Notes #1: Overview of Baryogenesis

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We are made of matter, not antimatter. While this might seem obvious, it is a mystery from the point of view of our current best theory of elementary particles, the Standard Model (SM). The SM treats matter and antimatter nearly identically, and it is not known how to explain why there is more matter than antimatter using SM physics alone.

Astrophysical and cosmological observations also point toward an excess of matter over antimatter. These indicate that regular matter makes up about 5% of the energy density of the universe, and consists nearly entirely of matter rather than antimatter. By mass (energy density), this excess is dominated by baryons, and thus the excess of matter over antimatter is often called the *baryon asymmetry*.

Baryogenesis refers to any process occurring in the early universe that creates a baryon asymmetry. The standard picture for the evolution of the cosmos is that inflation happens followed by reheating and (mostly) adiabatic expansion and cooling. Baryogenesis is expected to occur either during or after reheating, creating a small excess of matter over antimatter. As the universe eventually cools below the nucleon mass, $T \lesssim 1$ GeV, a net annihilation of baryons and antibaryons occurs until essentially only the excess baryons are left.

The mystery of baryogenesis is that there is no known mechanism for it within the SM. The baryon asymmetry is therefore a strong motivator for new physics beyond the SM. In these notes we will discuss the observational evidence for baryogenesis, the necessary ingredients needed for it to occur, and some proposed mechanisms for it. The notes to follow will cover two of the most promising baryogenesis mechanisms – leptogenesis and electroweak baryogenesis – in more detail. Some nice general reviews of baryogenesis can be found in Refs. [1, 2, 3, 4].

0 Notation

Throughout these lectures we will use natural units with $\hbar = c = k_B = 1$. This means that mass, energy, momentum, and temperature all have units of energy, and length and time have units of $(\text{energy})^{-1}$. We also use a mostly minus metric

$$\eta_{\mu\nu} = \eta^{\mu\nu} = \text{diag}\{1, -1, -1, -1\} \ , \tag{1}$$

together with $\epsilon^{01234} = +1$. When dealing with fermions, we will mostly use a 4-component notation in the chiral basis [5]. Finally, we will always use the reduced Planck mass $M_{\text{Pl}} = 1/\sqrt{8\pi\, G} \simeq 2.4 \times 10^{18}$ GeV.

1 Why Baryogenesis is Needed

Our local region of the universe is very clearly made of matter, with only very small amounts of antimatter consistent with creation by cosmic rays, nuclear decays, and high-energy particle colliders. This is good, because annihilating with a clump of antimatter would not be much fun. Going beyond our local region, there is extremely strong evidence that the baryon asymmetry persists throughout the entire visible universe based on observations of the cosmic microwave background (CMB), light element abundances, and cosmic rays.

1.1 Total Baryons

The best determination of the total baryon density (*i.e.* baryons and antibaryons) of the universe comes from measurements of the CMB temperature fluctuation power spectrum. Baryons remain coupled to photons until recombination, and this modifies how density perturbations evolve in time [6, 7]. The primary effect of this on the temperature power spectrum is to modify the heights of the peaks, enhancing the odd peaks relative to the even ones, as can be seen in the left panel of Fig. 1 (taken from Ref. [8]).[1] As a fraction $\Omega_B = \rho_B/\rho_c$ of the critical energy density $\rho_c = 8.0992\, h^2 \times 10^{-47}$ GeV^4, current CMB data combined with other data gives a baryon energy density of [10, 11]

$$\Omega_B h^2 = 0.0223 \pm 0.0002\ , \tag{2}$$

where $h = 0.679 \pm 0.006$ is the Hubble constant in funny units. The corresponding baryon number density is frequently expressed as a ratio of the photon number density,

$$\eta \;\equiv\; \frac{n_B}{n_\gamma} \;=\; 6.1 \times 10^{-10} \left(\frac{\Omega_B h^2}{0.0223}\right) \;\simeq (7.04) Y_B\ , \tag{3}$$

where $Y_B = n_B/s$ is the present number density relative to the entropy density.

A second determination of the CMB value of the total baryon density comes from primordial nucleosynthesis (BBN). As the early universe cooled below about 1 MeV, free nucleons began to bind into more complex nuclei. By extrapolating known nuclear reaction rates to the conditions expected in this era based on standard thermal cosmological evolution, the primordial abundances of several light elements (up to about ^{7}Li) can be predicted and compared to observations. The single input to this calculation is the baryon density η, and it is found that theory and experiment agree quite well for the value of η determined from the CMB. This is illustrated in the right panel of Fig. 1 (taken from Ref. [12]).[2]

Together, these and many other cosmological observations give us confidence in the standard cosmological model of inflation and reheating followed by ΛCDM with $\eta \simeq 6 \times 10^{-10}$. An obvious question, then, is how to explain this baryon density. Examining this question suggests that the total baryon density consists nearly entirely of baryons, with very few antibaryons.

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[1] See also Ref. [9] for a nice animation.

[2] Note that the observed abundance of ^{7}Li is a bit off from the prediction. It is not clear whether this is evidence for new physics or an indication of a discrepancy in the astrophysical interpretation of the data [13].

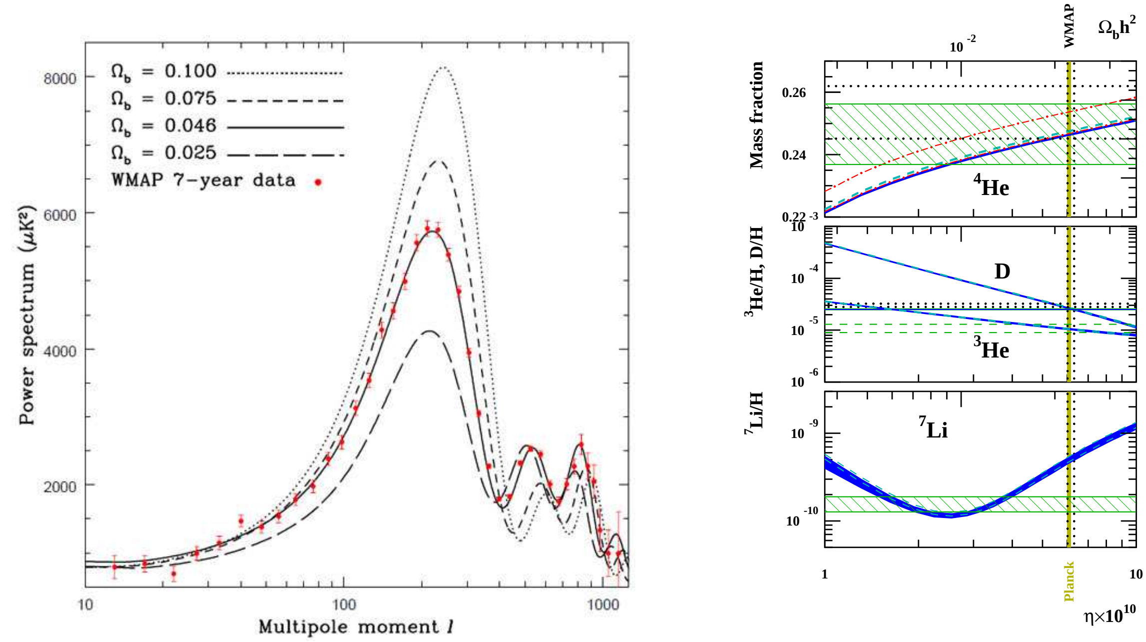

Figure 1: **Left:** Depedence of the CMB temperature power spectrum on the total baryon density (from Ref. [8]). **Right:** Predictions for the light element abundances from primordial nucleosynthesis as a function of η compared to the values deduced from observations (from Ref. [12]).

1.2 Asymmetric Baryons

Suppose there were equal numbers of baryons and antibaryons early on in the cosmos, at temperature $T \gg m_p \sim$ GeV. Assuming standard thermal evolution, a net annihilation of baryons and antibaryons would begin near the QCD phase transition at $T \sim \Lambda_{QCD} \sim$ 200 MeV. The annihilation would continue until it became too slow to keep up with the Hubble expansion, similar to the freezeout of dark matter. Tracking the freezeout, one finds that it occurs at $T \sim 20$ MeV and yields $\eta \simeq 5 \times 10^{-19}$ [14], nearly ten orders of magnitude too small! This most basic picture of baryon evolution is clearly incorrect.

Instead, let us assume that at some point prior to BBN an excess of baryons over antibaryons was created. Repeating the argument above, a net annihilation will begin near the QCD phase transition. However, in contrast to the symmetric case, the annihilation will only continue until (nearly) all the antibaryons are used up. Once they are, the (effective) conservation of baryon number in the SM implies that the remaining baryons will be stable and remain until today. In this asymmetric scenario, the ultimate baryon density is set by the initial asymmetry rather than rate of baryon annihilation. Given the observed value of $\eta = n_B/n_\gamma$, the asymmetry needed is only about one part in 10^{10}.

The asymmetric scenario of baryon formation is also consistent with the observation that our local region of space (the Milky Way galaxy) consists of baryons. In fact, this conclusion can be extended to our local galaxy cluster (Virgo), corresponding to a radius of about 20 Mpc [15, 16].[3] Evidence for a nearly complete asymmetry of baryons over antibaryons can be extended even further based on observations of extragalactic gamma rays. Regions consisting mainly of antibaryons would lead to significant annihilation at the boundaries with regions of baryons, leading to gamma ray production. Using observations of the diffuse gamma ray spectrum, Ref. [17] concluded that such regions would have to be at least 1000 Mpc in size. For the rest of these lectures, we will focus exclusively on the asymmetric baryon scenario.

2 Baryogenesis Ingredients

Three basic ingredients are needed to create a baryon asymmetry in the early universe (up to a few potential loopholes). They are called the *Sakharov conditions* after the person who first formulated them [18]:

1. Baryon number (B) Violation
2. C and CP Violation
3. Departure for Thermodynamic Equilibrium

We discuss each of these conditions below, with a focus on how they are realized within the SM with ΛCDM cosmology.

2.1 B Violation

Baryon number (B) is the charge corresponding to an approximate global $U(1)_B$ symmetry of the SM under which quark fields have charge $B = 1/3$, and all other states have charge zero. After confinement, baryons have $B = 1$, antibaryons have $B = -1$, and all other states have $B = 0$. The very long lifetime of the proton, $\tau_p > 1.6 \times 10^{34}$ yrs [19], indicates that B is conserved to a very strong degree at the present time. Despite the apparent conservation today, a significant violation of baryon number in the early universe is typically needed for baryogenesis. The main reason for this is that we expect inflation to dilute away any pre-existing baryon number. Thus, we want to go from a universe with $B = 0$ to one in which $B > 0$, and this obviously requires the non-conservation of B number.

Many theories of physics beyond the SM predict new sources of B violation [20]. However, baryon number is also violated in the SM itself! The way this occurs is somewhat subtle; non-perturbative transitions between different vacuum states of the $SU(2)_L$ gauge theory change the combination of charges $(B + L)$. These transitions occur by quantum tunnelling at zero temperature with an unobservably slow rate, but can proceed much more quickly in the hot early universe at $T \gtrsim 100$ GeV through thermal fluctuations. The net result is

[3] For reference, the size of the observable universe today is about 10^4 Mpc.

rapid $(B+L)$ violation at high temperatures and approximate $(B+L)$ conservation at low temperatures.

Since the violation of $(B+L)$ in the SM is not very obvious, and because it plays an important role in every baryogenesis mechanism operative at temperatures $T \gtrsim 100$ GeV, it is worth spending a bit of time on the topic [2, 3, 4, 21, 22]. The first thing to notice is that the classical SM Lagrangian is invariant under independent $U(1)_B$ and $U(1)_L$ transformations. For these transformations to be a full quantum symmetry, the path integral measure used to define the theory must also be invariant, and it turns out that it is not. Instead, a one-loop calculation gives the result

$$\partial_\mu j^\mu_B \;=\; \partial_\mu j^\mu_L \;=\; \frac{n_g}{32\pi^2}\left(g^2 W^a_{\mu\nu}\widetilde{W}^{a\,\mu\nu} - g'^2 B_{\mu\nu}\widetilde{B}^{\mu\nu}\right)\;, \tag{4}$$

where $n_g = 3$ is the number of generations, $W^a_{\mu\nu}$ and $B_{\mu\nu}$ are the $SU(2)_L$ and $U(1)_Y$ field strength tensors, $\widetilde{W}^{a\,\mu\nu} = \frac{1}{2}\epsilon^{\mu\nu\alpha\beta}W^a_{\mu\nu}$ is the dual tensor (with $\epsilon^{0123} = +1 = -\epsilon_{0123}$), and the Noether currents of the classical symmetry (summed over generations) are

$$j^\mu_B \;=\; \frac{1}{3}\sum_i \left(\bar{Q}^i_L\gamma^\mu Q^i_L + \bar{u}^i_R\gamma^\mu u^i_R + \bar{d}^i_R\gamma^\mu d^i_R\right) \tag{5}$$

$$j^\mu_L \;=\; \sum_i \left(\bar{L}^i_L\gamma^\mu L^i_L + \bar{e}^i_R\gamma^\mu e^i_R\right)\;. \tag{6}$$

An important feature of the right-hand side of Eq. (4) is that it can be written as a total divergence:

$$\frac{g^2}{32\pi^2}W^a_{\mu\nu}\widetilde{W}^{a\,\mu\nu} = \partial_\mu K^\mu\;, \tag{7}$$

with

$$K^\mu_2 = \frac{g^2}{32\pi^2}\epsilon^{\mu\nu\alpha\beta}\left(W^a_{\nu\alpha}W^a_\beta - \frac{g}{3}\epsilon^{cab}W^c_\nu W^a_\alpha W^b_\beta\right)\;, \tag{8}$$

where $\epsilon^{cab} = f^{cab}$ is the $SU(2)_L$ structure constant. A similar expression with $f^{cab} \to 0$ applies for the $U(1)_Y$ term.

While Eq. (4) implies the non-conservation of $(B+L)$, it does not tell us how it actually occurs. To understand the underlying mechanism, we must make a brief detour and discuss the vacuum structure of pure non-Abelian gauge theories. It can be shown that such theories have many independent classical vacua that can be labelled by integers $N_{CS} \in \mathbb{Z}$, shown schematically in Fig. 2. These integers correspond to the topology of the mapping of the non-Abelian group space to the Euclidean spacetime boundary at infinity, and are called the *Chern-Simons number* of the mapping. Let us now define the time- and gauge-dependent quantity $N_{CS}(t)$ by

$$\begin{aligned} N_{CS}(t) \;&=\; \int d^3x\, K^0_2 \qquad (9)\\ &\to\; -\frac{g^3}{96\pi^2}\int d^3x\,\epsilon^{ijk}W^a_i W^b_j W^c_k \qquad (W^a_0 \to 0 \text{ gauge, } W^a_{\mu\nu} = 0) \end{aligned}$$

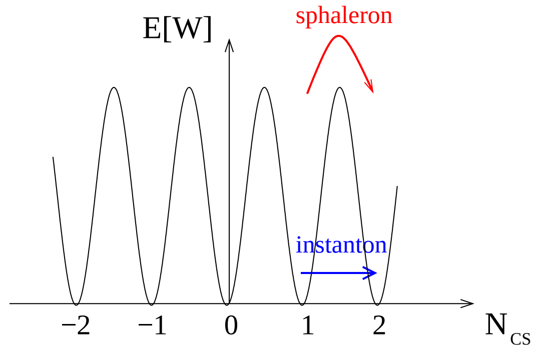

Figure 2: Cartoon of the vacuum structure of a non-Abelian gauge theory.

where W^a_μ refers to an $SU(2)$ gauge field. When this quantity is evaluated on a classical vacuum configuration with $W^a_0 = 0$ (which can be achieved by a choice of gauge) and $W^a_{\mu\nu}$ (so that it is a vacuum configuration), it becomes time-independent and coincides with the N_{CS} value of that vacuum. Beyond just having a non-trivial vacuum structure, there also exist solutions of the Euclidean space non-Abelian field equations with finite action that connect between vacua with different N_{CS} values [21, 22]. These solutions correspond to quantum tunnelling between the vacua, and they are called *instantons* because they have finite extent in both space and time.

Let us now return to Eq. (4) and compute the change in the charge $B = \int d^3x\, j^0_B$ between $t = -\infty$ and $t = +\infty$,

$$\begin{aligned} \Delta B &= \int_{-\infty}^{\infty} dt\, \partial_0 \int d^3x\, j^0_B \qquad (10) \\ &= \int_{-\infty}^{\infty} dt\, \int d^3x\, \left[\vec{\nabla}\cdot\vec{j}_B + \frac{n_g}{32\pi^2}\left(g^2 W^a_{\mu\nu}\widetilde{W}^{a\,\mu\nu} - g'^2 B_{\mu\nu}\widetilde{B}^{\mu\nu}\right)\right] \ . \end{aligned}$$

The spatial gradient term above gives zero because it reduces to a surface term involving the fermion fields. For the gauge terms above, let us focus on the $SU(2)_L$ piece and evaluate it using Eq. (7) in a gauge with $W^a_0 = 0$,

$$\begin{aligned} \int d^4x\, \frac{g^2}{32\pi^2} W^a_{\mu\nu}\widetilde{W}^{a\,\mu\nu} &= \int_{-\infty}^{\infty} dt \int d^3x\, (\partial_0 K^0_2 - \vec{\nabla}\cdot\vec{K}_2) \qquad (11) \\ &= \int d^3x\, K^0_2\Big|_{t\to\infty} - \int d^3x\, K^0_2\Big|_{t\to-\infty} \ + \ 0 \\ &= N_{CS}(t\to\infty) - N_{CS}(t\to-\infty) \ . \end{aligned}$$

Here, we have used the fact that field strengths must vanish at spacetime infinity to identify the spacial integral of K^0 with N_{CS} and set $K^i \to 0$ on the spatial boundary.[4] The same arguments show that the corresponding expression for the hypercharge gauge field vanishes.

[4] Any sensible field configuration has finite total energy (Euclidean action). Since the Hamiltonian density for a pure gauge theory is $\mathcal{H} = \Theta^{00} = F^{a\,0\mu}F^a_{0\mu} - \frac{1}{4}F^a_{\mu\nu}F^{a\,\mu\nu} = \frac{1}{2}(\vec{E}^a\cdot\vec{E}^a + \vec{B}^a\cdot\vec{B}^a) \geq 0$, a neccessary condition for finite energy is vanishing field strengths at spatial infinity [22, 23].

Combining Eq. (10) with Eq. (11), the net result is [24]

$$\Delta B = n_g \Delta N_{CS} \ . \tag{12}$$

Therefore baryon number violation corresponds directly to changes in the vacuum state of the $SU(2)_L$ non-Abelian gauge theory! As discussed above, this can occur through quantum tunnelling transitions called instantons. On very general grounds, it can be shown that the Euclidean action for a such a tunnelling event is $S_E \geq 8\pi^2|\Delta N_{CS}|/g^2$, implying a tunnelling rate proportional to

$$\Gamma \ \propto \ e^{-S_E} \ \leq \ e^{-8\pi^2/g^2} \ \simeq \ 10^{-160} \ . \tag{13}$$

This tiny dimensionless factor makes the instanton transitions much too slow to be observed, and consistent with the apparent stability of the proton.

The situation is very different at finite temperature. Instead of tunnelling through the barriers between vacua, thermal fluctuations can push the system over top of them. The rate for these thermal transitions is closely linked to energy of the lowest saddle point configuration in field space connecting a pair of classical $SU(2)_L$ vacua. This configuration involves both the vector fields W_μ^a and the Higgs field H, and is called the *sphaleron* [25]. When the Higgs vacuum expectation value (VEV) is non-zero, $\langle H \rangle = v$, the sphaleron energy is [26]

$$E_{sp}(T) \ \simeq \ \frac{8\pi \, v(T)}{g} \, f(m_h/m_W) \ , \tag{14}$$

for some function $f(m_h/m_W)$ of order unity. This leads to a sphaleron rate per unit volume at temperature T of [27, 28]

$$\Gamma_{sp} \ \simeq \ A \, (\alpha_W T)^4 \left[\frac{E_{sp}(T)}{T}\right]^7 e^{-E_{sp}(T)/T} \ , \tag{15}$$

where A is a constant, $\alpha_W = g^2/4\pi$, and the exponential can be understood as a Boltzmann suppression factor. At very high temperatures, $T \gg m_W$, the Higgs VEV is expected to vanish and a different estimate of the transition rate is needed based on the likelihood of certain N_{CS}-changing gauge field configurations to occur in the plasma. The best determination of the rate comes from lattice studies, which give [29, 30]

$$\Gamma_{sp} = (18 \pm 3) \, \alpha_W^5 T^4 \ , \tag{16}$$

which is notable for the lack of Boltzmann suppression.[5]

While our discussion has concentrated on $SU(2)_L$, a similar story applies to the strong group $SU(3)_c$; it has a non-trivial classicial vacuum structure and there exist transitions that connect these states. In contrast to $SU(2)_L$, however, B and L are not anomalous with respect to $SU(3)_c$, and thus $SU(3)_c$ instantons and "strong sphalerons" do not violate baryon or lepton number. Even so, we will see that these transitions can be important for baryogenesis.

[5]These transitions are still called sphaleron transitions (or just sphalerons), even though they do not actually involve the sphaleron gauge-Higgs field configuration.

2.2 C and CP Violation

The second Sakharov condition of C and CP violation is less obvious than B violation. These are needed to distinguish matter from antimatter in order to drive the B violating reactions to produce more baryons than antibaryons [3]. We will illustrate how this works in specific examples later on in the lectures. For now, it will be useful to review the action of C, P, and CP on the SM [5, 31, 32].

To begin, consider a theory with a complex scalar field ϕ and Lagrangian

$$\mathscr{L} = |\partial\phi|^2 - m^2|\phi|^2 \ . \tag{17}$$

In this form, we see that the theory has a global $U(1) \cong SO(2)$ global symmetry, with a corresponding Noether current

$$j_\mu = i\phi^\dagger \overleftrightarrow{\partial}_\mu \phi \ . \tag{18}$$

In terms of mode expansions in the free-field limit, we have

$$\phi(x) = \int \widetilde{dk} \left[a(\vec{k})e^{-ik\cdot x} + b^\dagger(\vec{k})e^{ik\cdot x}\right] \ , \tag{19}$$

$$Q = \int d^3x\, j^0 = \int \widetilde{dk} \left[a^\dagger(\vec{k})a(\vec{k}) - b^\dagger(\vec{k})b(\vec{k})\right] \ , \tag{20}$$

where $\widetilde{dk} = d^3k/2E_k(2\pi)^3$. Based on these mode expansions, we identify $a^\dagger(\vec{k})$ as the creation operator for a particle with positive charge and $b^\dagger(\vec{k})$ for a particle with negative charge. Since the underlying global $U(1)$ symmetry forces the masses of these particles to be equal and their charges to be opposite, we call the b-species the antiparticle of the a-species.[6] The natural definition of charge conjugation (C) on this theory is the exchange of particles and antiparticles in the sense

$$a(\vec{k}) \ \leftrightarrow \ b(\vec{k}) \ . \tag{21}$$

In terms of the field operators, this implies

$$\phi(x) \ \leftrightarrow \ \phi^\dagger(x) \ , \tag{22}$$

as well as $Q \to -Q$. It is clear that the scalar Lagrangian is invariant under this transformation.

For a Dirac fermion, written in 4-component notation in the chiral representation, the mode expansion in the free theory is

$$\psi(x) = \int \widetilde{dk} \sum_s \left[a_s(\vec{k})u(k,s)e^{-ik\cdot x} + b_s^\dagger(\vec{k})v(k,s)e^{ik\cdot x}\right] \ . \tag{23}$$

[6] Note that this only works to the extent that there is an underlying global symmetry. For example, adding a small explicit breaking term of the form $\mathscr{L} \to \mathscr{L} - \Delta m^2\left[(\phi)^2 + (\phi^*)^2\right]$ would lead to two real scalar mass eigenstates with mass-squared eigenvalues of $m^2 \mp \Delta m^2$ [33].

The standard free Dirac theory has a built-in $U(1)$ symmetry with current $j^\mu = \bar{\psi}\gamma^\mu\psi$ and charge

$$Q = \int \widetilde{dk} \sum_s \left[a_s^\dagger(\vec{k})a_s(\vec{k}) - b_s^\dagger(\vec{k})b_s(\vec{k})\right] \ . \tag{24}$$

We identify $a_s^\dagger(\vec{k})$ as the creation operator for a fermion of momentum $\vec{k}$ and spin s, and $b_s^\dagger(\vec{k})$ as the creation operator for an antifermion with the same momentum and spin, but opposite charge. The natural defintion of C on the mode operators is thus

$$a_s(\vec{k}) \leftrightarrow b_s(\vec{k}) \ . \tag{25}$$

The spinors appearing the mode expansions complicate the result of acting C on the field variables. Without going into details (which can be found in Refs. [5, 34]), the effect in the chiral representation is

$$\psi(x) \ \to \ (-i\gamma^2\gamma^0)\,\overline{\psi}^{\,t} \ , \qquad \bar{\psi}(x) \ \to \ \psi^t(-i\gamma^2\gamma^0) \ . \tag{26}$$

Based on this result, let us define $\Gamma_C = -i\gamma^2\gamma^0 = -\Gamma_C^t = -\Gamma_C^\dagger$. The action of C on fermion bilinears is then [34]

$$\overline{\psi}\Gamma_i\chi \ \to \ \overline{\chi}\,(\Gamma_C^{-1}\Gamma_i^t\Gamma_C)\,\psi \ = \ \eta_i\,\overline{\chi}\,\Gamma_i\,\psi \ , \tag{27}$$

where $\eta_i = +1$ for $\Gamma_i = 1$, $i\gamma^5$, $\gamma^\mu\gamma^5$ and $\eta_i = -1$ for $\gamma_i = \gamma^\mu$, $\sigma^{\mu\nu}$. In particular, the current $\overline{\psi}\gamma^\mu\psi$ is odd under C, as expected from the mode expansion. The free Dirac action is also invariant under this transformation.

The action of C on vector bosons can be deduced the result from the way the photon field couples to the electromagnetic current in QED,

$$\mathscr{L}_{QED} \ \supset \ -\,e\,A_\mu\,j^\mu_{em} \ = \ -\,e\,A_\mu\sum_i Q_i\overline{\psi}_i\gamma^\mu\psi_i \ . \tag{28}$$

Since the current is odd under C, $j^\mu_{em} \to -j^\mu_{em}$, the theory will be invariant under C if the photon field is odd as well,

$$A_\mu \ \to \ -\,A_\mu \ . \tag{29}$$

A similar story applies in QCD with one important modification [31]. The matter current is now

$$j^{\mu\,a}_{QCD} = \sum_i \overline{q}_i\gamma^\mu t^a q_i \ \to \ -\sum_i \overline{q}_i\gamma^\mu(t^a)^t q_i \ , \tag{30}$$

where t^a is a generator of the fundamental representation of $SU(3)_c$. It is possible to choose these generators such that they are either symmetric or antisymmetric for each value of $a = 1,\ldots,(N_c^2-1)$. This implies that under C

$$j^{\mu\,a}_{QCD} \ \to \ -\,\eta(a)j^{\mu\,a}_{QCD} \ , \tag{31}$$

where $\eta(a) = +1\,(-1)$ if t^a is symmetric (antisymmetric). With some work, one can show that this is consistent with the gluon self-interactions of QCD provided we assign

$$G^a_\mu \;\to\; -\,\eta(a) G^a_\mu \ . \tag{32}$$

With this, the QCD field strength $G^a_{\mu\nu}$ is also odd up to the $\eta(a)$ factor. With these definitions, both QED and QCD are invariant under C.

To apply C to the full SM, we must consider chiral fermions coupling to vector bosons. Recall that in the chiral representation, a Dirac fermion has chiral components

$$\psi = \left(\begin{array}{c}\psi_L\\ \psi_R\end{array}\right) \ . \tag{33}$$

Using the action of C defined above on the chiral components, one finds

$$\psi_L \;\to\; i\sigma^2\psi_R^* \ , \qquad \psi_R \;\to -i\sigma^2\psi_L^* \ . \tag{34}$$

We also define the action of C on the electroweak vectors bosons in the same way as for the photon and gluon above [31]. Together, this implies that the SM violates C since the chiral ψ_L and ψ_R components of the SM fermions (written as 4-component Dirac objects) come from different representations of the underlying $SU(2)_L \times U(1)_Y$ gauge group.

Turning next to parity (P), recall that it acts on spacetime according to

$$t \;\to t \ , \qquad \vec{x} \to -\vec{x} \ . \tag{35}$$

For scalar fields this produces

$$\phi(t,\vec{x}) \to \phi(t,-\vec{x}) \ , \tag{36}$$

while for Dirac fermions it gives (up to a possible phase)

$$\psi(t,\vec{x}) \to \gamma^0\psi(t,-\vec{x}) \ . \tag{37}$$

This implies that the spatial components of the QED and QCD matter currents are odd under P, and thus parity is maintained in these theories if

$$A^a_\mu \to -\eta(\mu) A^a_\mu(t,-\vec{x}) \ , \tag{38}$$

with $\eta(\mu = 0) = -1$ and $\eta(\mu = i) = 1$ For the SM with chiral fermions, the Dirac fermion transformation law implies for the chiral components that

$$\psi_L(t,\vec{x}) \;\to\; \psi_R(t,-\vec{x}) \ , \qquad \psi_R(t,\vec{x}) \;\to\; \psi_L(t,-\vec{x}) \ . \tag{39}$$

We also take the SM vector bosons to transform in the same way as the photon [31]. Thus, just like C, P is violated in the SM by the chiral fermion representations under the electroweak group.

While C and P are broken by the presence of chiral fermion reps in the SM, the combination CP is not. Combining the two operations, a CP transformation on a Dirac fermion ψ gives

$$\psi(t,\vec{x}) \;\to\; -\Gamma_C\psi^*(t,-\vec{x}) \;. \tag{40}$$

In terms of chiral components,

$$\psi_L(t,\vec{x}) \;\to\; -i\sigma^2\psi_L^*(t,-\vec{x}) \;, \qquad \psi_R(t,\vec{x}) \;\to\; i\sigma^2\psi_R^*(t,-\vec{x}) \;. \tag{41}$$

Since each chiral component transforms into itself, it is consistent with the chiral representations of the SM. Turning to the action, a CP transformation basically just replaces each operator by its complex conjugate while leaving the numerical coefficients unchanged. For example,

$$\overline{e}_R i\gamma^\mu(\partial_\mu - ig'B_\mu)e_R \;\to\; \overline{e}_R i\gamma^\mu(\partial_\mu - ig'B_\mu)e_R \tag{42}$$

$$(y_u)_{ij}\overline{Q}_{L\,i}Hu_{R\,j} + (y_u^\dagger)_{ji}\overline{u}_{R\,j}H^\dagger Q_{L\,i} \;\to\; (y_u^t)_{ji}\overline{u}_{R\,j}H^\dagger Q_{L\,i} + (y_u^*)_{ij}\overline{Q}_{L\,i}Hu_{R\,j} \tag{43}$$

The operator in Eq. (42) is clearly invariant, while that of Eq. (43) is invariant if the Yukawa coupling matrix is real.

The result of Eq. (43) is more general than just the SM: CP violation is typically associated with the presence of complex couplings in the action. However, while the presence of complex couplings in the action can give rise to CP violation, it does not guarantee it. Let us see how this works in the SM. After electroweak symmetry breaking, the fermion fields are usually transformed by making field redefinitions in flavor space to obtain diagonal mass matrices with positive mass eigenvalues [32]. The only remnant of these transformations in the interactions of the theory is the Cabbibo-Kobayashi-Maskawa (CKM) matrix. To be precise, the flavor transformations take the form

$$u_L \to V_{u_L}u_L \;, \quad u_R \to V_{u_R}u_R \;, \quad d_L \to V_{d_L}d_L \;, \quad d_R \to V_{d_R}d_R \;, \tag{44}$$

which gives the CKM matrix

$$V_{CKM} = V_{u_L}^\dagger V_{d_L} \;. \tag{45}$$

This is a 3×3 unitary matrix, and can be parametrized by three rotation angles and six phases. The phases in the CKM matrix are a potential source of CP violation. However, not all of these phases are physical since we can make further field redefinitions that do not alter the positive diagonal fermion mass matrices but do change the CKM matrix. These have the form

$$V_{u_L} = diag\{e^{i\alpha_1}, e^{i\alpha_2}, e^{i\alpha_3}\} = V_{u_R} \;, \qquad V_{d_L} = diag\{e^{i\beta_1}, e^{i\beta_2}, e^{i\beta_3}\} = V_{d_R} \;. \tag{46}$$

There are six phases here, and these can be used to remove five of the six independent phases in the CKM matrix.[7] This leaves a single irreducible phase in the CKM that leads to

[7] The universal overall phase with $\alpha_i = \beta_i$ coincides with baryon number and does not change V_{CKM}.

physical CP violation in the SM. This counting also implies that a version of the SM with n_g generations would have a CKM matrix with $(n_g - 1)(n_g - 2)/2$ physical CP violating phases. Experiments have measured CKM phase CP violation in Kaon and B-meson mixing and decay [11, 35].

There is an additional potential source of CP violation in the SM that presents a serious puzzle. It corresponds to the allowed operator [31, 36]

$$\mathscr{L}_{SM} \supset \frac{g_3^2}{32\pi^2}\Theta\, G^a_{\mu\nu}\widetilde{G}^{a\,\mu\nu} . \tag{47}$$

This operator violates P (and T but not C) and therefore CP. Its origin is connected to the non-trivial vacuum structure of $SU(3)_c$, and represents a quantum superposition of the classical N_{CS} vacua [21, 22]. The potentially observable effect of a non-zero Θ is a T-violating permanent electric dipole moment (EDM) of the neutron. Current limits on the neutron EDM imply

$$\Theta \lesssim 10^{-10} . \tag{48}$$

The *strong CP problem* refers to the mystery of why this dimensionless parameter (with natural range $\Theta \in [0, 2\pi)$) is so small. Note as well that Θ term can also be written for $SU(2)_L$. However, unlike for $SU(3)_c$ (with massive quarks), the $SU(2)_L$ term is not observable because it can be removed by making a $(B + L)$ transformation without altering any other part of the SM action.

While C, P, and CP are all broken explicitly in the SM, the combination CPT is respected. This is a very general requirement for any reasonable Lorentz quantum field theory [37]. Very bad things tend to happen when CPT is violated, and the limits on CPT breaking from data are extremely severe [38].

2.3 Departure from Thermodynamic Equilibrium

Thermodynamic equilibrium corresponds to the state of a system left to settle down for an infinitely long amount of time. Expectation values are time independent in equilibrium, and therefore it would not be possible to from a $\langle B = 0\rangle$ state to a $\langle B \neq 0\rangle$ state if full thermodynamic were maintained throughout. The cosmological expansion of the universe implies a departure from thermodynamic equilibrium, but it is usually too small to be effective in realistic mechanisms of baryogenesis. Before illustrating explicitly how departures from thermodynamic equilibrium enter into baryogenesis, let us review here a few aspects of thermodynamics in the early universe.

Particles in the cosmological plasma can usually be described by distribution functions in momentum space corresponding to a dilute gas of temperature T and chemical potentials μ_i [14, 39]. If thermodynamic equilibrium is maintained, these take the form

$$f_i(\vec{p}; T, \mu_i) = \left[e^{(E_i-\mu_i)/T} \mp 1\right]^{-1} , \tag{49}$$

where i refers to the particle species, $E_i = \sqrt{m_i^2 + \vec{p}^2}$, and the minus (plus) sign corresponds to bosons (fermions). The distribution function is defined such that the mean number density is of the species i is

$$n_i(T,\mu) \ = \ g_i \int \frac{d^3p}{(2\pi)^3}\, f_i(\vec{p}) \ \simeq \ \begin{cases} g_i \xi_i^{(n)} \frac{\zeta(3)}{\pi^2} T^3 & ; \quad T \gg m_i,\, \mu_i \\ g_i \left(\frac{m_i T}{2\pi}\right)^{3/2} e^{-(m_i-\mu_i)/T} & ; \quad T \ll (m_i - \mu_i) \end{cases} \tag{50}$$

where g_i is the number of internal degrees of freedom, $\xi_i^{(n)} = 1$ $(3/4)$ for i a boson (fermion), and $\zeta(3) \simeq 1.2026$. Similarly, the mean energy density due to the species is

$$\rho_i(T,\mu) \ = \ g_i \int \frac{d^3p}{(2\pi)^3}\, E_i\, f_i(\vec{p}) \ \simeq \ \begin{cases} g_i \xi_i^{(\rho)} \frac{\pi^2}{30} T^4 & ; \quad T \gg m_i,\, \mu_i \\ m_i n_i & ; \quad T \ll (m_i - \mu_i) \end{cases} , \tag{51}$$

with $\xi_i^{(\rho)} = 1$ $(7/8)$ for i a boson (fermion). The partial pressure of species i is obtained by integrating with $\vec{p}^{\,2}/3m_i$, and is given by $\rho_i = \rho_i/3$ in the relativisitc limit and $p_i = n_i T \ll \rho_i$ in the non-relativistic limit.

Chemical potentials play an important role in baryogenesis, and it is worth looking at them in a bit more detail. In full thermodynamic equilibrium the only independent chemical potentials are those for conserved charges, rather than those for individual particle species. This is because interactions relate the species-specific chemical potentials μ_i to each other. If the process $A + B \leftrightarrow C + D$ is allowed, full equilibrium implies the relation

$$\mu_A + \mu_B = \mu_C + \mu_D\ . \tag{52}$$

Full equilibrium also implies that photons and gluons have $\mu = 0$ (*e.g.* $e^+e^- \to e^+e^-\gamma$ is allowed), and that particle-antiparticle pairs have equal and opposite chemical potentials since they can annihilate:

$$\mu_\psi = -\mu_{\bar{\psi}}\ . \tag{53}$$

This relation implies that μ_ψ sets the number density asymmetry between a particle ψ and its antiparticle $\bar{\psi}$,

$$\begin{aligned} n_\psi - n_{\bar{\psi}} \ &= \ g_\psi \int \frac{d^3p}{(2\pi)^3}\, [f_\psi(\vec{p};T,\mu_\psi) - f_\psi(\vec{p};T,-\mu_\psi)] \\ &\simeq \ \begin{cases} \xi_\psi^{(\Delta)} \frac{g_\psi}{6} T^3 \left[\left(\frac{\mu_\psi}{T}\right) + \frac{1}{\pi^2}\left(\frac{\mu_\psi}{T}\right)^3 + \ldots\right] & ; \ T \gg m_\psi \\ 2\sinh\left(\frac{\mu_\psi}{T}\right)\, g_\psi \left(\frac{m_\psi T}{2\pi}\right) e^{-m_\psi/T} & ; \ T \ll m_\psi \end{cases} \end{aligned} \tag{54}$$

where $\xi_i^{(\Delta)} = 2$ (1) for ψ a boson (fermion). Note that for $\mu_\psi \ll T$, the number asymmetry is approximately linear in the chemical potential.

The energy density of the cosmological plasma causes the universe to expand with scale factor $a(t)$. The expansion rate is controlled by the equation

$$H^2 = \left(\frac{\dot{a}}{a}\right)^2 = \frac{1}{3M_{\rm Pl}^2}\rho \ , \tag{55}$$

where ρ is the total energy density and $M_{\rm Pl} = 2.4 \times 10^{18}$ GeV is the reduced Planck mass. Extrapolating the SM back in time, the early universe was dominated by radiation with total energy density

$$\rho_r = \frac{\pi^2}{30} g_* T^4 \ , \tag{56}$$

with g_* an effective total number of "light" degrees of freedom given by

$$g_* \ \simeq \ \sum_i \xi_i^{(\rho)} g_i \left(\frac{T_i}{T}\right)^4 \Theta(T_i - m_i) \ , \tag{57}$$

where we have allowed for different individual temperatures T_i. The energy density dilutes with the scale factor as dictated by entropy conservation, which follows from the continuity equation. Specifically,

$$0 = \frac{d}{dt}(sa^3) \ , \tag{58}$$

where s is the entropy density defined by

$$s = \frac{(\rho + p)}{T} = \frac{2\pi^2}{45} g_{*S} T^3 \ , \tag{59}$$

with

$$g_{*S} \simeq \sum_i \xi_i^{(\rho)} g_i \left(\frac{T_i}{T}\right)^3 \Theta(T_i - m_i) \ . \tag{60}$$

For constant g_{*S} we see that $a(t) \propto 1/T$. When radiation is dominant, the Hubble rate is

$$H \ \simeq \ \sqrt{\frac{\pi^2 g_*}{90}} \frac{T^2}{M_{\rm Pl}} \ . \tag{61}$$

Compared to radiation with $\rho_r \sim a^{-4}$, the density of decoupled non-relativistic matter dilutes as $\rho_m \sim a^{-3}$. The era of radiation domination ended fairly recently, at redshift $z = a(t_0)/a(t_{eq}) - 1 \simeq 3600$, corresponding to a temperature close to $T_{eq} \simeq 1\,\text{eV}$.

Much of the evolution of the early universe can be described by near-equilibrium thermodynamics with an adiabatically changing temperature $T(t)$. The condition for equilibration is usually that the processes that lead to it occur much faster than the rate of expansion of the universe as determined by the Hubble "constant" $H(t) = \dot{a}/a$. In discussing thermodynamic equilibrium, it helpful to separate two differnet aspects of it: *kinetic equlibrium*, and *chemical* *equilibrium.* Kinetic equilibrium is usually maintained fairly easily among light particles through rapid scattering and decay processes that transfer energy efficiently among them, and implies that all the particles participating in these processes have the same temperature $T(t)$. This is also sufficient to ensure that the distribution functions of these particles take the form of Eq. (49) with chemical potentials $\mu_i(t)$. Chemical equilibrium occurs when the chemical potentials of these particles also satisfy the constraints of Eq. (52) for all possible reactions.

In practice, chemical equilibrium is more challenging to maintain in the early universe than kinetic equilibrium. Deviations from chemical equilibrium (with kinetic equilibrium) can described by *Boltzmann equations* for number densities. A possibly familiar example is the evolution of the number density of a massive dark matter species χ that can annihilate with itself into SM final states [14]:

$$\frac{dn_\chi}{dt} = -3Hn_\chi - \langle\sigma_{ann}v\rangle\left[n_\chi^2 - (n_\chi^{eq})^2\right] \ , \tag{62}$$

n_χ^{eq} is the number density in full equilibrium with $\mu_\chi = 0$ and $\langle\sigma_{ann}v\rangle$ is the thermally averaged annihilation cross section. For $n_\chi\langle\sigma_{ann}v\rangle \gg H$, the second term on the right-hand side of Eq. (62) dominates and $n_\chi \to n_\chi^{eq}$ enforcing $\mu_\chi \to 0$. Conversely, if $n_\chi\langle\sigma_{ann}v\rangle \ll H$ the number of density of χ simply dilutes with the expansion of the universe proportionally to $1/a^3$, and this usually requires a non-zero chemical potential $\mu_\chi(t)$. As a general rule of thumb, a reaction must have a rate greater than Hubble $H(t)$ to contribute to equilibrium. For inelastic processes contributing to chemical equilibrium, this means that we should only expect equalities of the form of Eq. (52) if their rates are greater than Hubble.

A particularly important set of processes for baryogenesis are the $(B + L)$ violating sphaleron transitions discussed above. When they are in equilibrium, they imply

$$0 = \sum_i(3\mu_{q_i} + \mu_{\ell_i}) \ , \tag{63}$$

where the sum runs over generations. To estimate when the sphaleron transitions are fast enough to maintain equilibrium, we should compare the rates of Eqs. (15,16) within a volume of size $V \sim 1/T^3$ (corresponding to a typical interparticle spacing) to the Hubble rate during radiation domination [2]. At high temperature $T \gg m_W$ Eq. (16) is relevant and one finds sphaleron equilibration for

$$T \lesssim 10^{12} \text{ GeV} \ . \tag{64}$$

This equilibrium is maintained until electroweak symmetry breaking occurs at $T \sim m_W$, at which point the sphaleron rate becomes exponentially suppressed and quickly turns off. A detailed study within the SM finds that this occurs for [30]

$$T \simeq (131.7 \pm 2.3) \text{ GeV} \ . \tag{65}$$

Thus, sphaleron transitions are in equilibrium in the SM for $T \in [132,\ 10^{12}]$ GeV.

3 Baryogenesis Mechanisms

Beyond just the Sakharov ingredients, explaining the baryon asymmetry requires a concrete mechanism for how it occurs. A number of potentially viable mechanisms for baryogenesis have been proposed, and they all require new physics beyond the SM in one form or another. We present here a toy model of baryogenesis that illustrates qualitatively how the ingredients combine to make baryons, and we give an overview of some of the most popular mechanisms on the market.

3.1 A Toy Model for Baryogenesis

Consider a theory containing a complex scalar X and a pair of two-particle decay modes, $X \to A_1$ and $X \to A_2$, where A_1 and A_2 refer to the two-particle final states [14, 2]. Assume further that state A_1 has baryon number B_1 and state A_2 has baryon number B_2. By CPT, the antiparticle X^* must have decays to $\overline{A}_1$ and $\overline{A}_2$, each with baryon numbers $-B_1$ and $-B_2$. Note that these decays violate B number if $B_1 \neq B_2$.

A baryon asymmetry can be created from the decays of X and X^*, even if they have no initial asymmetry, $\mu_X = \mu_{X^*} = 0$, provided there is C and CP violation. With both, the partial decay widths of X and X^* can differ, taking the form

$$\begin{aligned} \Gamma(X \to A_1) &= \Gamma_1 + \varepsilon\Delta\Gamma \ , \qquad \Gamma(X \to A_2) = \Gamma_2 - \varepsilon\Delta\Gamma \ , \\ \Gamma(X^* \to \overline{A}_1) &= \Gamma_1 - \varepsilon\Delta\Gamma \ , \qquad \Gamma(X^* \to \overline{A}_2) = \Gamma_2 + \varepsilon\Delta\Gamma \end{aligned} \tag{66}$$

Note that the total decay width of X is equal to that of X^*, $\Gamma_X = \Gamma_{X^*} = \Gamma_1 + \Gamma_2$, as required by CPT invariance, and the dimensionless coefficient ε is non-zero only with C and CP violation.

Next, let us assume that X and X^* decay in the early universe at time $t = \tau = 1/\Gamma_X$ when the yields of both states are

$$Y_X^{dec}(\tau) = \frac{n_X}{s} = Y_{X^*}^{dec}(\tau) \ , \tag{67}$$

where s is the entropy density at time τ. The baryon asymmetry created immediately after these decays is

$$\begin{aligned} Y_B^{dec} &= Y_X^{dec}\left(B_1\frac{\Gamma_1 + \varepsilon\Delta\Gamma}{\Gamma_X} + B_2\frac{\Gamma_2 - \varepsilon\Delta\Gamma}{\Gamma_X} - B_1\frac{\Gamma_1 - \varepsilon\Delta\Gamma}{\Gamma_X} - B_2\frac{\Gamma_1 + \varepsilon\Delta\Gamma}{\Gamma_X}\right) \\ &= 2\varepsilon\Delta\Gamma(B_1 - B_2)Y_X^{dec} \ . \end{aligned} \tag{68}$$

This result illustrates two of the three Sakharov conditions: $B_1 \neq B_2$ requires B violation; and $\varepsilon \neq 0$ requires C and CP violation.

To see the role of the third condition, departure from equilibrium, we must think a bit more carefully about what processes can occur in the early universe. On top of decays, there can also occur inverse decays like $A_1 \to X$ and $2 \to 2$ processes such as $A_1 \leftrightarrow A_2$ through an off-shell intermediate X boson. If these processes are active, they will tend to *wash out* the baryon asymmetry created by the decays. For example, the asymmetry here arises from an excess of A_1 over A_2 and a deficit of $\overline{A}_1$ over $\overline{A}_2$. This means that following the decays, the process $A_1 \to A_2$ will have a greater rate than $A_2 \to A_1$ and tend to cancel the asymmetry made by the decays. It is not hard to check that full chemical equilibration leads to zero final baryon asymmetry. Instead, if the decays of X and X^* occur very late at temperature $T \ll m_X$ with $n_X(\tau) = n_{X^*}(\tau) > n_X^{eq}(\tau)$, the inverse decay and $2 \to 2$ washout processes will be too slow relative to Hubble to restore chemical equilibrium, and the final baryon asymmetry today will be very close to the value at decay given in Eq. (68).

3.2 Some Baryogenesis Mechanisms

We give here a brief description of some of the most popular mechanisms of baryogenesis. In the lectures to follow, we will study leptogenesis and baryogenesis in more detail.

Leptogenesis

In leptogenesis, an asymmetry in lepton number L is created, and then reprocessed by electroweak sphaleron transitions into a non-zero B asymmetry. A very nice proposal to generate the L asymmetry is through the decay of a very heavy SM-singlet neutrino. Such neutrinos arise in many models for the masses and mixings of the neutrinos we observe, and they can give rise to explicit L violation. If their decays also involve C and CP violation and occur out of equilibrium, they can generate an initial lepton asymmetry L_i. Sphaleron transitions then transfer some of this lepton charge to baryon charge, creating the baryon asymmetry we observed today.

Electroweak Baryogenesis

Electroweak baryogenesis (EWBG) is a class of mechanisms to create baryons in the electroweak phase transition in the early universe, in which the Higgs field goes from $\langle H^\dagger H\rangle = 0$ to $\langle H^\dagger H\rangle = v^2 \simeq (174\text{ GeV})^2$. If the transition is strongly first order, it proceeds through the nucleation of bubbles of broken phase within the surrounding plasma of symmetric phase. These bubbles expand, collide, and coalesce until they eventually fill the entire universe. Baryon creation in EWBG occurs in the vicinity of the bubble walls, which provide a strong departure from thermodynamic equilibrium. Particle scattering off the bubble walls with C and CP violation can generate net chiral asymmetries (*e.g.* more left-handed quarks than right-handed quarks) which bias the sphaleron transitions outside the bubbles to create more baryons than antibaryons. Once the baryons are created, they are quickly swept up into the interior of the expansing bubbles where they are effectively stable.

GUT Baryogenesis

In many grand unified theories (GUTs) where $SU(3)_c \times SU(2)_L \times U(1)_Y$ gauge factors are embedded in a simple unified gauge group, the SM quarks and leptons are combined into irreducible representations of the GUT gauge group. For example, in minimal $SU(5)$ the matter representations are $\overline{\mathbf{5}} = (d_R^c,\, L_L)$ and $\mathbf{10} = (Q,\, u_R^c,\, e_R^c)$ for each generation [40]. These combinations imply that B and L are violated in the GUT, specifically by interactions involving exotic gauge and Higgs bosons [41]. The exotic gauge and Higgs states are assumed to get very large masses on the order of M_{GUT} from course of GUT symmetry breaking $SU(5) \to SU(3)_c \times SU(2)_L \times U(1)_Y$. If the exotic states decay out of equilibrium with C and CP violation, they can generate a net B asymmetry (similar to in our toy model) [1, 42].

A significant challenge for GUT baryogenesis scenarios is that many simple GUTs break $(B{+}L)$ but preserve $(B{-}L)$ [2]. We will see that this implies that the asymmetries created by heavy GUT particle decays will be washed out by sphaleron transitions. Furthermore, it can be difficult to create enough of the very heavy exotic states in standard models of inflation for the values of $M_{GUT} \gtrsim 10^{16}$ GeV typically implied by direct searches for nucleon decay [20]. Very high reheating temperatures can also lead to the overproduction of topological defects or gravitinos (in supersymmetric realizations) [43].

Affleck-Dine Baryogenesis

Affleck-Dine baryogenesis (ADBG) is a mechanism to generate baryons from the decays of an oscillating scalar field that carries baryon number [44, 45]. Consider a scalar field $\phi(x)$ with non-zero B (or L) charge Q_B. The contribution of the field to the baryon Noether charge is

$$\Delta j_B^\mu \;=\; iQ_B\phi^\dagger(\overleftarrow{\partial^\mu} - \overrightarrow{\partial^\mu})\phi \;. \tag{69}$$

Now suppose we expand the scalar as $\phi(x) = r(x)\exp[-i\theta(x)]/\sqrt{2}$ with real fields $r(x)$ and $\theta(x)$. The contribution to the baryon number density obtained from Eq. (69) is then

$$\Delta n_B \;=\; \Delta j_B^0 \;=\; Q_B\, r^2\, \frac{d\theta}{dt} \;. \tag{70}$$

Thus, an angular variation of the scalar field produces a non-zero baryon charge provided the radial amplitude is non-zero. Note that such an amplitude implies the spontaneous breaking of B, so we expect that the minimum of the potential for ϕ is located at $r = 0$. Even so, non-zero initial scalar field amplitudes arise readily in the early universe from inflationary effects If the scalar effective potential also contains terms that violate C, CP, and B, the subsequent time evolution will lead to a non-zero "kick" in the angular direction that produces a non-zero $d\theta/dt$ and thus a B charge. This charge can then be transferred to the SM fermions through B-conserving decays of the scalar field or scattering with the thermal background.

The ingredients needed for ADBG emerge very naturally in supersymmetric extensions of the SM [46, 47]. Scalar fields with non-zero B or L charge come from the superpartners of the SM fermions. Supersymmetry also puts severe constraints on the form of the scalar potential, and a common feature are *flat directions* in the scalar potential along which the scalar fields are only stabilized very feebly. A simple example is the H_u-L direction, with a potential of the form

$$V \;\simeq\; \frac{1}{2}(g^2+g'^2)(|H_u|^2 - |\widetilde{L}|^2)^2 + \frac{1}{M^2}(H\cdot\widetilde{L})^2 + (h.c.) + \ldots \tag{71}$$

The first term vanishes for $|H_u| \to (0,\, v)^t$ and $|L| \to (v,\, 0)^t$, and corresponds to the would-be flat direction, while the second term provides a higher-dimensional operator suppressed by a large mass M that lifts the flat direction and very weakly stabilizes the potential at the origin, and also violates L. When inflation and soft supersymmetry breaking are added, additional terms arise in the potential and these can violate C, CP, such as

$$V \;\supset\; (a\, m_{3/2} + b\, H)\frac{1}{M}(H_u\!\cdot\!\widetilde{L})^2 \tag{72}$$

where a and b can be complex and provide CP violation, $m_{3/2}$ is the scale of soft supersymmetry breaking, and H is the Hubble constant. Parametrizing the flat direction by the gauge invariant $L = -1$ scalar $\Phi^2 = (H_u \cdot \widetilde{L})$, the evolution of Φ can satisfy all the requirements needed for ADBG. In particular, the very flat potential for Φ implies that it can very easily be displaced during inflation, leading to a large fluctuation amplitude (*i.e.* r^2 in Eq. (70)) and a correspondingly large L charge.

Asymmetric Dark Matter and Dark Sector Baryogenesis

An interesting numerical observation is that the observed dark matter density is about a factor of five larger than the baryonic matter density. In most existing theories, dark matter and baryons are formed by completely independent processes, and from this point of view the similarity in their densities is a little puzzling. This might simply be an accident, but it could be suggestive of a link between the mechanisms of dark matter and baryon creation in the early universe.

Scenarios of asymmetric dark matter (ADM) scenarios can make such a link by connecting the dark matter density to the baryon asymmetry [48, 49]. The key feature of ADM models is that the DM density is set by an asymmetry between distinct DM and anti-DM particles, with efficient DM–anti-DM annihilation removing all but the excess of DM over anti-DM, in analogy to how the baryon asymmetry determines the baryon density. This stands in contrast to most WIMP DM scenarios, where the DM density typically arises from standard thermal freezeout. The DM density can also be connected to the baryon density if the DM charge is related to the baryon charge.

A simple example of ADM consists of new Dirac fermions X and Y and a complex scalar Φ that are all neutral under the SM together with the couplings (in 2-component fermion notation) [50, 51]

$$-\mathscr{L} \;\supset\; \frac{\lambda}{M^2} X u_R^c d_R^c d_R^c + \zeta\, XY\Phi + (h.c.) \tag{73}$$

where $u_R^c = \Gamma_C \bar{u}_R^t$. This coupling breaks standard B number, but it respects a generalized B number in which the charges are $[X] = 1$, $[Y] = y$, and $[\Phi] = -(1+y)$. Suppose an asymmetry is created in the generalized baryon number. This could arise from dynamics involving only the new SM-neutral *dark sector* states such as X, Y, and Φ, or through a more conventional baryogenesis mechanism in the visible sector. In either case, the first operator of Eq. (73) will share the asymmetry among the dark and visible sectors, corresponding to the equilibrium relations

$$\mu_X \;=\; \mu_{u_R} + 2\mu_{d_R} \;=\; -\,\mu_Y - \mu_\Phi\ . \tag{74}$$

If Y or Φ (or both) are stable, they can make up the dark matter. Moreover, if the annihilation of Y-$\overline{Y}$ and $\Phi\Phi^*$ is very efficient, the final DM density will be set by the asymmetry, related to the visible baryon asymmetry by Eq. (74).

References

[1] E. W. Kolb and M. S. Turner, Ann. Rev. Nucl. Part. Sci. **33**, 645 (1983). doi:10.1146/annurev.ns.33.120183.003241

[2] A. Riotto, "Theories of baryogenesis," hep-ph/9807454.

[3] W. Bernreuther, "CP violation and baryogenesis," Lect. Notes Phys. **591**, 237 (2002) [hep-ph/0205279].

[4] J. M. Cline, "Baryogenesis," hep-ph/0609145.

[5] M. E. Peskin and D. V. Schroeder, "An Introduction To Quantum Field Theory," *Reading, USA: Addison-Wesley (1995) 842 p.*

[6] W. Hu, "Lecture Notes on CMB Theory: From Nucleosynthesis to Recombination," arXiv:0802.3688 [astro-ph].

[7] V. A. Rubakov and A. D. Vlasov, "What do we learn from CMB observations," Phys. Atom. Nucl. **75**, 1123 (2012) doi:10.1134/S1063778812090116 [arXiv:1008.1704 [astro-ph.CO]].

[8] K. Garrett and G. Duda, "Dark Matter: A Primer," Adv. Astron. **2011**, 968283 (2011) doi:10.1155/2011/968283 [arXiv:1006.2483 [hep-ph]].

[9] W. Hu, "Intermediate Guide to the Acoustic Peaks and Polarization," http://background.uchicago.edu/ whu/intermediate/driving2.html, 2001.

[10] P. A. R. Ade *et al.* [Planck Collaboration], Astron. Astrophys. **594**, A13 (2016) doi:10.1051/0004-6361/201525830 [arXiv:1502.01589 [astro-ph.CO]].

[11] C. Patrignani *et al.* [Particle Data Group], "Review of Particle Physics," Chin. Phys. C **40**, no. 10, 100001 (2016);
http://pdg.lbl.gov/

[12] A. Coc, J. P. Uzan and E. Vangioni, "Standard big bang nucleosynthesis and primordial CNO Abundances after Planck," JCAP **1410**, 050 (2014) doi:10.1088/1475-7516/2014/10/050 [arXiv:1403.6694 [astro-ph.CO]].

[13] M. Pospelov and J. Pradler, "Big Bang Nucleosynthesis as a Probe of New Physics," Ann. Rev. Nucl. Part. Sci. **60**, 539 (2010) doi:10.1146/annurev.nucl.012809.104521 [arXiv:1011.1054 [hep-ph]].

[14] E. W. Kolb and M. S. Turner, "The Early Universe," Front. Phys. **69**, 1 (1990).

[15] G. Steigman, "Observational tests of antimatter cosmologies," Ann. Rev. Astron. Astrophys. **14**, 339 (1976). doi:10.1146/annurev.aa.14.090176.002011

[16] L. Canetti, M. Drewes and M. Shaposhnikov, "Matter and Antimatter in the Universe," New J. Phys. **14**, 095012 (2012) doi:10.1088/1367-2630/14/9/095012 [arXiv:1204.4186 [hep-ph]].

[17] A. G. Cohen, A. De Rujula and S. L. Glashow, "A Matter - antimatter universe?," Astrophys. J. **495**, 539 (1998) doi:10.1086/305328 [astro-ph/9707087].

[18] A. D. Sakharov, "Violation of CP Invariance, c Asymmetry, and Baryon Asymmetry of the Universe," Pisma Zh. Eksp. Teor. Fiz. **5**, 32 (1967) [JETP Lett. **5**, 24 (1967)] [Sov. Phys. Usp. **34**, 392 (1991)] [Usp. Fiz. Nauk **161**, 61 (1991)]. doi:10.1070/PU1991v034n05ABEH002497

[19] K. Abe *et al.* [Super-Kamiokande Collaboration], "Search for proton decay via $p \to e^+\pi^0$ and $p \to \mu^+\pi^0$ in 0.31megatonyears exposure of the Super-Kamiokande water Cherenkov detector," Phys. Rev. D **95**, no. 1, 012004 (2017) doi:10.1103/PhysRevD.95.012004 [arXiv:1610.03597 [hep-ex]].

[20] P. Langacker, "Grand Unified Theories and Proton Decay," Phys. Rept. **72**, 185 (1981). doi:10.1016/0370-1573(81)90059-4

[21] L. H. Ryder, "Quantum Field Theory," *Cambridge, UK: Cambridge University Press (1996) 508 p.*

[22] M. Srednicki, "Quantum field theory," *Cambridge, UK: Cambridge University Press (2007) 660 p.*

[23] C. Itzykson and J. B. Zuber, "Quantum Field Theory," *New York, Usa: Mcgraw-hill (1980) 705 p.(International Series In Pure and Applied Physics)*

[24] G. 't Hooft, "Computation of the Quantum Effects Due to a Four-Dimensional Pseudoparticle," Phys. Rev. D **14**, 3432 (1976) Erratum: [Phys. Rev. D **18**, 2199 (1978)]. doi:10.1103/PhysRevD.18.2199.3, 10.1103/PhysRevD.14.3432

[25] F. R. Klinkhamer and N. S. Manton, "A Saddle Point Solution in the Weinberg-Salam Theory," Phys. Rev. D **30**, 2212 (1984). doi:10.1103/PhysRevD.30.2212

[26] J. Kunz, B. Kleihaus and Y. Brihaye, "Sphalerons at finite mixing angle," Phys. Rev. D **46**, 3587 (1992). doi:10.1103/PhysRevD.46.3587

[27] P. B. Arnold and L. D. McLerran, "Sphalerons, Small Fluctuations and Baryon Number Violation in Electroweak Theory," Phys. Rev. D **36**, 581 (1987). doi:10.1103/PhysRevD.36.581

[28] S. Y. Khlebnikov and M. E. Shaposhnikov, "The Statistical Theory of Anomalous Fermion Number Nonconservation," Nucl. Phys. B **308**, 885 (1988). doi:10.1016/0550-3213(88)90133-2

[29] G. D. Moore, "Measuring the broken phase sphaleron rate nonperturbatively," Phys. Rev. D **59**, 014503 (1999) doi:10.1103/PhysRevD.59.014503 [hep-ph/9805264].

[30] M. D'Onofrio, K. Rummukainen and A. Tranberg, "Sphaleron Rate in the Minimal Standard Model," Phys. Rev. Lett. **113**, no. 14, 141602 (2014) doi:10.1103/PhysRevLett.113.141602 [arXiv:1404.3565 [hep-ph]].

[31] R. D. Peccei, "Discrete and global symmetries in particle physics," Lect. Notes Phys. **521**, 1 (1999) doi:10.1007/BFb010552, 10.1007/BFb0105521 [hep-ph/9807516].

[32] S. Willenbrock, "Symmetries of the standard model," hep-ph/0410370.

[33] D. Tucker-Smith and N. Weiner, "Inelastic dark matter," Phys. Rev. D **64**, 043502 (2001) doi:10.1103/PhysRevD.64.043502 [hep-ph/0101138].

[34] H. E. Haber and G. L. Kane, "The Search for Supersymmetry: Probing Physics Beyond the Standard Model," Phys. Rept. **117**, 75 (1985). doi:10.1016/0370-1573(85)90051-1

[35] Z. Ligeti, "TASI Lectures on Flavor Physics," doi:10.1142/9789814678766-0006 arXiv:1502.01372 [hep-ph].

[36] M. Dine, "TASI lectures on the strong CP problem," hep-ph/0011376.

[37] R. F. Streater and A. S. Wightman, "PCT, spin and statistics, and all that," Princeton, USA: Princeton Univ. Pr. (2000) 207 p.

[38] V. A. Kostelecky and M. Mewes, "Astrophysical Tests of Lorentz and CPT Violation with Photons," Astrophys. J. **689**, L1 (2008) doi:10.1086/595815 [arXiv:0809.2846 [astro-ph]].

[39] D. Baumann, "Cosmology," http://cosmology.amsterdam/education/cosmology/

[40] H. Georgi and S. L. Glashow, "Unity of All Elementary Particle Forces," Phys. Rev. Lett. **32**, 438 (1974). doi:10.1103/PhysRevLett.32.438

[41] S. Weinberg, "Baryon and Lepton Nonconserving Processes," Phys. Rev. Lett. **43**, 1566 (1979). doi:10.1103/PhysRevLett.43.1566

[42] E. W. Kolb and S. Wolfram, "Baryon Number Generation in the Early Universe," Nucl. Phys. B **172**, 224 (1980) Erratum: [Nucl. Phys. B **195**, 542 (1982)]. doi:10.1016/0550-3213(80)90167-4, 10.1016/0550-3213(82)90012-8

[43] M. Bolz, A. Brandenburg and W. Buchmuller, "Thermal production of gravitinos," Nucl. Phys. B **606**, 518 (2001) Erratum: [Nucl. Phys. B **790**, 336 (2008)] doi:10.1016/S0550-3213(01)00132-8, 10.1016/j.nuclphysb.2007.09.020 [hep-ph/0012052].

[44] I. Affleck and M. Dine, "A New Mechanism for Baryogenesis," Nucl. Phys. B **249**, 361 (1985). doi:10.1016/0550-3213(85)90021-5

[45] M. Dine and A. Kusenko, "The Origin of the matter - antimatter asymmetry," Rev. Mod. Phys. **76**, 1 (2003) doi:10.1103/RevModPhys.76.1 [hep-ph/0303065].

[46] M. Dine, L. Randall and S. D. Thomas, Phys. Rev. Lett. **75**, 398 (1995) doi:10.1103/PhysRevLett.75.398 [hep-ph/9503303].

[47] M. Dine, L. Randall and S. D. Thomas, "Baryogenesis from flat directions of the supersymmetric standard model," Nucl. Phys. B **458**, 291 (1996) doi:10.1016/0550-3213(95)00538-2 [hep-ph/9507453].

[48] D. E. Kaplan, M. A. Luty and K. M. Zurek, "Asymmetric Dark Matter," Phys. Rev. D **79**, 115016 (2009) doi:10.1103/PhysRevD.79.115016 [arXiv:0901.4117 [hep-ph]].

[49] K. M. Zurek, "Asymmetric Dark Matter: Theories, Signatures, and Constraints," Phys. Rept. **537**, 91 (2014) doi:10.1016/j.physrep.2013.12.001 [arXiv:1308.0338 [hep-ph]].

[50] H. Davoudiasl, D. E. Morrissey, K. Sigurdson and S. Tulin, "Hylogenesis: A Unified Origin for Baryonic Visible Matter and Antibaryonic Dark Matter," Phys. Rev. Lett. **105**, 211304 (2010) doi:10.1103/PhysRevLett.105.211304 [arXiv:1008.2399 [hep-ph]].

[51] H. Davoudiasl, D. E. Morrissey, K. Sigurdson and S. Tulin, "Baryon Destruction by Asymmetric Dark Matter," Phys. Rev. D **84**, 096008 (2011) doi:10.1103/PhysRevD.84.096008 [arXiv:1106.4320 [hep-ph]].