B_{ij} of (5.14), respectively, in the one singlet neutrino generation case. The sneutrinos can decay into fermionic final states that are the lepton and higgsino, $\tilde{N}_\pm \to \ell\tilde{H}$, and into bosonic final states that include the slepton and Higgs boson, $\tilde{N}_\pm \to \tilde{\ell}H$. The CP violation parameter can be defined as

$$\varepsilon = \frac{\sum_f [\Gamma(\tilde{N}_+ \to f) + \Gamma(\tilde{N}_- \to f) - \Gamma(\tilde{N}_+ \to \bar{f}) - \Gamma(\tilde{N}_- \to \bar{f})]}{\sum_f [\Gamma(\tilde{N}_+ \to f) + \Gamma(\tilde{N}_- \to f) + \Gamma(\tilde{N}_+ \to \bar{f}) + \Gamma(\tilde{N}_- \to \bar{f})]} \tag{5.16}$$

with $f = \tilde{H}\ell$ or $f = H\tilde{\ell}$. Because the mass difference is small $M_+ - M_- = |B| \ll M_\pm$, the largest contribution to (5.16) comes from the mixing diagrams [103]-[106] similar to Fig. 3.1(c) with singlet neutrinos $N_{1,2}$ now replaced by the sneutrinos $\tilde{N}_\pm$.

The soft supersymmetry breaking terms induce sneutrino mixing and thus the sneutrino mass eigenstates can generate lepton asymmetry. This scenario departs form non-supersymmetric case in that only one family of singlet neutrinos and their superpartners are needed. Another difference comes from the singlet neutrino mass scale M which in standard leptogenesis scenarios has to obey $M \gtrsim 10^9$ GeV [196] that is quite large compared to the upper bound of for the reheat temperature $T_R < 10^7$ GeV mentioned earlier. In soft leptogenesis the mass scale $M \lesssim 10^9$ GeV and so the reheat temperature and singlet neutrino masses no longer contradict each other.

5.3 Leptogenesis from sneutrino-antisneutrino asymmetry

Gauging $B - L$ in supersymmetry extends the model with the gauge group $U(1)_{B-L}$, and a new source for CP violation and leptogenesis is obtained

[197]. If the MSSM is augmented with singlet neutrinos and $B-L$ symmetry which is broken when SM singlets Δ and $\bar{\Delta}$ obtain VEVs, new heavy mass eigenstates emerge. These states can decay asymmetrically into sneutrinos and antisneutrinos thus providing lepton asymmetry before the standard soft leptogenesis or non-supersymmetric leptogenesis mechanisms take effect.

5.3.1 The model with MSSM Higgs sector

Gauging $B-L$ in the MSSM can be done with the help of new chiral superfields Δ, $\bar{\Delta}$, S and N_i with $B-L$ charges -2, +2, 0 and +1, respectively. The fields Δ and $\bar{\Delta}$ develop large vacuum expecation values and thus form $B-L$ breaking Majorana masses for the singlet neutrinos N_i. The superpotential is

$$\begin{aligned} W^{(B-L)} &= \lambda S(\Delta\bar{\Delta} - M^2) + \frac{1}{2} f_{ij} N_i N_j \Delta + Y_\nu^{\alpha i} L_\alpha N_i H_u \\ &\quad + \mu H_u H_d + M_1 S^2 + M_2 \Delta\bar{\Delta} + Y_1 S^3 \\ &\quad + Y_3 S H_u H_d, \end{aligned} \tag{5.17}$$

where the first line is from the setup considered in [197]. The second line is an addition we have made [3] and it couples the MSSM Higgs sector to the $B-L$ breaking sector via the superfield S.

Interactions between the chiral superfields Δ, $\bar{\Delta}$, S and the vector superfield $\mathcal{V}_B$ coming from the gauge group $U(1)_{B-L}$ are given by the Kähler Lagrangian

$$\begin{aligned} \mathcal{L}_D^{(B-L)} &= \int d^4\theta \Big(\Delta^\dagger e^{q_\Delta g_B \mathcal{V}_B} \Delta + \bar{\Delta}^\dagger e^{-q_\Delta g_B \mathcal{V}_B} \bar{\Delta} \\ &\quad + \sum_i \Phi_i^\dagger e^{q_i g_B \mathcal{V}_B} \Phi_i \Big). \end{aligned} \tag{5.18}$$

Here Φ_i denote the chiral superfields of MSSM and $q_{\Delta,i}$ denote the $B-L$ charges of the superfields Δ and Φ_i. By moving to unitary gauge

$$\begin{aligned} \Delta &= \frac{1}{\sqrt{2}}(|M| + \Delta_0) e^{q_\Delta g_B \Delta'} \\ \bar{\Delta} &= \frac{1}{\sqrt{2}}(|M| + \Delta_0) e^{-q_\Delta g_B \Delta' + i\phi} \end{aligned} \tag{5.19}$$

and making the supersymmetric transformations

$$\begin{aligned}\mathcal{V}_B &= \mathcal{V}_B^0 - \Delta' - \Delta'^\dagger &(5.20)\\ \Phi_i &\to e^{q_i g_B \Delta'}\Phi_i &(5.21)\end{aligned}$$

the Kähler Lagrangian becomes

$$\begin{aligned}\mathcal{L}_D^{(B-L)} &= \int d^4\theta \Big[(|M|^2 + |M|(\Delta_0 + \Delta_0^\dagger) + \Delta_0^\dagger\Delta_0)\cosh(q_\Delta g_B \mathcal{V}_B^0) \\ &\quad + \sum_i \Phi_i^\dagger e^{q_i g_B \mathcal{V}_B^0}\Phi_i\Big]. &(5.22)\end{aligned}$$

After integrating out the gauge vector superfield $\mathcal{V}_B^0$ the effective theory couplings for the interactions $(\Delta_0^\dagger + \Delta_0)(\sum_i \Phi_i^\dagger q_i \Delta_i)^2$, $\Delta_0^\dagger\Delta_0(\sum_i \Phi_i^\dagger q_i \Phi_i)^2$ and $(\Delta_0^2 + \Delta_0^{\dagger 2})(\sum_i \Phi_i^\dagger q_i \Phi_i)^2$ go as $1/|M|^3$, $1/|M|^4$ and $1/|M|^4$, respectively. Thus, interactions of MSSM superfields with Δ_0 and $\Delta_0^\dagger$ via exchange of the gauge superfield are suppressed by terms at least $1/|M|^3$.

Like any supersymmetric theory also this model has to exhibit supersymmetry breaking which is introduced by the soft supersymmetry breaking potential

$$\begin{aligned}V_{\text{soft}}^{(B-L)} &= [bH_uH_d + a_1 S + b_1 S^2 + b_2\Delta\bar{\Delta} + c_1 S^3 + c_2 S\Delta\bar{\delta} \\ &\quad + c_3 S H_u H_d + \frac{A_f f_{ij}}{2}\delta\tilde{N}_i\tilde{N}_j] + \text{h.c.} \\ &\quad + m_{H_u}^2|H_u|^2 + m_{H_d}^2|H_d|^2 + m_S|S|^2 \\ &\quad + m_\Delta|\Delta|^2 + m_{\bar{\Delta}}^2|\bar{\Delta}|^2. &(5.23)\end{aligned}$$

Supersymmetry breaking causes splittings in the masses of S, Δ and $\bar{\Delta}$ and shifts $\langle S\rangle$. Eq. (5.19) shows the transformation made to Δ and $\bar{\Delta}$ as we move to unitary gauge. After this transformation the mass spectrum for the scalar particles H_u, H_d, S and Δ_0 can be found from the scalar potential $V_F + V_D + V_{\text{soft}}$ by use of perturbation theory.

The mass spectrum consists of four Higgs bosons that correspond to the ones found in the MSSM and four heavier boson states originating from the $B-L$ sector. The mass eigenstates are denoted by X_i while the original states are $X_i' =$ (ReH_u,ImH_u, ReH_d, ImH_d, ReS, ImS, ReΔ_0, ImΔ_0) and

$$X_i' = \sum_{j=1}^{8} n_{ij}X_j. \tag{5.24}$$

The coefficients n_{ij} form the matrix that diagonalizes the mass matrix that is given in the non-diagonal basis X'_i. The real states X_6, X_7 and X_8 corresponding to the Higgs bosons of the MSSM receive corrections from the $B - L$ sector and thus the direct couplings Y_3 and c_3 have to be sufficiently small. The four states $X_2, ..., X_5$ are heavy with quasidegenerate masses $(M_{X_2}, M_{X_3}, M_{X_4}, M_{X_5}) \sim M$ and couple to the sneutrinos. The masses M_{X_i} and the diagonalizing matrix have been obtained by the use of perturbtion theory as described in the Appendix of [3].

5.3.2 New source for CP violation

CP violation arises from the decays of the states $X_2, ..., X_5$, Fig. 5.1, into sneutrinos and antisneutrinos, which induces an initial asymmetry in the abundances of sneutrinos and antisneutrinos. This influences the decays of sneutrinos and antisneutrinos into MSSM leptons, antileptons, sleptons and antisleptons which is the case of soft leptogenesis. As we have added new terms to the superpotential (5.17), CP violation arises not only from the soft potential but from (5.17) as well, this mechanism of generating lepton number via sneutrino-antisneutrino asymmetry is not soft leptogenesis.

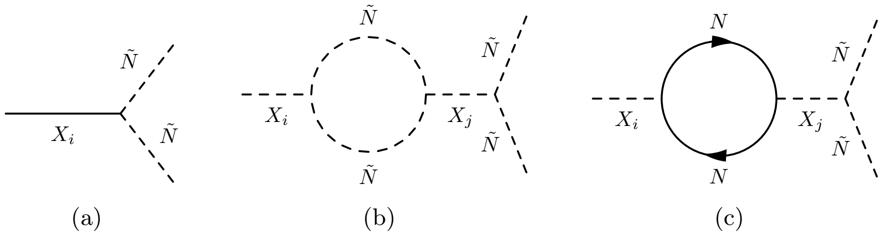

Figure 5.1: The relevant tree level and loop diagrams depicting the decay of the heavy mass states $X_2, ..., X_5$ to sneutrinos. Sneutrinos or their fermionic partners can run in the loop.

Monte Carlo analysis with constraints including the out-of-equilibrium condition, the correct hierarchy for the MSSM Higgs boson masses, correct magnitude for ε that produces (2.14) and an ad hoc limit to the degeneracy of the heavy boson masses reveals various aspects of the model. First, it shows that the decay channels $X_2 \to X_3^* \to \tilde{N}\tilde{N}$ and $X_3 \to X_2^* \to \tilde{N}\tilde{N}$ dominate with scan values shown in Table I of [3]. The mass difference $|M_{X_2} - M_{X_3}|$, Fig. 2(a) in [3] lies in the region $\sim 10^{-5} - 10^{-3}$ TeV whereas the other differences can be ~ 1 TeV (Fig. 2(b) in [3]), which can induce an enhancement in ε although the resonance region is not reached, Fig. 5 in [3]. Secondly, the removal of the constraint on the magnitude of ε highlights the difference between the extended model and the model in [197]. The addition of new couplings and the inclusion of the MSSM Higgs sector have introduced more freedom into the model. This is illustrated in the fact that the CP violation parameter can even become maximal $|\varepsilon| \lesssim 1/2$. The scenario in [197] includes only the first line of (5.17) and corresponding soft terms, and it is difficult to produce sufficient amounts of CP violation, Fig. 3 in [3].

The mass predictions for the right-handed neutrino, sneutrinos and MSSM Higgs bosons differ from those of [197]. As a consequence of the inclusion of the MSSM Higgs bosons, other possible terms for S and Δ fields and coupling the $B - L$ sector to the MSSM the singlet neutrino and superpartner sneutrino masses drop from the 1000 TeV scale to $\sim$1-10 TeV. Values favored by the lightest Higgs boson concentrate on the $\simeq 130$ GeV region while the two heavier Higgses have masses in the TeV region. So the masses of the particles N, $\tilde{N}$ and $\tilde{N}^*$ outside the MSSM are lowered in the extended scenario and the lightest Higgs mass tends to be in the lower end of the mass estimates.

5.3.3 Conclusions

A novel scenario for producing CP violation and lepton symmetry has been considered. In contrast to generating net lepton number via the asymmetric decays into MSSM (s)leptons and anti(s)leptons, the asymmetry arises between supersymmetric particles beyond the MSSM. By gauging $B - L$ and thus extending MSSM the authors in [197] have found a mechanism that places the abundances of sneutrinos and antisneutrinos into an imbalance that gives a nonzero lepton asymmetry without decays into MSSM particles. The initial asymmetry then influences the eventual decay of sneutrinos and antisneutrinos into MSSM (s)leptons and anti(s)leptons. One of the characteristics of the model in [197] is that producing sufficient CP violation is difficult. Once the scenario is augmented with the MSSM Higgs sector and additional superpotential terms involving the new chiral supermultiplets S, Δ and $\bar{\Delta}$ the system has more freedom with more parameters and larger amounts of CP violation can be obtained.

Besides the increased amount of CP violation, another thing that sets the model in [197] apart from [3] is the fact that singlet neutrino and superpartner sneutrino masses can be ∼1-10 TeV as opposed to ∼1000 TeV suggested by [197]. As these masses are reduced several orders of magnitude, observing signatures of these particles could become attainable in colliders during the coming years.

Increasing the freedom of the $B - L$ gauged MSSM with the MSSM Higgs sector and other allowed terms brings about flexibility that leads to significantly larger amounts of CP violation. With constraints imposed the higher end of the allowed interval of $|\varepsilon|$ is favored and once constraints are relaxed $|\varepsilon|$ can become $\sim 1/2$. Thus, it is essential to consider the complete model with all possible S, Δ, $\bar{\Delta}$ terms and the interaction $Y_3 S H_u H_d$ connecting the $B - L$ breaking sector and MSSM Higgs sector. The complete model where $B - L$ and MSSM sectors are included produces a conclusive picture of the novel source of leptogenesis.

Chapter 6

Summary and outlook

Two alternative leptogenesis scenarios have been considered. One incorporates flat extra dimensions with lepton asymmetry generation and the other has local $B - L$ symmetry built in an extension of the MSSM. Both models give rise to a new way of generating lepton asymmetry and consequent net baryon number. Extra dimensions, if observed, would revolutionize the way we understand space and time. Hints of extra dimensions are also likely to lead to the discovery of supersymmetry that lies at the heart of viable string theories. In this thesis, these frameworks serve as the background where leptogenesis is embedded. The other side of the story is provided by the LHC that is scheduled to confirm or disprove the existence of the Higgs boson of the SM by the end of 2012. If supersymmetry and extra dimensions are observed at the LHC [23, 198], the novel leptogenesis mechanisms could become viable explanations of the matter-antimatter asymmetry.

The considerations in this thesis have concentrated on finding new mechanisms of leptogenesis and verifying whether sufficient amounts of net baryon number can be derived from them. A more thorough treatment would be obtained if the Boltzmann equations for both systems were analyzed. This would give a more comprehensive picture of the decay and inverse decay processes that determine the washout factor. As the effectiveness of lepton number generation could become more apparent in a full treatment of the Boltzmann equations, also a better understanding of the parameter regions would result.

The main focus of this thesis has been in the origin of the 5 % portion of the energy density of the universe while the rest 95 % has remained untouched. However, the matter segments are not necessarily separate. The dark matter content of the universe is estimated to stand at 22-23 % [4] which gives $\rho_{\text{DM}} \simeq 4.5\rho_B$. Dark matter could then have a common origin with baryonic matter as their abundances are of the same order of magnitude. Moreover, the dark matter abundance could be the result of an asymmetry the same way net baryon number is. This type of dark matter is dubbed asymmetric dark matter [199]. To contrast this, one of the most common scenarios of dark matter is based on weakly interacting massive particles (WIMP) [23] whose density is determined by the freezeout of its annihilations to SM particles, and the source of the relic density is different from that giving the net baryon number of the universe. Unlike the WIMP approach asymmetric dark matter relies on the same baryogenesis mechanism that acts as the source of visible matter-antimatter asymmetry. It also predicts lighter dark matter particles with masses in the GeV region.

If baryogenesis is at the heart of both baryonic and dark matter, a natural question would be to ask whether leptogenesis is responsible for the existence of baryonic and dark matter sectors. This has been explored *e.g.* in [200] where the right-handed singlet neutrino not only couples to the SM Higgs and lepton doublets but also to the dark matter fermion and scalar field that carry lepton number. Thus, the same mechanism that generates lepton asymmetry can give rise to an asymmetry in the dark matter sector as well.

It is safe to say that numerous interesting descriptions of leptogenesis exist in the most studied extensions of the SM. Supersymmetry and extra dimensions both solve the hierarchy problem that ails the SM and give rise to new effects. A natural way of extending these scenarios further is to incorporate leptogenesis that can explain nonzero neutrino masses and act as an intermediary towards baryogenesis. Another prospect of leptogenesis briefly mentioned here is transcending the border between baryonic and dark matter and showing that these two could in fact have a common origin. Future observations put these scenarios to the test and show whether we have been on the right track on finding ways to describe nature.

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