Chapter 5

Implications of EDMs on Baryon Asymmetry in the Universe

5.1 General Introduction

The profound discovery of $CP$-violation in $K$ [1] and $B$ [2, 3] mesons is consistent with the Kobayashi-Maskawa model within the framework of the Standard Model of particle physics [4]. The $CP$-violation in SM is, thus, linked to a single physical phase of the CKM matrix which describes the mixing between the three known generations of quarks. The prediction and the observation of $CP$-violation arising through the CKM matrix is one of the greatest milestones of 20th century physics which indeed shaped the structure of the Standard Model (SM) of elementary particles. The SM is celebrated as one of the most precise models of elementary particles and their interactions so far. However, the two main classes of phenomena which require many orders of magnitude larger $CP$-violation than what is incorporated in the CKM model are; one, the EDM of the elementary particles within the reach of the current experimental sensitivities and the other, the observed baryon asymmetry of the Universe (BAU), i.e., the fact that, the observed Universe contains predominantly matter and either insignificant or almost no anti-matter.

Having described the physics of the EDMs of sub-atomic particles, especially that of the elec-
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To begin with, let us assume that the Universe began with symmetric initial conditions i.e., the net baryon density of the Universe was zero. In fact, the other possibility in which the Universe began with the non-zero baryon density is not, in general, preferred mostly for aesthetic reasons although it is not either experimentally or observationally ruled out and
also perhaps because the inflation washes out any such initial asymmetry even if it were existed before inflation. Following the former view, consider that the Universe has followed the well known physical laws of matter creation and annihilation during its evolution from the time of Big-Bang to the present day. We would anticipate, in that case, either the Universe with equal amounts of matter and anti-matter today or the Universe devoid of both matter and anti-matter because of the complete, however accidental, annihilation of the two. The number of planetary probes and satellite landings on various planets in the solar system and in addition, the solar wind samples have provided the valuable information in revealing the material content of our solar system. It is thus inferred that our solar system is made of entirely of matter, i.e., the same type of matter that we are made of. Furthermore, in order to determine the composition of our Galaxy and beyond up to the remote corners of the observable Universe one has to rely on the cosmic rays. The cosmic rays literally provide us the material samples directly from our Sun, our Galaxy and the galaxies far away in the Universe. The material content of the galactic cosmic rays (GCR) mostly include protons, electrons and atomic nuclei moving at extremely high velocities close to the speed of light. About 90% of the cosmic ray nuclei are protons, about 9% are alpha particles (helium nuclei) and the rest of the elements in the periodic table make up only 1% [6]. The cosmic rays from the extra galactic sources mostly include extremely energetic particles and $\gamma$-rays.

In contrast to what was anticipated, the cosmic ray detectors in space have only determined the anti-protons of about $10^{-4}$ for every proton on an average. This observed flux of anti-protons in the cosmic rays is almost consistent with the hypothesis that, these are produced occasionally as secondary particles when the high energy cosmic rays collide with the inter-stellar medium (ISM), mostly through the accelerator-like process, $p + p = 3p + \bar{p}$. Similarly, the observed ratio of the fluxes of anti-Helium ($\bar{He}^4$) to $He^4$ nuclei is $< 10^{-5}$ [7] which also is in agreement with the theory of their secondary production. The observed fluxes of these elements thus indicate that perhaps there is no primordial anti-matter in our Galaxy.

The indirect evidence for the presence of anti-matter in the outer Cosmos can be gathered by observing the characteristic $\gamma$-ray emission due to the annihilation of nuclei–anti-nuclei,
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i.e., for example, the annihilation of $p\bar{p}$ produces $\pi^0$ mesons which subsequently decay into two $\gamma$-rays. The annihilation products such as $\pi^\pm$ produce in flight the charged muons ($\mu^\pm$) and the muon neutrinos ($\nu_\mu$ or $\bar{\nu}_\mu$). The charged muons finally decay down to electrons and neutrinos of both electron- and muon-type. The spectra of all the annihilation products are very similar and they extend from a few MeV to a few hundred MeV and peaks between 100 – 200 MeV [8]. The energy released in the annihilations is carried away mainly by the $\gamma$-rays and the neutrinos but the latter are very difficult to detect due to their extremely weak interactions with matter.

When the galaxies made of matter and anti-matter co-exist in any dense cluster of galaxies and if they collide or interact with each other or merge together then one would expect a detectable background $\gamma$-radiation from these annihilations [8, 9, 10]. From the measured flux of annihilation products, the rates of annihilations can be estimated, from which the amount of interacting matter and anti-matter can also be derived. The absence of such a unique diffuse $\gamma$-ray background and an absence of any distortion of a cosmic microwave radiation in the CMBR spectrum suggest that the amount of anti-matter present either within the cluster of galaxies or at much larger scales is insignificant. Thus, the cosmic rays provide us the conclusive evidence that the universe consists of entirely of matter and the maximal asymmetry between baryons and anti-baryons exists on all scales up to the horizons of the observable Universe.

The observations mentioned above have been succeeded in putting the experimental upper bound on the amount of baryon asymmetry in the Universe. A convenient way to express the BAU is in terms of a dimensionless number given by the ratio of the number density of the net baryons ($n_B = n_b - n_\bar{b}$) to the number density of CMBR photons ($n_\gamma$) at the temperature $T_\gamma = 2.73$ K,

$$\eta = \frac{n_B}{n_\gamma} = \frac{n_b - n_\bar{b}}{n_\gamma} \approx 6.1^{+0.3}_{-0.2} \times 10^{-10} .$$

The most precise limit for the BAU given above is obtained by the Wilkinson Microwave Anisotropy Probe (WMAP) team in 2003 [11]. The parameter $\eta$ is very essential for the
determination of the primordial abundances of the light elements produced during the Big-Bang nucleosynthesis epoch during the evolution of the early Universe. The parameter $\eta$ is believed to have not been changed since the nucleosynthesis epoch. Two things are assumed in arriving at this conclusion; one, the net baryon density remains constant in the co-moving volume which moves together with the expansion of the Universe and the second, since the nucleosynthesis epoch, no process would have produced entropy to change the photon number and hence the number density of non-interacting photons also remains constant in the co-moving volume [12, 13].

The number density of photons did not remain constant during the most of the history of the Universe at least until the end of the radiation dominant era. However, when the Universe was in thermal equilibrium the entropy remained constant. Therefore, in the context of baryogenesis, a useful dimensionless quantity, which remained constant throughout the evolution history of the Universe, would be the ratio of the net baryon density to the entropy density of the Universe, which is given by [11, 13],

$$\eta^s = \frac{n_B}{s} = \frac{n_b - n_\gamma}{s} \approx 0.87 \pm 0.03 \times 10^{-10}. \quad (5.2)$$

The two quantities defined above, i.e., $\eta$ and $\eta^s$, to quantify the baryon asymmetry in the Universe, can only be related after the former remained constant, i.e., only after the time of $e^+e^-$ annihilation [14]. The expression relating the two is the following,

$$\eta \simeq 7 \eta^s = 7 \frac{n_B}{s}. \quad (5.3)$$

The unsettled question, *what could have created the imbalance between matter and antimatter in the Big-Bang theory with symmetric initial conditions?* is answered in the following section.

### 5.3 Sakharov’s Conditions for Baryogenesis

Soon after the path breaking discovery of $CP$-violation by Christenson et al in 1964 [1], it was put forth by A. D. Sakharov in his seminal paper [5] that the observed baryon asymmetry in the Universe would be generated from an initial baryon symmetric state only if
the following three conditions are met: (1) Baryon number violation - There should be baryon number violating interactions in the early Universe; otherwise, the present observed asymmetry would only reflect the asymmetric initial conditions. (2) C and CP-violation - In the absence of a preference for matter or anti-matter, B non-conserving reactions will produce baryon and anti-baryon excesses in the same rate, thereby maintaining zero net baryon number. So both C and CP-violation are necessary to supply a preferential arrow. (3) Out-of-thermal equilibrium conditions - In thermal equilibrium the phase space density of baryons and anti-baryons would necessarily be identical. So there should be departures from thermal equilibrium in order to generate a net baryon asymmetry in the early Universe.

Although, the first condition seems to be very trivial, there is no direct experimental evidence for the baryon number violating interactions so far. The proton decay experiments have yielded only the upper limit on the lifetime of the proton to be less than \((\text{a few } \times 10^{33})\) years [15]. In the present energy scale, whose dynamics is described by the electro-weak theory of elementary particles, a tiny non-conservation of baryonic charge is allowed through the higher order quantum corrections. Thus, the non-conservation of baryonic charge is highly suppressed at the low energies, however, at the energies or temperatures which existed in the early Universe, those interactions which violate baryonic charge could be very efficient and they could successfully produce a noticeable asymmetry by producing excess of baryons over anti-baryons. Thus, it is natural to expect that various grand unified theories (GUTs) inherently include certain heavy particles such as lepto-quarks which can transform quarks into leptons and thereby violating both baryon and lepton numbers. Alternatively, the Supersymmetric models which predict the existence of heavy gauge bosons or Higgs bosons could violate B, however at the energy scales lower than that of GUT energy scales, \(M \sim 10^{16}\) GeV. This would give rise to the proton decay and/or neutrino-anti-neutrino oscillations.

The violation of C and CP-symmetries have been known since several decades. However, the nature of CP-violation which, through the CPT-invariance implies T-violation, is still a subject of fundamental interest in particle physics. This, indeed, has been the central theme.
of our thesis which investigates the subtle connection between the EDM of the electron and the baryon asymmetry in the Universe.

We would like to address an important question about the second Sakharov’s condition required for the baryogenesis below:

*Does one need the breaking of both C- and CP-symmetries or is it sufficient to have CP-violation alone for the generation of cosmological matter-anti-matter asymmetry [17]?*

The short answer is, “Both are needed” !.

Let us assume that the Universe initially is charge symmetric, i.e., Universe is in \( C \) eigenstate.

\[
C|u\rangle = \xi|u\rangle, \quad \text{where } |\xi| = 1. \quad (5.4)
\]

Consider that \( C \) is conserved.

\[
\Rightarrow [C, H] = 0.
\]

The time evolution of the baryonic charge density \( B \) is governed by the equation:

\[
B(t) = \langle u|e^{-iHt} J_B^*e^{iHt}|u\rangle. \quad (5.5)
\]

On inserting a unity operator, \( I = C^{-1}C \) in the above equation we get,

\[
B(t) = \langle u|I e^{-iHt} J_B^*e^{iHt} I|u\rangle,
\]

\[
= -(u|e^{-iHt} J_B e^{iHt}|u),
\]

\[
= -B(t),
\]

\[
\Rightarrow B(t) = 0. \quad (5.6)
\]

It is considered here that the baryonic current is a \( C \)-odd operator, i.e., \( C J_B^* C^{-1} = -J_B^* \).

Thus, in the \( C \)-conserving state either baryonic charge or any other charge can not be generated. The similar arguments with \( CP \) leads to the conclusion that, with the conserved \( CP \) no charge asymmetry could be generated if the Universe is an eigenstate of \( CP \), i.e., \( CP|u\rangle = \zeta|u\rangle, \quad \text{where } |\zeta| = 1 \).
The baryon and lepton numbers of the Universe are $C$-odd and $CP$-odd quantities. Therefore, both $C$ and $CP$ should necessarily be violated in order for the early Universe to develop either baryon or lepton asymmetry.

5.4 Various Models of Baryogenesis

Among the three necessary conditions required for baryogenesis, which are discussed in the preceding section, the baryon number violation and $C$- and $CP$-violation can be investigated only within a given model of particle physics. However, there are many proposed models of particle physics, as of today, which thus provide many alternative mechanisms for baryogenesis. In addition, the departure from the thermal equilibrium condition can be satisfied in two ways; one, due to the presence of heavy gauge bosons which decay into baryons and leptons, however, with the decay rates much smaller than the expansion rate of the Universe and the other, due to the phase transitions which break some of the global or gauge symmetries [12]. Thus, instead of having a rather unique mechanism for baryogenesis, we will have many possible scenarios. With the present understanding of the subject and the experimental limitations, it is quite impossible to choose the correct model of the baryogenesis.

The subject of baryogenesis is one of the most active areas of current research. Voluminous amount of research papers have constantly been produced on the subject matter. Our motive here is to outline the gist of some of the attractive mechanisms proposed and discussed in the published literature. The elaborate discussions on various scenarios of baryogenesis can be found in the following references [7, 12, 13, 14, 18].

5.4.1 GUT Baryogenesis

The Grand Unified Theories (GUTs) unify the strong and the electro-weak interactions at the energy scales of $\approx 10^{16}$ GeV by a unique gauge group $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$. The quarks and leptons appear typically within the same gauge group and the super-heavy gauge bosons which mediate their interactions transform quarks into leptons and/or anti-quarks.
Therefore, at these energy scales, the baryon number violating interactions appear at the
tree-level itself and also such interactions will be extremely abundant. In addition, GUT
theories inherently include new sources of \( \mathcal{CP} \)-violation. The departures from thermal equi­
librium can be achieved naturally when the heavy gauge fields and Higgs fields drop out of
equilibrium due to their slow decay rates compared to the rate of expansion of the Universe.
Therefore, GUT scenarios of baryogenesis are quite popular.

Consider a simple model containing a charge symmetric collection of massive gauge bosons
\( X \) and its anti-particles \( \overline{X} \). The CPT theorem guarantees that, the mass, the lifetime and
the decay rates of \( X \) and \( \overline{X} \) should be the same. Therefore, the number densities of the
two species in thermal equilibrium would be equal, i.e., \( n_X = n_{\overline{X}} \). Let us consider that the
particle \( X \) has only two decay channels; \( qq (B = 2/3) \) and \( \overline{q}l (B = -1/3) \) where \( q \) and \( l \)
are quarks and leptons. The anti-particle \( \overline{X} \) decays to \( \overline{q} \overline{q} (B = -2/3) \) and \( q \overline{l} (B = 1/3) \).
The decays of \( X \) and \( \overline{X} \) violate both the baryon and the lepton numbers since the two
final states have different baryon and lepton numbers. Assuming that the Lagrangian is
Hermitian and \( \mathcal{CP} \) is conserved, the decay rates for \( X \) and \( \overline{X} \) are equal when \( \mathcal{C} \) or \( \mathcal{CP} \) is
conserved. Let \( r \) be the branching ratio for the decay channel, \( X \rightarrow qq \) and \( 1 - r \) for the
decay channel, \( X \rightarrow \overline{q}l \). The net baryon number produced by the decay of \( X \) would then
be, \( N_{B_X} = \frac{2}{3} r - \frac{1}{3} (1 - r) \). Similarly the net baryon number produced by the decay of \( \overline{X} \)
would be \( N_{B_{\overline{X}}} = -\frac{2}{3} \overline{r} + \frac{1}{3} (1 - \overline{r}) \). The total baryon number produced by the decays of \( X \)
and \( \overline{X} \) together would then be \( N_B = N_{B_X} - N_{B_{\overline{X}}} = r - \overline{r} \). The branching ratios \( r \) and \( \overline{r} \)
would exactly be equal which implies that any asymmetry in the net baryon number would
get washed off if \( \mathcal{C} \) or \( \mathcal{CP} \) is conserved. If \( \mathcal{C} \) or \( \mathcal{CP} \)-violation is invoked then this asymmetry
will persist all through, if one assumes that there are no other baryon violating interactions.

The out-of-thermal-equilibrium condition necessary for the baryogenesis could be achieved
in two ways; one, due to the presence of the heavy unstable particles which decay during
the expansion of the Universe and the other, due to the phase transitions which break some
global and/or some gauge symmetries [12]. Consider a baryon number violating heavy par­
ticle of mass \( m_X \). During the very early time when, \( T \gg m_X \), all the particles (here, only \( X \)
and $\overline{X}$) will be in the thermal equilibrium, therefore, $N_X \simeq N_{\overline{X}} \simeq N$, and $B = 0$. However, when $T \lesssim m_X$ the number densities, $n_X = n_{\overline{X}} \simeq (\frac{m_X}{T})^{3/2} e^{-\frac{m_X}{T}} \ll n_\gamma$. The expansion of the Universe defines a finite time scale, i.e., $\tau_U \sim H^{-1}$. If the rate of a reaction which diminishes the number of particles $X$ is smaller than the characteristic time of the Universe, $\Gamma_X \lesssim \tau_U$, then there is enough time for the process to occur and the particle $X$ is thermally coupled to the cosmic fluid. However, if $\Gamma_X > \tau_U$ then the particle $X$ will decouple and drop out of the thermal equilibrium.

The processes which are crucial in maintaining the equilibrium number densities of $X$ and $\overline{X}$ are: annihilation, decay (given by $\Gamma_D$), inverse decay (given by $\Gamma_{ID}$) and the $2 \leftrightarrow 2$ B-non-conserving scattering processes mediated by the exchange of $X$ and $\overline{X}$ bosons (given by $\Gamma_S$). However, the annihilation processes are self-quenching since $\Gamma_{Ann} \propto n_X$ and hence, they can be ignored in the following discussion. The rates for the other three processes are given below [7]:

$$\Gamma_D \simeq \alpha_W m_X \begin{cases} \frac{m_X}{T} & \text{for } T \gtrsim m_X, \\ 1 & \text{for } T \lesssim m_X. \end{cases}$$  \hspace{1cm} (5.7)

$$\Gamma_{ID} \simeq \Gamma_D \begin{cases} \frac{1}{\left(\frac{m_X}{T}\right)^{3/2}} e^{-\left(\frac{m_X}{T}\right)} & \text{for } T \gtrsim m_X, \\ \left(\frac{m_X}{T}\right)^{3/2} e^{-\left(\frac{m_X}{T}\right)} & \text{for } T \lesssim m_X. \end{cases}$$  \hspace{1cm} (5.8)

$$\Gamma_S \simeq n_\sigma \simeq T^3 \alpha^2 \frac{T^2}{(T^2 + m_X^2)^2}.$$  \hspace{1cm} (5.9)

The expansion rate of the Universe is given by,

$$H \simeq \frac{\sqrt{g} T^2}{m_{Pl}}.$$  \hspace{1cm} (5.10)

where $\alpha_W \sim \frac{\alpha}{4\pi}$ is the coupling strength of the $X$ boson and $m_{Pl}$ is the Planck mass. The decays are the most crucial mechanisms which regulate the number densities of $X$ and $\overline{X}$ bosons. During the thermal equilibrium, the decrease in the number density of $X$ bosons due to their natural decay will be compensated by the production of $X$ bosons due to inverse decays and $2 \leftrightarrow 2$ scattering processes. Whenever the interactions which create and destroy these heavy bosons occur slowly on the expansion time scale of the Universe then $X$ and $\overline{X}$
drop out of the thermal equilibrium. Hence, the unequal number densities of baryons and anti-baryons can be obtained. Thus, the departure from the thermal equilibrium felicitates the baryon asymmetry.

5.4.2 Electro-weak Baryogenesis

The Standard Model of electro-weak interactions of elementary particles satisfies all the Sakharov's conditions necessary for Baryogenesis. The experiments have confirmed the non-conservation of $C$ and $CP$ symmetries. The $CP$-violation in the SM appears either through the complex mass matrix with at least three generations of quarks or through the complex coupling constants of the Higgs fields. However, the B and L symmetries appear to be accidentally invariant at the current energy scales at which the SM is still precise. Quite surprisingly, the electro-weak theory allows for such violations at very high energies or temperatures which prevailed in the early Universe by the phenomenon called quantum chiral anomaly. The departure from the thermal equilibrium can arise at the electro-weak phase transition, a transition between the state in which $W$ and $Z$ bosons are massive and one in which they are massless. This transition can be of first order, providing an arrow of time. The baryogenesis scenarios at the EW scale are quite attractive because the predictions of which can be tested easily in the current or future particle accelerators.

The classical electro-weak interactions conserve baryonic current,

\[ \partial_{\mu} J^B_{\mu} = \sum_j \partial_{\mu} (\bar{q}_j \gamma_{\mu} q_j) = 0 , \]  \hfill (5.11)

However, the quantum corrections destroy this conservation law and hence the baryonic current, $J^B_{\mu}$ is not exactly conserved but rather satisfies [13]:

\[ \partial_{\mu} J^B_{\mu} = \frac{g^2 C}{16\pi^2} G_{\mu\nu} \tilde{G}_{\mu\nu} , \]  \hfill (5.12)

where $C$ is a constant and $\tilde{G}_{\mu\nu} = G_{\alpha\beta} \epsilon_{\mu\nu\alpha\beta}/2$ and the gauge field strength, $G_{\mu\nu}$ is given by,

\[ G_{\mu\nu} = \partial_{\mu} A_{\nu} - \partial_{\nu} A_{\mu} + g [A_{\mu} A_{\nu}] . \]  \hfill (5.13)
The anomalous current non-conservation is proportional to the total derivative of a vector operator, \( G_{\mu \nu} \tilde{G}_{\mu \nu} = \partial_\mu K_\mu \) where the anomalous current \( K_\mu \) is given by,

\[
K_\mu = 2 \epsilon_{\mu \nu \alpha \beta} \left( A_\nu \partial_\alpha A_\beta + \frac{2}{3} i g A_\nu A_\alpha A_\beta \right).
\]

(5.14)

To be more precise, the leptons also satisfy equations similar to those described above, i.e., Eqns. (5.11 - 5.14). The resultant quantity \( \partial_\mu K_\mu \) is gauge-invariant and in the temporal gauge \((G_0^\mu)\), this leads to \( \Delta N_{CS} \), where \( N_{CS} \) is Chern-Simons number which assigns a topological charge to a classical gauge field. A non-Abelian gauge theory will have infinite number of ground states whose vacuum gauge field configurations have different topological charges given by, \( \Delta N_{CS} = 0, \pm 1, \pm 2, \cdots \). It can be shown that \( \Delta B = \Delta L = N_g \Delta N_{CS} \) [14], where \( N_g \) is the number of generations of quarks or leptons, in the SM model, \( N_g = 3 \).

For small gauge field quantum fluctuations around the perturbative vacuum configuration, \( G_\mu^\nu = 0 \), B and L remain conserved because \( \Delta N_{CS} = 0 \) in that case. However, for large gauge fields, \( \Delta N_{CS} = \pm 1, \pm 2, \cdots \). This implies that B and L should differ by at least 3 units between the adjacent vacua. Thus, the transition from one vacuum state to another leads to a change in the baryonic and leptonic charges, i.e., B and L are violated by 3 units each, whereas, \( (B + L) \) is violated by 6 units, while, \( (B - L) \) is conserved. The transition between different vacua can happen either through tunneling (this is called instanton process) or through the classical cross over above the barrier (this is called sphaleron process). A schematic illustration of different vacua and the transitions between them are shown in Fig. 5.1.

The baryon number non-conservation through these non-perturbative tunneling effects is proportional to \( e^{- \frac{\Delta x}{\alpha_w}} \), where \( \alpha_W = \frac{2}{4\pi} \). Since \( \frac{4\pi}{\alpha_W} \gg 1 \), the instanton effects at the current energy scales are quite insignificant, therefore, the B-violating processes are exponentially suppressed (probability is \( \approx 10^{-160} \)), however, in the primordial Universe when the temperatures were quite high they were certainly important [7].

In the mechanism proposed by Kuzmin, Rubakov and Shaposnikov (KRS), the thermal fluctuations in the SU(2) gauge field \( (A_\mu^a) \) and symmetry breaking Higgs field \( (\phi^a) \) can drive appreciable number of transitions between different vacua over the potential barrier and
5.4.3: Baryogenesis through Leptogenesis

Figure 5.1: A schematic diagram of free energy vs. gauge field ($A^a_\mu$) or Higgs field ($\phi^a$) configurations, illustrating barriers separating different vacua; the tunneling path through the barrier is an instanton solution (I), whereas, the path crossing over the barrier is a sphaleron solution (S). The adjacent vacua are separated by $\Delta B = \Delta L = N_g$, where $N_g$ is the number of fermion generations taken to be 3. The Chern-Simons number, $N_{CS}$ for various ground states of the field configurations are also shown.

hence sufficient B- and L-violation can be achieved when the temperatures are approximately above 100 GeV or so, which is the typical electro-weak scale. However, at still higher temperatures above the electro-weak phase transition ($T_{EW} < T \lesssim 10^{12}$ GeV), the processes leading to B- and L-violation are quite rapid ($\propto T^5$) compared to the rate of expansion of the Universe ($\propto T^2/M_{Pl}$) which perhaps may wash out any asymmetry previously existed. Hence, the baryogenesis scenarios above $T_{EW}$ must be based on the particle physics models that violate (B-L) also [14]. If the electro-weak phase transition is second order, then the asymmetry can not be generated even below the phase transition since the thermal equilibrium remains undisturbed. Thus, the baryogenesis could proceed in a boundary region between the two phases [13].

The detailed calculations of the transition rates of either the sphaleron or instanton processes, either analytically or numerically, are quite involved and the published results are conflicting [13]. Hence, there is a scope for studying other baryogenesis scenarios.
5.4.3 Baryogenesis through Leptogenesis

The baryogenesis through leptogenesis mechanism is one of the alternate mechanisms proposed to explain the generation of baryon asymmetry in the Universe. This out-of-equilibrium-decay model assumes the existence of right handed heavy Majorona neutrinos with non-degenerate masses of the order of $\sim 10^{12} \text{GeV}$ in the early Universe [14]. At temperatures below their mass scales they fall out of thermal equilibrium. The interactions of these heavy leptons with other particles are quite feeble. As the name indicates, this mechanism involves two steps; first, the out-of-equilibrium decays of heavy Majorona neutrinos generate non-zero lepton number. In the second step, the lepton asymmetry is transformed into baryon asymmetry by $C$ and $CP$ conserving sphaleron processes in thermal equilibrium during the electro-weak stage. The sphalerons do not conserve baryonic and leptonic charges individually but they conserve the difference in their numbers $(B - L)$. Thus, initial $L$ could be redistributed in equilibrium in almost equal shares between $B$ and $L$. This mechanism first proposed in Ref. [19], has subsequently been developed over the years and the sufficient details have been given in the review article [20]. The observation of neutrino masses in the recent years makes this idea more attractive and plausible.

5.4.4 Spontaneous Baryogenesis

All those mechanisms proposed for baryogenesis incorporate the standard Sakharov conditions. However, there is one novel and a unique mechanism called spontaneous baryogenesis in which the explicit $CP$-violation and the non-equilibrium conditions are not required for generating the baryon asymmetry. However, this invokes a more rational approach in which the so called $CPT$-invariance is dynamically violated temporarily.

The novelty of this mechanism lies in the following term in the Lagrangian [7],

$$\mathcal{L} = \cdots + \frac{1}{f} \partial^\mu \phi J_B^\mu \hspace{1cm} (5.15)$$

where $\phi$ is a scalar field similar to an axion-like field called $\tilde{u}$ion [21] which couples to the baryon-number current and $f$ is the energy scale of the order of $10^{13} \text{GeV}$ or higher. The
CVT-invariance is violated when this term takes non-zero value. For a spatially constant $\phi$, we can write,

$$\frac{1}{f} \partial^\mu \phi J^\mu_B = \frac{1}{f} \dot{\phi} n_B. \quad (5.16)$$

The coefficient of the net baryon density, $n_B = n_b - n_\bar{b}$ can be viewed as a kind of chemical potential, $\mu_B$ for the baryon number and the evolution of $\phi$ thus provides a preferential direction for the production of baryons and anti-baryons by shifting the degeneracy of their energy levels. Since this arises dynamically without imposing any constraint, it is strictly not a chemical potential but is sometimes referred to as charge potential in the literature [18]. The non-zero baryon density obtained by spontaneous baryogenesis in thermal equilibrium is then given by [7],

$$n_B \approx -\frac{\dot{\phi} T^2}{f}, \quad (5.17)$$

and the baryon asymmetry parameter is given by,

$$\eta^* \approx -\frac{\dot{\phi}}{g_* f T}. \quad (5.18)$$

Thus, baryon asymmetry can be generated without explicit $C\bar{P}$-violation by invoking a dynamical violation of $CPT$-invariance, only temporarily, in thermal equilibrium.

### 5.5 Results and Discussions

Having described some of the attractive models of baryogenesis in the previous section, we would remark that the list of models one can consider for explaining baryogenesis is non-exhaustive. Since the amount of $C\bar{P}$-violation in the Standard Model which is described by the symmetry group $SU(2) \times U(1)$ is inadequate to explain some of the experimental data, let us turn on to one of the simplest extensions to the SM known as the two Higgs doublet model (2HDM) in which the particle content of the SM is extended by including an additional scalar Higgs doublet while keeping the fermion content of SM intact. Unlike the SM in which the $SU(2) \times U(1)$ symmetry is broken by the vacuum expectation value (VEV) of a single Higgs doublet, in the 2HDM the electro-weak symmetry is broken by the VEVs of the
Table 5.1: The particle content and their properties in the 2HDM model. The subscript ‘$\sigma$’ is the generation index and the superscript ‘c’ refers to Dirac conjugation.

<table>
<thead>
<tr>
<th>particles</th>
<th>$SU(2)_L \times U(1)_Y$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$(\nu_\sigma, l_\sigma)$</td>
<td>$(2, -\frac{1}{2})$</td>
</tr>
<tr>
<td>$(u_\sigma, d_\sigma)$</td>
<td>$(2, \frac{1}{6})$</td>
</tr>
<tr>
<td>$l_\sigma^c$</td>
<td>$(1, 1)$</td>
</tr>
<tr>
<td>$\nu_\sigma^c$</td>
<td>$(1, -\frac{2}{3})$</td>
</tr>
<tr>
<td>$d_\sigma^c$</td>
<td>$(1, \frac{1}{3})$</td>
</tr>
<tr>
<td>$(\phi_1^+, \phi_1^0)$</td>
<td>$(2, \frac{1}{2})$</td>
</tr>
<tr>
<td>$(\phi_2^+, \phi_2^0)$</td>
<td>$(2, \frac{1}{2})$</td>
</tr>
</tbody>
</table>

two independent scalar Higgs doublets. This model further considers that, one of the scalar doublets gives mass to the up-type quarks and the other to the down-type quarks and leptons.

The particles contained in the particle zoo of the 2HDM are listed in Table 5.1. The parenthesis in the second column contain $SU(2)_L$ multiplicity and $U(1)_Y$ quantum number; the latter is normalized by $Q = I_3 + Y$, where $Q$ is the electric charge, $I_3$ is the third-component of the isospin quantum number and $Y$ is the hypercharge.

The scalar potential of the most general re-normalizable 2HDM model is given by [22],

$$
V(\phi_1, \phi_2) = \lambda_1 [(\phi_1^+ \phi_1) - \nu_1^2]^2 + \lambda_2 [(\phi_2^+ \phi_2) - \nu_2^2]^2 + \lambda_3 [(\phi_1^+ \phi_1 - \nu_1^2) + (\phi_2^+ \phi_2 - \nu_2^2)]^2 + \lambda_4 [(\phi_1^+ \phi_1)(\phi_2^+ \phi_2) - (\phi_1^+ \phi_2)(\phi_2^+ \phi_1)] + \lambda_5 \left[\text{Re} (\phi_1^+ \phi_2) - \nu_1 \nu_2 \cos \xi\right]^2 + \lambda_6 \left[\text{Im} (\phi_1^+ \phi_2) - \nu_1 \nu_2 \sin \xi\right]^2,
$$

(5.19)

with the non-zero vacuum expectation values of the two neutral Higgs fields, $\phi_1^0$ and $\phi_2^0$, in the unitary gauge, are given by,

$$
\langle \phi_1^0 \rangle = (0, \nu_1) \text{ and } \langle \phi_2^0 \rangle = (0, \nu_2 e^{i\xi}),
$$

(5.20)
where \( v_1 \) and \( v_2 \) are the zero temperature VEVs corresponding to the scalar fields \( \phi_1^0 \) and \( \phi_2^0 \) respectively and are constrained by \( v^2 = v_1^2 + v_2^2 \). Both the VEVs and the coupling constants \( \lambda_i \) are the real positive definite parameters. The phase angle \( \xi \) is a physical \( CP \)-odd phase. The last two terms of Eq. (5.19) with the coupling coefficients \( \lambda_5 \) and \( \lambda_6 \) contain explicit \( CP \)-violation. With the suitable choices of parameters, such as for example, (i) by setting either \( \lambda_5 = 0 \) or \( \lambda_6 = 0 \) (ii) by choosing either \( \xi = \frac{n \pi}{2} \) or \( \lambda_5 = \lambda_6 \), the Lagrangian can be made \( CP \)-invariant. Thus, these two terms govern the dynamics of \( CP \)-violation. The magnitude of \( CP \)-violation in this 2HDM model is defined as [23],

\[
\lambda_{CP} = (\lambda_5 - \lambda_6) \sin (2 \xi) .
\] (5.21)

The 2HDM consists of a charged Higgs particle and its anti-particle, \( H^\pm \) and the three neutral Higgs bosons, \( H_{1,2,3}^0 \) in its particle spectrum. The \( CP \)-violation in this model can happen through a phase in the Kobayashi-Maskawa (KM) matrix and also in the neutral as well as the charged Higgs boson exchange [24]. In the absence of \( CP \)-violation i.e., if \( \xi = 0 \), the neutral Higgs sector consists of two scalars which are \( CP \)-even and one pseudo-scalar boson which is \( CP \)-odd. In that case, the two scalars mix among themselves, however, the mixing with pseudo-scalar is ruled out. In the presence of \( CP \)-violation i.e., if \( \xi \neq 0 \), all three neutral states mix together giving rise to three mixing angles and three mass eigenstates, \( |\phi_{1,2,3}^0\rangle \) with their masses given by \( m_{1,2,3} \) respectively. In summary, these states couple to both scalar and pseudo-scalar currents and hence, they will no longer possess definite \( CP \)-symmetry. The \( CP \)-violation can appear, in the unitary gauge, in the propagators involving the two complex neutral Higgs fields, such as for example \( \langle T \{ \phi_1^0, \phi_2^0 \} \rangle \). Thus, the \( CP \)-violation needed for explaining the two phenomena considered in the present work; the electron EDM and the baryogenesis, is obtained only from the neutral Higgs sector. Thus, at least qualitatively they both are assumed to have the same source of \( CP \)-violation.

By considering the electro-weak baryogenesis scenario with the strong first order phase transition at some temperature \( T_c \) within the 2HDM, the details of which can be found in [23], the baryon asymmetry can be obtained as,

\[
\eta = \frac{45}{2\pi^2 N_{\text{eff}}} \kappa n_f \alpha_W^6 \sin^3 (2 \alpha) \lambda_{CP} \frac{m_1^2 \tau_2}{v_1^4 v_2},
\] (5.22)
where $N_{\text{eff}}$ is the number of effectively massless degrees of freedom at the EW phase transition, $\kappa$ is the numerical coefficient appearing in the equation of the time dependent rate of sphaleron transitions, $n_f$ is the number of fermion generations, $\alpha_W$ is the weak coupling constant, $\alpha = \tan^{-1}\left(\frac{m_1(T)}{m_2(T)}\right)$ is the mixing angle where $m_1(T)$ and $m_2(T)$ are the temperature dependent masses of the Higgs scalars and $m_t$ is the mass of the top quark.

The main thrust of the present work is to consider both the electro-weak baryogenesis scenario and the electron EDM within the framework of the two-Higgs doublet model and study the qualitative connection between the two intriguing issues and obtain constraints on the matter-antimatter asymmetry by using the best limit available on the electron EDM. We are motivated by the first published work of this kind, to the best of our knowledge, by Kazarian et al in Ref. [23] where they have obtained the cosmological lower bound on the EDM of the electron by using the value for the magnitude of CP-violation extracted from fitting the observed baryon asymmetry to their modeled value using the 2HDM. We would like to consider the EDM of the electron and the baryogenesis within their proposed model and re-examine their bounds in the light of the recent and improved experimental constraints on many of the model parameters. While doing so, we have derived the more general expression for the EDM of the electron within the framework of 2HDM.

The electron EDM in the 2HDMs arises from the two-loop diagrams as shown by Barr in Ref. [25], however, it was pointed out earlier by Weinberg for the neutron EDM [24] with higher order loop diagrams. The CP-violation necessary for the electron EDM comes from the exchange of the neutral Higgs bosons in their models which was based on the result in Ref. [26]. However, it was shown in Ref. [25] that by considering a heavy particle of mass greater than or equal to the Higgs mass, for example, a top quark, running around a loop which is shrunk to a point, i.e., effectively reducing a three loop diagram from Ref. [24] to a two loop diagram, can induce a substantially large EDM. Thus, Barr demonstrated that the two-loop contribution to the electron EDM is about 8 orders of magnitude larger than the one-loop contribution and thus, he obtained the electron EDM of the order of $\sim 10^{-26} \text{ ecm}$ which encouraged the experimentalists to look for the non-zero EDM of the elementary particles.
or composite systems without giving up a hope. It was also shown in [24, 25, 27] that, the EDM of the electron is proportional to the unknown dimensionless quantity, \( \text{Im}(Z_0) \) which, in turn, depends on the neutral Higgs boson propagator as shown below:

\[
d_e \simeq 1.3 \times 10^{-26} \text{Im}(Z_0) \text{ e cm},
\]

where \( \text{Im}(Z_0) \propto \frac{\langle \phi^0_1, \phi^0_2 \rangle}{v_1 v_2^2} \).

The neutral scalar propagator in the 2HDM is given by [25],

\[
\frac{\langle \phi^0_1, \phi^0_2 \rangle}{v_1 v_2^2} \equiv \sqrt{2} G_F \sum_{n=1}^{3} \frac{\text{Im}(Z_{0n})}{q^2 + m_n^2},
\]

where \( q \) is the momentum transfer and the index \( n \) refers to the different mass eigen states of the neutral Higgs bosons. From the above equation it is clear that, the EDM enhances when the least massive neutral Higgs boson mediates the interaction. Let us assume a mass hierarchy among the three mass eigen states of the neutral Higgs bosons as, \( m_1 < m_2 < m_3 \).

Taking into account the fact that the EDM form factor is calculated at zero-momentum transfer, the \( \text{Im}(Z_0) \) can be written as,

\[
\text{Im}(Z_0) \simeq \frac{m_1^2}{\sqrt{2} G_F} \frac{\langle \phi^0_1, \phi^0_2 \rangle}{v_1 v_2^2}. \tag{5.26}
\]

Let us define the complex neutral Higgs fields by including the real spin-0 physical quantum fields, \( R_i \) and \( S_i \) where \( R_i \) is a scalar field and \( S_i \) is a pseudo-scalar field:

\[
\phi^0_1 = [v_1 + \frac{1}{\sqrt{2}} (R_1 + i S_1)],
\]
\[
\phi^0_2 = e^{i \xi} [v_2 + \frac{1}{\sqrt{2}} (R_2 + i S_2)]. \tag{5.27}
\]

In terms of the quantum fields, let us define the following,

\[
G = \frac{v_1 S_1 + v_2 S_2}{\sqrt{2} \sqrt{v_1^2 + v_2^2}}; \quad \text{and} \quad A = \frac{-v_2 S_1 + v_1 S_2}{\sqrt{2} \sqrt{v_1^2 + v_2^2}}, \tag{5.28}
\]

where \( G \) is the Goldstone boson which gets absorbed after the symmetry breaking.

\[
\Rightarrow S_1 \approx -\sqrt{2} \frac{v_2}{v} A; \quad \text{and} \quad S_2 \approx \sqrt{2} \frac{v_1}{v} A. \tag{5.29}
\]
In terms of the quantum fields, the Higgs propagator given in Eq. (5.24) becomes,

\[
\langle \phi_1^0, \phi_2^0 \rangle \approx \frac{1}{2} \frac{\langle R_1 + i S_1, R_2 - i S_2 \rangle}{v_1 v_2},
\]

\[
\approx \frac{1}{2} \langle R_1, R_2 \rangle + \langle S_1, S_2 \rangle + i ((R_2, S_1) - (R_1, S_2)).
\] (5.30)

The first two terms in the above equation are CP-conserving, whereas, the imaginary terms are CP-violating which give rise to the EDM of the electron.

Now consider,

\[
\text{Im}\langle \phi_1^0, \phi_2^0 \rangle \approx \langle R_2, S_1 \rangle - \langle R_1, S_2 \rangle,
\]

\[
\approx \frac{1}{v} [-\langle R_2, A \rangle v_2 - \langle R_1, A \rangle v_1],
\]

\[
\approx \frac{1}{v} \langle (R_2 v_2 + R_1 v_1), A \rangle.
\] (5.31)

Therefore,

\[
\text{Im}(Z_0) \approx -\frac{m_1^2}{2\sqrt{2} G_F} \frac{1}{v} \frac{\langle (R_2 v_2 + R_1 v_1), A \rangle}{v_1 v_2}.
\] (5.32)

On substituting \(\phi_1^0\) and \(\phi_2^0\) from Eq. (5.27) into the Higgs potential defined in Eq. (5.19) and minimizing the potential with respect to each of the quantum fields we can obtain the mass squared matrix which are quadratic in the fields. The 4 × 4 matrix so obtained can be reduced to a 3 × 3 matrix by re-defining the pseudo-scalar fields in terms of \(G\) and \(A\) using Eq. (5.28). It is to be noted here that the elements of the matrix with Goldstone boson \(G\) vanish resulting in the effective 3 × 3 matrix in terms of the fields \(R_1, R_2\) and \(A\). The final mass squared matrix, \(M^2\), is given by,

\[
\begin{pmatrix}
2 v_2^2 (\lambda_1 + \lambda_3) + \frac{1}{4} v_2^2 (\lambda_5 + \lambda_6) & \frac{1}{4} v_1 v_2 (8 \lambda_3 + \lambda_5 + \lambda_6) & - \frac{1}{2\sqrt{2}} v_2 v \sin (2 \xi) (\lambda_5 - \lambda_6) \\
+ \frac{1}{4} v_2^2 \cos (2 \xi) (\lambda_5 - \lambda_6) & + \frac{1}{4} v_1 v_2 \cos (2 \xi) (\lambda_5 - \lambda_6) & - \frac{1}{2\sqrt{2}} v_2 v \sin (2 \xi) (\lambda_5 - \lambda_6) \\
\frac{1}{4} v_1 v_2 (8 \lambda_3 + \lambda_5 + \lambda_6) & 2 v_2^2 (\lambda_2 + \lambda_3) + \frac{1}{4} v_2^2 (\lambda_5 + \lambda_6) & - \frac{1}{2\sqrt{2}} v_2 v \sin (2 \xi) (\lambda_5 - \lambda_6) \\
+ \frac{1}{4} v_1 v_2 \cos (2 \xi) (\lambda_5 - \lambda_6) & + \frac{1}{4} v_2^2 \cos (2 \xi) (\lambda_5 - \lambda_6) & - \frac{1}{2\sqrt{2}} v_2 v \sin (2 \xi) (\lambda_5 - \lambda_6) \\
- \frac{1}{2\sqrt{2}} v_2 v \sin (2 \xi) (\lambda_5 - \lambda_6) & - \frac{1}{2\sqrt{2}} v_2 v \sin (2 \xi) (\lambda_5 - \lambda_6) & \frac{1}{4} v_2^2 (\lambda_5 + \lambda_6) + \frac{1}{4} v_2^2 \cos (2 \xi) (\lambda_6 - \lambda_5)
\end{pmatrix}
\]
5.5: Results and Discussions

We can diagonalize the $M^2$ matrix defined above by some unitary transformation matrix, $O$, i.e.,

$$O^T M^2 O = D^2,$$

$$\Rightarrow M^2 = O D^2 O^T,$$

$$M^{-2} = O D^{-2} O^T.$$  \quad (5.33)

Thus, we have,

$$\begin{pmatrix} R_1 \\ R_2 \\ A \end{pmatrix} = O \begin{pmatrix} H_1 \\ H_2 \\ H_3 \end{pmatrix}.$$  \quad (5.34)

Therefore, $R_1 = O_{1k} H_k$; $R_2 = O_{2k} H_k$; $A = O_{3k} H_k$,

$$\Rightarrow \langle R_1, A \rangle = \frac{O_{1k} O_{3k}}{(M^2)_{k}^{23}} \equiv (M^{-2})_{13},$$  \quad (5.35)

similarly, $\langle R_2, A \rangle \equiv (M^{-2})_{23}$.  \quad (5.36)

Therefore, the propagator defined in Eq. (5.32) becomes,

$$\text{Im} (Z_0) \simeq -\frac{m_f^2}{2 \sqrt{2} G_F} \frac{1}{v \text{Denom}} \left[ \frac{v_2 (M^{-2})_{23} + v_1 (M^{-2})_{13}}{v_1 v_2} \right].$$  \quad (5.37)

The $M^{-2}$ matrix is obtained by inverting the mass squared matrix given in Eq. (5.33). The (13) and the (23) components of the $M^{-2}$ matrix are, respectively, given below:

$$(M^{-2})_{13} = \frac{v_2 \left[ -v_1^2 \lambda_3 + v_1^2 (\lambda_2 + \lambda_3) \lambda_{CP} \right]}{\text{Denom}},$$

$$(M^{-2})_{23} = \frac{v_1 \left[ -v_2^2 \lambda_3 + v_2^2 (\lambda_1 + \lambda_3) \lambda_{CP} \right]}{\text{Denom}},$$  \quad (5.38)

where the denominator, $\text{Denom} = \sqrt{2} v \left[ v_1^4 (\lambda_1 + \lambda_3) \lambda_5 \lambda_6 + v_2^4 (\lambda_2 + \lambda_3) \lambda_5 \lambda_6 \\
+ 2 v_1^2 v_2^2 \left( \lambda_1 (\lambda_2 + \lambda_3) (\lambda_5 + \lambda_6) + \lambda_3 [-\lambda_5 \lambda_6 + \lambda_2 (\lambda_5 + \lambda_6)] \right) \\
- 2 v_1^2 v_2^2 \left[ \lambda_2 \lambda_3 + \lambda_1 (\lambda_2 + \lambda_3) \right] (\lambda_5 - \lambda_6) \cos (2 \xi) \right].$  \quad (5.39)
5.5: Results and Discussions

Using Eq. (5.37) in the expression for the electron EDM i.e., in Eq. (5.23), we obtain,

\[
d_e \simeq \frac{m_i^2}{2\sqrt{2}G_F} \frac{(v_2 (M^{-2})_{23} + v_1 (M^{-2})_{13})}{v_1 v_2 v} \cdot 1.3 \times 10^{-26} \text{e cm.}
\]  

(5.40)

Let us now estimate the magnitudes of both \(d_e\) and \(\eta\) parametrically within the framework of the 2HDM. Some of the constants/quantities appearing in Eqns. (5.22 and 5.40) take the following values:

\[
v_1 = v \cos \beta; \quad v_2 = v \sin \beta; \quad v = 174 \text{ GeV}; \quad G_F = 1.16637 \times 10^{-5} \text{ GeV}^{-2};
\]

\[
T_e = 100 \text{ GeV}; \quad m_t = 173.1 \text{ GeV}; \quad n_f = 3; \quad \alpha_W = 1/128; \quad N_{\text{eff}} = 100.
\]

(5.41)

We recall here that, the mixing angle \((\alpha)\) and the \(C\bar{P}\)-violating parameter \((\lambda_{\text{CP}})\) are defined as below:

\[
\alpha = \tan^{-1} \left( \frac{m_2^2}{m_1^2} \right); \quad \text{and} \quad \lambda_{\text{CP}} = (\lambda_5 - \lambda_6) \sin (2\xi).
\]

(5.42)

The variable parameters appearing in the calculations of \(d_e\) and \(\eta\), such as, the coupling constants in the Higgs potential \((\lambda_1-6)\), the phase angle \((\xi)\), neutral scalar Higgs masses, \((m_1,2)\), the numerical coefficient \((\kappa)\) and \(\beta\) are randomly varied within their allowed range of values listed below:

\[
\beta = \text{Random Real,} \{0, \pi/2\};
\]

\[
\xi = \text{Random Real,} \{0, \pi\};
\]

\[
\lambda_1 \text{ through } \lambda_6 = \text{Random Real,} \{0, 1\};
\]

\[
m_1 = \text{Random Real,} \{100, 200\};
\]

\[
m_2 = \text{Random Real,} \{200, 1000\};
\]

\[
\kappa = \text{Random Real,} \{0.1, 1\}
\]

(5.43)

In order to obtain statistically good results, the above parameters are randomly varied for over a large number of times. The scatterplot of the baryon asymmetry obtained as a function of the electron EDM, for the full range of those values, is shown in Fig. 5.2.

The best upper bound on the EDM of the electron, obtained by combining the atomic thallium EDM experimental result, given at 90% confidence level, by Regan et al [28] and our
accurate theoretical EDM enhancement factor obtained in the preceding chapter, is given by $d_e < 1.9 \times 10^{-27}$ e cm. Therefore, we have shown in Fig. 5.3 only those values of $\eta$, which are obtained using the constraint from the experimental bound on the electron EDM. The results indicate that, the magnitude of the $\mathcal{CP}$-violation in the two-Higgs doublet model may be sufficient enough to produce the electron EDM within the reach of the current experimental limits, however, it seems to suppress the baryon asymmetry by $6 - 7$ orders of magnitude smaller than their observed bound. It can be inferred that, the electron EDM provides a very stringent constraint on the amount of the observed baryon asymmetry generated through the mechanism proposed by Kazarian et al [23] and it rules out the proposed mechanism as the main source of $\mathcal{CP}$-violation required to explain the electron EDM and the baryon asymmetry together under the assumptions considered here. One has to look for other mechanisms, such as, for example, baryogenesis through leptogenesis and search for the possible subtle connection between the two disparate phenomena. One can not also ignore the possibility that, these two phenomena may not necessarily have been manifested with the same source of $\mathcal{CP}$-violation since the energy scales involved in the two are different.

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{scatterplot.png}
\caption{The scatterplot of the baryon asymmetry ($\eta$) vs. the electron EDM ($d_e$ in units of $10^{-26}$ e cm).}
\end{figure}
Figure 5.3: The scatterplot of the baryon asymmetry ($\eta$) for the allowed range of the electron EDM ($d_e$ in units of $10^{-26}$ e cm).
Bibliography


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