

Aspects of Dark Sector Assisted Leptogenesis

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by

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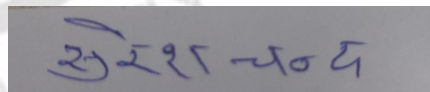


Declaration



This is to certify that the work contained in this thesis, entitled “**Aspects of Dark Sector Assisted Leptogenesis**” submitted by me to the Indian Institute of Technology, Guwahati, for the award of the degree of PhD, is a bona fide work carried out by me under the supervision of Dr. Debasish Borah and Dr. Pankaj K. Mishra and that this work has not been submitted elsewhere for a degree of diploma. I also wish to state that to the best of my knowledge and understanding nothing in this report amounts to plagiarism.

June 17, 2026

A rectangular box containing a handwritten signature in blue ink. The signature appears to be 'सुरेश चन्द' (Suresh Chand) in Devanagari script.

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Certificate



It is certified that the work contained in this thesis entitled “**Aspects of Dark Sector Assisted Leptogenesis**” by **Suresh Chand**, has been carried out under our supervision, and this work has not been submitted elsewhere for a PhD. degree.

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Abstract



The nature of dark matter (DM) and the origin of the baryon asymmetry of the Universe (BAU) remain two of the most compelling questions in contemporary particle physics and cosmology. Since the proposal of DM in 1933 by Fritz Zwicky, through its discovery in galaxy clusters and further validation in the cosmic microwave background (CMB) and large-scale structure observations, the hunt for a viable DM candidate has intensified. Simultaneously, the observed imbalance between matter and antimatter, manifested in the baryon asymmetry, has puzzled physicists for decades. Theoretical models, including those based on the Standard Model (SM) and its extensions, strive to address both phenomena, yet they remain unresolved.

This thesis addresses these two challenges, proposing a novel framework that connects low-scale leptogenesis to the dark matter sector. Specifically, it examines a minimal extension of the SM with three right-handed neutrinos, a charged scalar, and a neutral fermion, proposing that the decay of right-handed neutrinos, in combination with dark sector processes, can generate both the baryon asymmetry and the observed relic density of DM. The model highlights the role of dark matter interactions in generating the matter, antimatter asymmetry, exploring CP-violating effects from both visible and dark sectors.

Our work connects two previously disparate areas: leptogenesis and dark matter, revealing how they can be unified within a single theoretical framework. We demonstrate that through the feeble interactions of dark sector particles with right-handed neutrinos, low-scale leptogenesis can be realized even for light right-handed neutrino masses (< 2 TeV), without relying on the resonant mechanisms typically invoked at high scales. Moreover, this model also accommodates asymmetric dark matter, where the dark matter particles' interactions with the visible sector contribute to the final lepton asymmetry. This connec-

tion between the dark sector and leptogenesis provides a promising path to understanding both the observed BAU and DM relic density.

This thesis is organized as follows: In **Chapter 1**, we introduce the fundamental concepts and observational evidence for both DM and BAU. In **Chapter 2**, we explore the limitations of the Standard Model in explaining these phenomena, leading to the discussion of various beyond-the-SM (BSM) scenarios. **Chapter 3** outlines the main framework of our work, where we describe the proposed minimal extension of the SM and its potential to generate the baryon asymmetry and dark matter relic density. Finally, in **Chapter 4**, we examine a more detailed study of leptogenesis with CP phases induced by dark matter and their experimental implications, including predictions for collider experiments.

In **Chapter 5**, we present the summary and conclusion of the thesis. Finally, this chapter concludes with a brief outline of the future scope of the work presented in the thesis.

Through this work, we aim to offer a unified solution to two of the most profound problems in physics, potentially bridging the gap between early-Universe cosmology and high-energy physics in a way that can be tested in upcoming experiments, especially at the Large Hadron Collider (LHC).





Acknowledgements

*“ I am time, shaped into eight pahars.
I have come to see it as something vast enough
to hold an entire lifetime within itself.
Yet here, my mind gathers only a small passage
of that long journey—
the PhD years—
and names them these eight pahars.
Each pahar stands equal in weight and worth.
The first began with a turning point,
a new unfolding of life— my entry into the PhD.
Every moment along this path,
every measured pause of time,
was infinite in meaning,
and essential to the journey. “*

Life is full of challenges, and every journey is marked by countless tests. In the academic world, the PhD defense may be the ultimate one. Now, as I stand on the threshold of this final examination, I look back on the past eight pahars as a profound learning experience—one filled with both sweet and bitter memories. I would like to express my heartfelt gratitude to everyone I encountered along this journey, whose support, guidance, and encouragement made this long path not only possible but truly memorable.

Time has been my greatest supervisor, teaching me invaluable lessons, guiding me forward, and giving me the courage to navigate a challenging path and reach my destination



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 ... *To Empire of Nature* ... 

 ... *To those children who want to
protect the forests and restore the wooden
kingdom that fell victim to the axe, so that
the world may once again be filled with its
own fragrance.* ... 





Publications

1. Leptogenesis from a feebly interacting dark matter sector, **S. Chand**, Mariana Frank, and Poullose Poullose, [Phys. Rev. D **106**, 043030 \(2022\)](#).
2. Dark Matter-Driven Low-Scale Leptogenesis via Neutrino Portal, **S. Chand**, Avnish, and Poullose Poullose, [Phys. Rev. D **113**,116006 \(2026\)](#).



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Introduction

Two of the central open problems in particle physics are understanding the origin of the baryon asymmetry of the Universe (BAU) [1, 2] and identifying a viable dark matter (DM) candidate [3]. Most of the observations indicate that the present Universe is overwhelmingly matter dominated; antimatter is scarcely observed outside the particle accelerators [4]. The long-term survival of planetary probes implies that the solar system contains essentially no antimatter, while measurements of cosmic rays provide further evidence for matter dominance on galactic and extragalactic scales [5, 6].

Cosmic rays include roughly 10^{-4} as many antiprotons as protons [7]. The extremely low quantity of antiprotons in cosmic rays appears to lend credence to the theory that these particles are a consequence of cosmic ray collisions with the interstellar medium rather than direct proof of antimatter in the galaxy [8]. Anti- ^4He is around 10^{-5} as much as ^4He [9]. Cosmic rays are therefore an excellent means of demonstrating the distinction between matter and antimatter at the galactic level. The proof of asymmetry on greater scales is not as robust. X-rays from galaxy clusters indicate the presence of gas in the region between the galaxies [10]. Had matter and antimatter galaxies coexisted within the same cluster, the annihilation of nucleons and antinucleons would have produced a large amount of γ -ray emission [5, 6]. The absence of such γ -ray flux implies that nearby galaxy clusters such as Virgo are composed exclusively of baryons or antibaryons. The disparity between matter and antimatter is known as the matter–antimatter asymmetry, or BAU [11].

In Big Bang Nucleosynthesis, the baryon-to-photon ratio (which is the measure of baryon asymmetry of the universe) is taken as an input parameter. This asymmetry is the reason why BBN makes the right predictions for the primordial abundances of the light elements [12]. This picture is independently validated by measurements of the CMB,

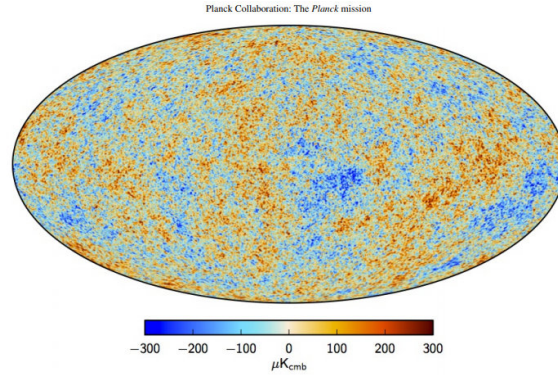


Figure 1.1: CMB temperature anisotropies observed across the sky. The color variations correspond to $\Delta T/T$, revealing the primordial density fluctuations that seeded large-scale structure. This figure is adapted from [14].

which we will discuss in a more elaborate way in the following subsections [Sec. 1.1.1 and Sec. 1.1.2] [13].

1.1 Evidence of Matter-Antimatter Asymmetry

The existence of matter-antimatter imbalance is well established by both experimental and observational evidence [13]. In particle Physics, CP violation constitutes a necessary ingredient for generating such an asymmetry, as formalized by the Sakharov conditions [11], and has been seen in the neutral kaon and B-meson systems [15, 16]. Moreover, neutrino oscillations and possible leptonic CP violation provide important clues toward understanding the matter antimatter asymmetry of the Universe [17, 18]. Evidence from the CMB, the dominance of baryonic matter throughout structure formation, and the extremely low abundance of antimatter in cosmic rays all point to the Universe's breaking of matter-antimatter symmetry [13, 19]. Taken together these considerations have motivated an early examination of the early Universe, where the observed quantities of light elements produced during the earliest moments following the Big Bang provide a sensitive window into the primordial asymmetries, connecting fundamental particle Physics with cosmological events [1, 4].

1.1.1 Matter–Antimatter Asymmetry from Big Bang Nucleosynthesis

Big Bang Nucleosynthesis (BBN) explains the formation of light elements during the first few minutes after the Big Bang [20, 21]. It predicts the primordial abundances of light

nuclei like deuterium (^2H), Helium-3 (^3He), Helium-4 (^4He), and Lithium-7 (^7Li) [12]. BBN proceeds through a sequence of nuclear reactions initiated by deuterium production via $n + p \leftrightarrow D + \gamma$, with the dominant reactions shown in the left panel of Fig. 1.2. The predicted abundances depend primarily on a single parameter, the baryon-to-photon ratio, which controls the nuclear reaction rates through the number densities of the participating nuclei [2].

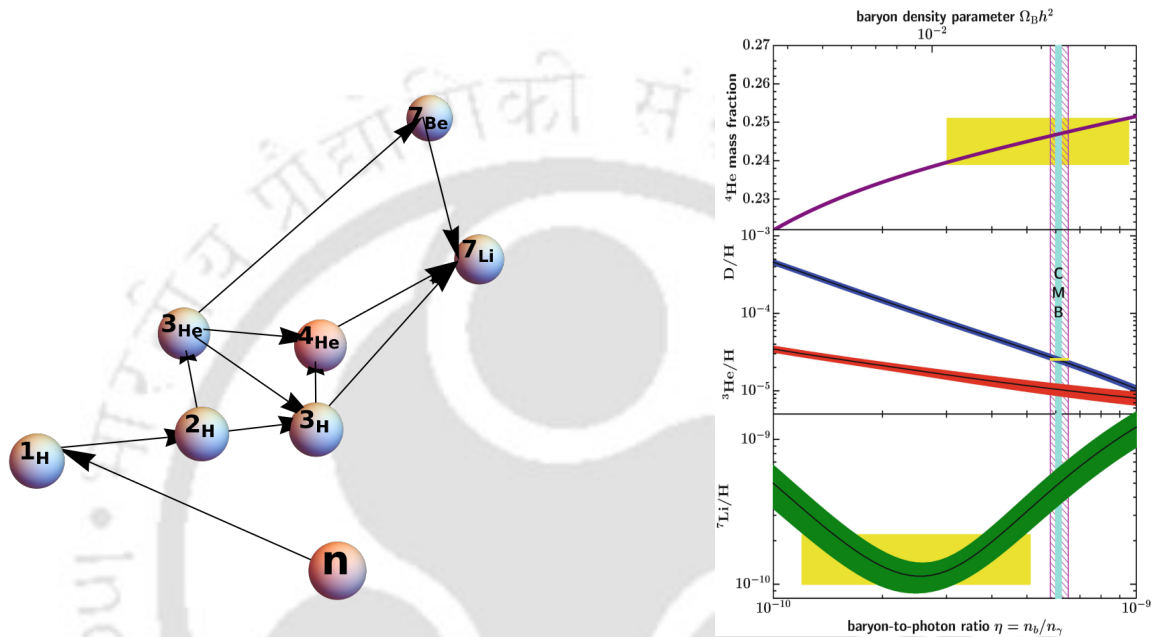


Figure 1.2: Left panel: Schematic illustration of the main nuclear reactions occurring during the Big Bang Nucleosynthesis (BBN). Right panel: Contour plots showing the predicted the observed quantities of light elements as a function of the baryon-to-photon ratio, η_B . The yellow boxes correspond to the measured astrophysical abundances. The narrow vertical band represents the baryon density inferred from CMB observations at 95% confidence level, while the wider band indicates the BBN concordance region at 95% confidence level. This figure is adapted from [22].

The right panel of Fig. 1.2 displays the predictions of BBN calculations for the primordial abundances of the light elements as a function of the baryon-to-photon ratio, η_B [23]. The corresponding observational constraints are indicated by the yellow bands. The primordial deuterium abundance is a decreasing function of it, η_B and thus deuterium is a particularly sensitive probe to the baryon density. The abundance of ^3He is a more complicated function of η_B because ^3He is both produced and destroyed in a number of competing nuclear reactions. The behaviors are derived from the chain of nuclear processes that determine the synthesis of light elements, such as the reactions $p + n \rightarrow d + \gamma$ and $p + d \rightarrow ^3\text{He} + \gamma$ [21]. These nuclear reactions, together with those responsible for the

production of ${}^4\text{He}$, proceed until the reaction rate Γ drops below the Hubble expansion rate H , marking the freeze-out nucleosynthesis [2].

Lithium production occurs through two distinct reaction pathways, each with a different dependence on η_B , resulting in a non-monotonic behavior of the primordial ${}^7\text{Li}$ abundance as a function of η_B [12]. The yellow shaded regions in Fig. 1.2 denote the observed primordial abundances of the light elements. The remarkable agreement between the predicted and observed abundances of ${}^4\text{He}$ and deuterium constrains η_B to a narrow range, providing strong support for the hot Big Bang model. At the 95% confidence level, this range is found to be $5.8 \times 10^{-10} \leq \eta_B \leq 6.5 \times 10^{-10}$ [13, 23]. Next we discuss the Matter-antimatter asymmetry inferred from cosmological observations, particularly those from the CMB.

1.1.2 Matter–Antimatter Asymmetry from Cosmic Microwave Background

The temperature of the CMB varies very little, by only about one part in 10^5 [13, 24]. We use the CMB power spectrum to quantify these tiny variations [25]. This power spectrum provides insights into fundamental questions, such as matter–antimatter asymmetry in the Universe, through precise determination of cosmological parameters including the baryon density [2, 13]. The CMB sky can be decomposed into spherical harmonics, and the angular power spectrum, as shown in Fig. 1.3, indicates the amount of power present at different angular scales across the sky [26]. The x-axis in Fig. 1.3 is represented by ℓ , the multipole moment. Small values of ℓ correspond to large angular scales, while large values of ℓ correspond to small angular scales [25]. The y-axis represents the power, measured in microkelvin. The power spectrum exhibits a distinct pattern of acoustic peaks, which arise from physical processes in the early Universe, specifically acoustic oscillations of the tightly coupled photon–baryon fluid prior to recombination [13, 26, 28].

Before recombination, photons were tightly coupled to the hot plasma of baryons and electrons through Thomson scattering [21, 29]. As the Universe cooled, photons decoupled from the plasma, making the Universe transparent [26]. The single most pronounced feature of the CMB temperature power spectrum is the first acoustic peak at $\ell \approx 220$, corresponding to angular scales of order 1° . Its position contains much information on the geometry of the Universe in space [13, 28]. The second peak, at $\ell \approx 540$, depends sensitively on the baryon density; a higher baryon density enhances compressional modes and suppresses rarefaction, resulting in a relatively smaller second peak [25, 26]. Together, the first and second peaks provide precise information about the amount of ordinary

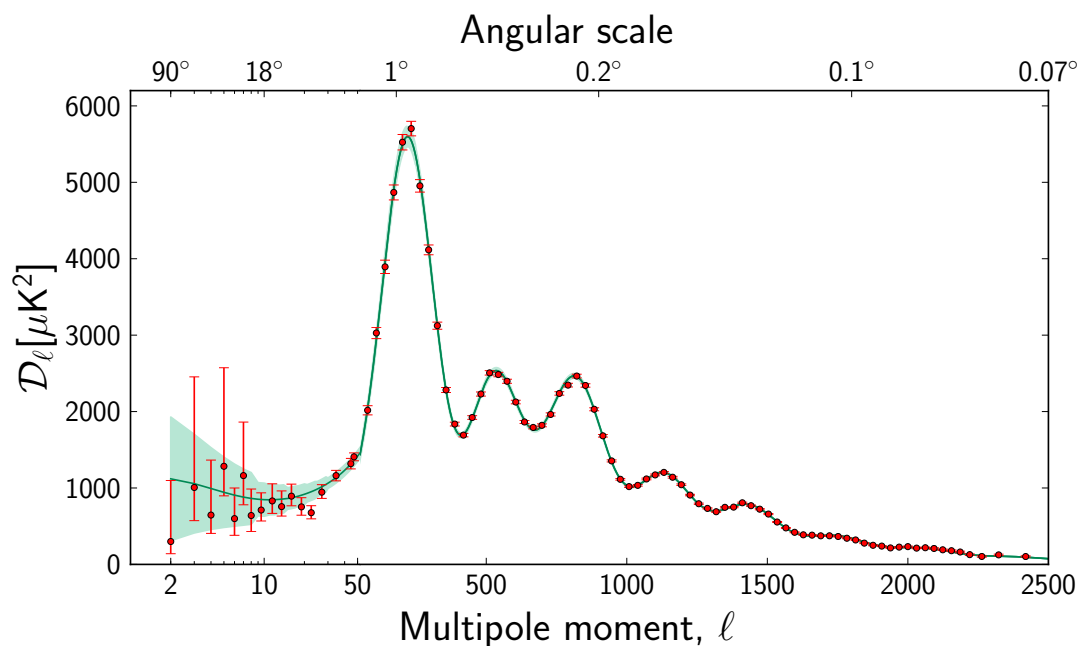


Figure 1.3: CMB temperature angular power spectrum from Planck, showing the measured anisotropies and the Λ CDM best-fit model. The positions and heights of the acoustic peaks are sensitive to the densities of baryonic matter, dark matter, and dark energy. Models without dark matter or with incorrect matter content fail to reproduce the observed peak structure. This figure is taken from [27].

(baryonic) matter in the Universe [13].

The third acoustic peak contains important information on the matter content of the Universe, and the ratio of the height of this peak to the height of the second peak is useful for constraining the dark matter density. In particular, the differences in the heights of the peaks are due to the interplay of the baryonic matter and the gravitational potentials dominated by dark matter. Later peaks are roughly harmonic series of acoustic oscillations and decay in amplitude with time. At large multipoles ℓ the power spectrum is strongly suppressed by photon diffusion (Silk damping) and the finite thickness of the last scattering surface with a prominent damping tail at small angular scales [13, 25, 26, 28, 30].

The small value of the baryon-to-photon ratio η_B indicates that photons greatly outnumber baryons in the Universe [25]. The numerical value of η_B is consistent with that independently obtained from Big Bang nucleosynthesis [12, 23]. This is a unique fingerprint of the early Universe, encoding precise information about cosmological parameters like the baryon density. From this the baryon to photon ratio η is inferred, suggesting the existence of a baryon asymmetry in the Universe [26]. Current observations, such as those from the Planck satellite, are remarkably well described by the Λ CDM model and

strongly support the existence of a matter–antimatter asymmetry in the Universe [2, 13].

The CMB power spectrum constrains the baryon-to-photon ratio η_B through the baryon density parameter $\Omega_b h^2$ [13, 26]. In general, the matter–antimatter asymmetry of the Universe is quantified by a parameter η , defined as the ratio of the difference between the baryon number density (n_B) and the antibaryon number density ($n_{\bar{B}}$) to the photon number density (n_γ) in the Universe, i.e. [2, 11, 21],

$$\eta = \frac{|n_B - n_{\bar{B}}|}{n_\gamma}. \quad (1.1)$$

The matter–antimatter asymmetry is frequently characterized by two quantitative measures [2, 12, 13, 23]:

$$\eta_B = \frac{|n_B - n_{\bar{B}}|}{n_\gamma} \Big|_0 = (6.21 \pm 0.16) \times 10^{-10}; \quad Y_B = \frac{|n_B - n_{\bar{B}}|}{s} \Big|_0 = (8.75 \pm 0.23) \times 10^{-11}. \quad (1.2)$$

Here n_B and $n_{\bar{B}}$ are the number densities of baryons and antibaryons, while n_γ and s represent the photon number density and entropy density of the Universe, respectively [21, 31].

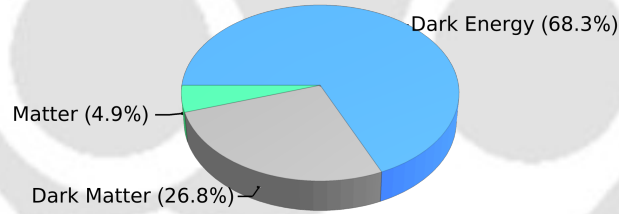


Figure 1.4: Relative contributions of dark energy, dark matter, and ordinary (baryonic) matter to the total energy density of the universe, as inferred from cosmological observations. Dark energy dominates the cosmic energy budget, while dark matter constitutes a significant fraction compared to ordinary matter, highlighting the need for non-baryonic dark matter in the Universe.

1.2 Evidence of Dark Matter

On the other hand the study of DM dates back to 1922 when Jacobus C. Kapteyn first proposed a statistical model for the Milky Way. By counting stars and analyzing their motions, he estimated the galaxy’s mass density using the vertical movements of stars.

Although Kapteyn didn't detect a significant mass discrepancy, he suggested that, with improved theory, it might be possible to identify the amount of dark matter through its gravitational effects [32]. In the early 1930s, Kapteyn's graduate student Jan Oort advanced this research by employing better data and applying the Poisson equation. In 1932, Oort concluded that the local mass density was greater than that of visible matter, providing one of the first hints of dark matter's existence [33]. The concept gained further traction in 1933 when Fritz Zwicky applied the virial theorem to the Coma Cluster. His calculations, which showed that galaxy velocity dispersions of approximately 1000 km s^{-1} required a total mass about 400 times larger than the visible stellar mass, indicated the presence of unseen mass necessary to keep the cluster gravitationally bound. Zwicky coined the term *dunkle Materie* (dark matter) to describe this mysterious component [34]. In 1936, Sinclair Smith independently confirmed Zwicky's results by studying the Virgo Cluster, which showed a similar discrepancy in mass between visible matter and the total mass required to explain the gravitational binding of the cluster [35].

All these seminal observations further provided a solid evidence for the existence of dark matter. Radio observations of extended hydrogen line (HI) rotation curves revealed flat rotation profiles, which required the presence of massive dark halos to explain the observed dynamics [36–40]. Gravitational lensing and X-ray emission from hot intracluster gas independently mapped dark matter at the cluster scale [41–43, 10]. Additionally, the Cold Dark Matter (CDM) paradigm, supported by observations of CMB anisotropies [24], the acceleration of the Universe [44, 45], Baryon acoustic oscillations (BAO) [46], and the Bullet Cluster [47], has solidified dark matter as a fundamental component of cosmology. More recent precision cosmology datasets [48–51] have further confirmed this, while null results from direct detection experiments [52] and emerging tensions in the data [53] continue to refine our understanding of dark matter's properties.

The CMB acoustic peaks provided definitive evidence of DM's non-baryonic nature, which cannot be accommodated within the SM of particle physics. DM constitutes $\sim 27\%$ the Universe's total energy density, compared to $\sim 5\%$ ordinary baryonic matter. The remainder $\sim 68\%$ is attributed to dark energy, driving accelerated cosmic expansion. The present dark matter abundance, or relic density, has remained constant and is measured to be $\Omega h^2 = 0.120 \pm 0.001$ [54]. The evidence of DM is customary to focus on three main probes: the observation of individual spiral galaxies, clusters of galaxies, the CMB, and large-scale structures at cosmic scales.

1.2.1 Physical characteristics and evolution of individual Galaxies

The observation of how stars or gas rotate at various distances from the galactic center, and the resulting galactic rotation curves, provides crucial information about the distribution of matter within galaxies [38, 55, 56]. Deviations of observed rotation curves from those expected based on visible matter alone provide strong evidence for the presence of dark matter halos surrounding galaxies [57]. In this section we present a brief overview of these rotation patterns and their implications for the distribution of both the luminous and dark matter in the individual galaxies.

1.2.1.1 Rotation curves of spiral galaxies

The axis of rotation of a spiral galaxy is approximately perpendicular to the galactic disk. By measuring the Doppler shift of atomic and molecular spectral lines, one can determine the circular velocity of stars and other tracers—such as neutral hydrogen clouds and masers—as a function of their distance from the galactic center, yielding a galactic rotation curve as shown in Fig. 1.5 [38, 56, 57]. Assuming spherical symmetry and Newtonian gravity, Newton's second law, $F = ma$, predicts a simple relationship between the circular motion of a test particle of mass m (such as a star) and the total mass M enclosed within a radius r from the galactic center [56]:

$$\frac{v_{\text{circ}}^2(r)}{r} = \frac{GM(r)}{r^2} \Rightarrow v_{\text{circ}}(r) = \sqrt{\frac{GM(r)}{r}}. \quad (1.3)$$

It is therefore expected that the circular velocity follows a Keplerian decline, $v_{\text{circ}}(r) \propto r^{-1/2}$, when the mass distribution is dominated by a centrally concentrated luminous component [56]. A constant circular velocity, $v_{\text{circ}}(r)$, can instead be obtained by assuming a mass density profile $\rho(r) \propto r^{-2}$, extending to large radii, as expected for an isothermal dark matter halo [58]. In this case, the enclosed mass grows linearly with radius, $M(r) = 4\pi \int_0^r r'^2 \rho(r') dr' \propto r$, which naturally explains the observed flat galactic rotation curves [39, 57]. The rotation curve of the Andromeda galaxy (M31) was precisely measured in 1970 by Rubin and Ford [38]. Figure 1.5 shows that as one moves farther from the center of the galaxy, the rotational velocity remains approximately constant rather than exhibiting a Keplerian decline. This behavior implies that, for stars to remain gravitationally bound at large galactocentric radii, the enclosed mass must be significantly greater than the observed luminous mass. This discrepancy between dynamical and visible mass provides direct and compelling evidence for the presence of dark matter in galaxies [39, 56].

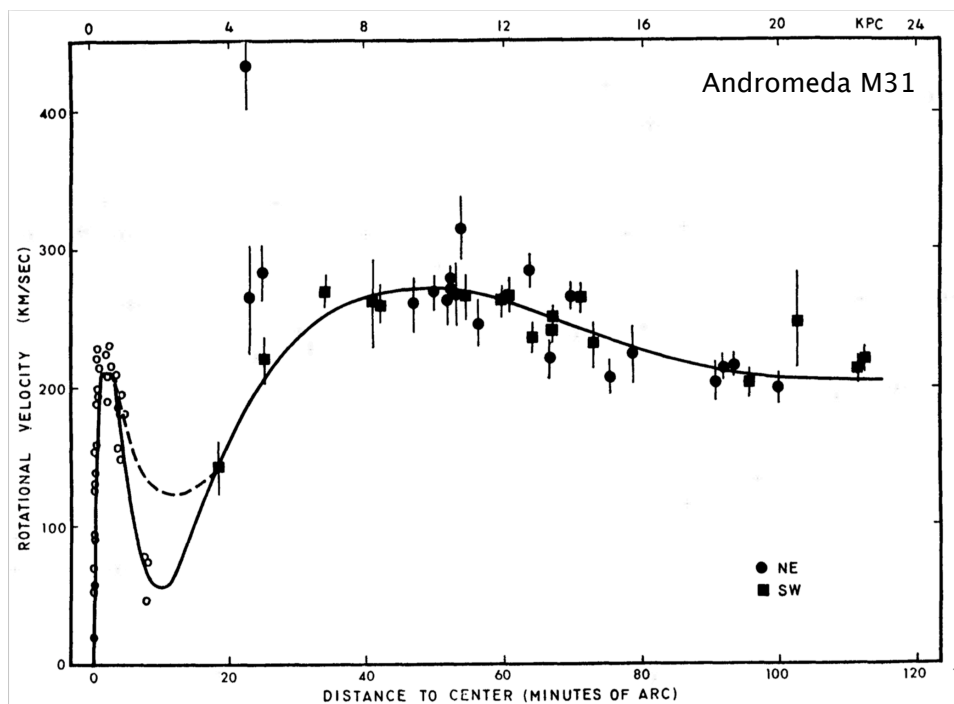


Figure 1.5: Observed rotation curve of the Andromeda galaxy (M31) as measured by Rubin and Ford (1970) [38]. The nearly flat rotation velocities at large radii indicate the presence of a substantial dark matter.

1.2.1.2 Additional evidence on a galactic scale

Numerous other pieces of evidence based on observations at galactic scales have been proposed to support the presence of dark matter (DM). First, other systems, especially dwarf galaxies, can benefit from the same kinematic tracer measurements. Dwarf spheroidal galaxies are particularly important in this context, as they are highly dark-matter dominated and their stellar velocity dispersions provide strong constraints on their mass distributions [59, 60]. Furthermore, dwarf galaxies themselves can be used as tracers to determine the mass distribution and rotation curves of larger spiral galaxies through the study of satellite kinematics [61]. This method supports the outcomes obtained from rotation curves of stars and gas and provides independent evidence for extended dark matter halos around galaxies [56].

Early numerical simulations and subsequent analytical interpretations in the 1970s demonstrated that spiral galaxies require a massive, approximately spherical dark matter halo to maintain disk stability [62]. In the absence of such a halo, simulated galactic disks were found to break apart within a few rotation periods and evolve into chaotic stellar distributions. In particular, instability occurs when the ratio of rotational kinetic

energy to gravitational potential energy exceeds approximately 14% [62]. A spherical dark matter halo increases the total gravitational binding energy of the system due to its geometry, thereby stabilizing the disk. More advanced studies, however, later showed that the presence of a dark matter halo can also have adverse effects on disk dynamics. In particular, dynamical friction between the rotating stellar bar and the dark matter halo can lead to the formation of a bar and subsequently slow down its pattern speed [63]. This behavior is in tension with observations indicating that many galactic bars rotate rapidly. As a result, the role of dark matter in disk stability and bar formation is complex and cannot be reduced to a single, simple conclusion [64].

It is also possible to demonstrate the presence of massive dark matter (DM) halos around galaxies through galaxy–galaxy gravitational lensing [65, 66]. In the same way that the Sun gravitationally deflects the light of background stars, the gravitational lensing effect of foreground galaxies causes distortions in the observed images of background galaxies [67, 68]. The deflection angle of a light ray passing a lens of mass M_{lens} at impact parameter b is given by $\delta\theta = \frac{4GM_{\text{lens}}}{b}$, where M_{lens} denotes the mass of the lensing galaxy [21]. According to geometric optics, a lens of mass M_{lens} at distance d_{lens} gravitationally focuses a background source at distance d_{source} into an arc or an Einstein ring with angular size $\theta_E = \sqrt{4GM_{\text{lens}} \left(\frac{1}{d_{\text{lens}}} - \frac{1}{d_{\text{source}}} \right)}$, when the lens and source are closely aligned along the line of sight; otherwise, multiple images may be formed [68, 66]. Although the deflection angles are typically small, $\delta\theta \ll 1$, this phenomenon is referred to as strong lensing only when the lens is sufficiently massive and the source lies close to the optical axis, producing arcs or multiple images [69]. More commonly observed is weak gravitational lensing, which results in subtle, coherent distortions of background galaxy shapes that do not produce visible arcs [70]. By statistically measuring these distortions, it is possible to infer the average mass of typical lensing galaxies. These measurements consistently indicate that the total mass inferred from lensing significantly exceeds the observed luminous mass [71, 70]. Furthermore, the same data allow the reconstruction of the typical density profiles of dark matter halos, which are found to be extended and significantly flattened compared to the distribution of visible matter [66].

1.2.1.3 An Observational Study of a Galaxy with Unusually Low Dark-Matter Content

Van Dokkum et al. [72] reported in 2018 that the ultra-diffuse galaxy NGC1052–DF2, located near the massive elliptical galaxy NGC1052, appears to contain little or no dark matter. Using the velocities of a small population of globular clusters as dynamical

tracers, its total mass was estimated to be less than twice its stellar mass, inferred to be $\sim 2 \times 10^8 M_\odot$. Given that galaxies of this type typically exhibit mass-to-light ratios of several hundred, this result was highly unexpected. After a careful reassessment of uncertainties, the same team identified a second galaxy with similar properties, NGC1052–DF4, in 2019 [73].

If confirmed, these findings raise important questions regarding galaxy formation and the dark matter paradigm. Such systems may represent rare statistical outliers or the outcome of unusual evolutionary processes in which dark matter was stripped through strong interactions with a massive host galaxy. While some numerical simulations suggest that dark-matter-deficient galaxies can form through tidal interactions and stripping mechanisms [74], other studies have found it difficult to reproduce such objects within standard Λ CDM scenarios [75]. Moreover, these galaxies also pose a challenge to modified-gravity theories, since such frameworks generally predict observable deviations from Newtonian gravity in the presence of baryonic matter at sufficiently large radii [76].

1.2.2 Mesoscale Structures: Galaxy Clusters

Fritz Zwicky is widely credited with first clearly articulating the evidence for dark matter (“Dunkle Materie”) in 1933 through his study of the Coma cluster [34]. As discussed by Bertone and Hooper [77], however, the underlying concept and terminology were already present in the astronomical community, with earlier contributions from researchers such as Smith, Poincaré, Thomson (Lord Kelvin), Öpik, Kapteyn, Lundmark, and Oort. Nonetheless, Zwicky’s work remains particularly significant due to its historical impact and the remarkably direct and transparent nature of its reasoning when viewed in retrospect.

By analyzing the velocity dispersion of galaxies in the Coma cluster, Zwicky inferred the presence of additional, unseen matter required to gravitationally bind the system [34]. Galaxy clusters—the largest gravitationally bound structures in the Universe, spanning several megaparsecs and containing hundreds to thousands of galaxies—serve as valuable probes of the large-scale matter distribution [78]. Although they no longer provide the most precise measurements of the cosmic dark matter density, studies of clusters still imply a contribution of approximately $\Omega_{\text{DM}} \simeq 0.2$ [77].

Zwicky’s conclusion was based on the virial theorem, which for gravitational systems relates the average kinetic and potential energies through $\langle K \rangle = -\frac{1}{2}\langle V \rangle$. This relation allows the total mass of a system with many constituents to be inferred from its characteristic size and velocity, providing a simple yet powerful method to estimate the mass

required to gravitationally bind large astrophysical systems [34, 79, 77].

$$N \frac{mv^2}{2} = \frac{1}{2} \frac{N^2 G m^2}{R} \Rightarrow mN = \frac{2Rv^2}{G}. \quad (1.4)$$

Applying these considerations to the Coma cluster, Zwicky concluded that its total mass far exceeded the visible matter, implying the presence of additional dark matter [34, 79]. Although this interpretation was not immediately embraced, it was taken seriously, with the general view that more observational data were needed to fully understand such systems [77].

Since the 1980s, X-ray observations have become a powerful method for measuring both baryonic and dark matter in galaxy clusters (see [10, 78] for reviews). Clusters contain large amounts of ionized hydrogen and helium that are heated to $T_{\text{gas}} \sim 10^8$ K (~ 10 keV) through shocks and adiabatic compression as the gas settles into the cluster potential well. Despite these high temperatures, nuclear fusion is negligible due to the low densities. The hot gas emits X-rays primarily via thermal bremsstrahlung, and under the assumptions of spherical symmetry and hydrostatic equilibrium, its pressure gradient can be related to the gradient of the gravitational potential.

$$\frac{1}{\rho_{\text{gas}}(r)} \frac{d\phi_{\text{gas}}}{dr} = -\frac{d\phi}{dr} = -\frac{GM(r)}{r^2}, \quad (1.5)$$

where $M(r)$ is the total gravitational mass (i.e., dark matter and hydrogen/helium gas) inside radius r and $\rho_{\text{gas}}(r)$ is the gas density. Assuming an ideal gas, the relationship between pressure and density is

$$p_{\text{gas}} = \frac{\rho_{\text{gas}} k T_{\text{gas}}}{\mu m_p}, \quad (1.6)$$

where $\mu \approx 0.6$ is the mean molecular weight of an admixture of $\sim 25\%$ helium and $\sim 75\%$ hydrogen [10, 78]. The intensity and spectrum of the X-ray emission can be used to determine the density ρ_{gas} and the temperature T_{gas} . This allows the reconstruction of the total cluster mass and its profile using eq. (1.6) [80]. The findings verify that the gas only makes up a fraction of the galaxy cluster's overall mass. In a sense, using X-ray emissions is a microscopic version of Zwicky's method, where the total gravitating mass that maintains the galaxy cluster's gravitational boundness is inferred from the kinetic energy of the gas molecules (i.e., their temperature) rather than the galaxies [34, 79].

Another way X-ray data are highly helpful is in the analysis of special configurations, including galaxy cluster collisions, which we will discuss next [81].

1.2.2.1 Dark Matter via Weak Gravitational Lensing: Cosmic Shear and the Bullet Cluster

The observations of a pair of colliding clusters known as the bullet cluster, located 3.7 Gyr away and cataloged as 1E0657-558 (or 1E0657-56), which was first observed in detail in 2006, as well as from similar systems, provide some of the most striking evidence for the existence of dark matter on the length scales of galaxy clusters today [82–84]. The bullet cluster’s X-ray emissions reveal the dispersion of the heated gas that makes up the majority of its baryonic mass. Weak lensing was used to measure the distribution of the entire mass, including both visible and dark matter [82, 85].

The visible matter and dark matter are spatially separated, which is a unique property of the bullet cluster system (see fig. 1.6). The following is the interpretation: Each of the two galaxy clusters was once a typical system, with DM and visible matter coexisting. 150 million years ago, the two objects collided. The heated gas from the two clusters suffered a collisional shock wave because visible matter interacts with itself in a substantial way. The DM clouds of the two systems, however, merely passed through one another because DM had very few collisions with regular matter and with itself. As shown in fig. 1.6, this resulted in the current separation of the visible and dark matter components. Numerous comparable systems have been examined since the bullet cluster was discovered. The results on 72 of them are reported by Harvey et al. (2015) [86], who conclude that the existence of DM may be demonstrated with a significance of greater than 7σ .

Such results severely constrain alternate explanations in which a change of gravity takes the place of DM. The anomalous lensing signal would follow the distribution of visible matter because such alterations cannot be spatially isolated from normal matter (unless they also introduce something that essentially behaves as DM).

The particle physics characteristics of DM are also revealed by the Bullet Cluster and related systems. An upper bound $\wp \sim \sigma n r_{\text{cl}} \lesssim 1$ on the probability $\wp$ that a DM particle scatters with another DM particle on the typical scale of the system is implied by the fact that the DM halos did not slow down in comparison to the collisionless trajectories. In this case, r_{cl} is the cluster size, n is the DM number density, and σ is the DM self-interaction cross section. Data show $r_{\text{cl}} \sim 150$ kpc and $\rho \sim 0.5$ GeV/cm³ in the Bullet Cluster instance. One derives a bound on σ/M , where M is the DM mass, as observations give information on the DM density $\rho = Mn$ (rather than directly on the DM number density n). Various investigations yield somewhat varying numbers [84, 83, 86], roughly within a factor of 2.

$$\frac{\sigma}{M} \lesssim 1 \frac{\text{cm}^2}{\text{g}} = \frac{1.8 \text{ mb}}{\text{GeV}} = \frac{4580}{\text{GeV}^3}, \quad (1.7)$$



Figure 1.6: Galaxies contain stars, gas, dust, and dark matter bound by gravity. The Bullet Cluster represents the collision of two enormous galaxy clusters, each held together by its own gravitational potential. Figure taken from [Bullet Cluster \(Webb and Chandra Image\)](#).

Under the assumption of point interactions. For example, a typical QCD cross section for pp scattering is 50 mb. Therefore, if DM is composed of particles that are either somewhat heavier than protons or somewhat less strongly interacting than protons, the bound in eq. (1.7) is satisfied. Astrophysics and cosmology also constrain the DM/matter cross section to be roughly less than QCD size at $M \sim 1$ GeV, while there are several other, much stricter constraints [87–89].

Weak lensing shear due to large-scale structure Additionally, the observation of cosmic shear is evidence for dark matter (DM) at length scales that are somewhat intermediate between galactic/cluster and truly cosmic scales [90, 91]. When light from very distant galaxies is deflected by the gravitational pull of foreground mass concentrations, this phenomenon is known as cosmic shear. These structures take the form of considerably broader and more diffuse formations, such as massive filaments and loose clumps, rather than the compact DM halos of individual galaxies or clusters (as in galaxy–galaxy lensing or cluster collisions).

Multi-wavelength surveys are used to statistically measure the shapes of millions of distant galaxies; such surveys must be both very deep (to detect high-redshift sources) and very wide (to cover large fractions of the sky) [92, 93]. Cosmic shear distorts the major-to-minor axis ratios of background galaxies by approximately 2%, which is detectable only through statistical analysis [90]. Observations consistently indicate a dark matter density parameter of $\Omega_{\text{DM}} \approx 0.25$ [13].

Cosmic shear provides far more than just confirmation of the existence of DM. By reconstructing the matter distribution along the line of sight, one can effectively trace the growth history of the large-scale DM structures that form the backbone of the Universe [94].

In a similar vein, the so-called **CMB lensing** phenomenon entails searching for distortions in maps of the CMB induced by foreground mass concentrations [95]. This represents, in a sense, the culmination of weak gravitational lensing as a probe of DM. Because the CMB is the earliest observable radiation, it offers a unique opportunity to study large-scale DM structures at high redshifts ($z \gtrsim 2$) [96].

The existence of DM is strongly supported by the detection of CMB lensing at high statistical significance in data from PLANCK and other experiments [96–98]. Gravitational lensing remaps the primary CMB anisotropies, modifying both temperature and polarization fields. In particular, it converts E -mode polarization into B modes and smooths the acoustic peaks at large multipoles ℓ in the CMB power spectrum [95]. As a result, lensing acts as a contaminant for measurements of primordial CMB signals, necessitating “delensing” procedures to recover the fundamental CMB fluctuations [99].

1.2.3 Global structure and evolution of the Universe

At present, the largest accessible scales—the entire observable Universe—provide the most compelling and precise evidence for the existence of dark matter (DM). The central qualitative argument is that, without the presence of DM, the Universe would not have evolved into the structure we observe today: galaxies would not be clustered in the observed manner, and the temperature anisotropies of the CMB would exhibit markedly different characteristics [13, 100, 101]. Dark matter plays a crucial role as a catalyst for structure formation. It drives the evolution of the Universe from an initial state of near-perfect homogeneity, with only tiny primordial density perturbations, into a richly structured cosmos spanning a wide range of scales [102, 103]. In particular, DM enables the gravitational growth of primordial inhomogeneities, ultimately leading to the formation of galaxies and large-scale structures observed today [25, 104]. Ordinary (baryonic)

matter alone cannot fulfill this role, primarily because of its strong coupling to radiation prior to recombination, which suppresses the growth of density perturbations in the baryon–photon fluid [28, 30]. Only after recombination can baryons fall into the gravitational potential wells already established by dark matter, thereby completing the process of structure formation.

1.2.3.1 Dark Matter Signatures in CMB Acoustic Peaks

Temperature anisotropies in the cosmic microwave background (CMB) provide a clear imprint of the presence of dark matter (DM). These anisotropies are commonly characterized through the angular power spectrum, whose lowest multipoles include the dipole moment, as discussed in the previous section in the context of matter–antimatter asymmetry [25, 105]. The detailed shape of the CMB temperature power spectrum is governed by the gravitational influence of dark matter, which provides the dominant potential wells that drive acoustic oscillations in the tightly coupled baryon–photon fluid prior to recombination [106, 107]. As a result, the relative heights and positions of the acoustic peaks are highly sensitive to the total dark matter content of the Universe [26, 101]. Consequently, the abundance of dark matter plays a crucial role in shaping the observed CMB anisotropies and encodes essential information about the composition of the Universe and the symmetry properties of matter on cosmological scales [13].

Figure 1.7 shows the angular power spectrum of temperature anisotropies in the CMB. The green curve represents the best-fit Λ CDM model to the CMB data, which provides an excellent description of the observed acoustic peak structure [13, 105]. When dark matter (DM) is replaced by baryonic matter alone, the resulting spectrum—shown by the red curve—fails to reproduce the observed peak heights and relative amplitudes, owing to the absence of deep gravitational potential wells prior to recombination [101, 109]. If DM is instead replaced by spatial curvature, the corresponding spectrum (blue curve) exhibits incorrect peak positions and angular scales, reflecting the strong dependence of the acoustic peak locations on the geometry of the Universe [110]. Finally, replacing DM with dark energy produces the black curve, which also fails to fit the data, since dark energy does not cluster on sub-horizon scales and therefore cannot support the formation of the acoustic oscillation pattern observed in the CMB [13, 111]. This comparison clearly demonstrates that the presence of dark matter is essential for achieving a successful fit to the observed CMB power spectrum.

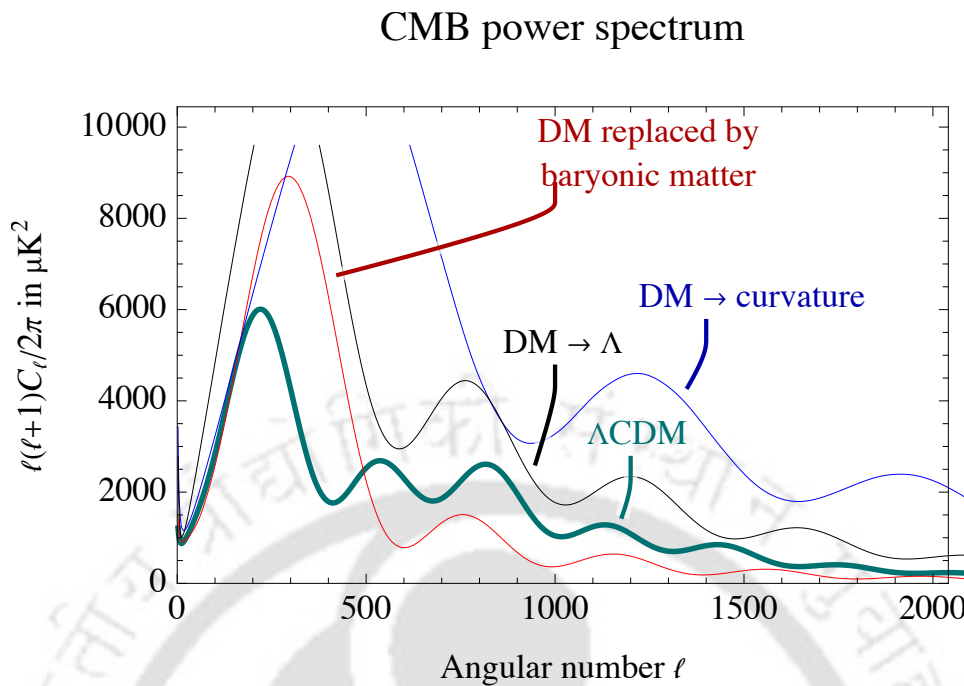


Figure 1.7: CMB temperature power spectrum measured by Planck compared to the best-fit Λ CDM model. The presence of dark matter is required to match the observed peak amplitudes and positions. Replacing the dark matter density with alternative components fails to reproduce the characteristic acoustic peak structure, providing strong evidence for the existence of non-baryonic dark matter. This figure is adapted from [108].

1.3 The Standard Model: Achievements and Limitations

Proposed by Weinberg and Salam in 1967–1968 [112, 113], the SM of particle physics provides a comprehensive framework for describing visible matter in the Universe. The mathematical framework of the SM [112–118] is rooted in Quantum Field Theory (QFT). As experimental observations brought us closer to an understanding of nature, increasingly sophisticated mathematical frameworks were developed to describe natural phenomena. The pursuit of unification in the language of mathematics began in the 1860’s with Maxwell’s equations, which unified electricity and magnetism [119]. This unifying principle later motivated the search for a unified description of all fundamental forces. The SM represents a major step in this ongoing journey of force unification [112, 113, 120]. Within the SM, the electromagnetic and weak interactions are unified into the electroweak force, Electromagnetic + Weak $\rightarrow$ Electroweak, while the strong interaction is incorporated separately. Gravity is not included in the SM [115, 121, 122]. The theory is invariant under Poincaré symmetry and a local gauge symmetry [115, 122–124] described by

Fields	$SU(3)_c$	$SU(2)_L$	$U(1)_Y$
$Q_L^\alpha = \begin{pmatrix} u_\alpha \\ d_\alpha \end{pmatrix}$	3	2	$\frac{1}{6}$
u_R^α	3	1	$\frac{2}{3}$
d_R^α	3	1	$-\frac{1}{3}$
$\ell_L^\alpha = \begin{pmatrix} \nu_\alpha \\ e_\alpha \end{pmatrix}$	1	2	$-\frac{1}{2}$
e_R^α	1	1	-1
$H = \begin{pmatrix} H^+ \\ H^0 \end{pmatrix}$	1	2	$\frac{1}{2}$

Table 1.1: Particle content of the SM and their representations under the gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$. The index α labels the three fermion generations ($\alpha = 1, 2, 3$). The subscripts L and R denote the left-handed and right-handed components, respectively.

$SU(3)_c \times SU(2)_L \times U(1)_Y$, where the subscripts c , L , and Y denote color, left-handed weak isospin, and hypercharge, respectively. It accounts for the fundamental particles and their electromagnetic, weak, and strong interactions. The present form of the SM is the result of decades of continued theoretical development and experimental verification.

The SM describes three fundamental classes of fields: A scalar Higgs field, which is responsible for electroweak symmetry breaking and mass generation; gauge bosons, which mediate the fundamental interactions; and fermions, which constitute matter. Table 1.1 summarizes the SM fields and their representations under the gauge group.

The SM Lagrangian, which includes all interactions of the fields in Table 1.1, is given by

$$\begin{aligned}
\mathcal{L}_{\text{SM}} = & \mathcal{L}_{\text{QCD}} - \frac{1}{4} \mathbf{W}_{\mu\nu} \cdot \mathbf{W}^{\mu\nu} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} + \overline{\ell}_L^\alpha i \not{D} \ell_L^\alpha + \overline{Q}_L^\alpha i \not{D} Q_L^\alpha + \overline{e}_R^\alpha i \not{\partial} e_R^\alpha + \overline{u}_R^\alpha i \not{\partial} u_R^\alpha \\
& + \overline{d}_R^\alpha i \not{\partial} d_R^\alpha - (y_{\alpha\beta}^\ell \overline{\ell}_L^\alpha H e_R^\beta + y_{\alpha\beta}^d \overline{Q}_L^\alpha H d_R^\beta + y_{\alpha\beta}^u \overline{Q}_L^\alpha H^c u_R^\beta + \text{h.c.}) \\
& + |D_\mu H|^2 - \mu_H^2 H^\dagger H - \lambda_H (H^\dagger H)^2.
\end{aligned} \tag{1.8}$$

where $V(H) = \mu_H^2 H^\dagger H + \lambda_H (H^\dagger H)^2$, with $\mu_H^2 < 0$ and $\lambda_H > 0$ and $\mathcal{L}_{\text{QCD}}$ represent the Lagrangian for Quantum Chromodynamics (QCD), the theory of strong interactions, typically written as $\mathcal{L}_{\text{QCD}} = -\frac{1}{4} G_{\mu\nu}^a G^{a,\mu\nu}$, where $G_{\mu\nu}^a$ are the gluon field strengths, and the quark fields are included in the covariant derivatives, but in this notation, the quark kinetic terms are separated into the electroweak parts below.

The definitions of the covariant derivatives are $\partial'_\mu = \partial_\mu - ig_Y \frac{Y}{2} B_\mu$ for the right-handed fields that only couple to hypercharge, and $D_\mu = \partial_\mu - ig_W \frac{\sigma}{2} \cdot \mathbf{W}_\mu - ig_Y \frac{Y}{2} B_\mu$ for the left-handed fields under $SU(2)_L \times U(1)_Y$. The charge-conjugated Higgs field is defined as $H^c = i\sigma_2 H^*$ to allow for up-type quark Yukawa couplings. The gauge field strength tensors are given by $B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu$ for the $U(1)_Y$ hypercharge and $W_{\mu\nu}^a = \partial_\mu W_\nu^a - \partial_\nu W_\mu^a + g_W \epsilon^{abc} W_\mu^b W_\nu^c$ for the $SU(2)_L$ weak isospin, where B_μ denotes the hypercharge gauge boson and W_μ^a ($a = 1, 2, 3$) are the gauge bosons associated with the $SU(2)_L$ symmetry. The Levi-Civita tensor ϵ^{abc} represents the structure constants of the $SU(2)_L$ gauge group, enabling the non-Abelian self-interactions of the weak bosons.

For clarity and organizational purposes, the SM Lagrangian in Eq. (1.8) can be decomposed as $\mathcal{L}_{\text{SM}} = \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{fermions}} + \mathcal{L}_{\text{Higgs}} + \mathcal{L}_{\text{Yukawa}} + \mathcal{L}_{\text{GF}} + \mathcal{L}_{\text{ghosts}}$, where each term corresponds to a distinct sector of the theory. The gauge sector $\mathcal{L}_{\text{gauge}}$ includes the kinetic terms and self-interactions of the gauge bosons, encompassing QCD, $SU(2)_L$, and $U(1)_Y$; the fermion sector $\mathcal{L}_{\text{fermions}}$ covers the kinetic and gauge interaction terms for leptons and quarks; the Higgs sector $\mathcal{L}_{\text{Higgs}}$ describes the scalar potential and kinetic term for the Higgs doublet, leading to symmetry breaking; and the Yukawa sector $\mathcal{L}_{\text{Yukawa}}$ handles the fermion-Higgs couplings. Meanwhile, $\mathcal{L}_{\text{GF}}$ and $\mathcal{L}_{\text{ghosts}}$ arise from the gauge-fixing procedure in quantization, ensuring consistent perturbation theory by introducing gauge-fixing terms and ghost fields to handle gauge redundancies. The Yukawa Lagrangian $\mathcal{L}_{\text{Yukawa}}$ encodes the interactions between the Higgs field and the fermions and is responsible for generating fermion masses after electroweak symmetry breaking, as the Higgs vacuum expectation value induces mass terms through these couplings.

After spontaneous symmetry breaking, the Higgs field acquires a vacuum expectation value (VEV), denoted as $v = \sqrt{-\frac{\mu_H^2}{\lambda_H}} = 246 \text{ GeV}$. This process fundamentally alters the dynamics of the electroweak sector in the SM.

The Higgs kinetic term then generates mass terms for the $W^\pm$ and Z gauge bosons through their interactions with the Higgs field. The resulting masses are given by

$$m_{W^\pm} = \frac{1}{2} g_W v, \quad m_Z = \frac{1}{2} \frac{g_W}{\cos \theta} v, \quad (1.9)$$

where g_W is the weak coupling constant and θ is the weak mixing (Weinberg) angle, which

parameterizes the relative strengths of the neutral currents.

In this broken symmetry phase, the physical fields emerge as mixtures of the original gauge fields. Specifically, the photon A_μ remains massless, while the neutral Z_μ boson gains mass. Their relationships to the original hypercharge B_μ and weak isospin W_μ^3 fields are expressed as

$$\begin{pmatrix} A_\mu \\ Z_\mu \end{pmatrix} = \begin{pmatrix} \cos \theta_W & \sin \theta_W \\ -\sin \theta_W & \cos \theta_W \end{pmatrix} \begin{pmatrix} B_\mu \\ W_\mu^3 \end{pmatrix}, \quad (1.10)$$

where θ_W is the Weinberg angle, ensuring the photon decouples from the massive Z .

The fermion masses in the SM arise similarly from Yukawa interactions with the Higgs field. The mass terms related to the charged leptons and quarks are

$$\begin{aligned} m_\ell^{\alpha\beta} &= Y_{\alpha\beta}^\ell \frac{v}{\sqrt{2}}, \\ m_u^{\alpha\beta} &= Y_{\alpha\beta}^u \frac{v}{\sqrt{2}}, \\ m_d^{\alpha\beta} &= Y_{\alpha\beta}^d \frac{v}{\sqrt{2}}, \end{aligned} \quad (1.11)$$

where $Y^{\ell,u,d}$ are the Yukawa coupling matrices for charged leptons, up-type quarks, and down-type quarks, respectively. These matrices are generally non-diagonal in flavor space, reflecting the generational structure of fermions.

To align with observed mass eigenstates, we diagonalize these mass matrices using unitary transformations. For charged leptons and quarks, this yields

$$V_L^{e,\dagger} m_\ell^{\alpha\beta} V_R^e = \text{diag}(m_e, m_\mu, m_\tau), \quad (1.12)$$

$$V_L^{d,\dagger} m_d^{\alpha\beta} V_R^d = \text{diag}(m_d, m_s, m_b), \quad (1.13)$$

$$V_L^{u,\dagger} m_u^{\alpha\beta} V_R^u = \text{diag}(m_u, m_c, m_t), \quad (1.14)$$

where the diagonal entries are the physical masses of the electron, muon, tau, down, strange, bottom, up, charm, and top quarks. The matrices V_L and V_R rotate the left- and right-handed fields into the mass basis.

When we apply this basis transformation to the quark–gauge interaction terms, the resulting Lagrangian now includes terms that mix different quark flavors in the weak charged currents. These flavor-mixing effects are fully described by a single unitary matrix,

$$V = (V_L^u)^\dagger V_L^d, \quad (1.15)$$

known as the Cabibbo–Kobayashi–Maskawa (CKM)[125, 126] matrix. This matrix en-

codes the weak charged-current interactions between up-type and down-type quarks and contains one physical complex phase responsible for CP violation in the quark sector, enabling processes like matter-antimatter asymmetry in certain decays.

The SM of particle physics stands as one of the greatest achievements of modern science. It successfully classifies all known fundamental particles—six quarks, six leptons, and the gauge bosons that mediate their interactions—while providing a precise description of the electromagnetic, weak, and strong nuclear forces. The SM predicted and guided the discovery of the charm, bottom, and top quarks, the $W^\pm$ and Z bosons, the tau neutrino, and, most notably, the Higgs boson, whose discovery at CERN in 2012 confirmed the mechanism by which elementary particles acquire mass.

Theoretical predictions of the SM, ranging from electroweak unification to quantum chromodynamics, are in remarkable agreement with experimental observations, often matching measurements to an accuracy of up to ten decimal places. These include precision tests such as the anomalous magnetic moment of the muon and measurements of CP violation in the quark sector. By unifying the electromagnetic and weak interactions and explaining key phenomena such as asymptotic freedom and quark confinement, the SM provides the foundational framework for understanding the Universe at its most fundamental level. For more than five decades, it has driven progress in particle accelerator technology, cosmology, and numerous scientific and technological applications.

The success of the SM lies in its exceptional predictive power and extensive experimental validation [127]. It has accurately predicted the existence and properties of fundamental particles such as the charm quark [128], the bottom quark [129], and the top quark [130], as well as the $W^\pm$ and Z gauge bosons [131, 132], and the Higgs boson, which was discovered in 2012 [133, 134]. Moreover, the SM unifies the electromagnetic and weak interactions into a single electroweak framework [135, 112, 136], whose predictions are in remarkable agreement with precision measurements at the level of one part in a billion [137, 138].

The SM provides a robust description of the strong interaction through quantum chromodynamics (QCD), accounting for phenomena such as quark confinement and asymptotic freedom [139, 140]. In addition, it incorporates CP violation in the weak sector via the Cabibbo–Kobayashi–Maskawa (CKM) matrix [141, 142]. This theoretical framework enables a coherent understanding of matter–antimatter interactions and the particle dynamics of the early Universe [143]. Validated over a vast range of energy scales, from nuclear processes to experiments conducted at the Large Hadron Collider [144], the SM has withstood decades of increasingly stringent experimental scrutiny [145]. Beyond its theoretical achievements, the SM underpins numerous practical technologies, including

Issue	SM	Observed	BSM?
Gravity	Not included	Exists	Yes
Neutrino mass	$m_\nu = 0$	$m_\nu > 0$	Yes
Dark matter	—	$\Omega_{\text{DM}} h^2 \simeq 0.12$	Yes
Baryon asymmetry	$\eta \ll 10^{-20}$	$\eta \approx 6 \times 10^{-10}$	Yes
Hierarchy problem	m_H unstable	$m_H = 125 \text{ GeV}$	Yes

Table 1.2: Key reasons why the SM fails to account for the observed properties.

medical imaging, particle accelerators, and models of stellar nucleosynthesis [146–148], making it the most successful and comprehensive theory of fundamental particles and interactions developed to date [149].

Despite its remarkable success, the SM remains fundamentally incomplete, as summarized in Table 1.2 [150]. Most notably, it does not incorporate gravity and lacks a consistent quantum description compatible with general relativity [151]. Moreover, the SM predicts neutrinos to be massless, in contradiction with experimental observations of neutrino oscillations, which require small but nonzero neutrino masses [152, 153].

The SM also fails to provide a viable candidate for dark matter, which constitutes approximately 27% of the Universe’s total mass–energy content [154], nor does it offer an explanation for dark energy, the unknown component driving the accelerated expansion of the Universe [155]. In addition, the amount of CP violation arising from the Cabibbo–Kobayashi–Maskawa (CKM) matrix is far too small to account for the observed matter–antimatter asymmetry of the Universe [156, 157].

Another major theoretical challenge is the hierarchy problem: the Higgs boson mass of approximately 125 GeV is unnaturally small compared to the Planck scale, $M_{\text{Pl}} \sim 10^{19} \text{ GeV}$, requiring severe fine-tuning to remain stable under quantum corrections [158, 159]. Further unresolved issues include the strong CP problem, wherein quantum chromodynamics permits CP violation that is experimentally absent [160], as well as the presence of 25 free parameters in the SM—such as fermion masses and mixing angles—whose numerical values lack a deeper theoretical explanation [161].

The table 1.2 summarizes key shortcomings of the SM of particle physics. It lists seven major issues—ranging from the absence of gravity to the strong CP problem—comparing SM predictions with experimental observations and indicating the need for physics beyond the Standard Model (BSM) in each case, with specific BSM solutions noted where applicable (e.g., axions, SUSY/GUTs).

Despite its remarkable success, the SM of particle physics does not provide a viable

dark matter candidate [162, 163]. All electrically neutral particles in the SM either interact too strongly or are too light to account for the observed dark matter abundance. In particular, SM neutrinos constitute hot dark matter and are ruled out by structure formation constraints [164, 165]. Furthermore, cosmological observations from Big Bang nucleosynthesis and measurements of the CMB indicate that dark matter must be predominantly non-baryonic, which cannot be accommodated within the SM [166, 167]. These considerations strongly motivate physics beyond the Standard Model.

The SM Lagrangian conserves both baryon number (B) and lepton number (L), and therefore does not violate these quantum numbers. After spontaneous symmetry breaking, the vacuum expectation value of the Higgs scalar field becomes nonzero. As discussed in [168], the group $SU(2) \otimes U(1)$ maps into $U(1)$, which can be associated with an integer-valued winding number known as the Higgs winding number [169, 170]. For a given vacuum configuration, the Higgs winding number is equal to its Chern–Simons number. Transitions between different vacuum configurations change the Chern–Simons number, which is directly related to changes in the baryon number [171].

A conserved charge arises from a continuous symmetry through Noether’s theorem. In the SM, the baryon number (B) and lepton number (L) are defined as

$$B = \int d^3x J_0^B(x), \quad L = \int d^3x J_0^L(x), \quad (1.16)$$

where J_μ^B and J_μ^L are the baryon and lepton currents, respectively. Explicitly, these currents are given by

$$J_\mu^B = \frac{1}{3} \sum_\alpha (\bar{Q}_L^\alpha \gamma_\mu Q_L^\alpha + \bar{u}_R^\alpha \gamma_\mu u_R^\alpha + \bar{d}_R^\alpha \gamma_\mu d_R^\alpha), \quad (1.17)$$

$$J_\mu^L = \sum_\alpha (\bar{\ell}_L^\alpha \gamma_\mu \ell_L^\alpha + \bar{e}_R^\alpha \gamma_\mu e_R^\alpha), \quad (1.18)$$

where α labels the fermion generations.

At the classical level, these currents are conserved. However, at the quantum level they are violated by chiral anomalies, leading to

$$\partial_\mu J_B^\mu = \partial_\mu J_L^\mu = \frac{N_f}{32\pi^2} \left(g_W^2 W_{\mu\nu}^a \tilde{W}^{a\mu\nu} - g_Y^2 F_{\mu\nu} \tilde{F}^{\mu\nu} \right), \quad (1.19)$$

where N_f is the number of fermion generations, $W_{\mu\nu}^a$ and $F_{\mu\nu}$ are the $SU(2)_L$ and $U(1)_Y$ field strength tensors, and $\tilde{X}^{\mu\nu} \equiv \frac{1}{2} \epsilon^{\mu\nu\rho\sigma} X_{\rho\sigma}$ denotes the dual field strength.

As a consequence, the anomaly-free combination satisfies

$$\partial_\mu (J_B^\mu - J_L^\mu) = 0, \quad (1.20)$$

while

$$\partial_\mu (J_B^\mu + J_L^\mu) = 2N_f \partial_\mu \kappa^\mu, \quad (1.21)$$

where the Chern–Simons current κ^μ is given by

$$\kappa^\mu = -\frac{g_W^2}{32\pi^2} \epsilon^{\mu\nu\rho\sigma} \left[W_\nu^a \left(\partial_\rho W_\sigma^a - \frac{g_W}{3} \epsilon^{abc} W_\rho^b W_\sigma^c \right) \right] + \frac{g_Y^2}{32\pi^2} \epsilon^{\mu\nu\rho\sigma} B_\nu B_{\rho\sigma}. \quad (1.22)$$

As first discussed by [172], the electroweak sector of the SM possesses an infinite set of degenerate vacuum states separated by a potential barrier. Each vacuum is characterized by a topological invariant known as the Chern–Simons number, N_{CS} . Transitions between different vacuum sectors correspond to changes in N_{CS} and are associated with nonperturbative gauge-field configurations.

Such transitions induce a change in baryon and lepton numbers according to

$$\Delta B = \Delta L = N_f \Delta N_{\text{CS}}, \quad (1.23)$$

where N_f is the number of fermion generations. These processes violate $B + L$ while conserving $B - L$.

At zero temperature, transitions between different vacua occur via quantum tunneling and are exponentially suppressed (instantons). At finite temperature, particularly around the electroweak phase transition, thermal fluctuations can overcome the barrier through sphaleron transitions. As a result, processes with

$$\Delta(B + L) = 2N_f \Delta N_{\text{CS}} = 6, 12, 18, \dots \quad (1.24)$$

are allowed in the SM, signaling changes in the vacuum topology. This has profound implications for baryogenesis, as any baryon asymmetry generated above the electroweak scale must carry a nonzero $B - L$ component in order to survive sphaleron washout.

At zero temperature, a first-order phase transition can occur, and CP violation is permitted within the SM. As discussed in Ref. [173, 174], the Euclidean action S_E is related to the Chern–Simons number N_{CS} via

$$S_E > \frac{8\pi^2}{g^2} N_{\text{CS}}. \quad (1.25)$$

The transition rate, which depends exponentially on the Euclidean action, is therefore also determined by the Chern–Simons number. For $N_{\text{CS}} = 1$, the transition rate at zero temperature is

$$\Gamma(T = 0) = e^{-2S_E} \approx 10^{-170}. \quad (1.26)$$

Such an extremely suppressed rate implies that baryon-number-violating processes are essentially unobservable at zero temperature, as these transitions proceed via quantum tunneling through a large potential barrier. In contrast, at finite temperature, transitions between different vacua can occur through thermal fluctuations that allow the system to overcome the barrier. As the temperature rises, the transition rate increases and also depends on the Higgs boson mass. To generate a sufficiently large baryon asymmetry η_B , the Higgs mass would need to be approximately 80 GeV, as reported in Ref. [175, 176]. However, this is inconsistent with the Higgs boson mass measured at the Large Hadron Collider. Therefore, although the SM formally satisfies all three Sakharov conditions (as discussed in Section 2.3), it fails to produce the observed magnitude of baryon asymmetry.

Taken together, these limitations show that the SM should be considered a highly successful low-energy effective theory. Its inability to address phenomena at cosmological scales and in the domain of quantum gravity strongly suggests the need for new physics beyond the SM.



Theoretical Frameworks of Dark Matter and Baryon Asymmetry

After presenting an extensive literature review of the Dark matter and Leptogenesis in the Chapter 1, in this Chapter, we present the theoretical frameworks underlying the origin of dark matter and the baryon asymmetry of the Universe, which constitute two of the most profound open problems in modern cosmology and particle physics. We begin by introducing the Boltzmann equations, which form the central tool used throughout this thesis to describe the thermal evolution of particle abundances in the early Universe. We then discuss the principal mechanisms for dark matter production, focusing on minimal models and the paradigms of freeze-out and freeze-in. This is followed by a detailed overview of theoretical models explaining matter–antimatter asymmetry, including baryogenesis, electroweak baryogenesis, Affleck-Dine baryogenesis, and scenarios arising from the decay of heavy particles such as grand unified theories and leptogenesis, with particular emphasis on the role of CP-violating asymmetry parameters.

2.1 Boltzmann equation

This thesis primarily focuses on matter–antimatter asymmetry and its possible connection to dark matter, which is explored in detail in the subsequent chapters. We do not restrict ourselves to the specific properties of dark matter particles, as one or multiple viable dark matter candidates may exist. If dark matter were present in sufficient abundance in the early Universe, it would have undergone annihilation into Standard Model (SM) particles. Following annihilation processes, a residual population of dark matter would have survived until the present epoch, giving rise to the observed dark matter relic density.

Alternatively, dark matter may be produced through the decay of heavier particles via the freeze-in mechanism, which also contributes to the present-day dark matter relic density. The evolution of any particle species in the expanding Universe is governed by the interplay between the cosmic expansion rate and the relevant interaction rates. These rates are systematically encoded in the Boltzmann equation, which determines the time evolution of the phase-space distribution function $f(x, p)$ for a particle species. The corresponding **Boltzmann equation** [2] is given by

$$\hat{L}[f] = C[f], \quad (2.1)$$

where the left-hand side describes the free streaming of particles in curved spacetime, while the right-hand side accounts for particle interactions through the collision operator $C[f]$. The Liouville operator $\hat{L}$ appearing in the Boltzmann equation is the covariant generalization of the classical phase-space flow operator. In a general curved spacetime it is given by

$$\hat{L} = p^\alpha \frac{\partial}{\partial x^\alpha} - \Gamma_{\beta\gamma}^\alpha p^\beta p^\gamma \frac{\partial}{\partial p^\alpha}, \quad (2.2)$$

where p^α denotes the four-momentum of the particle and $\Gamma_{\beta\gamma}^\alpha$ are the Christoffel symbols associated with the metric.

Specializing to a homogeneous and isotropic Friedmann–Robertson–Walker (FRW) universe, the Boltzmann equation can be simplified considerably. After integrating over momentum space and assuming spatial homogeneity, one obtains an evolution equation for the number density n of the particle species,

$$\frac{dn}{dt} + 3\frac{\dot{R}}{R}n = \frac{g}{(2\pi)^3} \int C[f] \frac{d^3p}{E}, \quad (2.3)$$

where $R(t)$ is the cosmic scale factor and the second term on the left-hand side accounts for dilution due to the expansion of the Universe.

The collision term on the right-hand side encodes all processes that create or destroy the particle species under consideration. For a general reaction of the form

$$X + a + b + \cdots \longleftrightarrow i + j + \cdots, \quad (2.4)$$

the integrated collision term can be written as

$$\begin{aligned} \frac{g}{(2\pi)^3} \int C[f] \frac{d^3 p_X}{E_X} = & - \int d\Pi_X d\Pi_a d\Pi_b \cdots d\Pi_i d\Pi_j \cdots \\ & \times (2\pi)^4 \delta^4(p_X + p_a + p_b \cdots - p_i - p_j \cdots) \\ & [|M|_{X+a+b+\dots i+j+\dots}^2 f_X f_a f_b \cdots (1 \pm f_i)(1 \pm f_j) \cdots \\ & |M|_{i+j+\dots X+a+b+\dots}^2 (1 \pm f_i)(1 \pm f_j) \cdots f_X f_a f_b \cdots] \end{aligned} \quad (2.5)$$

where the phase-space integrals, matrix elements, and the energy-momentum conserving delta function explicitly account for both the forward and inverse reactions.

To make further analytical progress, we assume that all particle species obey Maxwell-Boltzmann statistics. Under this assumption, the phase-space distribution function for a species x can be written as

$$f_x = \frac{n_x}{n_x^{\text{eq}}} e^{-E/T} = e^{-(E-\mu_x)/T}, \quad (2.6)$$

where μ_x denotes the chemical potential. In thermal equilibrium one has $\mu_x = 0$, and the equilibrium number density is given by

$$n_x^{\text{eq}} = g_x \int d\Pi_x e^{-E/T}, \quad (2.7)$$

$$n_x^{\text{eq}} = g_x \frac{m_x^3}{2\pi^2} \frac{K_2(z)}{z}. \quad (2.8)$$

Furthermore, since we are working in the dilute regime, quantum statistical effects such as Bose enhancement or Pauli blocking can be neglected, allowing the approximation $1 \pm f_x \simeq 1$. It is often convenient to rewrite the Boltzmann equation in terms of the comoving number density, $Y \equiv \frac{n}{s}$, where s is the entropy density of the Universe. This variable removes the explicit effect of cosmic expansion. During radiation domination, the Hubble parameter scales as $H(t) = 1/(2t)$, and the temperature evolves as $T \propto t^{-1/2}$. Introducing the dimensionless variable $x \equiv m/T$, the Boltzmann equation can finally be expressed in the compact form

$$\mathcal{H} s x \frac{dY}{dx} = \frac{g}{(2\pi)^3} \int C[f] \frac{d^3 p_X}{E_X}, \quad (2.9)$$

that serves as the starting point for studying the thermal evolution of particle abundances in the early Universe. The Boltzmann equation encapsulates two key pieces of information. First, it describes how the number density of a species changes due to the expansion of the

Universe. Second, it accounts for the effects of particle interactions—such as scattering, decay, and annihilation—on the number density. Using this equation, one can determine how long a species remains in thermal equilibrium, as well as when it undergoes freeze-out or freeze-in during cosmic evolution.

As discussed previously, changes in the number density of any species can be studied using the Boltzmann equation in Eq. 2.9. The right-hand side of this equation encodes information about the interactions among species. These interactions can occur in various ways, and based on such possibilities, we classify them into different regimes. For instance, consider a single species X that is produced via the decay of Y into X and Z , denoted as $Y \rightarrow XZ$. If the inverse decay process is forbidden, then in this scenario, the evolution of X can be analyzed using the Boltzmann equation from Eq. 2.9 as follows:

$$\mathcal{H} s x \frac{dY_X}{dx} = \gamma_Y. \quad (2.10)$$

where γ_Y is the thermally averaged decay width of decay $Y \rightarrow XZ$. With the help of first principle thermal averaged of decay width can be written as

$$\begin{aligned} \gamma(Y \rightarrow XZ) &= \int d\Pi_Y e^{-E_Y/T} \int d\Pi_X d\Pi_Z (2\pi)^4 \\ &\quad \delta^{(4)}(p_Y - p_X - p_Z) |\mathcal{M}(Y \rightarrow X + Z)|^2 \\ &= g_Y \int \frac{d^3 p_Y}{(2\pi)^3} e^{-E_Y/T} \frac{M_X}{E_X} \Gamma(Y \rightarrow X + Z) \\ &= n_Y^{\text{eq}} \frac{K_1(z)}{K_2(z)} \Gamma(Y \rightarrow X + Z). \end{aligned} \quad (2.11)$$

Where $n_Y^{\text{eq}} = g_Y \frac{m_Y^3}{2\pi^2} \frac{K_2(z)}{z}$ is the equilibrium number density of Y , $K_1(z)$ and $K_2(z)$ are the modified Bessel function. In this simple case, we assume that the species Y and Z remains in equilibrium through other interactions, and that the process respects CP symmetry. If Y is not in equilibrium, while all other conditions remain unchanged, this will impact the number density of X by altering the right-hand side of Eq. 2.10 accordingly, such that the Boltzmann equation becomes

$$\mathcal{H} s x \frac{dY_Y}{dx} = -\frac{Y_Y}{Y_Y^{\text{eq}}} \gamma_Y \quad (2.12)$$

$$\mathcal{H} s x \frac{dY_X}{dx} = \frac{Y_Y}{Y_Y^{\text{eq}}} \gamma_Y. \quad (2.13)$$

In this scenario, the number density of Y decreases, while that of X increases. The rates

at which Y decreases and X increases depend on the interaction rate γ_Y . In the first case, the evolution of the number density of X is independent of the number density of Y , whereas in the second case, the evolution of the number density of X is affected by the number density of Y . Now, consider the next scenario, where both decay and inverse decay are allowed. In this case, Eq. 2.10 becomes

$$\mathcal{H}sx \frac{dY}{dx} = \gamma_Y \left(1 - \frac{Y_X}{Y_X^{eq}} \right), \quad (2.14)$$

and Eq. 2.13 take the form as

$$\mathcal{H}sx \frac{dY_Y}{dx} = -\gamma_Y \left(\frac{Y_Y}{Y_Y^{eq}} - \frac{Y_X}{Y_X^{eq}} \right), \quad (2.15)$$

$$\mathcal{H}sx \frac{dY_X}{dx} = \gamma_Y \left(\frac{Y_Y}{Y_Y^{eq}} - \frac{Y_X}{Y_X^{eq}} \right). \quad (2.16)$$

The species Z is in thermal equilibrium. In this scenario, the forward process increases the abundance of X , while the inverse process reduces the abundance of Y . The rate at which the abundances of X and Y change depends on the corresponding interaction rates.

There are three species participating in the interaction $Y \rightarrow XZ$. If the species Y , X , and Z are all out of equilibrium, then the Boltzmann equation 2.13 becomes

$$\mathcal{H}sx \frac{dY_Y}{dx} = -\gamma_Y \left(\frac{Y_Y}{Y_Y^{eq}} - \frac{Y_Z Y_X}{Y_Z^{eq} Y_X^{eq}} \right), \quad (2.17)$$

$$\mathcal{H}sx \frac{dY_X}{dx} = \gamma_Y \left(\frac{Y_Y}{Y_Y^{eq}} - \frac{Y_Z Y_X}{Y_Z^{eq} Y_X^{eq}} \right), \quad (2.18)$$

$$\mathcal{H}sx \frac{dY_Z}{dx} = \gamma_Y \left(\frac{Y_Y}{Y_Y^{eq}} - \frac{Y_Z Y_X}{Y_Z^{eq} Y_X^{eq}} \right). \quad (2.19)$$

In this scenario, the evolutions of species Z , Y and X are somewhat different. The decay rate determines how quickly Y and X evolve over the history of the Universe.

In place of decay, if there are scattering processes $X, Y \rightarrow Z, N$, now we are interested in studying the abundance of X species. As mentioned, for all decay cases, the evolution of X during scattering can be studied using the Boltzmann equations.

$$\mathcal{H}sx \frac{dY_X}{dx} = -\frac{Y_X}{Y_X^{eq}} \gamma_{X,Y}^{Z,N} \quad \text{case 1} \quad (2.20a)$$

$$\mathcal{H} s x \frac{dY_X}{dx} = \gamma_{X,Y}^{Z,N} \left(1 - \frac{Y_X}{Y_X^{\text{eq}}} \right) \quad \text{case 2} \quad (2.21a)$$

$$\mathcal{H} s x \frac{dY_Y}{dx} = -\frac{Y_Y Y_X}{Y_Y^{\text{eq}} Y_X^{\text{eq}}} \gamma_{X,Y}^{Z,N} \quad \text{case 3} \quad (2.22a)$$

$$\mathcal{H} s x \frac{dY_X}{dx} = -\frac{Y_Y Y_X}{Y_Y^{\text{eq}} Y_X^{\text{eq}}} \gamma_{X,Y}^{Z,N} \quad (2.22b)$$

$$\mathcal{H} s x \frac{dY_Y}{dx} = \gamma_{X,Y}^{Z,N} \left(1 - \frac{Y_Y Y_X}{Y_Y^{\text{eq}} Y_X^{\text{eq}}} \right) \quad \text{case 4} \quad (2.23a)$$

$$\mathcal{H} s x \frac{dY_X}{dx} = \gamma_{X,Y}^{Z,N} \left(1 - \frac{Y_Y Y_X}{Y_Y^{\text{eq}} Y_X^{\text{eq}}} \right) \quad (2.23b)$$

$$\mathcal{H} s x \frac{dY_Z}{dx} = -\gamma_{X,Y}^{Z,N} \left(\frac{Y_Z Y_N}{Y_Z^{\text{eq}} Y_N^{\text{eq}}} - \frac{Y_Y Y_X}{Y_Y^{\text{eq}} Y_X^{\text{eq}}} \right) \quad \text{case 5} \quad (2.24a)$$

$$\mathcal{H} s x \frac{dY_N}{dx} = -\gamma_{X,Y}^{Z,N} \left(\frac{Y_Z Y_N}{Y_Z^{\text{eq}} Y_N^{\text{eq}}} - \frac{Y_Y Y_X}{Y_Y^{\text{eq}} Y_X^{\text{eq}}} \right) \quad (2.24b)$$

$$\mathcal{H} s x \frac{dY_Y}{dx} = \gamma_{X,Y}^{Z,N} \left(\frac{Y_Z Y_N}{Y_Z^{\text{eq}} Y_N^{\text{eq}}} - \frac{Y_Y Y_X}{Y_Y^{\text{eq}} Y_X^{\text{eq}}} \right) \quad (2.24c)$$

$$\mathcal{H} s x \frac{dY_X}{dx} = \gamma_{X,Y}^{Z,N} \left(\frac{Y_Z Y_N}{Y_Z^{\text{eq}} Y_N^{\text{eq}}} - \frac{Y_Y Y_X}{Y_Y^{\text{eq}} Y_X^{\text{eq}}} \right) \quad (2.24d)$$

Here, $\gamma_{X,Y}^{Z,N}$ denotes the thermally averaged reaction density for the scattering process $XY \rightarrow ZN$. From first principles, it can be written as

$$\gamma_{X,Y}^{Z,N} = \int d\Pi_X d\Pi_Y d\Pi_Z d\Pi_N (2\pi)^4 \delta^4(p_X + p_Y - p_Z - p_N) e^{-E_X/T} e^{-E_Y/T} |\mathcal{M}|^2. \quad (2.25)$$

Different scenarios associated with the scattering processes were discussed above. Cases 1 and 3 correspond to situations in which only the forward process is allowed, whereas Cases 2, 4, and 5 describe scenarios in which both the forward and reverse processes are permitted. In Cases 1 and 2, the species Y , Z , and N remain in thermal equilibrium. Consequently, a single Boltzmann equation is sufficient to describe the evolution of the abundance of species X . In Cases 3 and 4, only the species Z and N are in equilibrium, while in Case 5 all species, namely X , Y , Z , and N , are out of equilibrium.

The rate of the scattering process determines how rapidly or slowly the abundance of species X freezes out. The above discussion applies only when the processes considered do not violate CP symmetry. If either the decay or scattering processes violate CP symmetry, the Boltzmann equations must be modified accordingly. This scenario is studied in our leptogenesis work.

2.2 Possible Mechanism of Dark Matter

Based on their production mechanisms, dark matter models can be broadly classified into two categories: freeze-out and freeze-in. Each mechanism can be further subdivided according to the structure of the dark matter sector, such as single-component or multi-component dark matter. In the following discussion, we first examine the single- and multi-component WIMP dark matter scenarios, and then proceed to the single- and multi-component FIMP dark matter models.

2.2.1 A Simple Model of Dark Matter

There exist many models that address the dark matter problem. Here, however, we focus on the simplest realization in order to illustrate the underlying dark matter mechanism. The dark matter puzzle can be explained by a minimal extension of the SM that includes a real scalar field which is odd under a $\mathbb{Z}_2$ symmetry. With this extension, the Lagrangian of the model can be written as

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \frac{1}{2}(\partial_\mu S)(\partial^\mu S) - \left[\frac{1}{2}m_S^2 S^2 + \frac{1}{4}\lambda_S S^4 + \frac{1}{2}\lambda_{HS} S^2 H^\dagger H \right]. \quad (2.26)$$

The scalar particle S is stable due to the presence of a $\mathbb{Z}_2$ symmetry, which is imposed by hand and remains unbroken in the vacuum. As a result, the scalar S interacts with the SM particles only through the SM Higgs boson, a feature commonly referred to as the Higgs portal. Consequently, elastic scattering with nuclei, relevant for direct detection experiments, and dark matter annihilation processes in the early Universe are intrinsically related [].

The mass of the scalar dark matter particle is given by

$$M_S^2 = m_S^2 + \lambda_{HS} v^2, \quad (2.27)$$

where $v = 246$ GeV denotes the Higgs vacuum expectation value and m_S represents the bare mass of the scalar field S . This relation indicates that obtaining a light dark matter

mass M_S within this framework generally requires a cancellation between the two terms on the right-hand side, implying a degree of fine-tuning.

The stable scalar dark matter candidate can annihilate into SM particles via the Higgs portal. Within the WIMP paradigm, the particle S attains the observed dark matter relic abundance, $\Omega_{\text{DM}} h^2 \simeq 0.12$, through the thermal freeze-out mechanism in the early Universe. At high temperatures, S remains in thermal equilibrium with the SM plasma through processes such as

$$SS \rightarrow h^* \rightarrow \begin{cases} f\bar{f}, \\ W^+W^-, \\ ZZ, \\ hh. \end{cases} \quad (2.28)$$

The epoch at which dark matter decouples from the thermal bath depends on the interaction rate of the relevant processes. Taking into account all possible annihilation channels of dark matter into SM particles, the Boltzmann equation governing the evolution of the dark matter species S can be written as

$$\mathcal{H} s x \frac{dY_S}{dx} = (\gamma_{S,S}^{f\bar{f}} + \gamma_{S,S}^{W^+W^-} + \gamma_{S,S}^{ZZ} + \gamma_{S,S}^{hh}) \left[1 - \frac{Y_S^2}{(Y_S^{\text{eq}})^2} \right], \quad (2.29)$$

where $\gamma_{S,S}^{f\bar{f}}$, $\gamma_{S,S}^{W^+W^-}$, $\gamma_{S,S}^{ZZ}$, and $\gamma_{S,S}^{hh}$ denote the thermally averaged scattering cross sections for the corresponding processes, as defined in Eq. 2.25. The comoving number density Y_S can be obtained by solving Eq. 2.29. The present-day value of Y_S determines the dark matter relic density, which is given by

$$\Omega_S h^2 = \frac{s_0 h^2 m_S Y_S}{\rho_c}. \quad (2.30)$$

where $\rho_c = \frac{3H^2}{8\pi G}$ is the critical density. Here, $H = 100 h \text{ km s}^{-1} \text{ Mpc}^{-1}$ denotes the present value of the Hubble constant, h is the dimensionless reduced Hubble parameter, and G represents Newton's gravitational constant.

The next important question concerns the epoch at which dark matter decouples from the thermal bath. As the comoving number density Y_S approaches its equilibrium value Y_S^{eq} , dark matter remains in thermal contact with the plasma. The deviation from

equilibrium, defined as

$$\delta \equiv \left| \frac{Y_S - Y_S^{\text{eq}}}{Y_S^{\text{eq}}} \right|. \quad (2.31)$$

For thermal equilibrium $\delta \ll 1$. When the interaction rate drops below the Hubble expansion rate the abundance goes out of equilibrium and δ becomes order unity. This is the beginning of freeze-out. When this condition is satisfied, the dark matter particle S decouples from the thermal bath. This moment is referred to as the freeze-out point. At freeze-out, Eq. 2.29 can be approximated as

$$\mathcal{H} s x \frac{dY_S}{dx} \Big|_{x_f} = (\gamma_{S,S}^{f,\bar{f}} + \gamma_{S,S}^{W^+,W^-} + \gamma_{S,S}^{Z,Z} + \gamma_{S,S}^{h,h}) \Big|_{x_f}, \quad (2.32)$$

$$x_f = \frac{(\gamma_{S,S}^{f,\bar{f}} + \gamma_{S,S}^{W^+,W^-} + \gamma_{S,S}^{Z,Z} + \gamma_{S,S}^{h,h}) \Big|_{x_f}}{H(x_f) s(x_f) \Big|_{x_f} \frac{dY_S}{dx} \Big|_{x_f}}. \quad (2.33)$$

Here, we do not attempt to determine the allowed parameter space; rather, we focus on illustrating the underlying mechanism used to analyze dark matter production and to understand the conditions under which dark matter freezes out. In the subsection 2.2.1.1, we consider a simple example to study the evolution of the dark matter abundance.

The dark matter relic density is mainly controlled by two parameters: the Higgs-portal coupling λ_{HS} and the dark matter mass M_S . Typically, in order to obtain the observed dark matter relic abundance, the coupling is constrained within a range $n_1 < \lambda_{HS} < n_2$, where n_1 and n_2 represent the lower and upper bounds on the coupling constant, respectively. When the dark matter mass M_S approaches half of the Higgs boson mass, the dark matter annihilation cross section is resonantly enhanced, leading to a maximum annihilation rate.

In the context of direct detection, dark matter particles scatter off nucleons via t -channel processes, in particular through $SN \rightarrow SN$. The interaction rate for this process depends directly on the square of the coupling, λ_{HS}^2 . A more detailed discussion of the WIMP scenario, along with a concrete example, will be presented in subsection 2.2.1.1.

Another possibility for dark matter production is the FIMP scenario. In the simplest realization of this framework, the decay process $h \rightarrow SS$ can generate the dark matter abundance. Alternatively, dark matter may also be produced through the scattering of SM particles, such as $f\bar{f} \rightarrow SS$, $W^+W^- \rightarrow SS$, and $hh \rightarrow SS$. The resulting dark matter relic abundance depends on which of these processes dominates the production. The general Boltzmann equation governing the above-mentioned processes is given by

$$\mathcal{H} s x \frac{dY_S}{dx} = \left(\gamma_h^{SS} + \frac{Y_f Y_{\bar{f}}}{Y_f^{eq} Y_{\bar{f}}^{eq}} \gamma_{ff}^{SS} + \frac{Y_{W^+} Y_{W^-}}{Y_{W^+}^{eq} Y_{W^-}^{eq}} \gamma_{W^+ W^-}^{SS} + \gamma_{hh}^{SS} \right). \quad (2.34)$$

In this mechanism, the initial abundance of dark matter candidates is assumed to be negligible. As the Universe cools, DM particles are gradually produced through decays or scatterings of SM states. Since the interaction rates are feeble, DM never reaches thermal equilibrium with the SM plasma. A detailed discussion of the FIMP scenario will be presented in subsection 2.2.1.2, where we illustrate the mechanism with a concrete example.

2.2.1.1 Freeze Out: Origin of thermal Relic Dark Matter

As long as dark matter (DM) and the SM plasma remain in chemical equilibrium, they also maintain kinetic equilibrium as the temperature of the Universe decreases. In practice, kinetic equilibrium typically persists well after chemical decoupling. This is due to the abundance of light SM particles, which makes elastic scattering processes far more frequent than number-changing interactions. Nevertheless, this separation of timescale is not so universal. Exception can arise even in relatively simple models, and are expected to be more common in scenarios with richer dynamics, such as the decay of the heavier states or self-heating effect induced by semi-annihilations, which in turn can affect local thermal equilibrium.

Let us first consider the case of stable particle species (X). In this scenario, annihilation and the inverse annihilation process are possible. For example,

$$X, \bar{X} \rightarrow SM, SM \quad (2.35)$$

In this simple scenario, the Boltzmann equation are

$$\frac{dY_X}{dx} = \frac{\gamma^{sc}}{\mathcal{H} s} \left(1 - \left(\frac{Y_X}{Y_X^{eq}} \right)^2 \right) \quad (2.36)$$

where, γ^{sc} denotes the thermal average of the scattering cross section, $\mathcal{H}$ is the Hubble parameter, and s is the entropy density. For three different values of γ^{sc} , the corresponding evolution is shown in Fig. 2.1. For a large scattering rate, the species decouples later, at a larger value of x . In contrast, for smaller values of γ^{sc} , the particle X decouples from the thermal bath at earlier stages of the Universe. The freeze-out [178] value of this parameter,

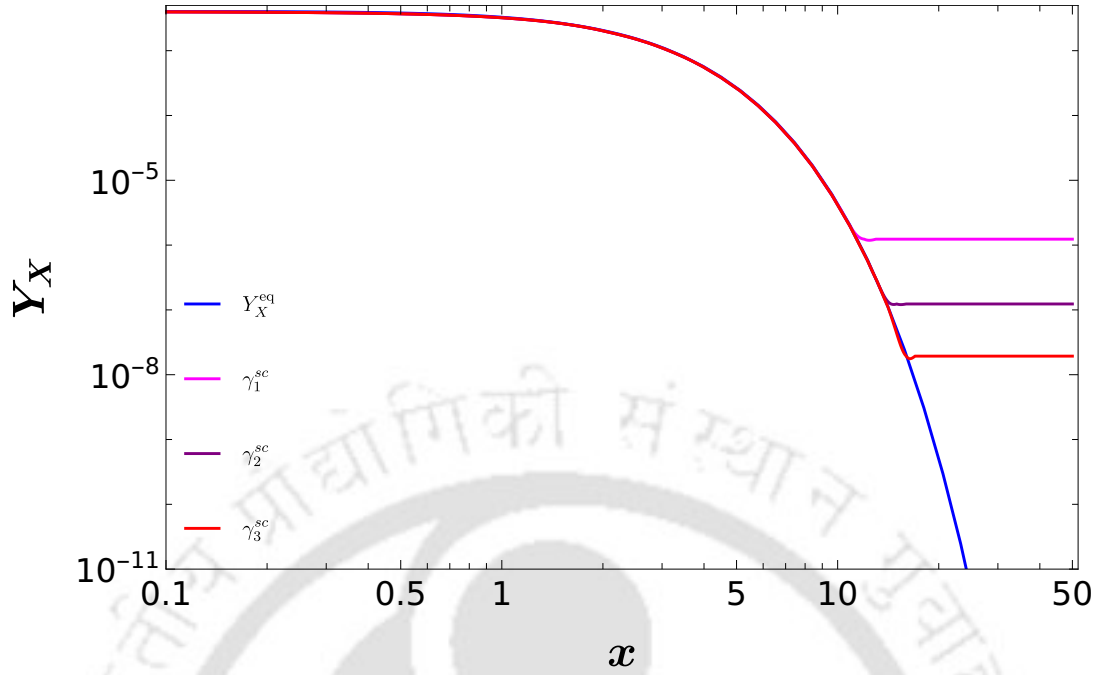


Figure 2.1: Freeze out the massive particle X species. The horizontal line actually represents the abundance after Freeze out. The thermal average of the scattering process γ^{sc} is increased from γ_1^{sc} to γ_3^{sc} .

x_f , can be calculated using Eq. 2.33. A species that decouples at the relativistic scale is often called a hot relic. The present relic density contribution by a hot relic of mass m_X is simply computed by using Eq. 2.30. The larger the annihilation cross section, the smaller the relic abundance. And the smaller the annihilation cross section, the larger the relic abundance. Interestingly, relic density X is inversely proportional to the annihilation cross section. The relation given in Eq. 2.33 makes it difficult to determine the exact value of the freeze-out x_f because the Hubble constant, entropy density, and thermal average of the scattering cross section all depend on x_f . The freeze-out point tells us the expansion rate, temperature, and entropy density of the universe at the time when DM decoupled from the thermal bath. The present relic density is directly proportional to the masses of the dark matter species.

2.2.1.2 Freeze In: Origin of Feebly Interacting Dark Matter

The freeze-in [179] mechanism admits many realizations, each characterized by distinct microscopic details and phenomenological signatures. Nevertheless, the underlying idea is simple and can be illustrated in a model-independent way. We consider a FIMP, denoted by X , which couples only very weakly to the thermal bath through a renormalizable

interaction at temperatures well above the weak scale. The mass of the heaviest particle at the interaction vertex is m_X , which typically assume to be close to the weak scale. There may be many particles from the thermal bath at the interaction vertex. The final amount of dark matter in freeze-in increases with smaller masses and weaker couplings. It is completely determined by physics at late times, not by conditions at very high temperatures. The standard freeze-out explanation is the opposite of this process, in which dark matter starts out in thermal equilibrium and then separates when annihilations stop working. Weaker interactions lead to a higher number of relics because annihilation happens sooner in freeze-out. Even though the two models depend on mass and coupling strength in different ways, they both say that the early Universe was warm and ran on low-energy processes.

For dark matter with masses near the weak scale, freeze-in can naturally replicate the observed relic density if the ratio of the weak scale to the Planck scale diminishes the interaction strength. In this way, freeze-in is like the parametric success of big particles with weak interactions, but it uses ultraweak couplings instead of electroweak ones.

Freeze-in models predict the presence of particles with exceptionally prolonged lifespans, often quantified in seconds. In numerous scenarios, the predominant abundance of dark matter arises from the disintegration of the lightest particle in the visible sector, which remains in thermal equilibrium. The appropriate relic density requires exceedingly small couplings, suggesting a high-energy physics origin in string theory, grand unification, or extradimensional frameworks.

Overall, freeze-in is a strong alternative to freeze-out. It could be tested through collider experiments, astrophysical signals, and cosmological observations. It also naturally links weak-scale dark matter masses to ultra-weak interactions that happen at very high energy levels.

At very high temperatures, we assume an initial X abundance is negligible. As the universe evolves, X can be produced through both decay and scattering processes. Let's take an example: the decay of a heavy particle A produces the DM particle X via the interaction $A \rightarrow XB$, where A is the heavy particle, which is in thermal equilibrium, and B is the SM particle. The evolution equation of the dark matter candidate X in this scenario can be written as

$$\frac{dY_X}{dx} = \frac{\gamma_A}{\mathcal{H}sx}, \quad (2.37)$$

Where γ_Y represents the thermal average of decay widths. The general solution of Eq. 2.37

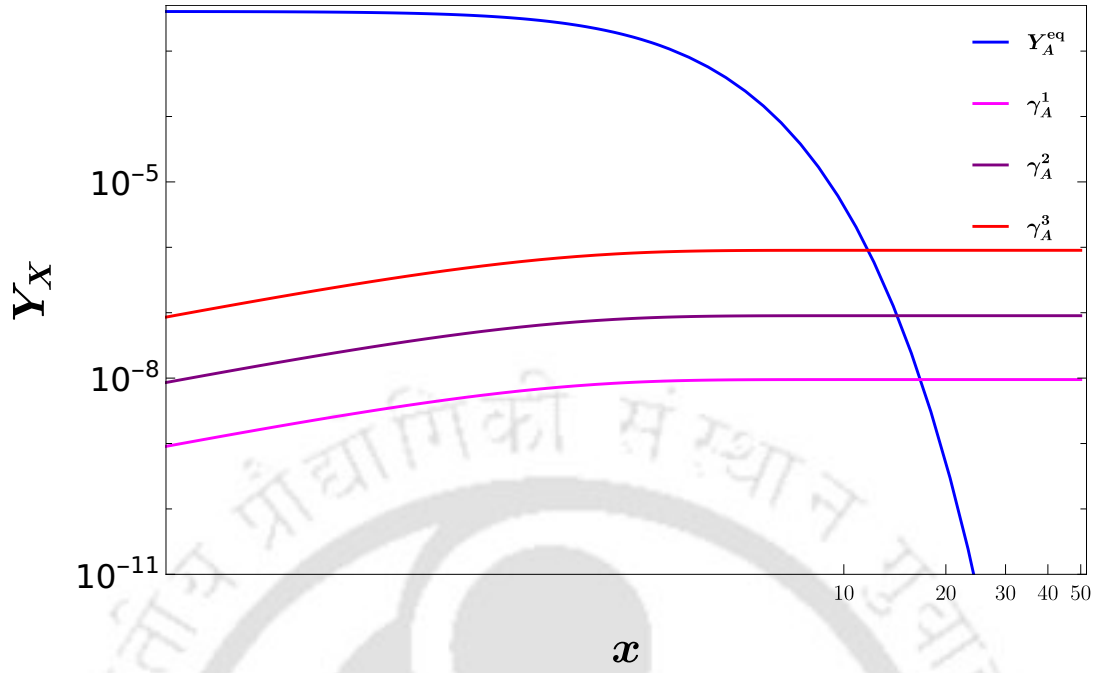


Figure 2.2: Freeze-In the massive particle X species. The thermal average γ_A of the decay/scattering process is increased from γ_A^1 to γ_A^3 .

is

$$Y_X = \int \frac{\gamma_A}{\mathcal{H}sx'} dx' + c \quad (2.38)$$

The evolution plot corresponding to Y_X vs x is shown in Fig. 2.2. It is clear that as the interaction rate increases, Y_X increases as well. In this scenario, the relic abundance of dark matter can be calculated using Eq. 2.30. There are so many possibilities of dark matter production. In the freeze-in scenario, the first possibility is that it can be produced via two-body decay, i.e., $A \rightarrow XB$, where A is the heavy particle, and B is the SM particle. It can also be produced via the $AB \rightarrow X$. Freezing in dark matter can also be produced by the $2 \rightarrow 2$ scattering process. Freeze-in production of dark matter is also possible via higher-order decays. In chapter 3, we will revisit the Freeze-in mechanism in detail.

2.3 Matter-Antimatter Asymmetry

The central question is how the asymmetry between matter and antimatter arose in the early Universe. The matter–antimatter asymmetry may either have been present from the very beginning of the universe or have emerged dynamically as the universe evolved. In the latter case, specific physical mechanisms would have been responsible for generating

the asymmetry. Any dynamical process capable of producing such an asymmetry during the evolution of the universe must satisfy a specific set of requirements, known as the **Sakharov conditions**:

(i) **Violation of Baryon/Lepton number conservation:** To generate a baryon/lepton asymmetry from a baryon-symmetric state of the Universe, certain interactions that contravene baryon/lepton number (B/L) conservation are necessary.

(ii) **Violation of C and CP symmetries:** A baryon number violation alone cannot cause baryon asymmetry. An asymmetry between matter and antimatter requires unequal decay and scattering rates. If interactions that break the baryon number still follow the CP symmetries, baryons and antibaryons would be made and destroyed at the same rate. As a result, there would be no net baryon asymmetry.

(iii) **Moving away from thermal equilibrium:** When a system is in chemical equilibrium, there can be no baryon asymmetry because the chemical potentials that are related to non-conserved quantum numbers go away. CPT symmetry also says that baryons and antibaryons must have the same mass. Thermal equilibrium, $f(p) = 1/(1 + \exp(\frac{p^2 + m^2}{T}))$, also means that baryons and antibaryons have the same phase-space distributions, which means that their number densities are the same, $n_B = n_{\bar{B}}$.

To demonstrate these criteria, **Charge symmetry** guarantees an equivalent reaction creating a baryon and antibaryon with the same probability, let us examine a straight forward example. Let's say that a heavy particle breaks down into lighter particles in a way that doesn't respect baryon or lepton number. So these process may generate the asymmetry in baryon or lepton sector. Means process must violate the charge conservation.

If the decay/scattering rate of this process is the same as that of its charge-conjugate process, any asymmetry that is created will be quickly erased. A net asymmetry may only occur if the decay rates are different, which means that CP symmetry must be broken. But CP violation by itself is not enough. If the system stays in thermal equilibrium, the decay will still create an asymmetry that the associated inverse processes will eliminate. To create a matter-antimatter imbalance, the forward and reverse reaction rates must be different. This means that the process must happen when the system is not in thermal equilibrium.

In conclusion, the creation of a cosmic asymmetry requires the simultaneous fulfillment of the Sakharov conditions. There are many models that satisfy the Sakharov conditions, which can be categorized by their underlying mechanisms, such as heavy particle decays, the electroweak phase transition, and other processes.

2.3.1 Models of baryogenesis

Numerous baryogenesis framework have been put forward. (Comprehensive reviews and specialized monographs on baryogenesis are available in [180–184]). Many fall into broader categories; the following is a selective overview:

2.3.1.1 Occurring during a first-order phase transition at electroweak symmetry breaking: Electroweak Baryogenesis

A first-order phase shift in the early universe around the electroweak scale could cause big changes in thermal equilibrium. This type of phase transition is fundamentally possible in numerous extended or alternative models, not confined to the SM. So, electroweak baryogenesis is seen as a general framework instead of a mechanism that only works with one model. Quantum anomalies in quantum field theory are what mostly cause baryon number violations. There are a lot of symmetries at the classical level that lead to conservation laws. Still, these same symmetries might not hold true when the theory is quantized. This breakdown is a natural result of how quantum theory is set up, not a mistake in the calculation. Chiral symmetry is a well-known example. For massless fermions, currents linked to chiral symmetry are conventionally considered conserved. However, this conservation is violated at the quantum level by the amalgamation of fermion fields in the presence of gauge fields. This is what "chiral anomaly" means. This peculiarity in the path-integral formalism arises from the integration measure of the fermion integral, rather than the Lagrangian. This measure does not stay the same when local chiral transformations happen, and this extra change is what breaks the conservation law. Consequently, a substantial disparity exists between quantum theory and classical symmetry. No specific gauge theory is solely associated with this chiral anomaly. When chiral fermions and non-Abelian gauge fields are present, global quantum numbers like baryon number and lepton number are not always conserved. Quantum processes can change the amounts of baryons and leptons, but some linear combinations, like baryon minus leptons, may still be kept. The existence of non-trivial configurations in the gauge field structure, including topological transitions, enhances the physical significance of these phenomena. In the early universe, when temperatures were high, these kinds of changes could happen on their own and actually change the baryon number. If a baryon-number violation happens at the same time as a first-order phase transition with a lot of CP-violation, the baryon asymmetry that comes out of it may stay the same as the universe changes. This process may pertain to various theories, including extensions of the SM, supplementary scalar fields, or altered gauge systems. Consequently, a comprehensive theoretical framework for baryogenesis is consti-

tuted by first-order phase transitions and quantum anomalies. This methodology is not confined to a singular model; it can be extensively applied to various theories, including gauge interactions, chiral fermions, and non-equilibrium dynamics in the primordial Universe. The SM [185], two-Higgs-doublet model [186], Minimal Supersymmetric SM [187], NMSSM [188], and nMSSM [188] are just a few of the many models that could cause matter-antimatter asymmetry. In conclusion, the SM [185] cannot create the necessary imbalance between matter and antimatter. Electroweak baryogenesis is now almost impossible in the MSSM [187]. The 2HDM [186], NMSSM [188], and nMSSM [189] can all make the necessary asymmetry. The main problems are dark matter and collider limits in the nMSSM, strict EDM and FCNC limits in the 2HDM, and changing parameters in the NMSSM.

2.3.2 Affleck–Dine Baryogenesis

The fundamental implementations of Grand Unified Theory (GUT) baryogenesis and electroweak baryogenesis face significant difficulties in elucidating the observable baryon asymmetry of the Universe. As a result, Affleck-Dine mechanism, have been looked at more closely. In this part, we talk about the main idea behind the Affleck-Dine mechanism. To accomplish this, we analyze a theoretical framework that incorporates a complex scalar field, whose dynamics are defined by a Lagrangian [190, 191] similar to that discussed as:

$$\mathcal{L}_\phi = (\partial^\mu \phi^\dagger)(\partial_\mu \phi) - m^2 \phi^\dagger \phi. \quad (2.39)$$

It is obvious that $\mathcal{L}_\phi$ is invariant under the phase redifination $\phi \rightarrow \phi' = e^{i\alpha} \phi$ with α being an arbitrary real constant. Under this redifination, conserved current is

$$J_B^\mu = i(\phi^\dagger \partial^\mu \phi - \phi \partial^\mu \phi^\dagger). \quad (2.40)$$

where the baryon number is identified as the global U(1) symmetry. It is significant to acknowledge that numerous supersymmetric theories incorporate scalar fields characterized by baryon and lepton numbers, exemplified by the supersymmetric counterparts of quarks and leptons [192]. The theory is CP-invariant because $\mathcal{L}_\phi$ is also invariant when you change ϕ to $\phi^\dagger$. Now, the interaction terms that Dine and Kusenko (2004) talked about are added.

$$\mathcal{L}_{\text{int}} = \lambda(\phi^\dagger \phi)^2 + [\epsilon \phi^\dagger \phi^3 + \delta \phi^4 + \text{h.c.}], \quad (2.41)$$

where the usual complicated numbers ϵ and δ are used. The ϵ and δ terms in Eq. (2.41) can thus be at odds with both the baryon number and CP invariance. In a uniform universe, the scalar field $\phi(x)$ changes according to

$$\frac{d^2\phi(t)}{dt^2} + 3H\frac{d\phi(t)}{dt} + \frac{\partial V(\phi)}{\partial\phi} = 0, \quad (2.42)$$

The scalar potential is represented by $V(\phi)$, and the Hubble parameter is represented by H . In the beginning, the scalar field was stuck at $\phi = \phi_0$, where $\phi_0 \gg 0$ was the expected value. This was because the condition $H \gg m$ was true. So, the initial baryon number must have been zero. Equation (2.42) demonstrates that the scalar field commenced oscillation when the Hubble parameter significantly decreased below m . When we take out the quartic term from the potential, we get $\phi(t) = \phi_0(mt)^{-3/2}\sin(mt)$ for the time when radiation was the most important force and $\phi(t) = \phi_0(mt)^{-1}\sin(mt)$ for the time when matter was the most important force. If we assume that ϕ_0 is real and that ϵ and δ are small coefficients, we can solve Eq. (2.42) to find the imaginary part of ϕ (Affleck and Dine, 1985):

$$\phi_I = a_r \frac{\text{Im}(\epsilon + \delta)\phi_0^3}{m^2(mt)^{3/4}} \sin(mt + \delta_r) \quad (2.43)$$

when radiation is the most important factor, with $a_r = 0.85$ and $\delta_r = -0.91$; and

$$\phi_I = a_m \frac{\text{Im}(\epsilon + \delta)\phi_0^3}{m^3t} \sin(mt + \delta_m) \quad (2.44)$$

Think about the matter-dominated case, where $a_m = 0.85$ and $\delta_m = 1.54$. By inserting Eq. (2.43) or Eq. (2.44) into the baryon-number current in Eq. (2.40), we get a baryon number that is not zero at large times:

$$n_B = 2a_r \frac{\text{Im}(\epsilon + \delta)\phi_0^2}{m^3t^2} \sin(\delta_r + \pi/8), \quad (\text{radiation}), \quad (2.45)$$

$$n_B = 2a_m \frac{\text{Im}(\epsilon + \delta)\phi_0^2}{m^3t^2} \sin(\delta_m), \quad (\text{matter}). \quad (2.46)$$

The baryon number in the scalar field will change into the baryon number of the SM particles through different decays of ϕ . For instance, the squark can break down into the SM quarks if it acts as the dynamic scalar field. The particulars of the inflationary model and the supersymmetric model dictate the final extent of the baryon number asymmetry [192, 190, 191].

Affleck-Dine baryogenesis is robust and efficient; however, its principal drawbacks in-

clude a significant reliance on supersymmetry, susceptibility to inflationary dynamics and initial conditions, and potential cosmological challenges such as iso curvature perturbations and Q-ball over production.

2.3.2.1 Arising from decays of heavy particles:

The decay of a heavy particle that satisfies all three Sakharov conditions can generate a matter–antimatter asymmetry. In fact, these conditions can be fulfilled by a wide range of models in high-energy physics. There are mainly two frameworks involved in this scenario:

(i) Grand Unified Theory (GUT) baryogenesis [193–198]; (ii) Leptogenesis [199];

GUT Baryogenesis The gauge groups $SU(3)_c \times SU(2)_L \times U(1)_Y$ are not unified in the SM due to the necessity of distinct coupling constants for each group. Georgi and Glashow (1974) presented the $SU(5)$ gauge group, marking the initial step towards unification. In 1975, Slansky presented the E_6 gauge group, while Fritzsch and Minkowski introduced the $SO(10)$ gauge group to explore subsequent advancements towards unification. Grand Unified Theories (GUTs) put quarks and leptons in the same multiplet, which is different from the SM.

$$\bar{5}_F = \begin{pmatrix} d_R^c \\ d_G^c \\ d_B^c \\ e_L \\ -\nu_L \end{pmatrix}, \quad 10_F = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & u_B^c & -u_G^c & u_R & d_R \\ -u_B^c & 0 & u_R^c & u_G & d_G \\ u_G^c & -u_R^c & 0 & u_B & d_B \\ -u_R & -u_G & -u_B & 0 & e^c \\ -d_R & -d_G & -d_B & -e^c & 0 \end{pmatrix}$$

The $\bar{5}_F$ representation has the right-handed down-type antiquarks d^c , the left-handed charged lepton e_L , and the left-handed neutrino ν_L . The antisymmetric tensor 10_F contains the right-handed charged lepton antiquark e^c , the right-handed up-type antiquarks u^c , and the left-handed quark doublet.

The breaking of the $SU(5)$ gauge symmetry happens in two steps. The SM gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$ is initially spontaneously broken from $SU(5)$. At lower energies, the electromagnetic symmetry $U(1)_Q$ is broken down even more to the electroweak symmetry $SU(2)_L \times U(1)_Y$.

This example shows that baryon number violation, which meets the first Sakharov condition, can happen on its own in a grand unified theory (GUT) model. The Kobayashi–Maskawa mechanism, which is similar to the SM [125], also causes CP violation. At the tree level, the decay rates of $X \rightarrow D + \ell$ and $X \rightarrow Q + U$ are the same

as those of their CP-conjugate processes. Nonetheless, loop-level corrections may induce CP-violating asymmetries. The third Sakharov condition, which says that the universe must not be in thermal equilibrium, is met if the heavy gauge bosons decay more slowly than the universe expands [178].

Baryogenesis based on GUTs eliminates any pre-existing asymmetry through the $(B + L)$ violating sphaleron process. To avert this problem, the $(B - L)$ symmetry must be violated in the GUT framework. This will stop the sphaleron process from erasing the asymmetry created, which will then be preserved [13] as the universe evolves.

2.3.3 Baryogenesis via Leptogenesis

Leptogenesis gives a simple and elegant explanation for the observed difference between matter and antimatter in the universe. One of its most interesting features is that it directly connects neutrinos' properties to the Universe's baryon asymmetry. In its simplest form, leptogenesis is driven by CP-violating interactions with the lightest heavy Majorana neutrino, which acts as a seesaw [200–202] partner for the SM neutrinos. The masses of both light and heavy neutrinos are constrained by the necessity of effective baryogenesis. In particular, all light neutrinos must have masses of less than about 0.1 eV.

The simplest extension of the SM with three right-handed neutrinos gives the tiny masses to the SM light neutrinos and can explain neutrino oscillations. Including the three right-handed neutrinos, the Lagrangian of our model becomes

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \sum_i \bar{N}_i i \gamma^\mu \partial_\mu N_i - \frac{1}{2} \sum_i m_{N_i} \bar{N}_i N_i - \left(\sum_i Y_N^{\alpha i} \bar{L}_\alpha \tilde{H} N_i + \text{h.c.} \right). \quad (2.47)$$

Also, the Type-I seesaw process can explain the mass of neutrinos by the presence of neutral fermions N_i . The mass matrix in the interaction basis is

$$\begin{pmatrix} 0 & m_D \\ m_D^T & m_N \end{pmatrix}, \quad (2.48)$$

In this 6×6 neutrino mass matrix, $m_D = v' Y_N$ represents the Dirac mass matrix, and m_N represents the Majorana mass matrix of the right-handed neutrinos, with $v' = \frac{v}{\sqrt{2}}$ and $v = 246$ GeV being the VEV of the Standard Model Higgs field. The SM light neutrino masses

$$m_\nu = m_D^T M^{-1} m_D, \quad (2.49)$$

The equation 2.49 shows how two Dirac mass insertions and the inverse of a heavy Majorana mass can make light neutrino masses. This implies that neutrinos are Majorana fermions, which naturally accounts for their lightness. Suppression is the most important part of the seesaw mechanism.

When we plug the Yukawa coupling into Eq. (2.49) for the Dirac mass matrix m_D , the light neutrino mass matrix becomes

$$m_\nu = Y_N^T M^{-1} Y_N v'^2. \quad (2.50)$$

For convenience, we define the new matrix κ that normalize by v'^2 . In general, the PMNS matrix can make κ diagonal, which it is:

$$\kappa = \frac{m_\nu}{v'^2}, \quad U^T \kappa U = \text{diag}(\kappa_1, \kappa_2, \kappa_3) \equiv D_\kappa, \quad (2.51)$$

The diagonal matrices D_κ in terms of Yukawa couplings becomes as:

$$D_\kappa = U^T Y_N^T D_M^{-1} Y_N U. \quad (2.52)$$

Where a diagonal mass matrix of D_M can be expressed as: $M = \text{diag}(m_{N_1}, m_{N_2}, m_{N_3}) \equiv D_M$. Rearranging the above equation in such a way that the combination of matrices satisfies the orthogonality condition. The simplest arrangement of the matrices is

$$1 = D_\kappa^{-1/2} U^T Y_N^T D_M^{-1} Y_N U D_\kappa^{-1/2}, \quad (2.53)$$

The Eq 2.53 can be rearrange as

$$1 = \left(D_M^{-1/2} Y_N U D_\kappa^{-1/2} \right)^T \cdot \left(D_M^{-1/2} Y_N U D_\kappa^{-1/2} \right). \quad (2.54)$$

This clearly meets the orthogonality condition $R^T R = 1$. So, we can find the orthogonal matrix by comparing Eq. (2.54) with $R^T R = 1$:

$$R = D_M^{-1/2} Y_N U D_\kappa^{-1/2}, \quad (2.55)$$

Solving Eq 2.55, for the Yukawa coupling matrix, we obtain

$$Y_N = D_M^{1/2} R D_\kappa^{1/2} U^\dagger. \quad (2.56)$$

This represents the most general parametrization of the neutrino Yukawa coupling

matrix Y_N . After substituting the value of U , R , D_M and D_κ in Eq 2.56, Y_N can be written as:

$$Y_N = \begin{pmatrix} a_1\sqrt{m_{N_1}}, a_2\sqrt{m_{N_1}}, a_3\sqrt{m_{N_1}} \\ b_1\sqrt{m_{N_2}}, b_2\sqrt{m_{N_2}}, b_3\sqrt{m_{N_2}} \\ c_1\sqrt{m_{N_3}}, c_2\sqrt{m_{N_3}}, c_3\sqrt{m_{N_3}} \end{pmatrix} \quad (2.57)$$

The parameters a_i , b_i , and c_i depend on the light neutrino masses, the elements of an arbitrary orthogonal matrix, and the elements of the PMNS matrix, where $i = 1, 2, 3$. These coefficients are explicitly defined in Eq. 2.58.

$$\begin{aligned} a_1 &= \sqrt{\kappa_1} R_{11} U_{11}^* + \sqrt{\kappa_2} R_{12} U_{12}^* + \sqrt{\kappa_3} R_{13} U_{13}^* \\ a_2 &= \sqrt{\kappa_1} R_{11} U_{21}^* + \sqrt{\kappa_2} R_{12} U_{22}^* + \sqrt{\kappa_3} R_{13} U_{23}^* \\ a_3 &= \sqrt{\kappa_1} R_{11} U_{31}^* + \sqrt{\kappa_2} R_{12} U_{32}^* + \sqrt{\kappa_3} R_{13} U_{33}^* \\ b_1 &= \sqrt{\kappa_1} R_{21} U_{11}^* + \sqrt{\kappa_2} R_{22} U_{12}^* + \sqrt{\kappa_3} R_{23} U_{13}^* \\ b_2 &= \sqrt{\kappa_1} R_{21} U_{21}^* + \sqrt{\kappa_2} R_{22} U_{22}^* + \sqrt{\kappa_3} R_{23} U_{23}^* \\ b_3 &= \sqrt{\kappa_1} R_{21} U_{31}^* + \sqrt{\kappa_2} R_{22} U_{32}^* + \sqrt{\kappa_3} R_{23} U_{33}^* \\ c_1 &= \sqrt{\kappa_1} R_{31} U_{11}^* + \sqrt{\kappa_2} R_{32} U_{12}^* + \sqrt{\kappa_3} R_{33} U_{13}^* \\ c_2 &= \sqrt{\kappa_1} R_{31} U_{21}^* + \sqrt{\kappa_2} R_{32} U_{22}^* + \sqrt{\kappa_3} R_{33} U_{23}^* \\ c_3 &= \sqrt{\kappa_1} R_{31} U_{31}^* + \sqrt{\kappa_2} R_{32} U_{32}^* + \sqrt{\kappa_3} R_{33} U_{33}^* \end{aligned} \quad (2.58)$$

The elements of the PMNS matrix are represented by U_{ij} , and the elements of the arbitrary orthogonal matrix are represented by R_{ij} . The complex and non-diagonal character of the Yukawa coupling is crucial for generating CP violation in the decay of N_1 . The next subsection 2.3.3.1 will talk about how to generate CP asymmetry.

2.3.3.1 The CP violating asymmetries parameter

Because direct CP violation can happen when decay amplitudes interfere with each other at the tree and loop levels, the difference between the decay rates $\Gamma(N_i \rightarrow \ell_\alpha + H)$ and $\Gamma(N_i \rightarrow \bar{\ell}_\alpha + H)$ goes away at the tree level but stays nonzero at the one-loop level. Fig. 2.3 shows the usual Feynman diagrams for $N_i \rightarrow \ell_\alpha + H$ decays. It is clear that the self-energy and vertex corrections are part of the one-loop contributions to the decay amplitudes.

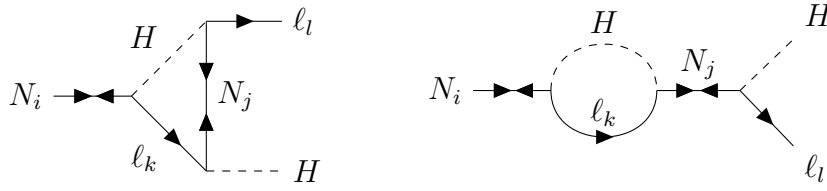


Figure 2.3: Feynman diagram corresponding to one-loop vertex and self-energy corrections to $N_i \rightarrow \ell_l H$.

The corresponding CP-violating asymmetries for the decays of N_i with each lepton flavor in the final states are characterized as

$$\epsilon_{i\alpha} \equiv \frac{\Gamma(N_i \rightarrow \ell_\alpha + H) - \Gamma(N_i \rightarrow \bar{\ell}_\alpha + H)}{\Gamma(N_i \rightarrow \ell_\alpha + H) + \Gamma(N_i \rightarrow \bar{\ell}_\alpha + H)}, \quad (2.59)$$

In this case, it is understood that the $SU(2)$ indices are added together. We have $\epsilon_i = \epsilon_{ie} + \epsilon_{i\mu} + \epsilon_{i\tau}$. It's easy to figure out N_i the overall decay rate at the tree level:

$$\Gamma_i \equiv \sum_{\alpha} \left[\Gamma(N_i \rightarrow \ell_\alpha + H) + \Gamma(N_i \rightarrow \bar{\ell}_\alpha + H) \right] = \frac{(Y_N^\dagger Y_N)_{ii}}{8\pi} m_{N_i}. \quad (2.60)$$

So, if there is no CP violation, $\Gamma(N_i \rightarrow \ell_\alpha + H) = \Gamma(N_i \rightarrow \bar{\ell}_\alpha + \bar{H})$.

In the initial examination of leptogenesis [200, 202], solely the vertex modifications were considered. Researchers found that the self-energy corrections are important [203, 204], and if the masses of heavy Majorana neutrinos are very close to each other [205], they may even be the most important ones. For CP violation to happen, the decay amplitudes must have at least two different weak phases and two different strong phases. If the lepton and Higgs doublets in the loops are on their mass shells, strong phases can happen. Weak phases, on the other hand, come from the complex Yukawa coupling matrix Y_N . This requirement is always met because N_i are heavy enough compared to the SM particles.

The CP violating asymmetries induced by the self-energy corrections:

$$\epsilon^s = \frac{1}{8\pi K_{11}} \sum_{j=2,3} \frac{m_{N_1} m_{N_j}}{m_{N_1}^2 - m_{N_j}^2} \text{Im}[(K_{ij})^2] \quad (2.61)$$

When the masses of two heavy Majorana neutrinos are the same, or $m_{N_i} = m_{N_j}$, the CP asymmetry parameter ϵ^s diverges. In this unusual case, we need to think about the heavy neutrinos N_j 's finite decay widths in their propagators. Even when it m_{N_i} gets very close to m_{N_j} , the CP violating asymmetries still work right because the decay widths Γ_j

act as regulators. When the mass differences $|m_{N_i} - m_{N_j}|$ are close to the decay widths Γ_i and Γ_j , the magnitudes of ϵ^s can be resonantly increased. This observation is the basis for the resonant leptogenesis mechanism [205, 206].

Next, we find the CP violating asymmetries that happen with the interferences of the one-loop vertex corrections and tree-level amplitude.

$$\epsilon^v = \frac{1}{8\pi K_{11}} \sum_{j=2,3} \text{Im}[K_{1j}^2] F_v \left(\frac{m_{N_j}^2}{m_{N_i}^2} \right) \quad (2.62)$$

Where the vertex function $F_v(m_{N_j}^2/m_{N_i}^2)$ is defined by $F_v(x) = \sqrt{x} [1 + (1+x)\text{In}(\frac{x}{1+x})]$. Overall CP violating asymmetry ϵ is simply a sum of the CP violating asymmetries induced by vertex and self-energy corrections

$$\epsilon = \epsilon^v + \epsilon^s \quad (2.63)$$

In type-I seesaw models, the aforementioned results are highly useful for examining flavor-dependent and flavor-independent leptogenesis scenarios. Heavy Majorana neutrinos N_i can generate a lepton-number asymmetry through their out-of-equilibrium decays that violate CP . The CP violating asymmetry ϵ_1 and the number density of heavy Majorana neutrinos N_i quantify the difference between the number densities of leptons Y_ℓ and antileptons $Y_{\bar{\ell}}$. That is to say, it $Y_\ell - Y_{\bar{\ell}}$ is directly related to ϵ_1 and $Y_{N_1} \equiv \frac{n_{N_1}}{s}$. Decay and inverse-decay processes, on the other hand, could change Y_{N_1} . Processes of decay and lepton-number-violating scattering can also change the values of Y_ℓ and $Y_{\bar{\ell}}$. At the loop level, the decay width N_1 is modified by the CP asymmetry parameter. Consequently, the decay rates of the forward process, the reverse process, and their CP conjugate processes are shifted according to the CP violating parameters.

$$\Gamma_{\vec{N}_1} = \frac{(1 + \epsilon_1)}{2} \Gamma_{N_1}, \quad \Gamma_{\vec{N}_1} = \frac{(1 - \epsilon_1)}{2} \Gamma_{N_1}, \quad \Gamma_{\vec{N}_1} = \frac{(1 + \epsilon_1)}{2} \Gamma_{N_1}, \quad \Gamma_{\vec{N}_1} = \frac{(1 - \epsilon_1)}{2} \Gamma_{N_1} \quad (2.64)$$

This parametrization is useful for deriving the evolution equation, which corresponds to the Boltzmann equation. To precisely compute the lepton and baryon number asymmetries, it is imperative to resolve the entire set of Boltzmann equations [2, 202, 207, 208].

2.3.3.2 Boltzmann equations

To study the asymmetries in the visible sector and the time evolution of the numerical densities of N_1 and L , we need to know how they change over time. When the system is not in thermal equilibrium, the Boltzmann equations tell us how the particle number densities change. In scenarios where the masses of the right-handed neutrinos (RHN) N_i satisfy the hierarchy $m_{N_3} \gg m_{N_2} \gg m_{N_1}$, any lepton asymmetry generated by the decays of the heavier neutrinos N_3 and N_2 at elevated temperatures is effectively nullified prior to the electroweak symmetry breaking (EWSB). Therefore, the disintegration of the lightest right-handed neutrinos, N_1 , constitutes the primary source of the lepton asymmetry that endures at lower temperatures. One of the ways that the number density of N_1 changes over time is through decay and scattering. In a comparable manner, processes that involve the creation (like N_1 decay and scattering processes) and destruction (like inverse decay and scattering processes) of leptons set their number density. enumerates the mechanisms relevant to the visible sector asymmetry. The interactions collectively dictate the quantity of N_1 and leptons contributing to the resultant lepton asymmetry. Sphaleron processes can subsequently transfer the produced lepton asymmetry into the detected baryon asymmetry of the Universe [209]. The relevant Boltzmann equations for deriving $Y_{N_1} = \frac{n_{N_1}}{s}$ and $Y_L = \frac{n_l - n_{\bar{l}}}{s}$ are as follows:

$$\mathcal{H}S z \frac{dY_{N_1}}{dz} = \left(1 - \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}}\right) (\gamma_D + 2\gamma_s + 2\gamma_{t_1} + 2\gamma_{t_2} + 2\gamma_V), \quad (2.65)$$

$$\begin{aligned} \mathcal{H}S z \frac{dY_L}{dz} = & \left(\frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} - 1\right) \epsilon \gamma_D - \frac{Y_L}{Y_L^{\text{eq}}} \left[\frac{\gamma_D}{2} + \gamma_{t_1} + \gamma_{t_2} + \gamma_{V_1} + \gamma_{V_3} + \gamma_{V_4} + \gamma_{V_5} \right. \\ & \left. + \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} (\gamma_s + \gamma_{V_2} + \gamma_{V_6}) \right], \end{aligned} \quad (2.66)$$

Where $\mathcal{H}$ denotes the Hubble constant, S is the entropy density, and γ_D is the thermally averaged decay rate for the process $N_1 \rightarrow \ell H$. For the processes $N_1 d \leftrightarrow u \ell$ and $N_1 \bar{d} \leftrightarrow \bar{u} \bar{\ell}$, γ_{t_1} is the thermal average, and γ_{t_2} is the same for the processes $N_1 d \leftrightarrow u \ell$ and $N_1 \bar{d} \leftrightarrow \bar{u} \bar{\ell}$. The sum of these contributions is $\gamma_t = \gamma_{t_1} + \gamma_{t_2}$. The s -channel processes $\ell N_1 \leftrightarrow d \bar{u}$ and $\bar{\ell} N_1 \leftrightarrow \bar{d} u$ are also represented by γ_s , while γ_{V_i} , where $i = 1, \dots, 6$, are the thermal averaged cross sections of scattering processes involving the gauge bosons: $V_\mu \ell$, $N_1 H \xleftrightarrow{s} N_1 \ell \leftrightarrow V_\mu H$, $N_1 V_\mu \xleftrightarrow{t} H \ell$, $N_1 V_\mu \xleftrightarrow{u} H \ell$, $N_1 H^\dagger \xleftrightarrow{t} V_\mu \ell$, $N_1 \ell \xleftrightarrow{t} V_\mu H^\dagger$, and their charged conjugate processes in that order, with $V_\mu = W_\mu, Z_\mu$.

The thermal average of the decay width of the process $N_1 \rightarrow \ell_\alpha H$ is given as

$$\begin{aligned}
 \gamma(N_i \rightarrow \ell_\alpha H) &= \int d\Pi_{N_i} e^{-E_{N_i}/T} \int d\Pi_{\ell_\alpha} d\Pi_H (2\pi)^4 \\
 &\quad \delta^{(4)}(p_{N_i} - p_{\ell_\alpha} - p_H) |\mathcal{M}(N_i \rightarrow \ell_\alpha + H)|^2 \\
 &= g_{N_i} \int \frac{d^3 p_N}{(2\pi)^3} e^{-E_{N_i}/T} \frac{M_i}{E_{N_i}} \Gamma(N_i \rightarrow \ell_\alpha + H) \\
 &= n_{N_i}^{\text{eq}} \frac{K_1(z)}{K_2(z)} \Gamma(N_i \rightarrow \ell_\alpha + H), \tag{2.67}
 \end{aligned}$$

Where $n_{N_1}^{\text{eq}} = g_{N_1} \frac{m_{N_1}^3}{2\pi^2} \frac{K_2(z)}{z}$ is the equilibrium number density of N_1 , $K_1(z) \equiv \frac{1}{z} \int_z^\infty dy e^{-y} y \sqrt{y^2 - z^2}$, and $K_2(z) \equiv \frac{1}{z^2} \int_z^\infty dy e^{-y} y^2 \sqrt{y^2 - z^2}$ are the modified Bessel functions. And the thermal average of scattering process $a, b \rightarrow c, d$ are For a $2 \rightarrow 2$ scattering process, where a and b are the incoming particles, the thermally averaged cross section γ_{ab} is given by

$$\gamma_{ab} = \frac{T}{64\pi^4} \int_{s_{\min}}^\infty ds 2s^{\frac{3}{2}} \lambda(1, \frac{m_a^2}{s}, \frac{m_b^2}{s}) \sigma_{ab} K_1\left(\frac{\sqrt{s}}{T}\right), \tag{2.68}$$

where T is the temperature, $\sqrt{s}$ is the center-of-mass energy, and K_1 is the modified Bessel function of the second kind of order one. The Källén function $\lambda(x, y, z)$ is given by

$$\lambda(x, y, z) = x^2 + y^2 + z^2 - 2(xy + yz + zx). \tag{2.69}$$

To ensure the kinematic feasibility of the scattering process, the lower limit of integration is

$$s_{\min} = \max[(m_{a_1} + m_{a_2})^2, (m_{a_3} + m_{a_4})^2], \tag{2.70}$$

The Boltzmann Eq. 2.65 and Eq. 2.66 given above fully describe how the system evolves and how the number density of leptons changes. The detailed study of leptogenesis is available in the next section. First, we will look at the analytical solution, and then we will come to the numerical analysis.

2.3.3.3 Analytic Solution

In the previous section, we gave the Boltzmann equations for the number density distribution functions for neutrinos, leptons, and antileptons in the model. We examine the alteration in the phase-space distribution of the lepton sectors to evaluate the asymmetry. N_1 has an effect on the number density of leptons. We solve the Boltzmann equation for

N_1 and the lepton sectors. Eqs. 2.65 and Eqs. 2.66 can be rewritten as

$$\frac{dY_{N_1}}{dz} + P_N(z)Y_{N_1} = Q_N(z), \quad (2.71)$$

$$\frac{dY_L}{dz} + P_L(z)Y_L = Q_L(z). \quad (2.72)$$

Where $P_N = \frac{\gamma_D + 2\gamma_s + 2\gamma_{t_1} + 2\gamma_{t_2} + 2\gamma_V}{\mathcal{H}SZ \frac{Y_{N_1}^{eq}}{Y_{N_1}^{eq}}}$, $Q_N = P_N Y_{N_1}^{eq}$, $P_L = \frac{1}{\mathcal{H}SZ \frac{Y_{N_1}^{eq}}{Y_{N_1}^{eq}}} \left[\frac{\gamma_{D_1}}{2} + \gamma_{t_1} + \gamma_{t_2} + \gamma_{V_1} + \gamma_{V_3} + \gamma_{V_4} + \gamma_{V_5} + \frac{Y_{N_1}}{Y_{N_1}^{eq}} (\gamma_s + \gamma_{V_2} + \gamma_{V_6}) \right]$ and $Q_L = \left(\frac{Y_{N_1}}{Y_{N_1}^{eq}} - 1 \right) \frac{\epsilon \gamma_D}{\mathcal{H}SZ}$ are the functions that incorporate information related to the evolution of N_1 and L . The number density of the daughter particles is affected by the rate at which the number density of parent particle per unit entropy changes with energy, as shown in Eq. 2.65. The number density of Y_L at that energy level stays the same in equilibrium, which means $Y_{N_1} = Y_{N_1}^{eq}$. This section analytically examines the impact of the coefficient multiplying $(Y_{N_1} - Y_{N_1}^{eq})$, which acts as a conversion factor into the leptonic sector, on the variation in lepton number density, as detailed in Eq. 2.66. This conversion factor also changes with energy. If the conversion factor is high at a certain energy, a bigger matter antimatter asymmetry can happen when Y_{N_1} is not equal to $Y_{N_1}^{eq}$. The last term in Eq. 2.66 shows the washout contribution. The following is a discussion of the analytical solutions to Eq. 2.71 and Eq. 2.72.

The analytic solution of the above equation 2.71, and 2.72 are given below

$$Y_{N_1} = \mu^{-1}(z) \left(\int \mu(z') Q_N(z') dz' + C_1 \right), \quad (2.73)$$

$$Y_L = \epsilon \kappa, \quad \text{Where} \quad (2.74)$$

$$\kappa = \frac{1}{\nu(z')} \left(\int \nu(z') Q_L(z') dz' + C_2 \right) \quad (2.75)$$

Where C_1 and C_2 are integration constants. They can be determined by imposing the initial conditions $Y_{N_1} = Y_{N_1}^{eq}$ and $Y_L = 0$ at $z = z_0$. The efficiency factor is represented by κ , and $\mu(z) = \exp(\int P_N(z') dz')$ and $\nu(z) = \exp(\int P_L(z') dz')$.

In this framework, there are two important variables that determine Y_L . The first is the difference between the heavy particle abundance and its equilibrium value, which is shown by the equation: $\delta = Y_N - Y_N^{eq}$. The interaction rate, which includes both scattering and decay processes, is second. The decay rate is the main thing that makes Y_N leave equilibrium. The maximum amount of asymmetry is reached when the product $\epsilon \delta$ is at its highest.

Further, based on the rate of the scattering process, it can be divided into two regimes: one weak washout regime and a second strong washout regime. In the weak-washout regime, scattering processes don't have a significant effect on asymmetry. Conversely, scattering mechanisms generally mitigate the asymmetry induced by decays in the strong-washout regime. The next part presents a numerical study demonstrating these effects.

2.3.3.4 Davidson–Ibarra Bound

The lower bound on masses of right handed neutrino to get required matter antimatter asymmetry has been explored by Davidson-Ibara [210]. The section 2.3.3.3 shows from Eq. 2.74 that the lepton asymmetry Y_L depends on the CP asymmetry parameter ϵ and the efficiency factor κ . The CP asymmetry parameter, which arises from the decay of the lightest right-handed neutrino N_1 , is discussed in Section 2.3.3.1, depends on the masses of the right-handed neutrinos and the Yukawa coupling Y_N .

To determine Y_N , we employ the Casas-Ibarra (CI) parametrization. In this framework, the Yukawa coupling Y_N depends on the masses of the right-handed neutrinos, the PMNS matrix, the light neutrino masses, and an arbitrary complex orthogonal matrix R . The central question is which values of ϵ and κ reproduce the observed value of the matter-antimatter asymmetry, i.e., the correct value of Y_L .

Numerical analysis shows that, within the CI parametrization, the efficiency factor κ is effectively independent of the right-handed neutrino masses. Moreover, Y_L reaches saturation for a specific value of κ . In both the strong washout regimes, this value is found to be of order

$$\kappa \simeq (2-4) \times 10^{-3}.$$

In what follows, we take $\kappa = 2.4 \times 10^{-3}$ for numerical estimates. The CP asymmetry parameters ϵ^s , ϵ^v , and ϵ are given in Eqs. 2.61, 2.62, and 2.63, respectively. Under the CI parametrization, the quantities K_{ij} and $(K_{1j})^2$ take the form

$$K_{ij} = \sqrt{m_{N_i} m_{N_j}} \sum_{k=1}^3 \kappa_k R_{ik}^* R_{jk}, \quad (2.76)$$

$$(K_{1j})^2 = m_{N_1} m_{N_j} \left(\sum_{k=1}^3 \kappa_k R_{1k}^* R_{jk} \right)^2. \quad (2.77)$$

Assuming a hierarchical spectrum of right-handed neutrino masses, $m_{N_j} \gg m_{N_1}$, the

CP asymmetry parameters ϵ^s , ϵ^v , and ϵ are bounded from above as

$$\epsilon^s \leq \frac{m_{N_1}}{8\pi} (\kappa_3 - \kappa_1) , \quad (2.78)$$

$$\epsilon^v \leq \frac{m_{N_1}}{16\pi} (\kappa_3 - \kappa_1) , \quad (2.79)$$

$$\epsilon \leq \frac{3m_{N_1}}{16\pi} (\kappa_3 - \kappa_1) . \quad (2.80)$$

Here, $\kappa_1 = m_{\nu_1}/v^2$ and $\kappa_3 = m_{\nu_3}/v^2$.

The lepton asymmetry is given by

$$Y_L = \epsilon \kappa . \quad (2.81)$$

Requiring Y_L to match the observed value,

$$1.72 \times 10^{-9} = \epsilon^{\max} \times 2.4 \times 10^{-3} , \quad (2.82)$$

we obtain

$$\epsilon^{\max} = 7.17 \times 10^{-7} . \quad (2.83)$$

Using Eq. 2.83 together with the upper bound on ϵ , we derive a lower bound on the mass of the lightest right-handed neutrino,

$$m_{N_1} \approx 1.2 \times 10^{-5} \frac{v^2}{\sqrt{\Delta m_{\text{atm}}^2}} \text{ GeV} . \quad (2.84)$$

The atmospheric mass-squared difference is $\sqrt{\Delta m_{\text{atm}}^2} \simeq 5 \times 10^{-11} \text{ GeV}$ [211]. This leads to the numerical lower bound $m_{N_1} \approx 7.3 \times 10^9 \text{ GeV}$. It is important to emphasize that this bound applies to the strong washout regime. In the weak washout regime, the value of κ changes, and consequently, the lower bound on the right-handed neutrino mass will be modified accordingly.

2.3.3.5 Numerical Analysis

As discussed in the previous subsection, the CP asymmetry parameter and the deviation Y_{N_1} from its equilibrium value $Y_{N_1}^{\text{eq}}$ are the key factors that determine the amount of asymmetry generated through decay processes. The left panel of Fig. 2.4 illustrates the deviation of Y_{N_1} from its equilibrium distribution $Y_{N_1}^{\text{eq}}$.

In this analysis, we consider four benchmark values for the mass of the lightest right-

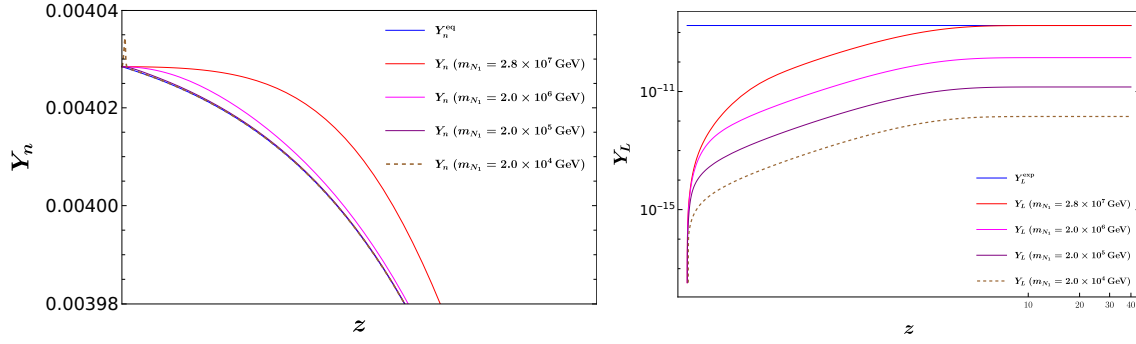


Figure 2.4: Evolution plots of Y_{N_1} and Y_L vs z for different numerical value of m_{N_1} .

m_{N_1} GeV	ϵ
2×10^4	$\approx 10^{-10}$
2×10^5	$\approx 10^{-9}$
2×10^6	$\approx 10^{-8}$
2.8×10^7	$\approx 10^{-7}$

Table 2.1: Values of the CP asymmetry parameter ϵ for different choices of the lightest right-handed neutrino mass m_{N_1} .

handed neutrino, namely $m_{N_1} = \{2.0 \times 10^4, 2.0 \times 10^5, 2.0 \times 10^6, 2.8 \times 10^7\}$ GeV. The masses of the heavier right-handed neutrinos are fixed at $m_{N_2} = 10^8$ GeV and $m_{N_3} = 2 \times 10^8$ GeV.

As the mass of the right-handed neutrinos increases, the deviation of Y_{N_1} from its equilibrium value also becomes more pronounced. When m_{N_1} reaches 2.8×10^7 GeV, the deviation of Y_{N_1} is sufficiently large to generate the required matter–antimatter asymmetry. This indicates that $m_{N_1} = 2.8 \times 10^7$ GeV represents the lower bound on the mass of the lightest right-handed neutrino in the weak washout regime. In contrast, in the strong washout regime, this lower bound is expected to increase.

For each value of the parameter m_{N_1} , the evolution of Y_L as a function of z is shown in the right panel of Fig. 2.4. The corresponding numerical values of the CP asymmetry parameter for each choice of m_{N_1} are listed in Table 2.1. In the weak washout regime, the parameter set $\epsilon = 10^{-7}$, $m_{N_1} = 2.8 \times 10^7$ GeV yields the required matter-antimatter asymmetry.

2.4 *Leptogenesis through the Scattering/annihilation process*

In this section, we explore leptogenesis via scattering processes. As discussed earlier, some scattering processes in conventional seesaw models can lead to lepton-number violation,

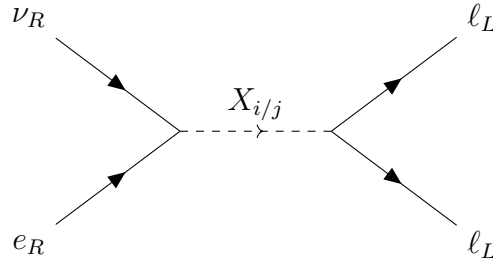


Figure 2.5: Feynman diagram for a two-to-two scattering process, where $i \neq j$.

typically acting as washout processes. To study leptogenesis through scattering in more detail, we consider a new model [212, 213], which provides a simple extension of the SM. This extension introduces two complex scalars, $X_{1,2} \equiv X_{\bar{1},2}$, and three right-handed neutrinos, ν_R . Within this framework, the Lagrangian of the model is given by:

$$\begin{aligned} \mathcal{L} = \mathcal{L}_{\text{SM}} &+ \overline{\nu_R^\alpha} i \not{\partial} \nu_R^\alpha + (D_\mu X_i)^\dagger (D^\mu X_i) - m_{X_i}^2 X_i^\dagger X_i \\ &- \left(\frac{1}{2} F_i^{\alpha\beta} \overline{(\ell_L^\alpha)^c} i \sigma_2 \ell_L^\beta X_i^\dagger + G_i^{\alpha\beta} \overline{(e_R^\alpha)^c} \nu_R^\beta X_i + \text{h.c.} \right) \end{aligned} \quad (2.85)$$

In this model, we do not discuss the mechanism by which neutrinos acquire their masses. Here, we focus solely on an example of leptogenesis through scattering processes. The interaction term $F_i^{\alpha\beta} \overline{(\ell_L^\alpha)^c} i \sigma_2 \ell_L^\beta X_i^\dagger$ violates lepton number by $\Delta L = 2$, while all other terms conserve lepton number. We do not present a full analysis of leptogenesis in this model; rather, we illustrate how matter–antimatter asymmetry can be generated specifically through the scattering process.

2.4.1 CP Asymmetry from $2 \rightarrow 2$ scattering

The scattering process $\nu_R e_R \rightarrow \nu_L e_L$ is mediated by the charged scalars $X_{1,2}$, as shown in Figure 2.5, and violates CP symmetry. We assume a hierarchical mass spectrum for the scalar particles, with $m_{X_1} \ll m_{X_2}$. In the center-of-mass frame, the CP asymmetry generated by this scattering process is given in Eq. 2.86. This asymmetry depends on the coupling constants and the denominators of the propagators; in other words, it is sensitive to the center-of-mass energy of the process.

$$\epsilon_{sc} = \frac{4 \text{Im}[F_2^* G_2^* F_1 G_1] \text{Im}[D_1^* D_2]}{|F_1|^2 |G_1|^2 |D_2|^2} \quad (2.86)$$

where, $D_i = s - M_{X_i^-}^2 + i m_i \Gamma_i$ denotes the denominator of the propagator. We will discuss the CP asymmetry generated by scattering processes in more detail in the upcoming chapter. In this simple scenario, this scattering process acts solely as a source term for the matter-antimatter asymmetry. To study the evolution of the lepton number, we will analyze the corresponding Boltzmann equations in the following subsection.

2.4.2 Time Evolution of Number Density of ν_R and $L = \ell - \bar{\ell}$

As discussed above, the scattering process violates CP symmetry and contributes in two distinct ways: first, as a source term generating the asymmetry, and second, as a washout term reducing it. While the SM electrons remain in thermal equilibrium, the right-handed neutrinos are not. The Boltzmann equations describing the evolution of the right-handed neutrino abundance ν_R and the net lepton number $L = \ell - \bar{\ell}$ for the simple scattering process $\nu_R e_R \rightarrow \nu_L e_L$ are:

$$\mathcal{H} s z \frac{dY_{\nu_R}}{dz} = \left(1 - \frac{Y_{\nu_R}}{Y_{\nu_R}^{eq}}\right) \gamma_{\nu e}^{LL}, \quad (2.87)$$

$$\mathcal{H} s z \frac{dY_L}{dz} = \left(1 + \frac{Y_{\nu_R}}{Y_{\nu_R}^{eq}}\right) \Theta_{\nu e}^{LL} - \frac{Y_L}{Y_L^{eq}} \gamma_{\nu e}^{LL}. \quad (2.88)$$

where, $\gamma_{\nu e}^{LL}$ and $\Theta_{\nu e}^{LL}$ denote the thermal averages defined in Eqs. 2.89 and 2.90, respectively, and λ is the Källén function given in Eq. 2.69. Initially, the abundance Y_{ν_R} remains in thermal equilibrium. The scattering rate $\sigma_{\nu_R e_R}^{\nu_L e_L}$ determines the extent to which ν_R departs from its equilibrium abundance $Y_{\nu_R}^{eq}$. This departure, in turn, affects the net lepton number density $L = \ell - \bar{\ell}$.

$$\gamma_{\nu e}^{LL} = \frac{T}{32\pi^4} \int_{s_{min}}^{\infty} ds s^{\frac{3}{2}} \sigma \lambda[1, \frac{m_\nu^2}{s}, \frac{m_e^2}{s}], \quad (2.89)$$

$$\Theta_{\nu e}^{LL} = \frac{T}{32\pi^4} \int_{s_{min}}^{\infty} ds s^{\frac{3}{2}} \epsilon_{sc} \sigma \lambda[1, \frac{m_\nu^2}{s}, \frac{m_e^2}{s}], \quad (2.90)$$

The thermal average of the scattering cross section σ is denoted by $\gamma_{\nu e}^{LL}$, while $\Theta_{\nu e}^{LL}$ represents the thermal average of $\epsilon_{sc} \sigma$. The quantity $\Theta_{\nu e}^{LL}$ acts as the source term responsible for generating the matter-antimatter asymmetry, whereas it $\gamma_{\nu e}^{LL}$ governs the washout rate. The analytic solutions of Eqs. 2.87 and 2.88 are presented below.

$$\begin{aligned}
Y_{\nu_R}(z) &= \exp \left[- \int_{z_i}^z dz' \frac{\gamma_{\nu e}^{LL}(z')}{\mathcal{H}(z') s(z') z' Y_{\nu_R}^{eq}(z')} \right] \left[Y_{\nu_R}(z_i) \right. \\
&\quad \left. + \int_{z_i}^z dz' \frac{\gamma_{\nu e}^{LL}(z')}{\mathcal{H}(z') s(z') z'} \exp \left(\int_{z_i}^{z'} dz'' \frac{\gamma_{\nu e}^{LL}(z'')}{\mathcal{H}(z'') s(z'') z'' Y_{\nu_R}^{eq}(z'')} \right) \right], \quad (2.91) \\
Y_L(z) &= \exp \left[- \int_{z_i}^z dz' \frac{\gamma_{\nu e}^{LL}(z')}{\mathcal{H}(z') s(z') z' Y_L^{eq}(z')} \right] \left[Y_L(z_i) \right. \\
&\quad \left. + \int_{z_i}^z dz' \frac{\Theta_{\nu e}^{LL}(z')}{\mathcal{H}(z') s(z') z'} \left(1 + \frac{Y_{\nu_R}(z')}{Y_{\nu_R}^{eq}(z')} \right) \exp \left(\int_{z_i}^{z'} dz'' \frac{\gamma_{\nu e}^{LL}(z'')}{\mathcal{H}(z'') s(z'') z'' Y_L^{eq}(z'')} \right) \right]. \quad (2.92)
\end{aligned}$$

In this framework, we study the lepton asymmetry generated through scattering processes. The resulting leptonic asymmetry $Y_L(z)$ depends on three key factors: first, the deviation of the right-handed neutrino abundance Y_{ν_R} from its equilibrium value; second, the CP asymmetry parameter; and third, the scattering cross section of the relevant process. In our work, we analyze in detail how each of these ingredients influences the generation of the lepton asymmetry through scattering.

2.5 Additional Framework Development for BAU and DM Issue Resolution

There are instances in which the origins of BAU and DM can be examined within a single framework, despite their potential lack of correlation. In this thesis, we investigate certain minimal extensions of the SM that can concurrently generate the reported amount of dark matter remnant and elucidate the observed baryon asymmetry through leptogenesis. The vanilla leptogenesis scenario, which is extensively studied in relation to canonical seesaw mechanisms of neutrino mass, posits that the scale of leptogenesis must be exceedingly high and beyond direct experimental accessibility, as established by the Davidson-Ibarra (DI) constraint [210]. In this thesis, we examine the potential for diminishing the scale of leptogenesis within a unified framework of dark matter and leptogenesis by: (a) Leptogenesis From a Feebly Interacting Dark Matter Sector [214]; (b) Dark Matter-Driven Low-Scale Leptogenesis via Neutrino Portal [215]; We show that some of these scenarios, or a mix of them, can lower the leptogenesis scale to a TeV, which would have interesting effects on both dark matter and Leptogenesis. In our scenario we developed the framework

that provide the direct connection between leptogenesis and dark matter.

2.6 Layout of the thesis

The layout of the subsequent chapters is as follows.

In **Chapter 3** we present a novel mechanism for low-scale leptogenesis in a minimally extended SM framework, incorporating three right-handed Majorana neutrinos, a charged scalar S^+ , a charged fermion χ^+ , and a neutral fermion DM candidate ψ . This setup explores two-body and four-body decay FIMP dark matter scenarios, where ψ interacts feebly with χ^+ and S^+ to match the observed DM relic density. Unlike conventional thermal leptogenesis, which requires high neutrino masses, mass degeneracy in right-handed neutrinos enables low-scale leptogenesis. Our model achieves both DM production via freeze-in and adequate CP violation for leptogenesis with TeV-scale right-handed neutrinos, without relying on resonant leptogenesis. This framework also satisfies the observed DM relic density in the sub-TeV mass regime.

Following this in **Chapter 4** we present a novel low-scale leptogenesis framework within a minimal SM extension, featuring three right-handed neutrinos, a neutral fermion, and a real scalar. The model generates CP asymmetry through two-body decays of right-handed neutrinos and dark sector scattering processes mediated by RHNs, which also transfer asymmetry from the dark sector to the visible sector. This leads to enhanced CP violation and successfully accounts for the observed baryon asymmetry with right-handed neutrino masses around 2 TeV. Our framework links leptogenesis and dark matter, predicting new dark-sector states within the LHC's reach and offering a testable connection between cosmology and collider physics.

Finally in **Chapter 5** we offer the main summary and conclusion of our Thesis with a brief outline of the future direction in which our work can further be extended.



Leptogenesis From a Feebly Interacting Dark Matter Sector

In this chapter, we present a study of leptogenesis from a feebly interacting dark matter sector within a minimal extension of the Standard Model. The model consists of two additional fermions, a charged state χ and a neutral state ψ , three right-handed neutrinos, and a charged gauge-singlet scalar S . The feeble interactions of the dark sector yield a relic abundance consistent with observational constraints, while the decay of right-handed neutrinos into leptons and S , together with virtual S -mediated contributions, generates an enhanced CP asymmetry. This framework successfully accounts for the observed baryon asymmetry of the Universe for right-handed neutrino masses without requiring the neutrino mass degeneracy. The work presented in this Chapter is published in Phys. Rev. D [214].

3.1 Introduction

The introduction related Leptogenesis has been quite extensively presented in Chapter 1 and Chapter 2. However, to make the Chapter self-contained here we provide a brief overview of Leptogenesis using the FIMP.

The origin of the baryon asymmetry of the Universe (BAU) and identifying a suitable dark matter (DM) candidate are two of the major challenges in the particle Physics. The BAU, defined by the ratio of the difference in baryon (n_B) and antibaryon ($n_{\bar{B}}$) number densities to photon density, is measured to be $\eta = (6.09 \pm 0.06) \times 10^{-10}$ [216, 217]. Sakharov's conditions [218] indicate that generating the BAU requires baryon number violation, CP violation, and out-of-equilibrium conditions, often modeled through elec-

troweak phase transitions [219, 220, 180]. An alternative mechanism, leptogenesis [200], transfers lepton number asymmetry to baryon asymmetry via sphaleron processes. Leptogenesis models often involve heavy Majorana neutrinos, with CP violation arising from their decay [200, 221, 202, 204, 222]. While standard models require these neutrinos to be heavier than 10^9 GeV [208, 223], alternative frameworks suggest low-scale leptogenesis [224]. Regarding DM, evidence for its existence is strong, but its nature is not well understood [225–228]. The most common DM candidates are weakly interacting massive particles (WIMPs), but they face challenges in satisfying both relic density and direct detection constraints [229, 230]. As an alternative, feebly interacting massive particles (FIMPs) have been proposed [231, 232], where the DM abundance is generated through the freeze-in mechanism. This approach naturally avoids direct detection constraints due to weak interactions [233–235]. Recent studies have combined FIMPs with low-scale leptogenesis [236, 232], showing that TeV-scale Majorana neutrinos can successfully generate BAU while satisfying DM constraints. This model, which incorporates Yukawa couplings for lepton flavor violation, provides a consistent FIMP DM scenario and low-scale leptogenesis without fine-tuning [237].

In this Chapter we present the leptogenesis, which is directly connected to single-component FIMP dark matter. In addition to the seesaw interaction, right-handed neutrinos also interact with the charge scalar and charged leptons. This additional interaction provides an extra phase to the CP asymmetry parameters, helping to obtain the required matter-antimatter asymmetry at a low scale.

Our organization of the subsequent part of the Chapter is as follows. We discuss the model in Section 3.2. In Sec. 3.3, we give a short explanation of how dark matter works in the FIMP mechanism. The framework also considers the possibility of achieving the correct relic density via two- or four-body decays in sec 3.3.1, and sec 3.3.2. Section 3.4 discusses how to make leptogenesis work in the suggested scenario, and Section 3.4.1 shows how the new scalar degree of freedom affects getting the right CP imbalance. In 3.4.2, we provide clear Boltzmann equations for how the number density changes over time when it is not in thermal equilibrium. Section 3.5 contains in-depth research and discussion about the evolution of the number density of right-handed neutrinos and the number density of leptons. In Section 3.5.0.5, we compare our findings with the standard scenario result for same set of parameters. Finally we conclude this Chapter in Sec. 3.6.

Field	Spin	Y	Z_2
S^+	0	+2	+
N_1, N_2, N_3	$\frac{1}{2}$	0	+
χ^+	$\frac{1}{2}$	+2	−
ψ	$\frac{1}{2}$	0	−

Table 3.1: Additional fields in the model, together with their hypercharges and Z_2 charges. All fields are $SU(2)_L$ singlets. The neutrinos N_i ($i = 1, 2, 3$) are Majorana fermions, χ and ψ are vector-like fermions, and S is a scalar.

3.2 A MODEL WITH FIMP DARK MATTER

We thus revisit the extended model with heavy neutrinos and additional particles added to the spectrum, as in [237]. The extension to the SM particle content includes a gauge singlet charged scalar field S^+ , plus one charged (χ^+) and one neutral (ψ) singlet fermions, interacting feebly with the SM particle content. With an additional Z_2 symmetry under which both χ^+ and ψ are odd, while all other particles even, ψ is a stable dark matter candidate. The additional particle spectrum together with their hypercharges and Z_2 charges are given in Table 3.1. Including the above particle content, the Lagrangian of the model is given by

$$\begin{aligned}
\mathcal{L}_m = & \mathcal{L}_{\text{SM}} + (D_\mu S)^\dagger D^\mu S + \bar{\chi} i \gamma^\mu D_\mu \chi + \bar{\psi} i \gamma^\mu \partial_\mu \psi + \sum_i \bar{N}_i i \gamma^\mu \partial_\mu N_i \\
& - m_\chi \bar{\chi} \chi - m_\psi \bar{\psi} \psi - \sum_{ij} m_{Nij} \bar{N}_i N_j \\
& - \left(y_1 \bar{\chi} S \psi + \sum_{ij} y_{2ij} \bar{N}_i S l_j + \sum_{ij} Y_{Nij} \bar{L}_i \tilde{H} N_j + \text{H.c.} \right) \\
& - (\mu_S^2 S^\dagger S + \lambda (S^\dagger S)^2 + \lambda_1 S^\dagger S H^\dagger H)
\end{aligned} \tag{3.1}$$

where H and $\tilde{H} = i\sigma^2 H^*$ represent the SM Higgs doublet, and L_i and l_i denote the SM left-handed lepton doublet and right-handed charged lepton singlet, respectively. The summation indices i, j run from 1 to 3, indicating the three flavors of leptons. Further, we consider hierarchical right-handed neutrinos in the basis with $m_N = \text{diag}(m_{N1}, m_{N2}, m_{N3})$. In the broken electroweak symmetry phase, the VEV of H , v induces a Dirac mass term through the Yukawa interaction of the right-handed neutrinos

to the SM leptons,

$$\mathcal{L}_Y \supset \sum_{ij} m_{Dij} \bar{\nu}_{Li} N_j + \text{H.c.} \quad (3.2)$$

where the Dirac mass matrix is given by $m_D = \frac{v}{\sqrt{2}} Y_N \equiv v' Y_N$. This leads to the neutrino mass terms,

$$\mathcal{L}_{\text{mass}} \supset (\bar{\nu}_L \bar{N}) \begin{pmatrix} 0 & m_D \\ m_D^T & m_N \end{pmatrix} \begin{pmatrix} \nu_L^c \\ N \end{pmatrix} \quad (3.3)$$

with the Type-I seesaw mechanism [238–243] generating light neutrino masses $m_\nu = m_D^T m_N^{-1} m_D$. The masses of the heavy right-handed neutrinos, on the other hand, remains the same as m_{Ni} with negligible correction, which we ignore. In terms of the mass eigenstates, we have

$$\mathcal{L}_{\text{mass}} \supset \sum_{\alpha} m_{\nu\alpha} \bar{\nu}_{L\alpha} \nu_{L\alpha}^c + \sum_i m_{Ni} \bar{N}_i N_i, \quad (3.4)$$

with the gauge eigenstates of the light neutrinos related to the mass eigenstates through mixing matrix U , known as the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix[244] as

$$\nu_{L\alpha} = \sum_i U_{\alpha i} \nu_{Li}, \quad (3.5)$$

where α indicating the mass eigenstates, while i the flavor eigenstates. Conversely, from the experimental results on the masses and mixings of the light neutrinos, we can reconstruct Y_N for given heavy neutrino masses. We adopt the Casas-Ibarra parametrization (CI) [245, 246] to represent Y_N as

$$Y_N = D_{\sqrt{M}} R D_{\sqrt{\kappa}} U^\dagger, \quad (3.6)$$

where $(D_{\sqrt{M}})_{ij} = \sqrt{m_{Ni}} \delta_{ij}$, $(D_{\sqrt{\kappa}})_{ij} = \frac{\sqrt{m_i}}{v} \delta_{ij}$, with m_{Ni} and m_i as the masses of the heavy right-handed neutrinos and the light neutrinos, respectively, and R is any arbitrary complex orthogonal matrix.

The VEV of H also contributes to the mass of the charged singlet scalar yielding

$$m_S^2 = \mu_S^2 + \frac{\lambda_1 v^2}{2}. \quad (3.7)$$

On the other hand, the masses of the dark vectorlike fermions, χ and ψ , arise purely through the parameters in the Lagrangian, m_χ and m_ψ , respectively.

3.3 FERMIONIC FIMP DARK MATTER

In this scenario, the neutral Z_2 -odd fermion, ψ emerges as a dark matter candidate, coupling with other particles solely through its Yukawa interaction with the vectorlike gauge singlet charged fermion χ and the charged scalar S . The charged scalar, on the other hand, decays primarily through $S \rightarrow Nl$ channel, which is controlled by the Yukawa coupling y_2 , by requiring the mass hierarchy, $m_S > m_N$ for on shell neutrino N . This, however, can be relaxed to accommodate $m_S < m_N$ so that S decay goes through an off shell neutrino N , $S \rightarrow lN^* \rightarrow l\nu H$. The χ decay has two possibilities. In the kinematic region with $m_\chi > m_S + m_\psi$, χ decays through $\chi \rightarrow \psi S$, whereas for $m_\chi < m_S + m_\psi$ the decay proceeds through, $\chi \rightarrow \psi S^* \rightarrow \psi N^* l \rightarrow \psi l \nu H$, further requiring $m_\chi > m_H + m_\psi + m_l$. While the two-body decay is controlled by the Yukawa coupling y_1 alone, the other decay is dictated by the combination of $y_1 y_2$. With the corresponding Yukawa couplings combination sufficiently large, the dark matter fermions can be produced copiously as to be in thermal equilibrium in the early stages of the Universe. Then, through the nonequilibrium processes near temperature $T \sim m_\psi$, the annihilation process ($\psi\psi \rightarrow SS$) brings down their number density to eventually satisfy the relic abundance at the decoupling. This requires the mass hierarchy $m_\psi > m_S$. On the other hand, if the relevant coupling combination is much weaker, the production process $\chi \rightarrow \psi + \dots$ can be sufficiently slow to enable the FIMP mechanism to generate the required dark matter abundance.

Viability of all the above scenarios are studied in Ref. [237], treating the Yukawa couplings y_2 as well as Y_N diagonal. Relaxing this requirement to include nonzero off-diagonal couplings in both cases is necessary to generate leptogenesis. This is the focus of the present work, where we explore the parameter space that could accommodate the required relic density. The off-diagonal couplings induce lepton flavor violation (LFV) through Y_N as well as y_2 . Since Y_N is involved in the seesaw mechanism, it is naturally required to be small. Furthermore, we shall restrict our s to the case where $y_2 \ll 1$, to avoid inducing large LFV. Thus in our model the LFV is negligibly small.

3.3.1 Freeze-in via two-body decay of χ

When the couplings involving the dark matter particle are very small, we may envisage a scenario where the initial number density of dark matter particles is negligible, and the observed abundance of dark matter is produced by the slow decay of the partner particles. In this case, the process could occur through the slow decay of $\chi \rightarrow S + \psi$ if kinematically allowed ($m_\chi > m_S + m_\psi$), and will continue until the Universe cools down

Parameter	Range (GeV)
m_χ	150–1000
m_S	150–1000
m_ψ	1–850

Table 3.2: Range of the masses in GeV, scanned to fix the coupling y_1 satisfying the observed relic density.

to a temperature $T < m_\chi$. Below this temperature, due to the Boltzmann suppression of the number density of χ ($n_\chi \propto \exp(-m_\chi/T)$), there is no further addition to the number density of the dark matter fermion ψ , leading to a constant comoving number density. This mechanism of generating the dark matter relic density is known in the literature as the freeze-in mechanism. Since the initial DM number density is negligible, the inverse decay is irrelevant, and the Boltzmann equation satisfied by the dark matter particle is given by

$$\frac{dY_\chi}{dz} = 2Y_\chi^{\text{eq}} \frac{z}{H} \frac{K_1(z)}{K_2(z)} \Gamma_{\chi \rightarrow S\psi}, \quad (3.8)$$

leading to the relic density, [247]

$$\Omega h^2 = 2.19 \times 10^{27} \frac{g_\chi}{g_{S*}} \sqrt{g_{\rho*}} \frac{m_\psi \Gamma_{\chi \rightarrow S\psi}}{m_\chi^2}. \quad (3.9)$$

With the Lagrangian in Eq. 3.1, the decay width becomes, in terms of the masses of the particles involved and the Yukawa coupling y_1

$$\Gamma_{\chi \rightarrow S\psi} = \frac{y_1^2}{16\pi m_\chi^3} [(m_\chi + m_\psi)^2 - m_S^2] \times [(m_\chi^2 - m_S^2 - m_\psi^2)^2 - 4m_S^2 m_\psi^2]^{1/2}. \quad (3.10)$$

Putting Eq. 3.10 back into Eq. 3.9, and demanding Ωh^2 to satisfy the observed relic density, the coupling y_1 can be constrained, for given masses of the dark matter and the partner particles. Notice that, in this scenario, the dark matter sector is decoupled from the neutrino sector and does not contribute to leptogenesis. Within this scenario, selecting the relic density to satisfy the observed value of $\Omega h^2 = 0.118 \pm 0.001$ [248], we scanned the parameter space spanned by the masses m_χ , m_ψ , and m_S to obtain the value of the coupling y_1 obeying the kinematic restriction for the decay, $m_\chi > m_S + m_\psi$. Focusing on light dark matter, and sub-TeV partner particles, we scanned over the parameter range as given in Table 3.2. Allowed values of the parameter space arising from the scan are presented in different planes in Fig. 3.1. The parameter m_S affects the relic density purely

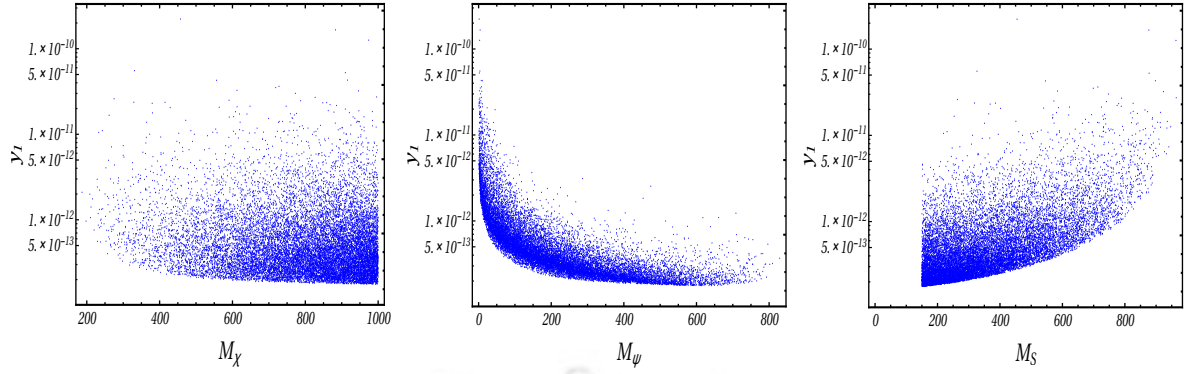


Figure 3.1: Yukawa coupling y_1 as a function of the masses of the dark matter (m_ψ), the charged fermionic partner particle (m_χ), and the charged scalar (m_S).

through the phase-space factor, and thus has a clear correlation with the couplings. On the other hand, m_χ is expected to have an anticorrelation effect, considering the dependence of the coupling and the mass of the decaying particle. This is indeed consistent with the plots in Fig. 3.1. However, the effect is somewhat subdued owing to the additional inverse dependence of the relic density on m_χ . In the case of the dark matter mass, the direct dependence of the relic density on the combination $m_\psi y_1^2$, where the coupling is coming from the decay width, yields a parabolic dependence between the two parameters. The influence of the phase-space factor is weak, as indicated by the spread of the points in the scan. We verified that the width of χ in this parameter range varies between 10^{-25} GeV and 10^{-21} GeV. The charged scalar would decay through $S \rightarrow Nl$, if kinematically allowed. However, as we shall consider the mass of the heavy neutrinos m_N to be in the range of 10 TeV or more, the two-body decay is disallowed. The three-body decay that follows through the heavily off shell neutrino N , along with the Yukawa coupling required for the seesaw mechanism being in the range of 10^{-8} slows down the decay of S . For the parameters considered here, the decay width is in the range of 10^{-24} to 10^{-19} GeV for m_S between 100 GeV and 1 TeV.

3.3.2 Freeze-in via four-body decay of χ

In the kinematic region where $m_\chi < m_S + m_\psi$, the two-body decay discussed above is kinematically disallowed. Noting that we are working with $m_S \ll m_N$, the leading allowed decay channel is the four-body decay shown in Fig. 3.2. This involves the couplings, y_1 , y_2 and Y_N . With $m_N \sim 10$ TeV, values of Y_N in the range of 10^{-8} to satisfy the seesaw condition the cross section for this process will require larger y_1 values. For consistency with the parameters chosen from leptogenesis, as discussed in Sec. 3.4, we restrict y_2 in

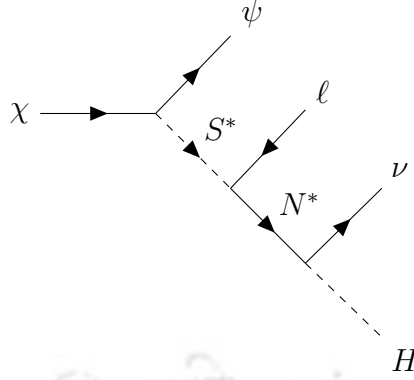


Figure 3.2: Feynman diagrams for the χ decay when $m_H + m_\psi < m_\chi < m_S + m_\psi$

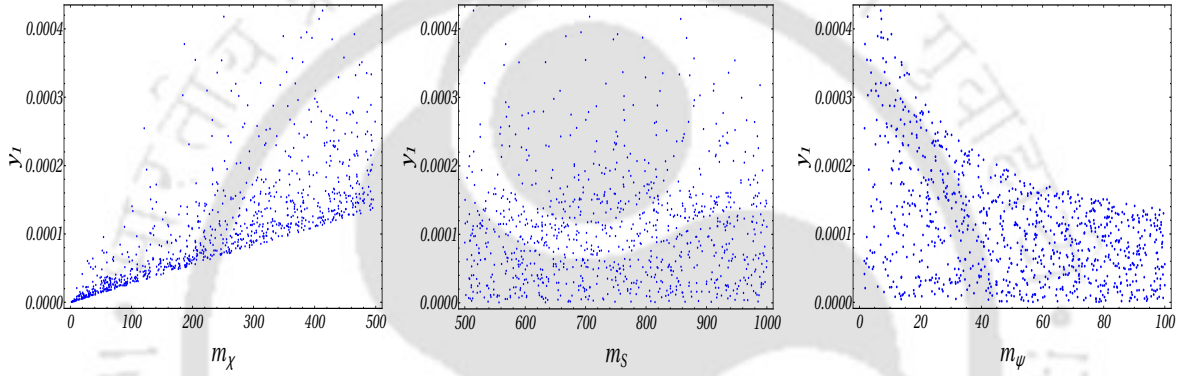


Figure 3.3: Yukawa coupling y_1 versus the mass of non-SM particles, from the requirement that they satisfy the relic density bound.

the range of 10^{-3} – 10^{-1} . With these restrictions, we scanned over the masses (m_χ , m_S , m_ψ) and used Eq. 3.9 replacing $\Gamma_{\chi \rightarrow S\psi}$ by the four-body width $\Gamma_{\chi \rightarrow \ell\nu H\psi}$, to obtain y_1 corresponding to the observed relic density. The scan range is as shown in Table 3.2 with the additional condition that $m_\chi < m_S + m_\psi$. The resulting parameter space points, projected on to $y_1 - m_i$ plane, where $i = \chi, S, \psi$, are presented in Fig. 3.3. Clearly, values of y_1 in the range of 10^{-4} are compatible with the dark matter observations.

3.4 LEPTOGENESIS

We move on to the analysis of leptogenesis in the scenario considered here, where the necessary lepton number violation and CP asymmetry are generated through the decay of one of the Majorana neutrinos present. In the standard scenario, leptogenesis is generated by the decay of neutrinos through the Higgs doublet, whereas in the scenario presented here, additional possibilities open through the Yukawa interaction term $y_{2ij}\bar{l}_i N_j S + \text{H.c.}$

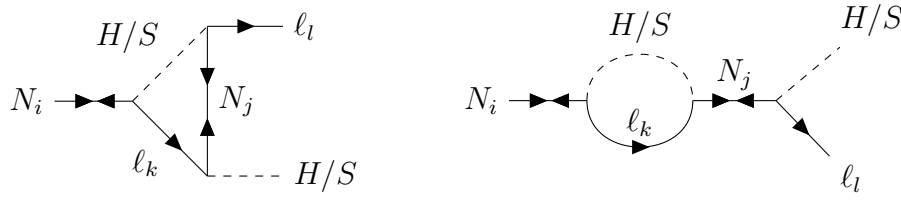


Figure 3.4: Feynman diagram corresponding to one loop vertex and self-energy corrections to $N_i \rightarrow \ell_l H/S$.

First, $N \rightarrow Sl$ and its CP-conjugate process provide new decay channels. Second, the self-energy and vertex corrections in each of the decays receive additional contributions. We now investigate possibilities in the low-energy leptogenesis enabled by these new interactions.

3.4.1 CP asymmetry

As in the standard leptogenesis, the CP asymmetry arises through the interference of the tree-level and the one-loop channels. We denote the CP-violating parameter arising from the standard process involving the Higgs boson as

$$\epsilon_1 = \frac{\Gamma(N_1 \rightarrow LH) - \Gamma(N_1 \rightarrow \bar{L}\bar{H})}{\Gamma(N_1 \rightarrow LH) + \Gamma(N_1 \rightarrow \bar{L}\bar{H})} \quad (3.11)$$

and the corresponding parameter arising from the decay involving the new charged singlet scalar as

$$\epsilon_2 = \frac{\Gamma(N_1 \rightarrow lS) - \Gamma(N_1 \rightarrow \bar{l}\bar{S})}{\Gamma(N_1 \rightarrow lS) + \Gamma(N_1 \rightarrow \bar{l}\bar{S})}. \quad (3.12)$$

Figure 3.4 shows the relevant one-loop diagrams involving the vertex corrections as well as the self-energy corrections. The CP-violating parameters in Eqs. (4.1) and (4.2) can be written explicitly in terms of the self-energy (ϵ_ξ^s) and vertex contributions (ϵ_ξ^v), with $\xi = H, S$, as

$$\epsilon_1 = \epsilon_H^s + \epsilon_H^v, \quad (3.13)$$

$$\epsilon_2 = \epsilon_S^s + \epsilon_S^v, \quad (3.14)$$

where

$$\epsilon_H^s = \frac{1}{8\pi K_{11}} \sum_{j=2,3} \left[\frac{m_{N1}}{m_{N1}^2 - m_{Nj}^2} \text{Im}(m_{Nj} K_{1j}^2 + m_{N1} \kappa_{1j} K_{1j} + m_{Nj} \kappa_{j1} K_{1j}) \right], \quad (3.15)$$

$$\epsilon_S^s = \frac{1}{8\pi\kappa_{11}} \sum_{j=2,3} \left[\frac{m_{N1}}{m_{N1}^2 - m_{Nj}^2} \text{Im}(m_{Nj}\kappa_{j1}^2 + m_{N1}K_{1j}\kappa_{j1} + m_{Nj}K_{j1}\kappa_{j1}) \right], \quad (3.16)$$

$$\epsilon_H^v = \frac{1}{8\pi K_{11}} \sum_{j=2,3} \text{Im}(K_{1j}^2) F\left(\frac{m_{Nj}^2}{m_{N1}^2}\right), \quad (3.17)$$

$$\epsilon_S^v = \frac{1}{8\pi\kappa_{11}} \sum_{j=2,3} \text{Im}(\kappa_{j1}^2) F\left(\frac{m_{Nj}^2}{m_{N1}^2}\right), \quad (3.18)$$

where $F(x) = \sqrt{x}[1 + (1+x)\ln\frac{x}{1+x}]$, and we have defined the relevant product of the Yukawa couplings as $K_{ij} = (Y_N^\dagger Y_N)_{ij}$ and $\kappa_{ij} = (y_2^\dagger y_2)_{ij}$. Unlike Y_N , the Yukawa coupling combination κ_{ij} connecting NIS do not play any role in the seesaw mechanism, but may influence the lepton flavor violating processes, as previously discussed. We shall therefore keep these couplings not larger than of order of 0.1, but otherwise unrestricted. Note that the standard leptogenesis is recovered in the limit $\kappa_{ij} \rightarrow 0$, when only the first terms in Eqs. 3.15 and 3.17 survive.

3.4.2 The Boltzmann equations

The time evolution of the number density when out-of thermal equilibrium is obtained using the Boltzmann equations. With $M_{N3} \gg M_{N2} \gg M_{N1}$, any lepton asymmetry generated by the decay of N_2 and N_3 at high temperature is washed out before the electroweak symmetry breaking and the surviving lepton asymmetry is generated by the decay of the lightest right-handed neutrino (N_1). The number density of N_1 depends on its decay, inverse decay, and scattering process. In the present setup, the following are the relevant processes

- (1) Decays: $N_1 \rightarrow HL$, $N_1 \rightarrow \bar{H}\bar{L}$, $N_1 \rightarrow Sl$, $N_1 \rightarrow \bar{S}\bar{l}$;
- (2) Inverse decays: $HL \rightarrow N_1$; $\bar{H}\bar{L} \rightarrow N_1$; $Sl \rightarrow N_1$; $\bar{S}\bar{l} \rightarrow N_1$;
- (3) Scattering processes: standard $\Delta L = 1$ s-channel processes: $lN_1 \rightarrow d\bar{u}$, $\bar{l}N_1 \rightarrow \bar{d}u$; standard $\Delta L = 1$ t-channel processes: $N_1 u \rightarrow d\bar{l}$, $N_1 \bar{u} \rightarrow \bar{d}l$, $N_1 d \rightarrow u\bar{l}$, $N_1 \bar{d} \rightarrow \bar{u}l$; standard $\Delta L = 1$ processes involving gauge boson A : $N_1 H \rightarrow AL$, $N_1 A \rightarrow HL$; new $\Delta L = 1$ processes involving S : $N_1 l \rightarrow SH$, $N_1 \bar{l} \rightarrow \bar{S}\bar{H}$; new $\Delta L = 1$ processes involving χ , ψ : $N_1 l \rightarrow \chi\psi$, $N_1 \bar{l} \rightarrow \bar{\chi}\bar{\psi}$.

Lepton number asymmetry is induced by all the above processes, and in addition, induced also by the $\Delta L = 2$ processes below:

- (4) Standard $\Delta L = 2$ processes: $ll \rightarrow \bar{H}\bar{H}$, $Hl \rightarrow \bar{H}\bar{l}$;
- (5) New $\Delta L = 2$ processes: $Sl \rightarrow \bar{S}\bar{l}$, $Hl \rightarrow \bar{l}\bar{S}$, $\bar{H}\bar{l} \rightarrow lS$.

As standard practice, we consider the evolution of the number density normalized by

the entropy density, denoted by Y_{N1} and the effective lepton number density, $Y_L = Y_l - Y_{\bar{l}}$. Within this framework the Boltzmann equations can be written in the following form,

$$\frac{dY_{N1}}{dz} = D(-Y_{N1} + Y_{N1}^{\text{eq}}), \quad (3.19)$$

$$\frac{dY_L}{dz} = \left(-\epsilon_1 \frac{\Gamma_{D1}}{Hsz} Y_{N1}^{\text{eq}} + \epsilon_2 \frac{\Gamma_{D2}}{Hsz} Y_{N1}^{\text{eq}} \right) (-Y_{N1} + Y_{N1}^{\text{eq}}) - SY_L, \quad (4.10) \quad (3.20)$$

where

$$\begin{aligned} D &= \frac{\Gamma_{D1} + \Gamma_{D2}}{Hsz} Y_{N1}^{\text{eq}} + S_N, \\ S_N &= \frac{1}{Hsz Y_{N1}^{\text{eq}}} \left(\Gamma_{s1} + \Gamma_{s2} + \frac{1}{2}(\Gamma_{s3} + \Gamma_{s4} + \Gamma_{s5} + \Gamma_{s6}) + 2\Gamma_{s,\text{new1}} + 2\Gamma_{s,\text{new2}} \right) \\ S &= \frac{\Gamma_{D1} + \Gamma_{D2} + 2\Gamma_{st}}{2Hsz Y_L^{\text{eq}}} + S_N \frac{Y_{N1}}{2Y_L^{\text{eq}}}. \end{aligned} \quad (3.21)$$

Here Γ_{D1} and Γ_{D2} are thermal averaged partial decay widths of $N_1 \rightarrow HL$ and $N_1 \rightarrow Sl$, respectively, with

$$\Gamma_{Di}(N_1 \rightarrow ab) = \int d\Pi_{N1} e^{-E_{N1}/T} \int d\Pi_a d\Pi_b (2\pi)^4 \delta^4(P_N - P_a - P_b) |M(N_1 \rightarrow ab)|^2, \quad (3.22)$$

where $d\Pi_i = d^3P_i/((2\pi)^3 2E_i)$, and M is the corresponding invariant amplitude. The thermal averaged scattering cross sections of the $\Delta L = 1$ standard processes $lN_1 \rightarrow d\bar{u}$, $lN_1 \rightarrow H\bar{A}$, $N_1 u \rightarrow d\bar{l}$, $N_1 d \rightarrow u\bar{l}$, $N_1 H \rightarrow A\bar{L}$, $N_1 \bar{A} \rightarrow H\bar{L}$ are denoted by Γ_{si} ; $i = 1, \dots, 6$, respectively, and the corresponding processes involving new particles, $N_1 l \rightarrow \chi\psi$ and $N_1 \bar{l} \rightarrow S\bar{H}$ are denoted by $\Gamma_{s\text{new1}}$ and $\Gamma_{s\text{new2}}$, respectively. In the case of $\Delta L = 2$ processes listed above, the thermal averaged scattering cross section is jointly denoted by Γ_{st} . In terms of the scattering cross section $\sigma(s)$ with a fixed centre of mass energy $\sqrt{s}$, the thermal averaged cross section is given by [249]

$$\Gamma_{si}(X + a \rightarrow 1 + 2) = \frac{T}{64\pi^4} \int_{s_{\min}}^{\infty} ds s^{1/2} \hat{\sigma}(s) K_1 \left(\frac{\sqrt{s}}{T} \right), \quad (3.23)$$

Where $\hat{\sigma} \equiv 2s\lambda[1, m_N^2/s, m_a^2/s]\sigma$ is the ‘reduced cross section’, $\lambda(x, y, z) = (x - y -$

$z)^2 - 4yz$, and $K_1(x)$ is the modified Bessel function of second kind.

Formally, we can write the solutions of Eqs. 3.19 and 3.20 in integral form as

$$Y_{N1} = e^{-\int Ddz} \left(\int Y_{N1}^{\text{eq}} D e^{\int Ddz} dz' + \text{const} \right), \quad (3.24)$$

$$Y_L = \epsilon_1 \zeta_1 + \epsilon_2 \zeta_2, \quad (3.25)$$

where the efficiency factors ζ_i are given by

$$\zeta_i = e^{-\int Sdz} \left(\int \frac{\Gamma_{Di}}{Hsz} Y_{N1}^{\text{eq}} \left(-e^{-\int Ddz} \left(\int Y_{N1}^{\text{eq}} D e^{\int Ddz} dz' \right) + Y_{N1}^{\text{eq}} + \text{const} \right) e^{\int Sdz} dz'' + \text{const}' \right). \quad (3.26)$$

Notice that the CP-asymmetry parameter ϵ_1 and the corresponding efficiency factor ζ_1 yield the standard contribution to the lepton number asymmetry, while ϵ_2 and ζ_2 add to the lepton number asymmetry arising through the new dynamics. As expected, in the case of hierarchical right-handed neutrinos with TeV scale masses, the contributions from the standard processes are negligible. We shall explore the parameter region of the model considered where this deficit is compensated by the contribution from the new channels.

3.5 NUMERICAL ANALYSIS

In this section, we perform the numerical analysis for the lepton asymmetry generated through a combination of CP asymmetry arising through the standard decay and the decay to the newly introduced charged scalar, as well as considering the influence of the washout effect arising through different scattering processes mentioned in Sec. 3.4.2. The CP parameters depend on the coupling constants Y_N and y_2 , where, as we noted before, Y_N is constrained by the light neutrino masses and mixings, whereas y_2 is largely unconstrained. Considering the CI parametrization as in Eq. 3.6, Y_N depends on the masses and mixing matrix elements of the light neutrinos, on the masses of the heavy neutrinos, and on the elements of an arbitrary orthogonal matrix, R . We assume normal hierarchy in the light neutrino sector, with the lightest neutrino taken to be massless, and the other two masses set in agreement with constraints from the neutrino oscillation experiments [211, 250]. Accordingly, we chose $m_1 = 0$ eV, $m_2 = 0.0083$ eV and $m_3 = 0.051$

eV. The complex orthogonal matrix, R is parametrized as

$$R = \begin{pmatrix} \cos \theta & \sin \theta & 0 \\ -\sin \theta & \cos \theta & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (3.27)$$

with θ complex. We scan over the heavy neutrino masses keeping the normal hierarchy $m_{N1} \ll m_{N2} \ll m_{N3}$, with the lightest mass in the 1–100 TeV range, and the heavier ones differing from it by at least one order of magnitude. The dependence of Y_N and y_2 enter the CP asymmetries through their combination K_{1j} and κ_{1j} with $j = 2, 3$, respectively, as given in Eqs. 3.15–3.18. The complex nature of these parameters is important in determining the amount of CP violation.

First we consider K_{1j} . With the standard Yukawa couplings given by the CI parametrization as in Eq. 3.6,

$$K_{1j} = (Y_N^*)_{k1}(Y_N)_{kj} = \sum_{k,\alpha,\beta} \frac{m_{Nk}}{v^2} \sqrt{m_\alpha m_\beta} R_{k\alpha}^* R_{k\beta} U_{1\alpha} U_{j\beta}^*, \quad (3.28)$$

which, for the structure of R and the choice of $m_1 = 0$, takes a simpler form,

$$K_{1j} = \frac{m_{N1}m_2}{v^2} |R_{12}|^2 U_{12} U_{j2}^* + \frac{m_{N2}m_2}{v^2} |R_{22}|^2 U_{12} U_{j2}^* + \frac{m_{N3}m_3}{v^2} U_{13} U_{j3}^*. \quad (3.29)$$

Thus, any phase of the elements of R is irrelevant, and we can consider it a real orthogonal matrix. Taking U as the PMNS matrix [211], the only phases that enter in K_{12} are those of U_{13} and U_{22} , while in K_{13} the phases of U_{13} and U_{32} give rise to the relevant CP phase. However, noticing that for the hierarchical case considered in the heavy neutrino sector, the third term, proportional to M_{N3} , dominates over the others, the relevant phase in both K_{12} and K_{13} is that of U_{13} . We write this as

$$K_{12} \sim \frac{m_{N3}m_3}{v^2} \sin \theta_{23} \cos \theta_{13} \sin \theta_{13} e^{-i\delta_{CP}}, \quad (3.30)$$

$$K_{13} \sim \frac{m_{N3}m_3}{v^2} \cos \theta_{23} \cos \theta_{13} \sin \theta_{13} e^{-i\delta_{CP}}, \quad (5.3) \quad (3.31)$$

while the real matrix element K_{11} can be written as

$$K_{11} \sim \frac{m_{N3}m_3}{v^2} \sin^2 \theta_{13}. \quad (5.4) \quad (3.32)$$

We use the present experimental values for the mixing angles and the δ_{CP} parameter [250]

$$\theta_{13} = 8.57^{+0.13^\circ}_{-0.12^\circ}, \quad \theta_{23} = 49.0^{+1.1^\circ}_{-1.4^\circ}, \quad \delta_{CP} = 195^{+51^\circ}_{-25^\circ}. \quad (3.33)$$

Taking $m_{N3} \sim 10^3$ TeV, we have $K_{1j} \sim 10^{-8}$. The standard contribution to the self-energy term in the CP -violating parameter can then be read from Eq. 2.86 as

$$\epsilon_H^s|_{\text{std}} = \frac{1}{8\pi K_{11}} \left(\frac{m_{N1}}{m_{N2}} \text{Im}(K_{12}^2) + \frac{m_{N1}}{m_{N3}} \text{Im}(K_{13}^2) \right) \sim \frac{1}{8\pi K_{11}} \frac{m_{N1}}{m_{N2}} \text{Im}(K_{12}^2) \quad (3.34)$$

and the new contribution to the standard decay channel is dominantly

$$\epsilon_H^s|_{\text{new}} \sim \frac{1}{8\pi K_{11}} \frac{m_{N1}}{m_{N2}} \text{Im}(\kappa_{21} K_{12}). \quad (3.35)$$

With a suitably chosen κ_{21} , it may then be possible to lift up the CP -violating parameter to the value required in leptogenesis, even when the standard contribution is a few orders of magnitude smaller. In addition, there is an effect from the new decay channel $N_1 \rightarrow lS$, which contributes to the self-energy Eq. 3.16,

$$\epsilon_S^s \sim \frac{1}{8\pi \kappa_{11}} \frac{m_{N1}}{m_{N2}} \text{Im}(\kappa_{21}^2). \quad (3.36)$$

Notice that the contribution to the standard channel from the virtual effects of the new scalar particle, S is present even if κ_{21} is real, whereas the contribution coming from the new channel is subdominant in this case, as this will now be

$$\epsilon_S^s \sim \frac{1}{8\pi \kappa_{11}} \frac{m_{N1}}{m_{N2}} \text{Im}(\kappa_{21} K_{21}). \quad (3.37)$$

If κ_{1j} is real, the effect of S is completely absent in the vertex correction contribution to the CP -violating parameter, Eqs. 3.17 and 3.18. However, once κ_{1j} has an imaginary part, it could be suitably chosen to get the required CP asymmetry, which now has a dominant contribution from the new scalar S

$$\epsilon_S^v \sim \frac{1}{8\pi \kappa_{11}} \frac{m_{N3}}{m_{N1}} \text{Im}(\kappa_{31}^2). \quad (3.38)$$

As it is proportional to m_{N3} , the vertex contribution dominates over the self-energy contribution for κ_{1j} of the same order. With these considerations, we now under-take the numerical analysis, scanning the parameter space to find regions that are compatible with the observed baryon asymmetry. We consider four distinct possibilities for κ_{1j} , i.e., (1) Real and positive κ_{1j} ; (2) Real and negative κ_{1j} ; (3) Imaginary and negative

$$m_\chi = 200 \text{ GeV}, \quad m_S = 175 \text{ GeV}, \quad m_\psi = 60 \text{ GeV}, \quad 10^{-4} \leq \kappa_{11} \leq 10^{-1}$$

Scenario		couplings	Mass of Heavy Neutrinos		
		κ_{12}, κ_{13}	m_{N_1} TeV	m_{N_2} TeV	m_{N_3} TeV
Case 1	Real positive κ	$10^{-3} \leq \kappa_{1j} \leq 10^{-1}$	(10, 100)	$(10^3, 10^4)$	$(10^5, 10^6)$
Case 2	Real negative κ	$-10^{-1} \leq \kappa_{1j} \leq -10^{-3}$	(10, 100)	$(10^3, 10^4)$	$(10^5, 10^6)$
Case 3	Imaginary κ	$-i10^{-3} \leq \kappa_{1j} \leq -i10^{-1}$	(10, 100)	$(10^3, 10^4)$	$(10^5, 10^6)$
Case 4	Complex κ	$(1 - i)10^{-3} \leq \kappa_{1j} \leq (1 - i)10^{-1}$	(10, 100)	$(10^3, 10^4)$	$(10^5, 10^6)$

Table 3.3: The possible leptogenesis scenarios, yielding correct signs for the CP asymmetry, and the limits on the couplings, and neutrino masses, where applicable, from the asymmetry parameters for each case.

κ_{1j} ; (4) Complex κ_{1j} with a negative imaginary part. The signs are chosen so that in each situation, we get the correct sign for the CP asymmetry, and these are the only viable scenarios for leptogenesis in our model. Throughout In our analysis, we allow the right-handed neutrino masses to vary in order to clearly understand the influence of the couplings (κ) on the CP asymmetries ϵ_1 and ϵ_2 . We then the generated baryon asymmetry to determine if each scenario is compatible with the experimental data. The scans are performed over the mass and coupling parameters in the range specified in Table 3.3 for each scenario. We now consider specific cases with κ_{1j} taken to be real, purely imaginary or complex and perform separate s for each case.

3.5.0.1 Case 1: Real positive κ

If κ is constrained to be real and positive, varying the right-handed neutrino masses, a scan will reveal the effects of the masses of right-handed neutrinos as well as of κ_{1j} . Here the mass of the lightest right-handed neutrino, m_{N_1} , is varied from 10 TeV to 100 TeV; the mass m_{N_2} is varied from 10^3 TeV to 10^4 TeV and the mass m_{N_3} is varied from 10^5 TeV to 10^6 TeV. The parameters ϵ_1 and ϵ_2 depend only on κ_{11} , κ_{12} , and κ_{13} . In this case, the vertex correction does not contribute to ϵ_2 , and thus the CP asymmetry is dominated by ϵ_S^s as in Eq. 3.37. In the standard channel, the influence of S to the self-energy correction will also be affected by κ_{12} being real, as in Eq. 3.35. The efficiency factors ζ_1 and ζ_2 depend on the mass of right-handed neutrinos. Since ϵ_1 and ϵ_2 depend more sensitively on κ_{12} than on any other κ_{ij} 's, we focus on variations with κ_{12} . in Fig. 3.5 we show the dependence of ϵ_1 and ϵ_2 on κ_{12} and on the masses m_{N_1} and m_{N_2} . We note that throughout, varying the right-handed neutrino masses, ϵ_1 remains negative, while ϵ_2 is small and positive, showing that this is a robust prediction of real positive values of κ . The CP asymmetry parameter ϵ_2 is very small due to its inverse dependence on κ_{11} , which is orders of magnitude larger

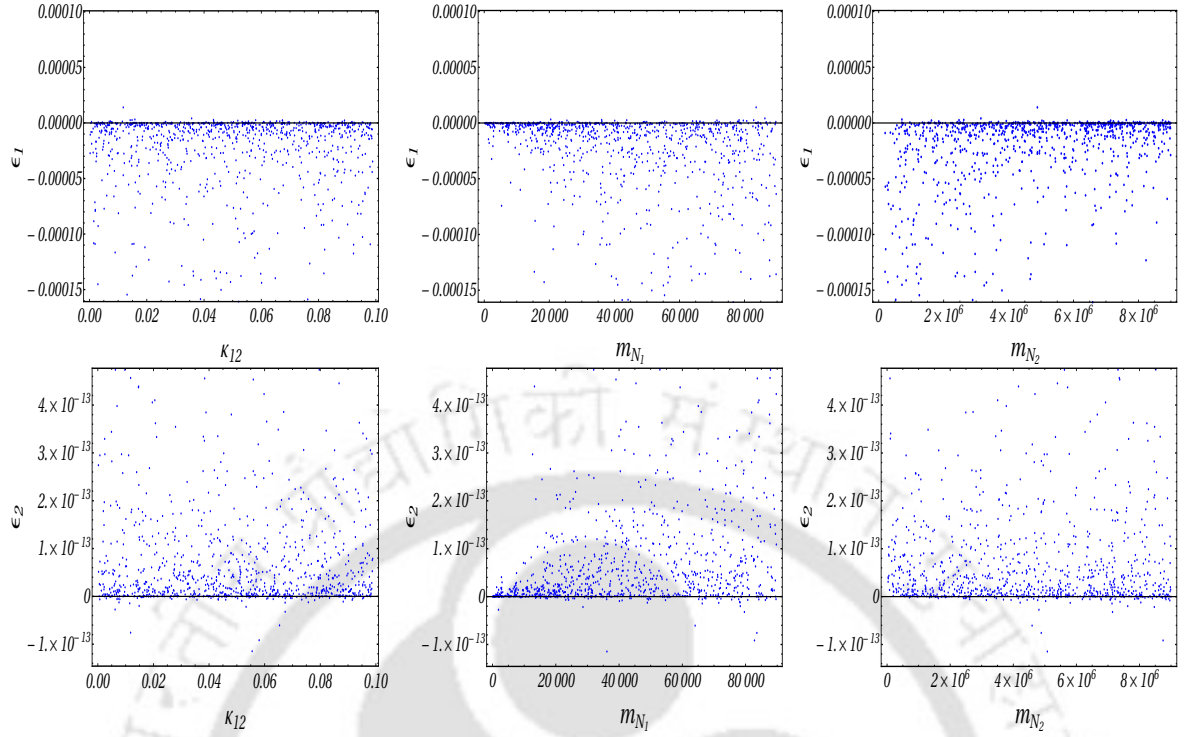


Figure 3.5: CP asymmetry against the relevant coupling κ_{12} (left panels) and the masses, m_{N1} (middle panels) and m_{N2} (right panels) for Case 1, with real κ and masses m_{Nj} in GeV. Top panels: Variation of ϵ_1 with κ_{12} , m_{N1} and m_{N2} . Bottom panels: same, but for ϵ_2 .

than K_{11} , while ϵ_1 is of order 10^{-5} . To further elucidate the resulting baryon asymmetry generated in this case, in Fig. 3.6 we plot the lepton asymmetry and the corresponding generated baryon asymmetry η against $z = m_{N1}/T$ for a selected set of parameter values, as given in Table 3.4. Clearly, the baryon-antibaryon asymmetry obtained is many orders of magnitude smaller than the value required to generate the observed matter-antimatter asymmetry.

3.5.0.2 Case 2: Real negative κ

For negative values of κ_{1j} , the standard channel yields positive contributions (ϵ_1) to the CP asymmetry, while the contributions due to the new channels (ϵ_2) are negative, as is clear from Eqs. 3.35 and 3.37. These are plotted in Fig. 3.7.

However, with ϵ_2 being many orders smaller than ϵ_1 , the overall effect is such that the sign of η is negative, indicating a non acceptable case of excess of antibaryons. Note however that the sign prediction for ϵ_1 is not robust. Due to the large uncertainty in the measurement of δ_{CP} Eq. 3.33, the sign of K_{12} is flipped for the lower range of δ_{CP}

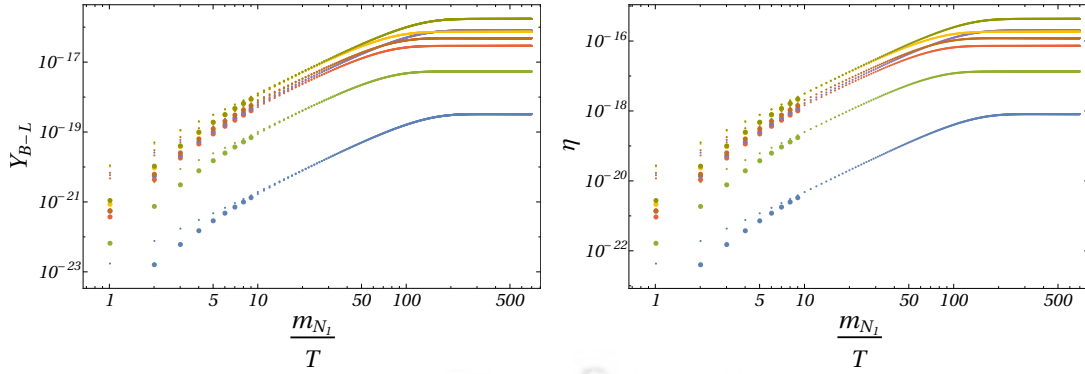


Figure 3.6: The lepton asymmetry, Y_{B-L} (left) and the baryon asymmetry, η (right) versus $z = m_{N1}/T$ corresponding to real κ for the set of selected parameter values as given in Table 3.4.

κ_{11}	κ_{12}	m_{N1} (TeV)	m_{N2} (TeV)
0.024	0.097	29.8	6937.1
0.028	0.071	49.8	5802.6
0.025	0.071	61.9	4564.2
0.097	0.086	83.7	7585.8
0.013	0.087	19.6	7373.1
0.022	0.090	75.8	1459.1
0.025	0.022	19.8	7849.5
0.026	0.008	88.6	1060.7
0.031	0.008	11.3	8482.2
0.025	0.090	73.3	7837.5

Table 3.4: Parameter values for the Y_{B-L} and η plots in Fig. 3.6. The values of Y_{B-L} and η corresponding to these parameters are $Y_{B-L} \in (10^{-19} - 10^{-16})$ and $\eta \in (10^{-18} - 10^{-15})$.

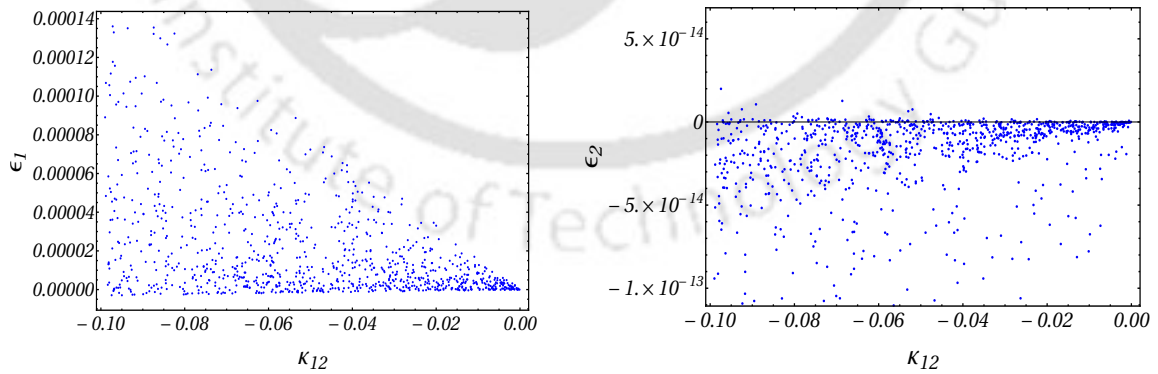


Figure 3.7: CP-asymmetry parameters ϵ_1 (left) and ϵ_2 (right) versus κ_{12} for Case 2, for real negative κ values and right-handed neutrino masses, m_{N_j} , in a range as given in Table 3.3.

Table 3.5: Parameter values for the Y_{B-L} and η plots in Fig. 3.9.

κ_{11}	κ_{12}	m_{N1} (TeV)	m_{N2} (TeV)
0.085	-0.00063 i	41.9	3205.6
0.026	-0.00076 i	16.4	1464.6

compared to the central value and upper range. In particular, if we consider $\delta_{CP} < 180^\circ$, we can obtain the desired sign for ϵ_1 . However, even in that case, the numerical values for the CP asymmetry are found to be far too small to provide the required leptogenesis.

3.5.0.3 Case 3: Imaginary κ

We consider next the case where κ_{12} and κ_{13} are purely imaginary. The imaginary part of K_{12} plays the role of the phase factor for real positive κ_{12} and real negative κ_{12} , but in this case the real part of K_{12} also contributes to the asymmetry. The order of the real and imaginary parts of K_{12} are approximately same. Unlike the case where κ_{12} was real, here ϵ_1 and ϵ_2 have the same sign, meaning that their numerical values are negative for imaginary values of κ_{12} and κ_{13} . The numerical values of ϵ_2 are very small in comparison to ϵ_1 , so ϵ_1 and the efficiency factor ζ_1 are not able to give required matter-antimatter asymmetry, as shown below. Varying neutrino masses populates the plot of ϵ_1 versus κ_{12} for ϵ_1 values closer to 0, while affecting ϵ_2 somewhat less significantly. The variations of ϵ_1 and ϵ_2 with κ_{12} and the masses, m_{N1} and m_{N2} , are shown in Fig. 3.8. Investigating the lepton number asymmetry generated in this case, sample points from across the parameter region considered could not provide the required baryon asymmetry. In Fig. 3.9 we show the evolution of Y_{B-L} and η for a couple of selected sets of parameter values, as given in Table 3.5, to illustrate this point.

3.5.0.4 Case 4: Complex κ

We now the case where κ_{12} has both real and imaginary parts. The parameters K_{12} and K_{13} are complex but K_{11} is real. The real part of κ_{12} contributes to the asymmetry, together with the combination of the imaginary part of K_{12} , the imaginary part of κ_{12} , along with the real part of K_{12} . There are two phases in this case; one in the coupling constant Y_N and another in κ_{12} . We find that ϵ_1 and ϵ_2 can both acquire negative numerical values for complex κ_{12} and κ_{13} . More importantly, the numerical values of ϵ_2 and ϵ_1 are of the same order. In Fig. 3.10 we plot the dependence of ϵ_1 (top panels) and ϵ_2 (bottom panels) with the parameter κ_{12} and κ_{13} , separately for their real and imaginary parts, while in Fig. 3.11 we plot the dependence of ϵ_1 and ϵ_2 on the relevant right-handed

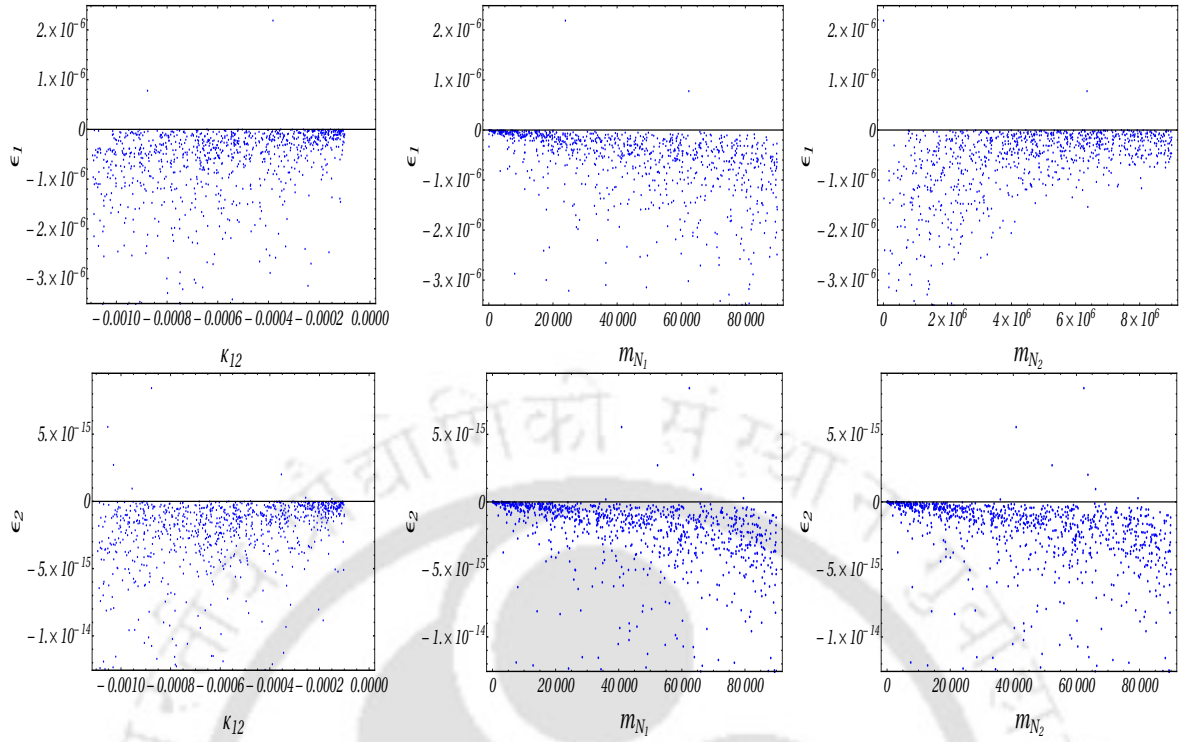


Figure 3.8: CP asymmetry against the relevant coupling combination κ_{12} (left panels) and right-handed neutrino masses, m_{N1} (middle panels) and m_{N2} (right panels) for ϵ_1 (top panels) and ϵ_2 (bottom panels), for Case 3, with imaginary values for κ_{12} .

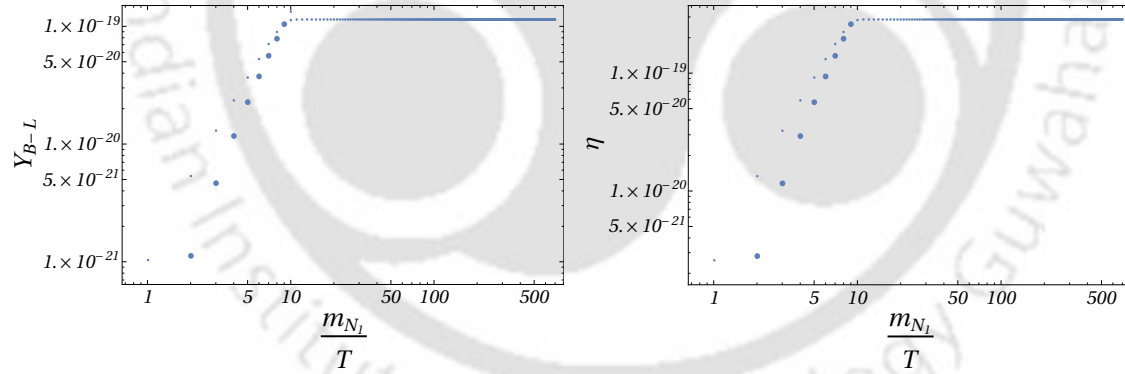


Figure 3.9: The lepton asymmetry, Y_{B-L} (left) and the baryon asymmetry, η (right) versus $z = m_{N1}/T$ corresponding to imaginary κ with parameter values as in Table 3.5.

neutrino masses m_{N1} , m_{N2} and m_{N3} . We observe that by varying neutrino masses, regions of ϵ_1 negative and close to zero are evenly populated, while ϵ_2 is allowed to be positive. Note the similarities between the dependence of ϵ_1 and ϵ_2 on $\text{Re } \kappa_{12}$ and $\text{Re } \kappa_{13}$; while the plots for ϵ_2 versus $\text{Im } \kappa_{12}$ and $\text{Im } \kappa_{13}$ favor values close to zero, the variation of ϵ_1 on the same parameters is slightly different.

To study the asymmetry, in this case in Fig. 3.13 we plot the dependence of the lepton

asymmetry, Y_{B-L} (left) and the baryon asymmetry, η (right), with $z = m_{N1}/T$ for the case with complex κ for a representative set of parameter values, as indicated in the caption. One can see that, unlike scenarios where we consider κ_{12} purely real or purely imaginary, this case is compatible with the requirement of yielding sufficient baryon asymmetry in the Universe.

3.5.0.5 Comparison of the standard leptogenesis with the new model effects

In the previous subsections we showed that, with suitably chosen parameter values, our model could explain the observed baryon asymmetry of the Universe. The presence of the new charged scalar that interacts directly with the right-handed neutrinos boosts the CP asymmetry, inducing the required baryon asymmetry. We further illustrate this by comparing the standard leptogenesis case (in the absence of the new degrees of freedom) with the scenario considered here, including all possibilities described above. To highlight the differences, we fix the mass of the lightest right-handed neutrino to be $m_{N1} = 10$ TeV, and other masses and relevant parameter values as given in Table 3.6. Choosing m_{N1} so light means sufficient leptogenesis cannot be generated in seesaw Type-I models without neutrino mass degeneracy, whereas in our model there is an additional channel that violates lepton number in the decay process, and both channels are important in the analysis. Numerical solutions of Boltzmann equations for this case are plotted in Fig. 3.14, giving the evolution of the lepton asymmetry and the corresponding induced baryon asymmetry. The left side plot shows the lepton asymmetry Y_{B-L} , with curves labeled Y_{B-L}^{S0} and Y_{B-L}^S representing the yields in the standard case without and with the washout effects induced by scattering included, respectively, and Y_{B-L}^0 and Y_{B-L} , respectively, representing the total yields including the new contributions. The curve Y_{B-L}^{Exp} represents the lepton asymmetry needed to generate the correct matter- antimatter asymmetry. Similarly, in the right side plot, η^0 and η^{S0} correspond to the standard contribution alone without and with scattering effects, respectively, and the top curves η^0 and η correspond to the total contribution in our model, without and with the effects of washout scattering processes. The flat curve labeled η^{Exp} represents the observed baryon asymmetry. These plots show that while standard leptogenesis alone is unable to generate required matter-antimatter asymmetry at low masses, the addition of the new degrees of freedom makes up the deficit.

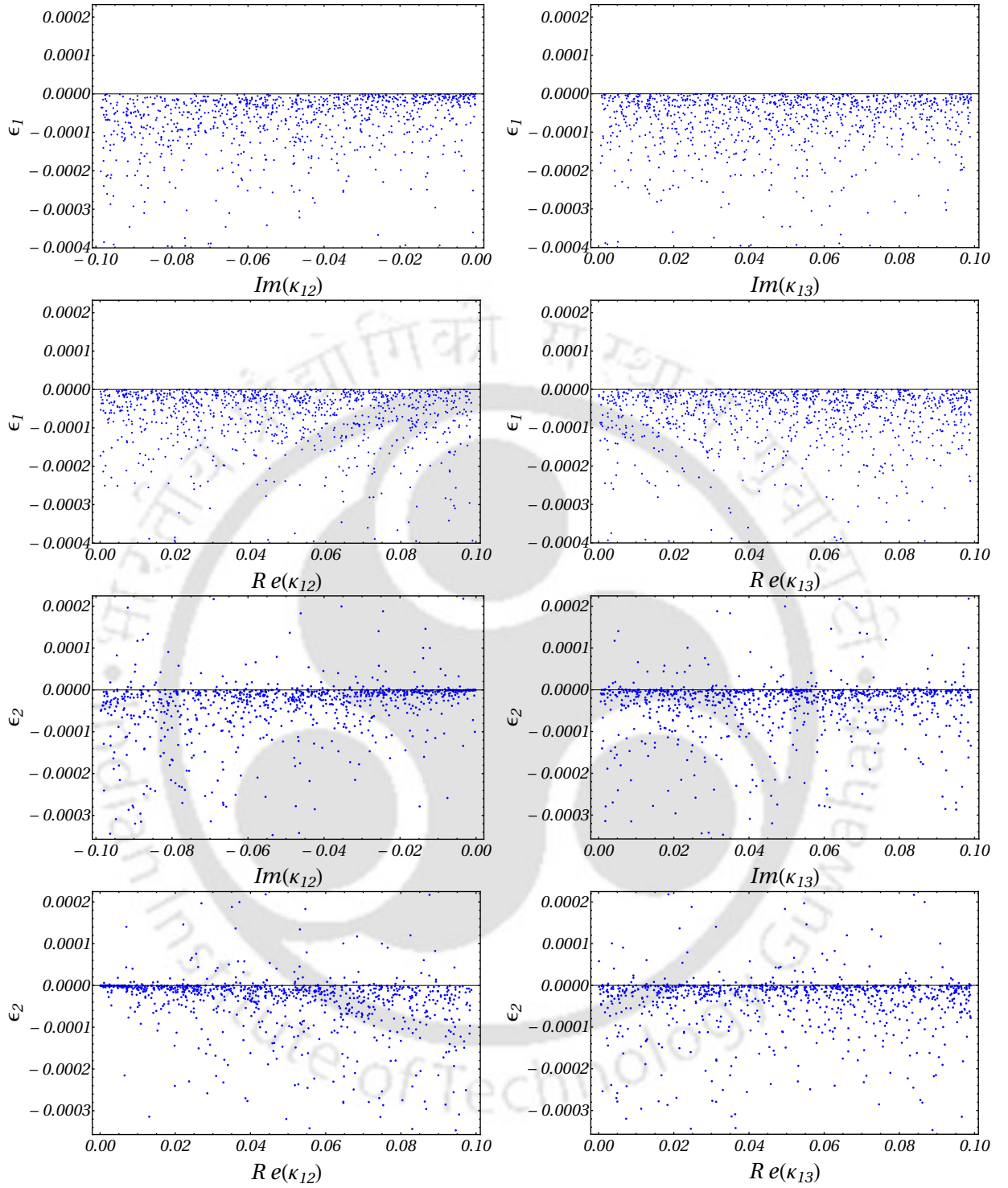


Figure 3.10: CP asymmetry ϵ_1 (top panels) and ϵ_2 (bottom panels) plotted against the real and imaginary parts of the relevant coupling combination for Case 4, with complex κ and varying m_{Nj} . We show the dependence on $\text{Im } \kappa_{12}$ (left panels), $\text{Im } \kappa_{13}$ (left middle panels), $\text{Re } \kappa_{12}$ (right middle panels) and $\text{Re } \kappa_{13}$ (right panels).

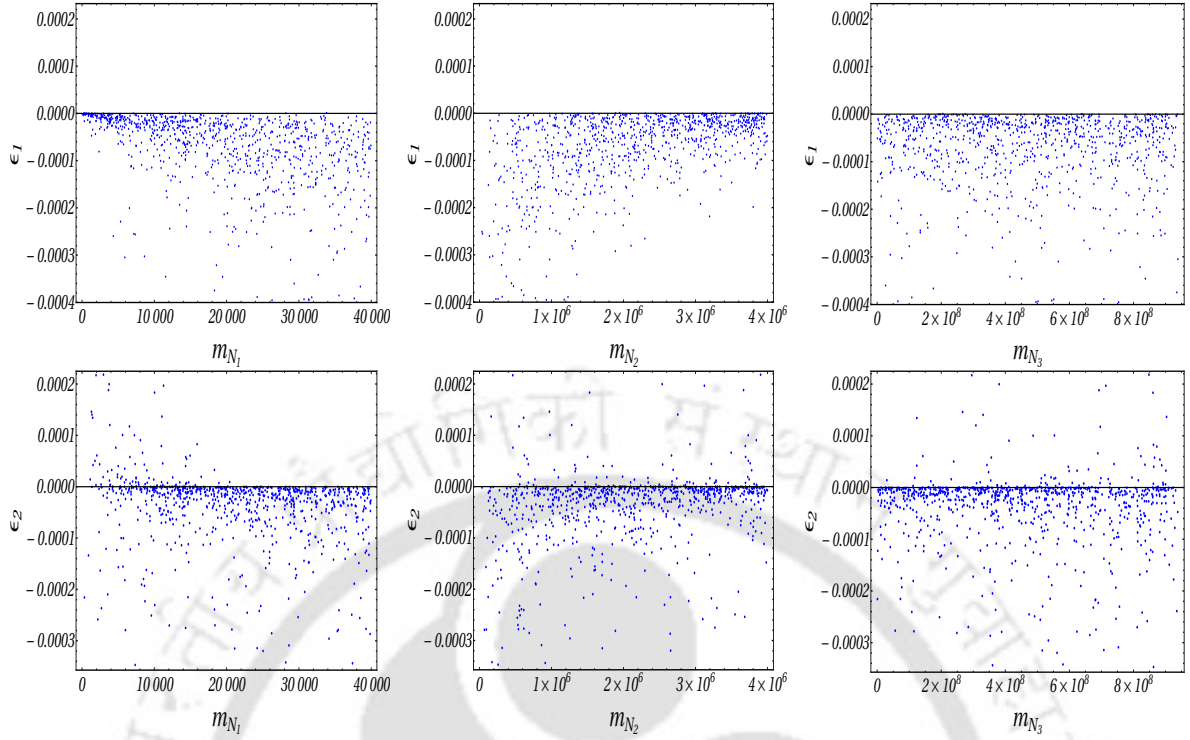


Figure 3.11: CP-asymmetry ϵ_1 (top panels) and ϵ_2 (bottom panels) plotted against the masses, m_{N_1} (left panels), m_{N_2} (middle panels) and m_{N_3} (right panels), for Case 4b, with complex κ .

Table 3.6: Benchmark point for the comparison of standard leptogenesis case with the new scenario.

m_{N_1} (TeV)	m_{N_2} (TeV)	m_{N_3} (TeV)	m_χ (GeV)	m_ψ (GeV)	m_S (GeV)	κ_{11}	κ_{12}	κ_{13}
10	1020	8.76×10^5	200	60	175	0.018	$(2 + 0.055i) \times 10^{-3}$	$(4.51 - 5.5i) \times 10^{-4}$

3.6 CONCLUSION

In this Chapter we have presented the analysis of a simple extension of the SM to account for both a dark matter candidate and a mechanism for leptogenesis which could produce sufficient asymmetry to generate the matter-antimatter discrepancy in the Universe. The extension includes, in addition to three Majorana right-handed neutrinos, a gauge singlet charged scalar field S^+ , plus a charged (χ^+) and a neutral (ψ) singlet fermions, with the latter two are odd under an additional Z_2 symmetry. In this scenario, ψ is a stable dark matter candidate, interacting with the other particles through its Yukawa coupling alone. When this coupling is very small, the abundance of dark matter is due to the slow two-body decay of $\chi^+ \rightarrow S^+ \psi$, if kinematically allowed, or the four-body decay $\chi^+ \rightarrow H \nu \psi$, in the regions where the two-body decay is forbidden. We implement

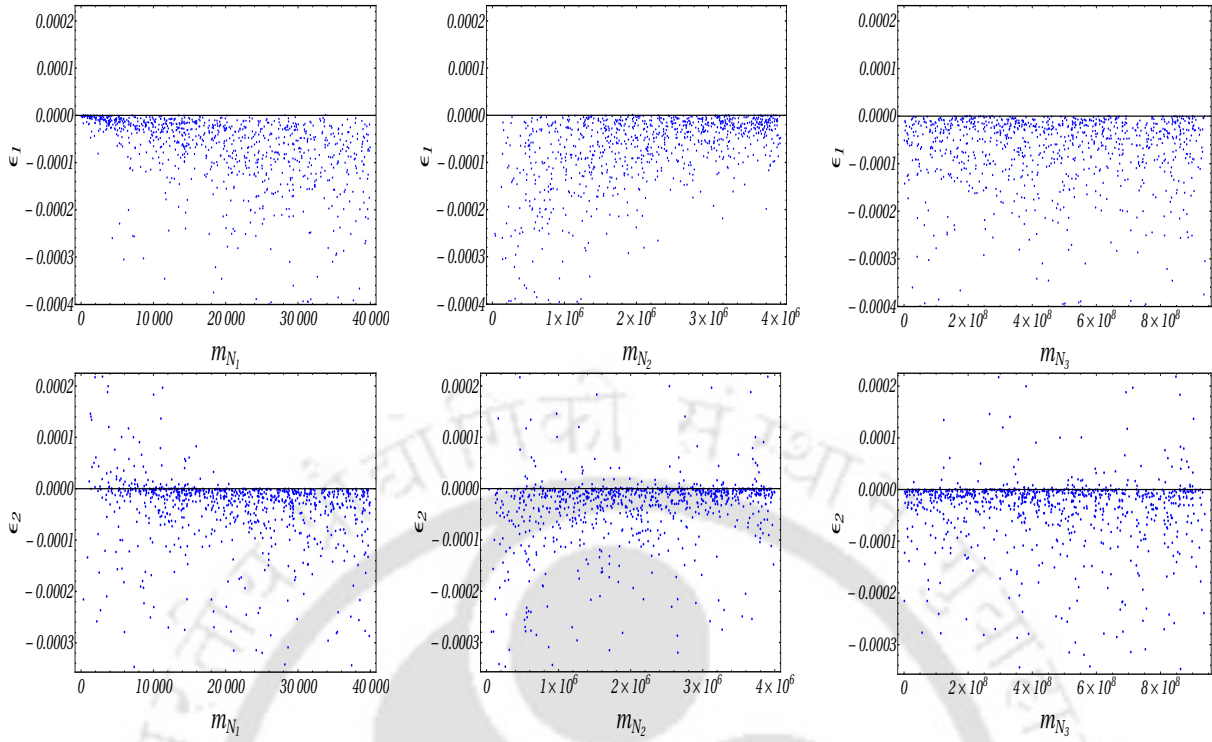


Figure 3.12: CP asymmetry ϵ_1 (top panels) and ϵ_2 (bottom panels) plotted against the masses, m_{N1} (left panels), m_{N2} (middle panels) and m_{N3} (right panels), for Case 4, with complex κ .

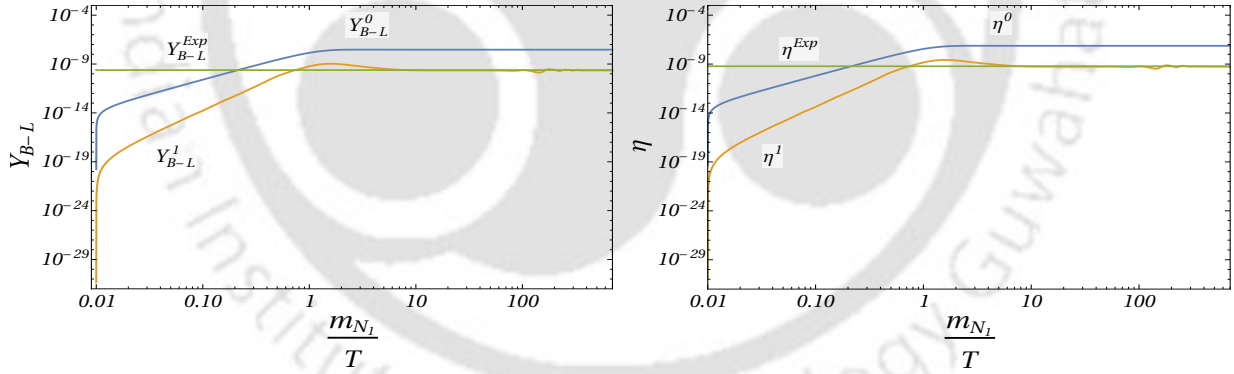


Figure 3.13: The lepton asymmetry, Y_{B-L} (left) and the baryon asymmetry, η (right) versus $z = m_{N1}/T$ for the case with complex κ and $m_{N1} = 10$ TeV, $m_{N2} = 10^3$ TeV, $\kappa_{13} = 10^{-1}$, $\kappa_{12} = 2(1-i)10^{-3}$. The superscripts 0 and 1 correspond to the cases without and with the washout terms induced by the scattering processes, respectively.

the freeze-in mechanism and show that, for a wide range of masses and couplings, the relic abundance is consistent with the experimental data. We then scan this parameter space to find regions favorable to decays of heavy right-handed neutrinos into leptons (or antileptons) plus doublet Higgs bosons (standard leptogenesis), or into the new charged

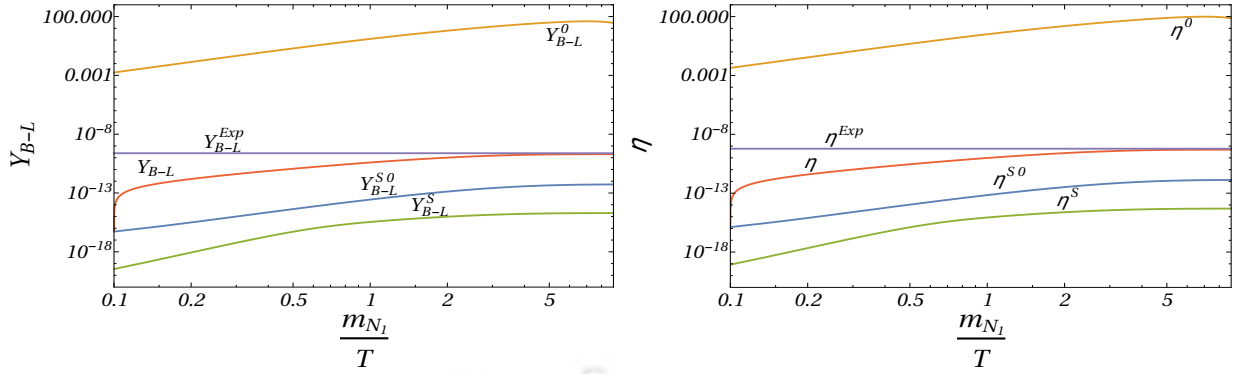


Figure 3.14: Plot for Y_{B-L} and η showing the contribution of the standard leptogenesis and the total contribution in the new scenario for the benchmark point given in Table 3.6. (Left panel): the lepton asymmetry Y_{B-L} . (Right panel): the induced baryon asymmetry η . Here the superscript S indicates the standard contribution, while superscripts 0 and 1 represent contributions without and with scattering, respectively.

scalar S^+ and leptons or antileptons (new contributions). Leptogenesis is generated by the CP-violating interference between the tree-level process and one-loop contributions from self-energy and vertex corrections, facilitated by the non-equilibrium conditions when the temperature of the Universe is of the same order as the mass of the decaying heavy neutrino.



Dark Matter-Driven Low-Scale Leptogenesis via Neutrino Portal

In Chapter 3 we developed a minimal framework in which leptogenesis is realized through feebly interacting dark sector dynamics, while in this chapter we extend this idea to a novel low-scale scenario where both decays and scattering processes involving the dark sector play a central role. In particular we propose a novel framework for low-scale leptogenesis within an extension of the Standard Model (SM) that includes three $SU(2)$ singlet right-handed neutrinos, a singlet vector-like neutral fermion, and a real scalar field. In this setup, the CP asymmetry arises through a rich interplay of mechanisms, including two-body decays of the lightest right-handed neutrino into leptons and Higgs boson or into dark-sector particles, as well as multiple $2 \rightarrow 2$ scattering processes involving visible and dark sector particles. Crucially, the CP -violating phases originate not only from conventional vertex and self-energy corrections of the decay of the right-handed neutrinos but also from novel interference effects in the scattering processes mediated by the dark sector, which significantly enrich the sources of asymmetry. A distinctive feature of our framework is the direct connection between the dark sector and leptogenesis, providing a unified explanation for both the matter–antimatter asymmetry and DM abundance. Moreover, the viable parameter ranges, including the mass spectrum, are within the range that LHC can explore. This Chapter is based on the work published in Phys. Rev. D [251].

4.1 Introduction

The introduction related Leptogenesis has been quite extensively presented in Chapter 1 and Chapter 2. However, to make the Chapter self-contained here we provide a brief overview of the series of work related to the Low-scale Leptogenesis.

The Baryon Asymmetry of the Universe (BAU) and the nature of dark matter are two unsolved problems in cosmology. While the Standard Model (SM) satisfies some conditions for BAU, it cannot fully explain the observed matter-antimatter asymmetry, primarily due to insufficient CP violation and a smooth electroweak phase transition [218, 190, 252–254]. Leptogenesis, a widely studied mechanism, proposes that lepton asymmetry is generated through L -violating decays, which can later be converted into baryon asymmetry via $B + L$ -violating sphaleron processes [219, 200, 219]. Meanwhile, the puzzle of dark matter, which comprises about 85% of the Universe’s matter, has led to various beyond-the-SM theories, with the WIMP [255] and FIMP [256, 179, 214] scenarios being prominent. In particular, FIMPs, which interact weakly with SM particles and are produced non-thermally, evade direct detection constraints [230].

In this Chapter, we propose a unified framework that connects leptogenesis, dark matter, and neutrino mass generation. Our model introduces new fields, such as a singlet vector-like fermion, a neutral scalar, and right-handed neutrinos, and shows how dark sector interactions can enhance CP violation, leading to successful BAU generation at the TeV scale, while simultaneously accounting for dark matter and neutrino masses [241, 257–259, 208, 260–263].

The structure of subsequent part of this Chapter is as follows. In Section 4.2, we present the theoretical framework of the model, including the particle content, symmetries, and relevant interactions. Section 4.3.0.2 is devoted to the dark matter phenomenology, where the WIMP realizations are discussed, and the parameter space consistent with the observed relic abundance is identified. In Section 4.4.2.1, we analyze the mechanism of leptogenesis, highlighting the novel sources of CP violation arising from neutrino–portal interactions with the dark sector, including decay and scattering contributions. The numerical solutions of the coupled Boltzmann equations and a detailed discussion of the generated asymmetries are presented in Section 4.5. Finally, we summarize our results and outline future directions in Section 4.6.

4.2 The Model

The model framework under consideration extends the Standard Model (SM) by introducing three right-handed gauge-singlet neutrinos, denoted by N_i , a singlet vector-like fermion ψ , and a singlet real scalar field ϕ . Both ϕ and ψ are assigned to be odd under a discrete Z_2 symmetry, while all other fields remain even. This discrete symmetry is introduced to ensure the stability of the dark matter (DM) candidate. The gauge symmetry of the model is given by the SM gauge group supplemented with an additional discrete Z_2 symmetry. Depending on the mass hierarchy between ϕ and ψ , the lighter of the two becomes stable and can serve as a viable DM candidate. In addition, we assign a generalized lepton number (global $U(1)_{L'}$ symmetry) to the fermion ψ , the heavy neutrinos N_i , the SM lepton doublet L and the SM lepton singlet e_R , while the scalar ϕ and rest of the Standard Model particles are taken to be lepton-number neutral. This symmetry governs the generation of asymmetries by distinguishing ψ from its antiparticle $\bar{\psi}$. The relevant particle content of this extended model, along with their transformation properties under the SM gauge symmetry, $U(1)_{L'}$ symmetry, and the Z_2 symmetry, is summarized in Table 4.1.

Particles	Spin	$SU(3)_c$	$SU(2)_L$	$U(1)_Y$	$U(1)_{L'}$	Z_2
$N_{1,2,3}$	$\frac{1}{2}$	1	1	0	1	1
ψ	$\frac{1}{2}$	1	1	0	1	-1
ϕ	0	1	1	0	0	-1
L	$\frac{1}{2}$	1	2	$-\frac{1}{2}$	1	1
e_R	$\frac{1}{2}$	1	1	-1	1	1

Table 4.1: The relevant particle content of the model. All the SM particles are Z_2 even. The SM leptons have $U(1)_{L'}$ charge equal to +1, while all the other SM particles have zero $U(1)_{L'}$ charge.

The Lagrangian for this model framework is given by

$$\begin{aligned}
\mathcal{L}_m = & \mathcal{L}_{\text{SM}} + (\partial^\mu \phi)^\dagger (\partial_\mu \phi) + \bar{\psi} i \gamma^\mu \partial_\mu \psi + \bar{N}_i i \gamma^\mu \partial_\mu N_i - m_\psi \bar{\psi} \psi - m_{N_{ij}} \bar{N}_i N_j \\
& - \left(y_i \bar{\psi} \phi N_i + Y_{ij}^N \bar{L}_i \tilde{H} N_j + \text{h.c.} \right) - \lambda' \phi^2 H^\dagger H - \frac{1}{2} m^2 \phi^2 - \zeta \phi^4.
\end{aligned} \tag{4.1}$$

where $L_i = (\nu_{e_i}, e_i)_L^T$ are the left-handed lepton doublets with $e_i = e, \mu, \tau$, and H denotes the SM Higgs doublet field. After electroweak symmetry breaking (EWSB), the

Higgs field acquires a vacuum expectation value (VEV),

$$\langle H \rangle = \begin{pmatrix} 0 \\ \frac{v}{\sqrt{2}} \end{pmatrix} \quad (4.2)$$

while the scalar field ϕ does not develop a VEV, ensuring that the Z_2 symmetry remains an exact symmetry. The Higgs VEV contributes to the mass of the scalar field ϕ , yielding its physical mass square

$$m_\phi^2 = m^2 + \lambda' v^2 \quad (4.3)$$

The exact Z_2 symmetry prevents any mixing between ϕ and the SM Higgs boson. The only interaction connecting the dark scalar to the SM Higgs sector arises from the quartic term $\lambda' \phi^2 H^\dagger H$, which can open an additional Higgs decay channel $H \rightarrow \phi\phi$ for sufficiently light ϕ . However, the coupling λ' can be chosen to be very small, independently of the scalar mass, thereby evading constraints from the invisible Higgs decay searches.

The mass of the vector-like fermion ψ is not generated through symmetry breaking and is entirely determined by the parameter m_ψ . Although ψ does not couple directly to the Standard Model fields, it communicates with the visible sector through interactions mediated by the heavy right-handed neutrinos N_i and the scalar field ϕ . As we shall demonstrate in the subsequent sections, these neutrino-portal interactions play a crucial role in linking the dynamics of the dark sector to the generation of the matter-antimatter asymmetry of the Universe. This interplay between dark matter and leptogenesis constitutes the central focus of the present work.

The model addresses the origin of neutrino masses via the standard type-I seesaw mechanism [264], induced by the Yukawa interaction governed by the couplings Y_{ij}^N . For simplicity, we assume the Majorana mass matrix of the right-handed neutrinos to be diagonal, $m_{N_{ij}} = m_{N_i} \delta_{ij}$, so that the physical masses of the heavy neutrinos are given by m_{N_i} . The seesaw Yukawa couplings Y_{ij}^N are fixed using the Casas–Ibarra parameterization [246].

The right-handed neutrinos N_i , being Majorana fermions, naturally provide the necessary ingredients for leptogenesis through their out-of-equilibrium decays. In the standard type-I seesaw framework, leptogenesis arises solely from the Yukawa interactions between N_i , the SM leptons, and the Higgs doublet. In the present model, this conventional picture is extended by the additional interaction term $\bar{\psi} N_i \phi$, which directly connects the right-handed neutrinos to the dark sector. The associated Yukawa couplings y_i , being unconstrained by low-energy experiments, introduce new sources of CP violation in addition to those present in the standard leptogenesis scenario. As a result, the decay of N_i can

generate asymmetries simultaneously in the visible and dark sectors. Moreover, this interaction opens up a variety of scattering processes involving leptons, dark-sector particles, and heavy neutrinos, which can act both as sources of lepton asymmetry and as channels for transferring asymmetry between the two sectors. In addition, when kinematically allowed, the decay of the scalar field ϕ can itself generate a dark-sector asymmetry, which may subsequently influence the evolution of the lepton asymmetry through these scattering processes. The combined effect of these decay and scattering channels establishes a nontrivial interplay between dark matter dynamics and leptogenesis, which constitutes the central focus of this work.

4.2.1 Collider Constraints

Major collider constraints on the present framework arise from searches for heavy neutral leptons (HNLs), corresponding to the right-handed neutrinos N_i , at the Large Hadron Collider. At the LHC, HNLs can be produced through their mixing with Standard Model leptons, typically via processes such as $pp \rightarrow W^* \rightarrow \ell N$, followed by N decays into leptons and electroweak gauge bosons. Both of the ATLAS and the CMS collaborations have carried out extensive searches for HNLs in prompt and displaced decay channels [265–269]. These searches are sensitive to the active-sterile mixing along the masses of HNLs. For HNL masses between 100 GeV to 20 TeV, the mixing is limited to $< (0.1 - 1)$ [267]. We shall restrict our study to heavy neutrinos with masses $m_{N_i} \gtrsim 2$ TeV and active-sterile mixing angles smaller than $|U_{\ell N}|^2 \lesssim 1.2 \times 10^{-5}$, well below the current sensitivity of LHC searches. Within this regime, collider constraints are naturally evaded, allowing us to focus on the cosmological implications of the model.

We first analyze the framework to determine the regions of parameter space consistent with the observed dark matter relic density. Subsequently, we investigate the viability of leptogenesis within this restricted parameter space, with particular emphasis on the novel contributions arising from neutrino-portal interactions. In this setup, the interplay between the dynamics of the dark sector and lepton number generation plays a central role, enabling successful low-scale leptogenesis while simultaneously satisfying constraints on dark matter.

4.3 Dark Matter

Both ϕ and ψ are odd under an exact discrete Z_2 symmetry, which guarantees their stability and renders them viable dark matter (DM) candidates. When kinematically

allowed, the heavier of the two particles can decay into the lighter one, effectively yielding a single-component DM scenario. On the other hand, if such decays are kinematically forbidden or sufficiently suppressed, a two-component DM framework arises naturally, with both ϕ and ψ contributing to the total dark matter relic abundance.

As discussed in Section 4.2, the quartic coupling between ϕ and the standard Higgs field, λ' is taken to be very small, so as to avoid collider constraints. Thus, the only interaction ϕ and ψ have with the visible sector is through the Yukawa interaction $\bar{\psi}\phi N_i$, controlled by the respective couplings, y_i . For this coupling being in the range of $\mathcal{O}(y_i) < 10^{-10}$, we have FIMP dark matter which follows a non-thermal genesis of the dark matter relic abundance. We have explored this case and found that for a large range of dark matter masses in the GeV to TeV range, the compatible parameters lead to $y_i \sim 10^{-12}$. However, such small couplings do not lead to any significant effect in the leptogenesis. Therefore, we shall not discuss this scenario any further in this work.

On the other hand, the WIMP scenario is of particular interest in this framework, as the presence of sizable couplings can lead to a significant impact of dark matter interactions on leptogenesis. The viable dark matter candidates in this case are dictated by the mass hierarchy among the exotic particles. Accordingly, we consider the different possibilities separately. In the following, we first demonstrate why a single-component dark matter scenario is not viable within this framework. We then systematically analyze the two-component scenarios, identify the corresponding viable regions of parameter space, and discuss their phenomenological consequences.

4.3.0.1 Single-Component Dark Matter

Depending on the mass hierarchy, we may have either a scalar DM scenario or a scotino (fermionic) DM scenario. In the first case of scalar DM scenario, the dark scalar, ϕ is lighter than the new fermion, ψ . The scalar DM candidate, ϕ will remain in thermal equilibrium with the SM particles in the early Universe through its Higgs portal interactions. This makes it a potential WIMP candidate. The annihilation of DM will be controlled by the quartic coupling, $\lambda'\phi^2 H^\dagger H$. However, the same coupling also decides the DM-nucleus scattering in the direct detection experiments. Therefore, it turned out to be quite challenging to deplete the DM sufficiently to avoid an overabundance in DMRD while satisfying the measurements at the direct detection experiments. This is a typical issue with scalar DM and Higgs portal interactions, and the present model framework is no exception.

In the scotino DM scenario, the dark fermion ψ remains in thermal equilibrium with

the Standard Model (SM) particles through processes such as $\psi\psi \rightarrow N_i N_j$, $\phi\phi$. However, for $m_\psi < (m_\phi, m_N)$, no tree-level annihilation channels are available that can efficiently deplete the ψ abundance to reproduce the observed dark matter relic density during thermal freeze-out. Annihilation processes of the type $\psi\psi \rightarrow \text{SM}$ may arise at higher orders; however, these contributions are strongly suppressed by loop factors, the smallness of the Higgs-portal coupling λ' required by collider constraints, and the bounds on the seesaw Yukawa couplings.

In view of the above constraints, we do not pursue the single-component WIMP scenario any further in this work.

4.3.0.2 Co-existing Fermion-Scalar Two-component Dark Matter

In kinematic regions where ϕ and ψ are mass-degenerate or nearly degenerate, the co-annihilation channel $\phi\psi \rightarrow \text{SM}$ becomes significant. In contrast, the annihilation processes $\phi\phi \rightarrow \text{SM}$ and $\psi\psi \rightarrow \text{SM}$ are highly suppressed, such that the dark matter relic abundance predominantly determined by the co-annihilation channel.

The process mediated by the right-handed neutrinos, $\phi\psi \rightarrow \ell H$, is the main co-annihilation channel. The Feynman diagram corresponding to these processes is depicted in Fig. 4.1.

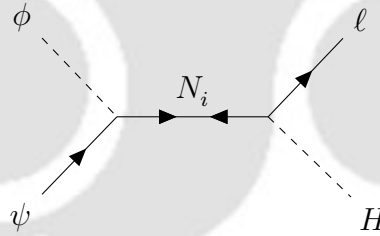


Figure 4.1: Feynman diagram showing the tree-level co-annihilation process $\phi\psi \rightarrow \ell H$.

The temperature evolution of the yields Y_ϕ and Y_ψ corresponding to the fields ϕ and ψ , respectively; is governed by the coupled Boltzmann equations,

$$\mathcal{H}Sx \frac{dY_\phi}{dx} = \gamma^{\text{sc}} \left(1 - \frac{Y_\phi}{Y_\phi^{\text{eq}}} \frac{Y_\psi}{Y_\psi^{\text{eq}}} \right), \text{ and}$$

$$\mathcal{H}Sx \frac{dY_\psi}{dx} = \gamma^{\text{sc}} \left(1 - \frac{Y_\phi}{Y_\phi^{\text{eq}}} \frac{Y_\psi}{Y_\psi^{\text{eq}}} \right),$$

where γ^{sc} denotes the thermally averaged cross section for the processes $\phi\psi \rightarrow \ell H$, $\phi\psi \rightarrow \bar{\ell} H^\dagger$, $\phi\bar{\psi} \rightarrow \ell H$, and $\phi\bar{\psi} \rightarrow \bar{\ell} H^\dagger$ (see Appendix A.3). We have assumed that

the scotino ψ and its antiparticle $\bar{\psi}$ have equal number densities. The slight asymmetry, while important in the leptogenesis, is not relevant in the dark matter dynamics.

As discussed later in the context of leptogenesis, the complex Yukawa coupling y_i associated with the interaction term $\bar{\psi} \phi N_i$ can generate a small asymmetry between the number densities of ψ and $\bar{\psi}$. However, this effect is subdominant for the dark matter dynamics, which is primarily governed by the WIMP freeze-out process. Nevertheless, such an asymmetry in the fermionic dark sector can have important implications for leptogenesis, and this constitutes one of the key features of our framework. A detailed discussion of this aspect is presented in the leptogenesis section.

The remaining parameters in Eq. 4.4 have their standard meanings: $x = m_\phi/T$ denotes the scaled inverse temperature, $\mathcal{H}$ is the Hubble parameter, S is the entropy density of the Universe, and $Y_{\phi,\psi}^{\text{eq}}$ are the equilibrium yields. Since the equilibrium number densities of ϕ and ψ differ primarily by their internal degrees of freedom, the ratios Y/Y^{eq} are approximately equal for the two species. Nevertheless, for completeness, we solve the coupled equations in Eq. 4.4 explicitly.

While the Boltzmann equations determine the late-time comoving yields $Y_{\phi,\psi}$, cosmological observations constrain the present-day dark matter abundance in terms of the dimensionless relic density parameter Ωh^2 . Accordingly, the relic density contributions are expressed as

$$\Omega_{\psi,\phi} h^2 = \frac{S h^2 m_{\psi,\phi}}{\rho_c} Y_{\psi,\phi}, \quad (4.4)$$

where $\rho_c = \frac{3\mathcal{H}^2}{8\pi G}$ is the critical density, with $\mathcal{H} = 100 h \text{ km s}^{-1} \text{ Mpc}^{-1}$ and G denoting Newton's constant. The total dark matter relic abundance is then given by

$$\Omega_{\text{tot}} h^2 = 2 \Omega_\psi h^2 + \Omega_\phi h^2. \quad (4.5)$$

Here $\Omega_\psi h^2$ denotes the contribution from either ψ or $\bar{\psi}$ separately. Since the Boltzmann equation is solved for a single scotino species with $Y_\psi = Y_{\bar{\psi}}$, the total fermionic dark matter contribution carries an explicit factor of two in Eq. 4.5.

Notice that, for a fixed m_{N_i} , the relevant seesaw Yukawa couplings Y_{ij}^N are fixed. Consequently, the co-annihilation cross section for the process $\phi \psi \rightarrow \ell H$ depends only on the common dark matter mass $m_\phi = m_\psi$ and the Yukawa coupling y_i . For a fixed relic abundance (or, equivalently, for a fixed thermally averaged effective annihilation rate) an increase in the dark matter mass must be compensated by a decrease in the coupling y_i , leading to an anti-correlation between $m_{\phi,\psi}$ and y_i . As a result, the perturbative limit $y_i = \sqrt{4\pi}$ corresponds to the minimum value of $m_{\phi,\psi}$ that reproduces the observed dark

matter relic abundance. Further, with m_{N_1} in the TeV range and $m_{N_{2,3}} \gg m_{N_1}$, the effect of N_2 and N_3 mediated processes are negligible in the co-annihilation process. Our analysis indicates that the DM candidate must be in the range of $\mathcal{O}(10^2)$ GeV to satisfy the observed relic density in this scenario. To remind the reader, large couplings y_i required here will not influence the detection prospects at the direct detection experiments, as the process is mediated by heavy neutrinos, which do not have any direct interaction with the nucleus.

Considering benchmark scenarios with m_{N_1} in the TeV range, Table 4.2 presents the lower bound on the dark matter mass corresponding to the perturbative limit of the coupling y_1 . We keep the value of the quartic coupling, λ' fixed to 1.0×10^{-10} all through this study. As mentioned above, this evades the dark matter direct detection constraints, as well as the Higgs sector collider constraints. As an illustrative example, for $m_{N_1} = 2$ TeV one finds $m_{\phi,\psi} \geq 200$ GeV. The sensitivity of the relic abundance to variations in m_{N_1} is also illustrated in the Table 4.2. A decrease in m_{N_1} enhances the co-annihilation cross section for the process $\phi\psi \rightarrow \ell h$, resulting in a smaller relic density. This effect can be compensated either by reducing the $\bar{\psi}\phi N_1$ Yukawa coupling y_1 or by lowering the common dark matter mass $m_\phi = m_\psi$. Conversely, increasing m_{N_1} suppresses the co-annihilation cross section. To maintain the observed relic abundance in this case, the minimum allowed value of $m_\phi = m_\psi$ shifts to larger values, since the coupling cannot exceed the perturbative limit. As mentioned above, for $m_{N_1} \ll m_{N_{2,3}}$, the contributions

BPs	$m_\phi = m_\psi$ [GeV]	m_{N_1} [GeV]	y_1	$\Omega_{\text{tot}} h^2$
BP 1	200	2000	$2\sqrt{\pi}$	0.12
BP 2	200	1000	$2\sqrt{\pi}$	0.04
BP 3	200	3000	$2\sqrt{\pi}$	0.21
BP 4	200	1000	$\sqrt{1.56}$	0.12
BP 5	187	1000	$2\sqrt{\pi}$	0.12
BP 6	209	3000	$2\sqrt{\pi}$	0.12

Table 4.2: Benchmark values of model parameters $m_\phi = m_\psi$, m_{N_1} , and y_1 and the corresponding DM relic abundance $\Omega_{\text{tot}} h^2$. The heavier right-handed neutrino masses are fixed at $m_{N_2} = 10^9$ GeV and $m_{N_3} = 2.6 \times 10^{13}$ GeV, and the quartic coupling is fixed at $\lambda' = 10^{-10}$.

of N_2 and N_3 to the dark matter dynamics are negligible, and their precise masses are therefore not relevant for the present discussion. However, these masses play an important role in leptogenesis, since the corresponding seesaw Yukawa couplings enter the generation of the lepton asymmetry. Accordingly, we fix $m_{N_2} = 10^9$ GeV and $m_{N_3} = 2.6 \times 10^{13}$ GeV.

In the subsequent analysis, we fix $m_{N_1} = 2$ TeV as an illustrative benchmark point, keeping in mind the sensitivity of both the dark matter dynamics and leptogenesis, as discussed later, to this parameter space. As shown above, this choice requires a minimum dark matter mass of $m_{\phi,\psi} \simeq 200$ GeV when the coupling saturates its perturbative limit, $y_1 = \sqrt{4\pi}$. We find that such large couplings generically lead to a lepton asymmetry that is much below the observed value.

Our analysis further indicates that the generated lepton asymmetry is more sensitive to the coupling y_1 than to the dark matter mass, since y_1 enters both the CP asymmetry and the relevant reaction rates. In practice, viable leptogenesis typically requires y_1 in the range 0.01–0.1. Ideally, one would perform a comprehensive scan over the dark matter mass and coupling to identify the region consistent with the relic density, and independently carry out a similar scan for leptogenesis. The overlap of these regions would then determine the parameter space compatible with both dark matter constraints and successful leptogenesis. However, such a global scan is computationally prohibitive.

We therefore adopt a more pragmatic strategy. We fix the dark matter mass to $m_{\phi,\psi} = 250$ GeV, motivated by its proximity to the minimum mass required for the chosen neutrino-sector parameters. The coupling y_1 is then scanned over a coarse range to identify regions corresponding to underabundant and overabundant dark matter. For these values of y_1 , the resulting lepton asymmetry is subsequently computed to assess the viability of this parameter region. The results are presented in Table 4.3. Let us now

BPs	$m_\phi = m_\psi$ [GeV]	m_{N_1} [GeV]	y_1	$\Omega_{\text{tot}} h^2$
BP 7	250	2000	0.0042	0.12
BP 8	250	2000	0.01	0.032
BP 9	250	2000	0.1	0.0033
BP 10	250	2000	0.5	0.0028

Table 4.3: Dark matter relic abundance $\Omega_{\text{tot}} h^2$ as a function of the coupling y_1 for a fixed DM mass of 250 GeV. The heavier right-handed neutrino masses are fixed at $m_{N_2} = 10^9$ GeV and $m_{N_3} = 2.6 \times 10^{13}$ GeV.

briefly examine the sensitivity of the above discussion to variations in m_{N_1} . As shown in Table 4.2, increasing m_{N_1} while keeping it within the few-TeV range modifies the minimum allowed dark matter mass by no more than $\mathcal{O}(10\%)$. For dark matter masses above these minima, one can always identify a corresponding value of the coupling y_1 that reproduces the observed relic abundance. We also note that an underabundant dark matter contribution is phenomenologically acceptable, provided the total dark matter

relic density is not required to be saturated by this model alone.

4.4 Leptogenesis

We now turn our attention to the details of leptogenesis within the framework proposed in this study, where the required lepton number violation and CP asymmetry arise from the decay of heavy Majorana neutrinos $N_{1,2,3}$. In the conventional thermal leptogenesis scenario, such asymmetry is generated via the out-of-equilibrium decays of right-handed Majorana neutrinos into the SM leptons and the Higgs doublet. However, CP violation arises here through the quantum corrections to the decay process, both self-energy corrections and vertex corrections. In the standard scenario, these corrections are decided by the type-1 seesaw interaction [200, 204]. In our framework, $y_i \bar{\psi} \phi N_i$ provides the additional self-energy and vertex correction contributions. This additional interaction modifies the dynamics of leptogenesis in several important ways. Notably, the vertex and self-energy loop corrections contributing to CP violation in the decay processes receive new contributions from diagrams involving the fields ϕ and ψ . These additional contributions can potentially enhance the CP asymmetry compared to the standard scenario, and allow a viable leptogenesis at lower energy scales. Notice that, apart from the standard decay channel, $N_1 \rightarrow \ell H$, the heavy neutrino N_1 decays to the dark sector, $N_1 \rightarrow \phi \psi$ when kinematically allowed. In this study, learning from the dark matter analysis, we restrict ourselves to a degenerate $\phi - \psi$ dark matter, which is always considerably lighter than N_1 . The decay of a heavy neutrino into dark-sector particles also provides a source of CP asymmetry at one loop, very similar to the standard scenario. This asymmetry, however, does not induce any lepton number asymmetry directly. At the same time, this can generate significant asymmetry in the dark sector, which in turn can influence the thermal evolution of the lepton number asymmetry. While strict mass degeneracy of ψ with ϕ disallows ϕ decay, making it part of the dark matter spectrum, sufficient non-zero mass-splitting could allow ϕ decay through 3-body or 4-body final states. In principle, these 3-body and 4-body decays of ϕ could also generate additional CP violation. However, in the rest of this study, we impose strict mass degeneracy of ϕ and ψ to prohibit such 3/4-body decays of ϕ . Other sources of CP violation arise from scattering processes. Scattering channels mediated by right-handed neutrinos can also violate CP symmetry through the interference of amplitudes involving different heavy neutrinos, N_i and N_j with $i \neq j$. The relevant scattering processes in this case are $\phi \psi \rightarrow \ell H$, $\phi \bar{\psi} \rightarrow \ell H$, $\ell \phi \rightarrow H \psi$, $\bar{\ell} \phi \rightarrow H^\dagger \psi$, $\bar{\ell} \psi \rightarrow H^\dagger \phi$, and $\ell \psi \rightarrow H \phi$.

In the standard leptogenesis, scattering processes typically contribute to washout ef-

fects, often significantly reducing the lepton asymmetry generated by CP -violating decays. In the present framework, however, scattering processes play three distinct roles: they act as source terms that generate lepton asymmetry through their intrinsic CP violation, as washout terms that suppress the generated asymmetry, and as asymmetry-transfer processes that redistribute the asymmetry between the dark sector and the leptonic sector, and vice versa.

Before proceeding to the thermal evolution of the yields of the various particle species, we first analyze the CP asymmetries arising in the different processes.

4.4.1 CP Asymmetry

The CP asymmetry arising from the decay of right-handed neutrinos N_i , as discussed in Refs. [200, 270, 223], originates from the interference between tree-level and loop-level diagrams. In the standard case, the loop-level contributions to the decay of right-handed neutrinos arise from diagrams with leptons and the Higgs boson running in the loop. In contrast, within the framework considered here, the CP asymmetry in the decay of N_1 receives contributions not only from the standard lepton–Higgs loops but also from additional quantum corrections involving the dark-sector fields ψ and ϕ . While the general idea of linking dark sector dynamics with leptogenesis has been explored previously, notably in Ref. [271–273], the scenario studied here is substantially different. Our framework does not rely on asymmetric dark matter, but instead considers a multi-component WIMP scenario in which the dark matter abundance is set by thermal freeze-out. Furthermore, scattering processes play a qualitatively new role in our analysis, contributing not only to washout but also directly to the generation and transfer of lepton asymmetry, hence an important distinction of our study in comparison to Ref. [271]. We discuss the various contributions to the CP asymmetry in the following.

4.4.1.1 CP Asymmetry from the decay of $N_1 \rightarrow \ell H$

The CP asymmetry arising from the decay of N_1 is defined as

$$\epsilon_1 = \frac{\Gamma(N_1 \rightarrow \ell H) - \Gamma(N_1 \rightarrow \bar{\ell} H^\dagger)}{\Gamma(N_1 \rightarrow \ell H) + \Gamma(N_1 \rightarrow \bar{\ell} H^\dagger)}, \quad (4.6)$$

where $\Gamma(N_1 \rightarrow \ell H)$, and $\Gamma(N_1 \rightarrow \bar{\ell} H^\dagger)$ are the decay widths of N_1 into leptons and antileptons, respectively, given explicitly by

$$\Gamma_{N_1 \rightarrow \ell H} = \Gamma_{N_1 \rightarrow \ell H}^{tree} (1 + \epsilon_1) \quad (4.7)$$

$$\Gamma_{N_1 \rightarrow \bar{\ell} H^\dagger} = \Gamma_{N_1 \rightarrow \bar{\ell} H^\dagger}^{tree} (1 - \epsilon_1) \quad (4.8)$$

where the tree-level decay width is given by

$$\Gamma_{N_1 \rightarrow \ell H}^{tree} = \Gamma_{N_1 \rightarrow \bar{\ell} H^\dagger}^{tree} = \frac{(Y^{N^\dagger} Y^N)_{11}}{8\pi m_{N_1}^3} (m_{N_1}^2 - m_H^2)^2, \quad (4.9)$$

where m_{N_1} and m_H are the masses of the heavy neutrino and the SM Higgs boson, respectively. The quantum corrections to the decay process at one loop level leading to CP violation are shown in Fig. 4.2. Notice that, in addition to the standard contributions arising through the presence of H and ℓ in the loop, in the present case, the quantum corrections get contributions from the interaction of N_i with the dark sector particles, ϕ and ψ . These additional contributions are proportional to the new Yukawa couplings, y_i . While in the DM dynamics, only the magnitudes of these couplings $|y_i|$ play a role; the complex nature of the couplings y_i is vital for leptogenesis. The CP asymmetry, including

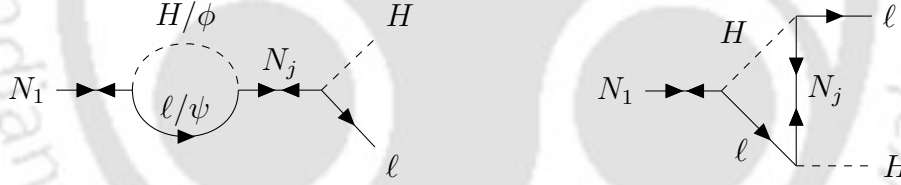


Figure 4.2: Feynman diagrams representing the one-loop corrections to the decay $N_1 \rightarrow \ell H$ including the dark sector contributions to the self-energy.

the contributions arising from the dark sector dynamics, is given below, with total CP asymmetry ϵ_1 getting contributions from the self-energy (ϵ^s) and vertex correction (ϵ^v):

$$\epsilon_1 = \epsilon^s + \epsilon^v, \text{ where} \quad (4.10)$$

$$\epsilon^s = \frac{1}{8\pi K_{11}} \sum_{j=2}^3 \frac{m_{N_1}}{m_{N_1}^2 - m_{N_j}^2} \text{Im} (m_{N_j} K_{1j}^2 + m_{N_1} \kappa_{j1} K_{1j} + m_{N_j} \kappa_{1j} K_{1j}), \text{ and} \quad (4.11)$$

$$\epsilon^v = \frac{1}{8\pi K_{11}} \sum_{j=2}^3 \text{Im} (K_{1j}^2) \mathcal{F} \left(\frac{m_{N_j}^2}{m_{N_1}^2} \right), \quad (4.12)$$

where $K_{ij} = (Y^{N\dagger} Y^N)_{ij}$, and $\kappa_{ij} = (y_i^\dagger y_j)$, and the function

$$\mathcal{F}(x) = \sqrt{x} \left[1 + (1+x) \ln \left(\frac{x}{1+x} \right) \right]$$

is related to vertex loop factor. In the self-energy correction contribution, ϵ^s given in Eq. 4.11 has three different terms. The first term is related to the standard self-energy part of CP asymmetry, whereas the second and the third terms are the novel contributions due to ϕ and ψ in the loop. These additional terms are crucial for low-scale leptogenesis, without requiring mass degeneracy of right handed neutrinos.

4.4.1.2 CP Asymmetry from decay of $N_1 \rightarrow \phi\psi$

As discussed before in Sec. 4.4.1.1, the CP asymmetry generated by the decay of N_1 into the visible sector is influenced by its interaction with the dark sector. At the same time, the decay of N_1 into the dark sector also generates the CP asymmetry, leading to asymmetric generation of ψ and its antiparticle, $\bar{\psi}$. The CP asymmetry in this case is defined as

$$\epsilon_{DM} = \frac{\Gamma(N_1 \rightarrow \phi \psi) - \Gamma(N_1 \rightarrow \phi \bar{\psi})}{\Gamma(N_1 \rightarrow \phi \psi) + \Gamma(N_1 \rightarrow \phi \bar{\psi})},$$

where ϕ is its own antiparticle. The decay widths are explicitly given by

$$\Gamma_{N_1 \rightarrow \phi \psi} = \Gamma_{N_1 \rightarrow \phi \psi}^{tree} (1 + \epsilon_{DM}) \quad (4.13)$$

$$\Gamma_{N_1 \rightarrow \phi \bar{\psi}} = \Gamma_{N_1 \rightarrow \phi \bar{\psi}}^{tree} (1 - \epsilon_{DM}) \quad (4.14)$$

where the tree-level contributions is

$$\Gamma_{N_1 \rightarrow \phi \psi}^{tree} = \Gamma_{N_1 \rightarrow \phi \bar{\psi}}^{tree} = \frac{|y_1|^2}{16\pi m_{N_1}^3} [m_{N_1}^2 + m_\psi^2 - m_\phi^2] \sqrt{(m_{N_1}^2 - m_\phi^2 - m_\psi^2)^2 - 4 m_\phi^2 m_\psi^2},$$

while the interference between the tree-level amplitude and the one-loop contributions gives rise to the CP asymmetry parameter ϵ_{DM} . The Feynman diagrams associated with the one-loop processes are shown in Fig. 4.3. As it is clear, the CP asymmetry in this decay channel potentially induce asymmetry in the number density of ψ and $\bar{\psi}$. Since the lepton number asymmetry is related to the number density of N_1 , it is important to investigate the possible impact of dark sector asymmetry on the visible sector (lepton–antilepton) asymmetry. We will address it in the numerical discussion in section 4.5 where we shall demonstrate that the dark sector asymmetry feeding into the lepton–antilepton asymmetry has a significant role in certain regions of the parameter space. At the same time, we assert

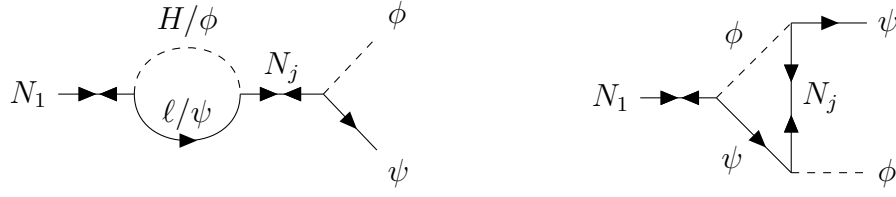


Figure 4.3: Feynman diagrams representing the one-loop corrections to the decay $N_1 \rightarrow \phi \psi$.

that the effect of this asymmetry in dark matter relic density is negligible (in the relevant parameter region) as illustrated in the previous section. Thus, unlike the asymmetric dark matter scenario discussed in the literature [271], we have a thermal WIMP dark matter framework.

The CP asymmetry arising in the decay $N_1 \rightarrow \phi \psi$ is given by

$$\epsilon_{DM} = \epsilon_{DM}^s + \epsilon_{DM}^v, \text{ where} \quad (4.15)$$

$$\epsilon_{DM}^s = \frac{1}{8\pi \kappa_{11}} \sum_{j=2}^3 \frac{m_{N_1}}{m_{N_1}^2 - m_{N_j}^2} \text{Im} (m_{N_j} \kappa_{1j}^2 + m_{N_1} K_{j1} \kappa_{1j} + m_{N_j} K_{1j} \kappa_{1j}), \text{ and} \quad (4.16)$$

$$\epsilon_{DM}^v = \frac{1}{8\pi \kappa_{11}} \sum_{j=2}^3 \text{Im} (\kappa_{1j}^2) \mathcal{F} \left(\frac{m_{N_j}^2}{m_{N_1}^2} \right), \quad (4.17)$$

respectively.

4.4.1.3 CP Asymmetry from $2 \rightarrow 2$ Scattering process $\phi \psi \rightarrow \ell H$

A new contribution to the total lepton asymmetry comes from the scattering process, $\psi \phi \rightarrow H \ell$ mediated by heavy neutrinos $N_{1,2,3}$. Notice that, while Ref. [273] discuss the asymmetry in scattering processes induced through self-energy corrections of the propagator, in our framework the CP asymmetry is generated at the tree-level through interference of two different propagators. The corresponding Feynman diagram is shown in Fig. 4.4. The importance of these CP asymmetry sources is that they affect the dark sector as well as visible sector asymmetries both. Therefore, these channels turn into a bridge between the dark sector and the visible sector asymmetries. An asymmetry generated in the dark sector is transferred to the visible sector and vice-versa. The CP asymmetry generated by scattering process $\psi \phi \rightarrow H \ell$ is parameterized as

$$\epsilon_2 = \frac{\sigma(\psi \phi \rightarrow H \ell) - \sigma(\bar{\psi} \phi \rightarrow H^\dagger \bar{\ell})}{\sigma(\psi \phi \rightarrow H \ell) + \sigma(\bar{\psi} \phi \rightarrow H^\dagger \bar{\ell})}, \quad (4.18)$$

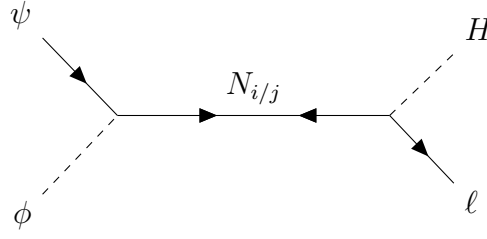


Figure 4.4: Feynman diagrams of the scattering process $\phi \psi \rightarrow \ell H$ ($i \neq j$).

where σ 's denote the cross sections of corresponding processes.

At the tree level, interference between two different scattering $\phi \psi \rightarrow \ell H$ processes that are mediated by different heavy right-handed neutrinos provides the additional sources of CP asymmetry. This phenomenon illustrates how multiple pathways in particle scattering can lead to significant effects in terms of CP violation. Here, we take an approximation where $m_{N_1} \ll m_{N_2} \ll m_{N_3}$. In this scenario, the dominant contribution to the denominator arises from the process that is mediated by the lightest heavy right-handed neutrino N_1 , whereas ϵ_2 itself gets contributions from the interference between N_1 and N_2 mediated channels. In the center of mass frame of the process (with the center of mass energy $\sqrt{s}$), the generated asymmetry is given by

$$\epsilon_2 = \frac{4 \operatorname{Im}[y_1^* y_2 Y_{1\alpha}^N (Y_{2\alpha}^N)^*] \operatorname{Im}[D_1^* D_2]}{|y_1|^2 K_{11} |D_2|^2}, \quad (4.19)$$

where $K_{11} = (Y^{N\dagger} Y^N)_{11}$. The parameters $y_{1,2}$ and $Y_{(1,2)\alpha}^N$, with $\alpha = 1, 2, 3$, denote the Yukawa couplings of N_1 to the dark sector and to the visible sector, respectively. Besides that $D_i = s - m_{N_i}^2 + i m_{N_i} \Gamma_i$, with $i = 1, 2$, denotes the denominators associated with the propagators of different right-handed neutrinos with Γ_i being the total decay width of the mediating heavy neutrino N_i . After substituting D_i into Eq. 4.19, ϵ_2 becomes:

$$\epsilon_2 = \frac{4 \operatorname{Im}[y_1^* y_2 Y_{1\alpha}^N (Y_{2\alpha}^N)^*] [(s - m_{N_1}^2) m_{N_2} \Gamma_2 - (s - m_{N_2}^2) m_{N_1} \Gamma_1]}{|y_1|^2 K_{11} [(s - m_{N_2}^2)^2 + m_{N_2}^2 \Gamma_2^2]}. \quad (4.20)$$

It is clear from Eq. 4.20 that ϵ_2 depends on the center of mass energy $\sqrt{s}$, the Yukawa couplings y and Y , the masses of right-handed neutrinos, m_{N_i} , and the total decay widths of right-handed neutrinos. Here, for m_{N_3} being very heavy in comparison to m_{N_1} and m_{N_2} , the contribution from N_3 is highly suppressed, and therefore has been neglected in our analysis.

4.4.1.4 CP Asymmetry from $2 \rightarrow 2$ Scattering process $\phi \bar{\psi} \rightarrow \ell H$

Another scattering process $\phi \bar{\psi} \rightarrow \ell H$ also violates the CP asymmetry. The relevant Feynman diagram is shown in Fig. 4.5. Similar to the scattering process $\phi \psi \rightarrow \ell H$, this channel also works as a bridge between the dark sector and the visible sector asymmetries. The CP asymmetry generated by the scattering process $\phi \bar{\psi} \rightarrow \ell H$ is parameterized as

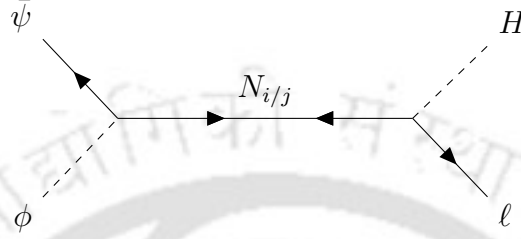


Figure 4.5: Feynman diagrams for the scattering process $\phi^\dagger \bar{\psi} \rightarrow \ell H (i \neq j)$.

$$\epsilon'_2 = \frac{\sigma(\bar{\psi} \phi \rightarrow H \ell) - \sigma(\psi \phi \rightarrow H^\dagger \bar{\ell})}{\sigma(\bar{\psi} \phi \rightarrow H \ell) + \sigma(\psi \phi \rightarrow H^\dagger \bar{\ell})}, \quad (4.21)$$

where σ denotes the cross section of the associated process. The CP asymmetry associated with this process ϵ'_2 originates from the interference of two different N_i mediated processes shown in Fig. 4.5. For being insignificant, the contributions from N_3 has been ignored in this analysis. After putting the cross-sections of relevant processes in Eq. 4.21, ϵ'_2 in the center of mass frame of the process becomes,

$$\epsilon'_2 = \frac{4 \operatorname{Im}[y_1 y_2^* Y_{1\alpha}^N (Y_{2\alpha}^N)^*] \operatorname{Im}[D_1^* D_2]}{|y_1|^2 K_{11} |D_2|^2}, \quad (4.22)$$

where $\alpha = 1, 2, 3$ is summed over, and D_i denotes the denominator of the corresponding propagator. Substituting the explicit expression for this D_i factor, ϵ'_2 is given by

$$\epsilon'_2 = \frac{4 \operatorname{Im}[y_1 y_2^* Y_{1\alpha}^N (Y_{2\alpha}^N)^*] [(s - m_{N_1}^2) m_2 \Gamma_2 - (s - m_{N_2}^2) m_{N_1} \Gamma_1]}{|y_1|^2 K_{11} [(s - m_{N_2}^2)^2 + m_{N_2}^2 \Gamma_2^2]}. \quad (4.23)$$

Notice that, as the total decay width of N_1 depends on the mass spectrum of the dark sector, therefore ϵ_2 and ϵ'_2 indirectly depend on the masses of DM candidates.

4.4.1.5 CP Asymmetry from $2 \rightarrow 2$ Scattering t-channels

Apart from the above CP asymmetry generating $2 \rightarrow 2$ s-channel scattering processes, there are several $2 \rightarrow 2$ t-channel scattering processes that violate CP symmetry. The

relevant Feynman diagrams are listed in Fig. 4.6. Notice that both ϕ and ψ are DM candidates. Certain t-channel scattering processes allow the transfer of asymmetry generated in the dark sector to the visible sector and vice-versa, whereas in a few other t-channel processes, this asymmetry transfer does not occur.

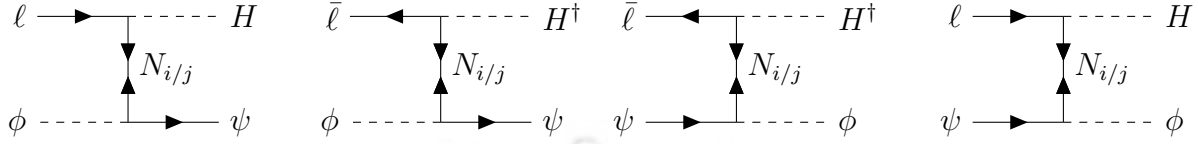


Figure 4.6: Feynman diagrams for the scattering process $\phi \ell \rightarrow \psi H$, $\phi \bar{\ell} \rightarrow \psi H^\dagger$, $\psi \bar{\ell} \rightarrow \phi H^\dagger$ and $\psi \ell \rightarrow \phi H$, where $i \neq j$.

The CP asymmetry parameter generated by scattering processes $\ell \phi \xrightarrow{t} H \psi$, $\bar{\ell} \phi \xrightarrow{t} H^\dagger \psi$, $\bar{\ell} \psi \xrightarrow{t} H^\dagger \phi$ and $\ell \psi \xrightarrow{t} H \phi$ are given by

$$\epsilon_1^t = \frac{\Gamma(\ell \phi \xrightarrow{t} H \psi) - \Gamma(\bar{\ell} \phi \xrightarrow{t} H^\dagger \bar{\psi})}{\Gamma(\ell \phi \xrightarrow{t} H \psi) + \Gamma(\bar{\ell} \phi \xrightarrow{t} H^\dagger \bar{\psi})} = -\frac{\Gamma(\bar{\ell} \psi \xrightarrow{t} H^\dagger \phi) - \Gamma(\ell \bar{\psi} \xrightarrow{t} H \phi)}{\Gamma(\bar{\ell} \psi \xrightarrow{t} H^\dagger \phi) + \Gamma(\ell \bar{\psi} \xrightarrow{t} H \phi)}, \quad (4.24)$$

$$\epsilon_2^t = \frac{\Gamma(\bar{\ell} \phi \xrightarrow{t} H^\dagger \psi) - \Gamma(\ell \phi \xrightarrow{t} H \bar{\psi})}{\Gamma(\bar{\ell} \phi \xrightarrow{t} H^\dagger \psi) + \Gamma(\ell \phi \xrightarrow{t} H \bar{\psi})} = -\frac{\Gamma(\ell \psi \xrightarrow{t} H \phi) - \Gamma(\bar{\ell} \bar{\psi} \xrightarrow{t} H^\dagger \phi)}{\Gamma(\ell \psi \xrightarrow{t} H \phi) + \Gamma(\bar{\ell} \bar{\psi} \xrightarrow{t} H^\dagger \phi)}, \quad (4.25)$$

respectively. The CP asymmetry parameters ϵ_1^t and ϵ_2^t in the center of mass frame of the process are given as follows

$$\epsilon_1^t = \frac{4 \operatorname{Im}[y_1^* y_2 Y_{1\alpha}^N (Y_{2\alpha}^N)^*] \operatorname{Im}[D_1^{t*} D_2^t]}{|y_1|^2 K_{11} |D_2^t|^2}, \text{ and} \quad (4.26)$$

$$\epsilon_2^t = \frac{4 \operatorname{Im}[y_1^* y_2 (Y_{1\alpha}^N)^* Y_{2\alpha}^N] \operatorname{Im}[D_1^{t*} D_2^t]}{|y_1|^2 K_{11} |D_2^t|^2}, \quad (4.27)$$

where $y_{1,2}$ and $Y_{(1,2)\alpha}^N$ with $\alpha = 1, 2, 3$ are Yukawa couplings. With the propagator factor, $D_i^t = t - m_{N_i}^2 + i m_{N_i} \Gamma_i$, the expressions for ϵ_1^t and ϵ_2^t are given by

$$\epsilon_1^t = \frac{4 \operatorname{Im}[y_1^* y_2 Y_{1\alpha}^N (Y_{2\alpha}^N)^*] [(t - m_{N_1}^2) m_{N_2} \Gamma_2 - (t - m_{N_2}^2) m_{N_1} \Gamma_1]}{|y_1|^2 K_{11} [(t - m_{N_1}^2)^2 + m_{N_1}^2 \Gamma_1^2]}, \text{ and} \quad (4.28)$$

$$\epsilon_2^t = \frac{4 \operatorname{Im}[y_1^* y_2 (Y_{1\alpha}^N)^* Y_{2\alpha}^N] [(t - m_{N_1}^2) m_{N_2} \Gamma_2 - (t - m_{N_2}^2) m_{N_1} \Gamma_1]}{|y_1|^2 K_{11} [(t - m_{N_2}^2)^2 + m_{N_2}^2 \Gamma_2^2]}, \quad (4.29)$$

respectively, where m_{N_i} are the right handed neutrino masses of N_i , and Γ_i are the total decay widths of N_i . The imaginary parts of coupling products and denominator of propagator products are responsible for generating the CP violating phase. It is clear that

the complex couplings, mass splittings, and decay widths of intermediate particles play a crucial role in generating CP violating phases. The dependency of ϵ_1^t on right-handed neutrino masses and couplings is the same as that in the case of ϵ_2 .

In total, there are ten CP asymmetry sources in this framework. Two are associated with the decays of N_1 . Another six CP asymmetry sources are connected to the scattering processes (two from s-channel processes and rest four from t-channel processes). The last two CP asymmetry sources associated with three-body decay and four-body decay of ϕ . These last two sources of CP asymmetry are not relevant in our analysis due to mass degeneracy of m_ϕ and m_ψ .

4.4.2 Time Evolution of the Number Densities of Particles

The time evolution of the number densities of N_1 , ϕ , ψ , and $\bar{\psi}$ is crucial for studying the asymmetries in both the dark sector and the visible sector. The evolution of the particle number densities when the system is out of thermal equilibrium is governed by the Boltzmann equations. In scenarios where the masses of the right-handed neutrino (RHN) N_i satisfy the hierarchy $m_{N_3} \gg m_{N_2} \gg m_{N_1}$, any lepton asymmetry generated by the decays of the heavier neutrinos N_3 and N_2 at high temperatures is efficiently washed out before the EWSB. As a result, the lepton asymmetry that survives to lower temperatures is predominantly produced by the decay of the lightest right-handed neutrinos, N_1 . The evolution of the number density of N_1 is influenced by several processes, including its decay and scattering processes. Similarly, the number density of the leptons is decided by processes involving their production (like N_1 decay and scattering processes), and annihilation (like the inverse decay and scattering processes). The processes relevant to the visible sector asymmetry and the dark sector asymmetry in our case are presented in Appendix A.1. These interactions collectively determine the abundance of N_1 , ϕ , ψ , $\bar{\psi}$ and the resulting lepton asymmetry. The generated lepton asymmetry can subsequently be changed into the observed baryon asymmetry of the Universe via sphaleron processes [209]. The relevant Boltzmann equations for yields $Y_{N_1} = \frac{n_{N_1}}{s}$, $Y_L = \frac{n_l - n_{\bar{l}}}{s}$ and $Y_{\Delta\psi} = \frac{n_\psi - n_{\bar{\psi}}}{s}$ are

as follows:

$$\mathcal{H}S z \frac{dY_{N_1}}{dz} = \left(1 - \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}}\right) (\gamma_{D_1} + \gamma_{D_2} + 2\gamma_s + 2\gamma_{t_1} + 2\gamma_{t_2} + 2\gamma_V), \quad (4.30)$$

$$\begin{aligned} \mathcal{H}S z \frac{dY_L}{dz} = & \left(\frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} - 1\right) \epsilon_1 \gamma_{D_1} + 2\Theta_{\epsilon_2}^{\phi\psi} + 2\Theta_{\epsilon_2'}^{\phi\bar{\psi}} - 2\Theta_{\epsilon_1'}^{\phi\ell} - 2\Theta_{\epsilon_1'}^{\psi\bar{\ell}} + 2\Theta_{\epsilon_2'}^{\phi\bar{\ell}} + 2\Theta_{\epsilon_2'}^{\psi\ell} - \gamma_D^{\text{os}} \epsilon_1 \\ & + \frac{Y_{\Delta\psi}}{Y_\psi^{\text{eq}}} \left(\frac{\gamma_{\phi\psi}}{2} + \frac{\gamma'_{\phi\psi}}{2} + \frac{\gamma_{\phi\ell}^t}{2} - \frac{\gamma_{\phi\bar{\ell}}^t}{2} \right) \\ & - \frac{Y_L}{Y_L^{\text{eq}}} \left(\frac{\gamma_{D_1}}{2} + \gamma_{t_1} + \gamma_{t_2} + \gamma_{V_1} + \gamma_{V_3} + \gamma_{V_4} + \gamma_{V_5} + \frac{\gamma_{\phi\psi}}{2} + \frac{\gamma'_{\phi\psi}}{2} + \frac{\gamma_{\phi\ell}^t}{2} + \frac{\gamma_{\phi\bar{\ell}}^t}{2} + \gamma_{\psi\bar{\ell}}^t \right. \\ & \left. + \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} (\gamma_s + \gamma_{V_2} + \gamma_{V_6}) - \gamma_D^{\text{os}} \right), \text{ and} \end{aligned} \quad (4.31)$$

$$\begin{aligned} \mathcal{H}S z \frac{dY_{\Delta\psi}}{dz} = & \left(\frac{Y_N}{Y_N^{\text{eq}}} - 1\right) \epsilon_{DM} \gamma_{D_2} - 2\Theta_{\epsilon_2}^{\phi\psi} + 2\Theta_{\epsilon_2'}^{\phi\bar{\psi}} + 2\Theta_{\epsilon_1'}^{\phi\ell} + 2\Theta_{\epsilon_1'}^{\psi\bar{\ell}} + 2\Theta_{\epsilon_2'}^{\phi\bar{\ell}} + 2\Theta_{\epsilon_2'}^{\psi\ell} + \gamma_D^{\text{os}} \epsilon_1 \\ & + \frac{Y_L}{Y_L^{\text{eq}}} \left(\frac{\gamma_{\phi\psi}}{2} - \frac{\gamma'_{\phi\psi}}{2} + \frac{\gamma_{\phi\ell}^t}{2} - \gamma_{\phi\bar{\ell}}^t - \gamma_D^{\text{os}} \right) \\ & - \frac{Y_{\Delta\psi}}{Y_\psi^{\text{eq}}} \left(\frac{\gamma_{D_2}}{2} + \frac{\gamma_{\phi\psi}}{2} + \frac{\gamma'_{\phi\psi}}{2} + \frac{\gamma_{\phi\ell}^t}{2} + \frac{\gamma_{\psi\bar{\ell}}^t}{2} + \frac{\gamma_{\phi\bar{\ell}}^t}{2} + \frac{\gamma_{\psi\ell}^t}{2} \right), \end{aligned} \quad (4.32)$$

where γ_{D_1} and γ_{D_2} are the thermal averaged decay widths of $N_1 \rightarrow \ell H$ and $N_1 \rightarrow \phi \psi$, respectively. γ_{t_1} is the thermal average of the processes $N_1 u \leftrightarrow d \bar{\ell}$ and $N_1 \bar{u} \leftrightarrow \bar{d} \ell$, and γ_{t_2} is the same for the processes $N_1 d \leftrightarrow u \ell$ and $N_1 \bar{d} \leftrightarrow \bar{u} \bar{\ell}$. We denote the sum of these contributions as $\gamma_t = \gamma_{t_1} + \gamma_{t_2}$. Similarly, γ_s corresponds to the s -channel processes $\ell N_1 \leftrightarrow d \bar{u}$, and $\bar{\ell} N_1 \leftrightarrow \bar{d} u$, while γ_{V_i} , where $i = 1, \dots, 6$, are the thermal averaged cross sections of scattering processes involving the gauge bosons: $N_1 H \xleftrightarrow{s} V_\mu \ell$, $N_1 \ell \xleftrightarrow{s} V_\mu H$, $N_1 V_\mu \xleftrightarrow{t} H \ell$, $N_1 V_\mu \xleftrightarrow{u} H \ell$, $N_1 H^\dagger \xleftrightarrow{t} V_\mu \ell$, $N_1 \ell \xleftrightarrow{t} V_\mu H^\dagger$, and their charged conjugate processes in that order, with $V_\mu = W_\mu, Z_\mu$. The quantities $\gamma_{\phi\psi}$, $\gamma'_{\phi\psi}$, $\gamma_{\phi\ell}^t$, $\gamma_{\phi\bar{\ell}}^t$, $\gamma_{\psi\bar{\ell}}^t$, γ_{LL} , γ_{HL} and $\gamma_{\psi\ell}^t$ represent the thermally averaged scattering cross-sections of the processes $\phi\psi \rightarrow \ell H$, $\phi\bar{\psi} \rightarrow \ell H$, $\phi\ell \rightarrow \psi H$, $\phi\bar{\ell} \rightarrow \psi H^\dagger$, $\psi\bar{\ell} \rightarrow \phi H^\dagger$, $\ell\ell \rightarrow H^\dagger H^\dagger$, $H\ell \rightarrow H^\dagger \bar{\ell}$ and $\psi\ell \rightarrow \phi H$, respectively (as given in the Appendix [Eq. A.4]). While $\Theta_{\epsilon_\alpha}^{ab}$, given in Appendix [Eq. A.6], represent the thermal average of $\epsilon_\alpha \sigma^{ab}$ corresponding to the respective scattering process. In addition, $\gamma_D^{\text{os}} = \gamma_D \mathcal{B}_1 \mathcal{B}_2$, where γ_D is the thermal average of the total decay width of N_1 , whereas $\mathcal{B}_1$ and $\mathcal{B}_2$ correspond to the branching fractions for $N_1 \rightarrow \ell H$ and $N_1 \rightarrow \phi\psi$, respectively. These terms represent the on-shell contributions explicitly subtracted from the scattering processes $\phi\psi \rightarrow \ell H$ and $\phi\bar{\psi} \rightarrow \ell H$, as well as the corresponding conjugate channels, in order to avoid double counting of the real intermediate N_1 state already included through decay and inverse-

decay contributions. For the range of parameters considered in this work, these on-shell subtraction terms are numerically negligible compared to the other reaction rates. In Eq. 4.31, the terms proportional to $Y_{\Delta\psi}$ indicate the transfer of dark-sector asymmetry into the lepton asymmetry. Similarly, in Eq. 4.32, the terms involving Y_L represent the transfer of visible sector asymmetry into the dark sector. The strength of the transfer of asymmetry from the dark sector to the visible sector and vice versa depends on the choice of parameter values. In a simplified scenario, if the scattering processes involving dark matter occur in such a way that the numbers of dark sector particles ψ and $\bar{\psi}$ remain equal during the evolution, then no dark sector asymmetry is generated, nor can any dark-sector asymmetry be transferred to the visible sector. Under this condition, only two equations are required to analyze the resulting lepton asymmetry. As described below, this simplified case can be solved analytically in a schematic manner.

4.4.2.1 Case when asymmetry in dark sector is absent

We shall explore the specific case when dark sector asymmetry is negligible. Consequently, the abundances of ψ and $\bar{\psi}$ remain equal ($Y_\psi(z) = Y_{\bar{\psi}}(z)$). Under this condition, Eq. 4.30 and Eq. 4.31 simplify to

$$\frac{dY_{N_1}}{dz} + P_N(z)Y_{N_1} = Q_N(z), \quad (4.33)$$

$$\frac{dY_L}{dz} + P_L(z)Y_L = Q_L(z), \quad (4.34)$$

where $P_N(z)$, $Q_N(z)$, $P_L(z)$ and $Q_L(z)$ are defined in Eq. 4.35, 4.36, 4.37 and 4.38 respectively.

$$P_N(z) = \frac{\gamma_{D_1} + \gamma_{D_2} + 2\gamma_s + 2\gamma_{t_1} + 2\gamma_{t_2} + 2\gamma_V}{\mathcal{H}S z Y_{N_1}^{\text{eq}}}, \quad (4.35)$$

$$Q_N(z) = Y_{N_1}^{\text{eq}} P_N(z), \quad (4.36)$$

$$P_L(z) = \frac{1}{\mathcal{H}S z Y_L^{\text{eq}}} \left[\frac{\gamma_{D_1}}{2} + \gamma_{t_1} + \gamma_{t_2} + \gamma_{V_1} + \gamma_{V_3} + \gamma_{V_4} + \gamma_{V_5} + \frac{\gamma_{\phi\psi}}{2} + \frac{\gamma'_{\phi\psi}}{2} + \frac{\gamma_{\phi\ell}^t}{2} + \frac{\gamma_{\phi\bar{\ell}}^t}{2} + \gamma_{\psi\bar{\ell}}^t + \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} (\gamma_s + \gamma_{V_2} + \gamma_{V_6}) - \gamma_D^{\text{os}} \right], \text{ and} \quad (4.37)$$

$$Q_L(z) = \frac{\epsilon_1 \gamma_{D_1}}{\mathcal{H}S z} \left(\frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} - 1 \right) + \frac{2}{\mathcal{H}S z} \left[\Theta_{\epsilon_2}^{\phi\psi} + \Theta_{\epsilon'_2}^{\phi\bar{\psi}} - \Theta_{\epsilon_1^t}^{\phi\ell} - \Theta_{\epsilon_1^t}^{\psi\bar{\ell}} + \Theta_{\epsilon_2^t}^{\phi\bar{\ell}} + \Theta_{\epsilon_2^t}^{\psi\ell} \right] - \gamma_D^{\text{os}} \epsilon_1$$

$$= \epsilon_1 Q_1(z) + Q_{\epsilon_2}(z) + Q_{\epsilon'_2}(z) + Q_{\epsilon_1^t}(z) + Q'_{\epsilon_1^t}(z) + Q_{\epsilon_2^t}(z) + Q'_{\epsilon_2^t}(z). \quad (4.38)$$

In Eq. 4.38, the individual source terms $Q_i(z)$ are defined implicitly through decomposition. This simple form, allows us to solve Eq. 4.33 and Eq. 4.34 formally with the solutions

$$Y_{N_1}(z) = \frac{1}{\mu(z)} \left[Y_{N_1}^{\text{eq}}(z_0) + \int_{z_0}^z \mu(s) Q_N(s) ds \right], \text{ and} \quad (4.39)$$

$$Y_L(z) = \epsilon_1 \kappa_1 + \kappa_2 + \kappa_3 + \kappa_4 + \kappa_5 + \kappa_6 + \kappa_7, \quad (4.40)$$

respectively, where $\mu(z) = \exp \left(\int_{z_0}^z P_N(z') dz' \right)$. The definitions of the various efficiency factors κ_i are given by

$$\begin{aligned} \kappa_1 &= \frac{1}{\nu(z)} \int_{z_0}^z \nu(z') Q_1(z') dz', & \kappa_2 &= \frac{1}{\nu(z)} \int_{z_0}^z \nu(z') Q_{\epsilon_2}(z') dz', \\ \kappa_3 &= \frac{1}{\nu(z)} \int_{z_0}^z \nu(z') Q_{\epsilon'_2}(z') dz', & \kappa_4 &= \frac{1}{\nu(z)} \int_{z_0}^z \nu(z') Q_{\epsilon_1^t}(z') dz', \\ \kappa_5 &= \frac{1}{\nu(z)} \int_{z_0}^z \nu(z') Q'_{\epsilon_1^t}(z') dz', & \kappa_6 &= \frac{1}{\nu(z)} \int_{z_0}^z \nu(z') Q_{\epsilon_2^t}(z') dz', \\ \kappa_7 &= \frac{1}{\nu(z)} \int_{z_0}^z \nu(z') Q'_{\epsilon_2^t}(z') dz', & & \end{aligned} \quad (4.41)$$

where each κ_i represents the accumulated contribution of the corresponding source term $Q_i(z)$, while being weighted by the integrating factor $\nu(z)$ given by

$$\nu(z) = \exp \left(\int_{z_0}^z P_L(z') dz' \right). \quad (4.42)$$

These efficiency factors encode the different sources of matter-antimatter asymmetry. Specifically, κ_1 contributes to the asymmetry only when Y_{N_1} deviates from its equilibrium value $Y_{N_1}^{\text{eq}}$. The factors κ_2 and κ_3 arise from the s -channel $2 \rightarrow 2$ scattering processes $\phi\psi \rightarrow \ell H$ and $\phi\bar{\psi} \rightarrow \ell H$, respectively; their contributions are maximized when the CP -violating combinations $\epsilon_2\sigma^{\phi\psi}$ and $\epsilon'_2\sigma^{\phi\bar{\psi}}$ are largest. Similarly, κ_4 and κ_5 originate from the t -channel processes $\ell\phi \rightarrow H\psi$ and $\bar{\ell}\psi \rightarrow \phi H^\dagger$, respectively, while κ_6 and κ_7 stem from the related t -channel processes $\bar{\ell}\phi \rightarrow H^\dagger\psi$ and $\ell\psi \rightarrow H\phi$, respectively. The contributions from κ_4 , κ_5 , κ_6 , and κ_7 are maximized when the respective CP -violating scattering rates $\epsilon_1^t\sigma^{\phi\ell}$, $\epsilon_1^t\sigma^{\bar{\ell}\psi}$, $\epsilon_2^t\sigma^{\phi\bar{\ell}}$, and $\epsilon_2^t\sigma^{\ell\psi}$ are the largest. Finally, the CP asymmetry parameter ϵ_1 and all efficiency factors $\kappa_{1,2,\dots,7}$ depend on the Yukawa couplings $y_{1,2,3}$, the heavy neutrino masses $m_{N_{1,2,3}}$, and the DM masses m_ϕ and m_ψ . Consequently, the lepton asymmetry Y_L depends solely on these free parameters.

In Eq. 4.31 and Eq. 4.32, the parameters ϵ_2 , ϵ'_2 , ϵ_1^t , and ϵ_2^t appear linearly. The CP asymmetry parameters ϵ_1 and ϵ_{DM} are associated with the decays of RHNs, which behave as a source term for the visible sector and dark sector asymmetry, respectively. While the scattering processes serves as sources to generate asymmetry in both the dark sector and the visible sector, they also acts as a washout processes, as well as transferring asymmetry from the dark sector to the visible sector and vice-versa. The CP asymmetry sources ϵ_2 and ϵ'_2 contribute inversely to each other in the leptonic sector.

4.5 Numerical Analysis and results

In this section, we perform a numerical analysis of the lepton asymmetry (visible sector) and DM asymmetry generated by decays and scattering processes involving the heavy right-handed neutrinos. These asymmetries depend on the free parameters y_i , m_{N_i} (for $i = 1, 2, 3$), m_ϕ , and m_ψ . The CP asymmetry parameters discussed in Sec. 4.4.1 play a crucial role in the generation of these asymmetries. In this framework, there are eight such parameters and we have considered only ϵ_1 , ϵ_{DM} , ϵ_2 , ϵ'_2 , ϵ_1^t , and ϵ_2^t in our investigation as the CP asymmetries from three-body and four-body decays of ϕ are not allowed here due to mass degeneracy in the dark sector ($m_\phi = m_\psi$). The CP asymmetry parameter ϵ_1 and ϵ_{DM} depend on y_i and m_{N_i} , while other CP asymmetry parameters ϵ_2 , ϵ'_2 , ϵ_1^t

B.P.	m_{N_1} [GeV]	m_{N_2} [GeV]	m_{N_3} [GeV]	y_1	y_2	y_3
(a)	2000	10^9	26×10^{12}	$0.0059(1 - i)$	$0.20(1 + i)$	$0.3\sqrt{2\pi}(1 - i)$
(b)	2000	10^9	26×10^{12}	$\frac{0.1}{\sqrt{2}}(1 - i)$	$0.20(1 + i)$	$0.3\sqrt{2\pi}(1 - i)$
(c)	2000	1.5×10^7	26×10^{12}	$0.12(1 - i)$	$0.20(1 + i)$	$0.3\sqrt{2\pi}(1 - i)$

Table 4.4: Fixed parameter points for η vs z plot Fig 4.7.

and ϵ_2^t depend on the decay width of N_1 in addition to y_i , and m_{N_i} . We discuss the visible sector asymmetry study in sub-section 4.5.1 and dark sector asymmetry study in sub-section 4.5.2 below.

4.5.1 Numerical Analysis of visible sector asymmetry

As explained in the previous sections, in the present framework, the lepton asymmetry is influenced by the dark sector in multiple ways. Our numerical study elaborated in this section will categorically establish that this influence is significant enough to have sufficiently large leptogenesis with TeV-scale N_1 . We shall illustrate this with the results presented in Fig. 4.7, where the baryon number asymmetry, η is given as a function of $z = \frac{m_{N_1}}{T}$. Parameter values used for this study are given in Table 4.4. The couplings, y_2 and y_3 as well as the masses of the heavy neutrinos, N_2 and N_3 are fixed in both the cases described in the first two plots. In Fig. 4.7(a), the observed baryon asymmetry is achievable for $m_\phi = m_\psi = 200$ GeV for a corresponding $y_1 = 0.0059(1 - i)$. However, notice that this case correspond to over-abundant dark matter scenario. The red curve corresponding to $m_\phi = m_\psi = 250$ GeV for the same coupling produces larger than observed baryon asymmetry, while being consistent with the dark matter abundance. Fig. 4.7(a) also demonstrates the significance of dark matter asymmetry (η_{DM}) in leptogenesis. As expected, the when $\eta_{DM} = 0$, the baryon asymmetry is insensitive to the dark matter mass. One can conclude that, for $m_{N_1} = 2$ TeV, and for the other right-handed neutrino masses and seesaw Yukawa's fixed as in this case, the dark matter mass around or below 200 GeV are not consistent with both dark matter observations and required baryon asymmetry. Keeping the dark matter mass a little higher, at 250 GeV, it is possible to adjust the coupling y_1 to get the right value of η , while being consistent with the dark matter observations, as illustrated in Fig. 4.7(b). The coupling required in this case is $y_1 = \frac{0.1}{\sqrt{2}}(1 - i)$. Whereas, Fig. 4.7(c) demonstrates that the dark matter asymmetry is quite sensitive to the value of m_{N_2} . At $m_{N_2} = \mathcal{O}(10^7)$ GeV, the effect of η_{DM} is negligible. However, suitable couplings provide the required baryon asymmetry while being consistent with the dark matter observations for a dark matter mass of 250 GeV.

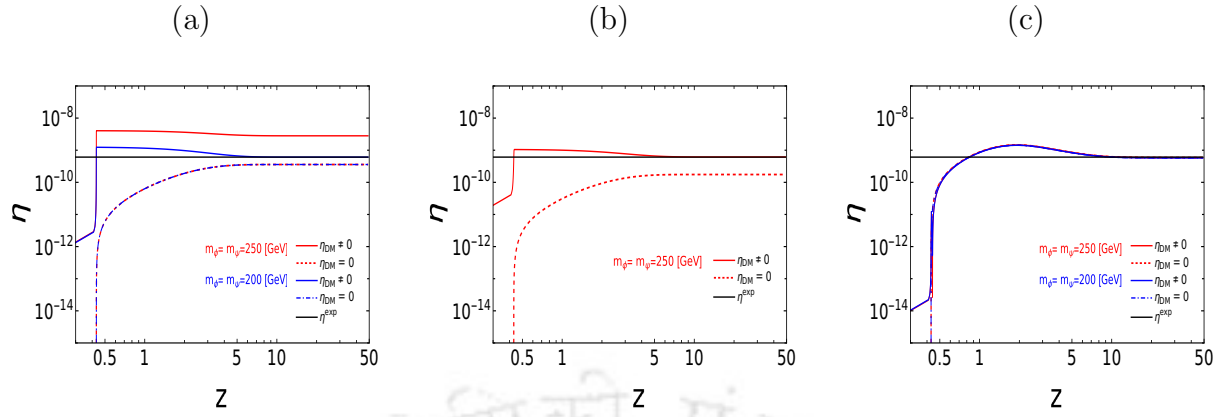


Figure 4.7: The baryon asymmetry parameter η as a function of the dimensionless variable $z = \frac{m_{N_1}}{T}$. The parameter values considered for these studies are given in Table 4.4.

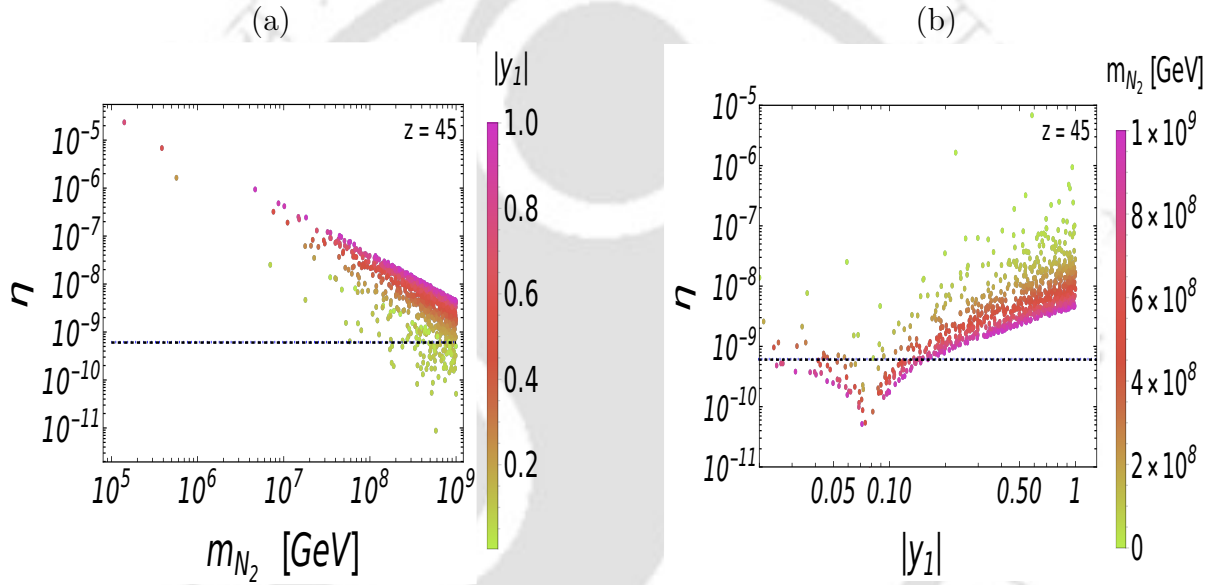


Figure 4.8: Scan of the baryon asymmetry parameter $\eta[z = 45]$: (a) as a function of the right-handed neutrino mass m_{N_2} , with color variation representing the Yukawa coupling $|y_1|$, (b) as a function of $|y_1|$, with color variation representing m_{N_2} . The parameter values considered for the scan are given in Table 4.5.

Although leptogenesis depends on the parameters $y_1, y_2, y_3, m_\phi, m_\psi, m_{N_1}, m_{N_2}, m_{N_3}$, the resulting asymmetry η is particularly sensitive to just two of them: y_1 and m_{N_2} . Consequently, a broad range of (y_1, m_{N_2}) values can reproduce the observed η . This is explicitly shown in the scans over the parameter space as presented in Fig. 4.8, Fig. 4.9, and Fig. 4.10.

The scans were carried out over the ranges $\frac{10^{-3}(1-i)}{\sqrt{2}} < y_1 < \frac{(1-i)}{\sqrt{2}}$ and $8 \times 10^4 \text{ GeV} < m_{N_2} < 10^9 \text{ GeV}$, while other parameter were fixed at values given in Table 4.5. For each point in this parameter space, we solved the Boltzmann Eqs. 4.30–4.32 and extracted the

m_{N_1} [GeV]	m_{N_3} [GeV]	m_ϕ [GeV]	m_ψ [GeV]	y_2	y_3
2000	26×10^{12}	250	250	$0.200781(1+i)$	$0.3\sqrt{2\pi}(1-i)$
$8 \times 10^4 < m_{N_2} < 10^9$			$\frac{10^{-3}}{\sqrt{2}}(1-i) < y_1 < \frac{1}{\sqrt{2}}(1-i)$		

Table 4.5: Parameter values considered for the scan results presented in Fig. 4.8, Fig. 4.9, and Fig. 4.10.

corresponding value of η at $z = 45$. The left panel in Fig. 4.8 shows η as a function of m_{N_2} with the color scale encoding $|y_1|$. The right panel in Fig. 4.8 shows η as a function of $|y_1|$, with the color scale encoding m_{N_2} . Both panels include horizontal lines marking the observed value of the BAU η_{obs} in 3σ limit. Corresponding to a large value of m_{N_2} , a combination with a small value of y_1 can provide the observed value of η . It is evident that $\eta \propto \frac{1}{m_{N_2}}$ with $\eta > 0$. This behavior typically arises from the CP asymmetry parameters. For a large value of m_{N_2} , a smaller value of y_1 leads to a smaller value of η , whereas a larger value of y_1 results in a larger value of η . For small values of $|y_1|$, η remains suppressed, whereas it increases significantly at larger values of y_1 .

4.5.2 Numerical Analysis of Dark Sector Asymmetry

We investigate the numerical behavior of the dark sector asymmetry and its possible impact on the visible-sector asymmetry. Our analysis indicates that, for certain choices of parameters, the dark sector asymmetry can significantly affect the visible sector asymmetry. This is particularly so when the dark sector asymmetry is sizable at early times (small values of z). For other regions of parameter space, however, the visible sector asymmetry remains largely insensitive to the dark-sector contribution. To study this interplay, we compute the dark sector asymmetry $Y_{\Delta\psi}$ at three representative values of the evolution parameter z : $z = 1$, $z = 10$, and $z = 45$, as shown in Fig. 4.9. The parameter values used in this analysis are listed in Table 4.5. For each point in parameter space, we numerically solve Eq. 4.32 to obtain the corresponding values of $Y_{\Delta\psi}$ at the three chosen values of z . We then perform parameter scans of $Y_{\Delta\psi}$ as a function of m_{N_2} and $|y_1|$ at $z = 1$, $z = 10$, and $z = 45$, as illustrated in Fig. 4.9.

The dark-sector asymmetry at $z = 1$ and $z = 10$ is both sizable and nearly identical. In contrast, at later times ($z = 45$) the asymmetry becomes strongly suppressed. For all values of z , the asymmetry $Y_{\Delta\psi}$ increases as the mass of the right-handed neutrino m_{N_2} decreases. The dependence of $Y_{\Delta\psi}$ on the Yukawa coupling $|y_1|$, however, varies with z and exhibits distinct behaviors at different stages of the evolution. In this analysis, the mass of the lightest right-handed neutrino is fixed at $m_{N_1} = 2000$ GeV. When the

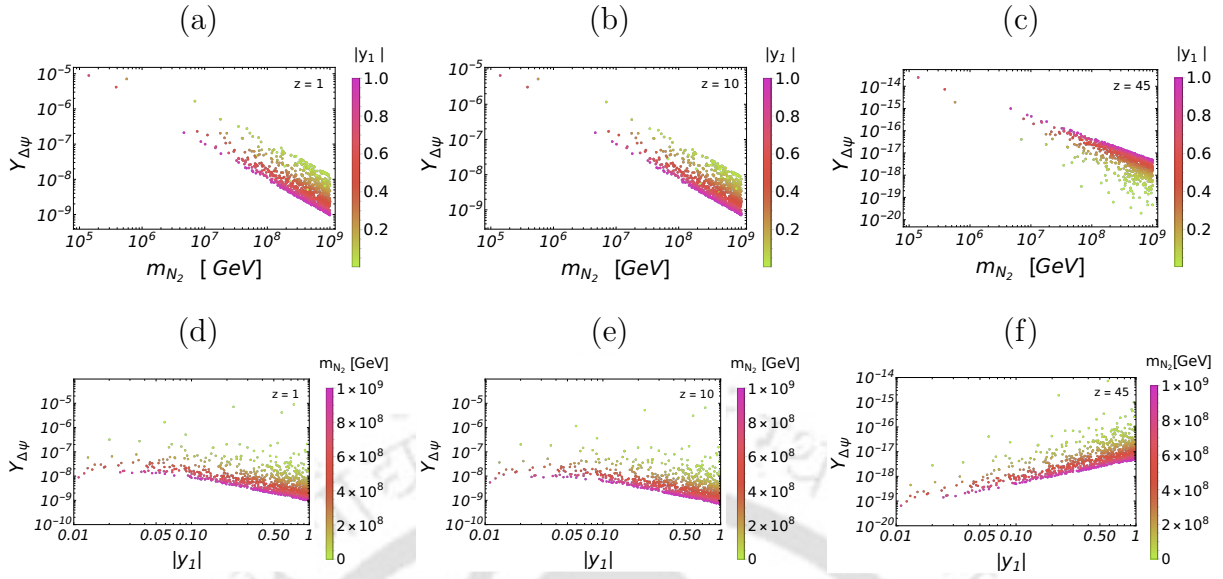


Figure 4.9: Dark-sector asymmetry $Y_{\Delta\psi}$: (a)–(f) as described in the text, for the benchmark point given in Table 4.5.

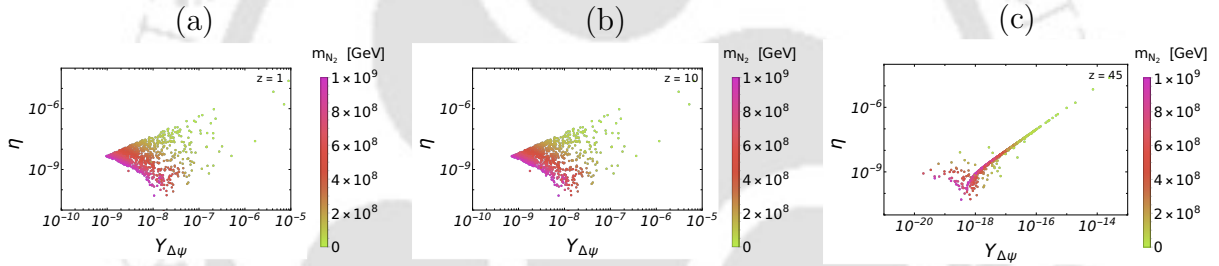


Figure 4.10: Scan of the dark sector asymmetry $Y_{\Delta\psi}$ as a function of η for an alternative benchmark point given in Table 4.5.

mass m_{N_2} approaches m_{N_1} , the CP -asymmetry parameter ϵ_{DM} increases sharply, leading to an enhancement of the dark sector asymmetry. As m_{N_2} moves further away from m_{N_1} , the dark sector asymmetry decreases steadily, as shown in subfigures (a), (b), and (e) of Fig. 4.9. An increase in the Yukawa coupling y_1 enhances both the decay rate of the process $N_1 \rightarrow \phi\psi$ and the associated scattering cross sections. As a result, the dark sector asymmetry initially grows because the parameter ϵ_{DM} itself increases with y_1 . Beyond a critical value of y_1 , however, washout processes become dominant and suppress the asymmetry. This leads to a characteristic turnover, after which the dark sector asymmetry decreases, as seen in subfigures (d), and (e) of Fig. 4.9.

Although the dark sector asymmetry can affect the visible sector asymmetry, as illustrated in Fig. 4.7 and further supported by the scatter plot in Fig. 4.10, the transfer of asymmetry from the dark sector to the visible sector is most efficient at early times,

namely at $z = 1$ and $z = 10$. At later times ($z = 45$), the transfer becomes highly suppressed and has a negligible impact on the dark matter relic density and visible sector asymmetry. This is because dark matter decouples at a later epoch, by which time the number densities of ψ and $\bar{\psi}$ are approximately equal.

4.6 Conclusion

In this Chapter, we have presented an analysis of a simple extension of the SM that accommodates both a dark matter candidate and a viable mechanism for leptogenesis in the early Universe. The model introduces three Majorana right-handed neutrinos, a Z_2 -odd real scalar field ϕ , and a vector-like neutral fermion ψ . In this setup, both ϕ and ψ are stable and serve as stable dark matter particles. To distinguish the scotino (ψ) from its antiparticle, we introduce a global $U(1)_{L'}$ corresponding to a generalized lepton number. The fermionic dark matter candidate ψ interacts with other fields only through its Yukawa couplings. In contrast, the scalar field ϕ couples through quartic scalar couplings as well. However, the quartic coupling involved must be extremely small in order to satisfy dark matter direct detection constraints. As a consequence, ϕ cannot serve as a viable single-component dark matter candidate in this scenario, when considered as a WIMP. Since ψ pair annihilation is prohibited when ϕ and N_1 are heavier, the only phenomenologically viable dark matter configuration in this model is a two-component fermion–scalar dark matter scenario, where the correct relic abundance is achieved predominantly through co-annihilation channels. For simplicity, we have considered an exact mass degeneracy of ψ and ϕ in our study. The Yukawa coupling of ϕ , ψ interaction and the masses of dark matter candidate, for the given masses of right handed neutrinos is constrained from the observed dark matter relic density. Minimum masses of dark matter candidate depends on the masses of right handed neutrinos. For a lightest right-handed neutrino mass of 2000 GeV, the lightest dark matter mass can naturally lie around 200 GeV, while for $M_{N_1} = 1000$ GeV, this is slightly reduced to 187 GeV.

Canonical thermal leptogenesis with hierarchical right-handed neutrinos typically requires a very heavy right handed neutrino in order to generate the observed baryon asymmetry of the Universe, unless at least two of them are nearly degenerate in mass. In our model, however, additional possibilities arise due to the Yukawa interactions between the right-handed neutrinos and the dark matter candidate, which contribute to the leptogenesis mechanism. In addition, the decay of N_1 into the dark matter candidate produces a primordial dark sector asymmetry. We have shown that this asymmetry can be efficiently transferred to the visible sector, thereby contributing to the observed baryon asymmetry.

In the early Universe, the dark sector asymmetry is significantly larger and capable of sourcing the visible sector asymmetry through the transfer mechanism. However, in the present epoch, the residual dark sector asymmetry becomes negligible, ensuring that the dark matter relic density remains unaffected.



Conclusion and Future Outlook

In this chapter, we summarize the thesis, which addresses a key feature that CP -violating phases arise not only from standard vertex and self-energy corrections but also from novel interference effects mediated by the dark sector and additional charged scalars, substantially enhancing the available sources of asymmetry. It successfully accounts for the observed baryon asymmetry of the Universe even when the lightest right-handed neutrino masses lie in the TeV range. It also stabilized the connection between dark matter and leptogenesis. The different interactions of right-handed neutrinos provide the main contribution of our thesis work. It connects low-scale leptogenesis with dark matter.

In Chapter 1, we discuss the observational evidence from astrophysics and cosmology suggesting the existence of matter-antimatter asymmetry and dark matter in the Universe. It is followed by Chapter 2, where we briefly discuss these two observed phenomena within the SM of particle physics, followed by a summary of popular BSM scenarios explaining the origin of matter-antimatter asymmetry and dark matter with its constraint that is provided by experiment. A new possibility, involving an additional interaction of a Majorana neutrino with a charged lepton and a charged scalar, opens up a new framework for low-scale leptogenesis, the full extension and details of which are discussed in Chapter 3. We have explored the scenario in which a feebly interacting dark sector can contribute to generating baryon asymmetry via leptogenesis at the TeV scale through the out-of-equilibrium decay of right-handed neutrinos. This additional decay channel gives a new phase to the second decay and also opens up new possibilities with a new decay channel. This framework allows low-scale leptogenesis at 10 TeV without requiring mass degeneracy of right-handed neutrinos. Another extension of the SM that allows for the interaction of right-handed neutrinos with a DM candidate, establishing a connection between DM and leptogenesis. Full extension and study are available in Chapter 4.

In this work, not only decay but also scattering provides sources of CP asymmetry. The transfer of dark sector asymmetry contributes to the visible sector asymmetry. This study establishes a connection between the mass of dark matter and the lepton asymmetry. This connection helps to get the matter-antimatter asymmetry at 2 TeV without requiring a mass degeneracy.

There could be several future possibilities that can be explored on the basis of different interactions of right-handed neutrinos. Out of all, one can explore the extension of our proposed model in Chapter 4 with the charged scalar S^+ , which leads to several interesting features relevant for both DM and leptogenesis. In this scenario, the lightest right-handed neutrino N_1 can decay into a Higgs and a lepton, into S^+ , a charged lepton, and into particles in the dark sector. These three decay modes can act as independent sources of CP violation while also generating additional phases that influence one another and CP violating parameters. Furthermore, certain scattering processes violate CP asymmetry in ways that transfer asymmetry between the dark sector and the visible sector in both directions. Overall, this framework can provide a viable region of parameter space that is compatible with FIMP dark matter production and successful leptogenesis for a common parameter space. Additionally, the new charged scalars can be probed at collider experiments.



Appendix A

In this appendix we present the fundamental formulas and derivations used in the analysis of particle interactions in the early universe as presented in the Chapter 4. Specifically, we summarize the treatment of decay and scattering processes that contribute to the evolution of species abundances, including the thermal average of the decay width for processes of the type $X \rightarrow ab$, the thermal average of the decay width and thermally averaged cross section for scattering reactions $ab \rightarrow cd$ and the thermal averaging of general $2 \rightarrow 2$ scattering processes. These expressions form the basis for computing interaction rates under finite-temperature conditions and are referenced throughout the calculations in Chapter 4.

A.1 The relevant Decay and scattering processes

The relevant processes to the Boltzmann equations 4.30, 4.31, and 4.32 of the yields, Y_{N_1} , $Y_L = Y_\ell - Y_{\bar{\ell}}$, $Y_{\Delta\psi} = (Y_\psi - Y_{\bar{\psi}})$, are listed as follows:

1. Decay and inverse decay:

$$N_1 \leftrightarrow \ell H, \quad N_1 \leftrightarrow \bar{\ell} H^\dagger, \quad N_1 \leftrightarrow \phi \psi, \quad N_1 \leftrightarrow \phi \bar{\psi}$$

1. $\Delta\ell = 1$ scattering processes:

standard s -channel processes :

$$\ell N_1 \leftrightarrow d \bar{u}, \quad \bar{\ell} N_1 \leftrightarrow \bar{d} u,$$

standard t -channel processes :

$$N_1 u \leftrightarrow d \bar{\ell}, \quad N_1 \bar{u} \leftrightarrow \bar{d} \ell,$$

$$N_1 d \leftrightarrow u \ell, \quad N_1 \bar{d} \leftrightarrow \bar{u} \bar{\ell}$$

standard processes involving gauge boson ($V_\mu = W_\mu$ and B_μ):

$$N_1 H \xleftrightarrow{s} V_\mu \ell, \quad N_1 H^\dagger \xleftrightarrow{s} V_\mu \bar{\ell},$$

$$N_1 \ell \xleftrightarrow{s} V_\mu H, \quad N_1 \bar{\ell} \xleftrightarrow{s} V_\mu H^\dagger,$$

$$N_1 V_\mu \xleftrightarrow{t} H \ell, \quad N_1 V_\mu \xleftrightarrow{t} H^\dagger \bar{\ell}$$

$$N_1 V_\mu \xleftrightarrow{u} H \ell, \quad N_1 V_\mu \xleftrightarrow{u} H^\dagger \bar{\ell}$$

$$N_1 H^\dagger \xleftrightarrow{t} V_\mu \ell, \quad N_1 H \xleftrightarrow{t} V_\mu \bar{\ell},$$

$$N_1 \ell \xleftrightarrow{t} V_\mu H^\dagger, \quad N_1 \bar{\ell} \xleftrightarrow{t} V_\mu H,$$

Non-standard scattering processes:

$$\ell H \xleftrightarrow{s} \phi \psi, \quad \ell H \xleftrightarrow{s} \phi \bar{\psi},$$

$$\ell \phi \xleftrightarrow{t} H \psi, \quad \bar{\ell} \psi \xleftrightarrow{t} H^\dagger \phi,$$

$$\bar{\ell} \phi \xleftrightarrow{t} H^\dagger \psi, \quad H \phi \xleftrightarrow{t} \ell \psi,$$

$$\phi \psi \xleftrightarrow{s+t} \phi \bar{\psi}, \quad \phi \bar{\psi} \xleftrightarrow{s+t} \phi \psi, \quad \phi \phi \xleftrightarrow{t} \psi \bar{\psi}, \quad \phi \phi \xleftrightarrow{t} \psi \psi, \quad \phi \phi \xleftrightarrow{t} \bar{\psi} \bar{\psi};$$

2. $\Delta\ell = 2$ scattering processes: $\ell \ell \xleftrightarrow{t+u} H^\dagger H^\dagger, \quad H \ell \xleftrightarrow{s+t} H^\dagger \bar{\ell},$
and all the reverse reactions.

As it is clear, some of the decay and scattering processes mention above are new, while others are studied in standard thermal leptogenesis, [223, 274, 249].

A.2 Thermal average of decay width of $X \rightarrow a b$

Consider a particle X that decays into particles a and b . Let Γ_X denote the decay width of the particle X . The thermally averaged decay width of X is given by

$$\begin{aligned} \gamma_X &= \int d\Pi_X e^{-E_X/T} \int d\Pi_a d\Pi_b (2\pi)^4 \delta^4(p_X - p_a - p_b) |M_{X \rightarrow ab}|^2 \\ &= n_X^{\text{eq}} \frac{K_1(z)}{K_2(z)} \Gamma_{X \rightarrow ab}, \end{aligned} \quad (\text{A.1})$$

where n_X^{eq} is the equilibrium number density of X , and $\Gamma_{X \rightarrow ab}$ is the decay width of $X \rightarrow ab$ and $K_1(z)$ and $K_2(z)$ are the modified Bessel functions of the second kind of order one and two, respectively. Besides that remaining details are same as usual. The modified Bessel functions of the second kind of order one $K_1(z)$ and two $K_2(z)$ are as follows:

$$K_1(z) = \frac{1}{z} \int_z^\infty e^{-y} \sqrt{y^2 - z^2} dy, \quad \text{and} \quad (\text{A.2})$$

$$K_2(z) = \frac{1}{z^2} \int_z^\infty y e^{-y} \sqrt{y^2 - z^2} dy, \quad (\text{A.3})$$

respectively.

A.3 Thermal averaged cross-section of $a b \rightarrow c d$ scattering processes

The thermal average cross-section γ for a $2 \rightarrow 2$ scattering process with a and b being incoming particles, is given by

$$\gamma^{ab} = \frac{T}{64\pi^4} \int_{s_{\min}}^{\infty} ds \sqrt{s} \hat{\sigma}^{ab}(s) K_1 \left(\frac{\sqrt{s}}{T} \right), \quad (\text{A.4})$$

where:

- T is the temperature,
- $\sqrt{s}$ is the center-of-mass energy,
- K_1 is the modified Bessel function of the second kind of order one,
- $\hat{\sigma}(s)$ is the reduced cross-section, defined as

$$\hat{\sigma}^{ab}(s) = 2s \sigma^{ab}(s) \lambda \left(1, \frac{m_a^2}{s}, \frac{m_b^2}{s} \right), \quad (\text{A.5})$$

where $\sigma^{ab}(s)$ is the cross-section for the process, and $\lambda(x, y, z) = x^2 + y^2 + z^2 - 2(xy + yz + zx)$.

- m_a, m_b are the masses of the incoming particles.

The lower limit of integration, $s_{\min} = \max[(m_a + m_b)^2, (m_c + m_d)^2]$, ensures that the process is kinematically allowed.

Note: For scattering processes that violate CP asymmetry, the definition of thermally averaged cross-section of scattering process in Eq. A.4 slightly modifies to Eq. A.6. In this case, the reduced cross section $\hat{\sigma}$ is replaced with $\epsilon \hat{\sigma}$ (or $\hat{\sigma} \rightarrow \epsilon \hat{\sigma}$), where ϵ denotes the CP -asymmetry parameter. Finally, the thermally averaged cross-section of scattering processes that violate CP asymmetry $\Theta_{\epsilon_\alpha}^{ab}$ is given by

$$\Theta_{\epsilon_\alpha}^{ab} = \frac{T}{64\pi^4} \int_{s_{\min}}^{\infty} ds \sqrt{s} \epsilon_\alpha \hat{\sigma}^{ab}(s) K_1 \left(\frac{\sqrt{s}}{T} \right), \quad (\text{A.6})$$

where other details remain same as in Eq. A.4.

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