LEPTOGENESIS: A LINK BETWEEN
THE MATTER-ANTIMATTER ASYMMETRY AND
NEUTRINO PHYSICS

J. ORLOFF

Laboratoire de Physique Corpusculaire, Université Blaise Pascal, 24 rue des Landais, F-63177 Aubière

We review the experimental evidence for a net baryon density in cosmology, and the theoretical mechanism for producing it, called leptogenesis, which relies on the creation of a lepton asymmetry at an intermediate step. The naturality of this mechanism and its possible relations with neutrino oscillations are outlined.

1 Facts and Fancy about the Matter Asymmetry

Matter is so tightly connected to our everyday experience, that the fascinating prediction of anti-matter in Dirac’s theory first raised skepticism, which turned into solid confidence after discovery of the positron. Clearly, anti-matter on earth exists only briefly after the high-energy collisions we use to study fundamental interactions. At first sight, this fugacity seems a natural explanation for the absence of anti-matter around us. In fact, it is rather a consequence of the domination of matter which we tend to take for granted by a kind of anthropic argument. Quantitatively explaining this domination (or asymmetry) of protons is the purpose of baryogenesis. Let us first see how to quantify this baryon asymmetry.

Even when turning off accelerators on earth, we can collect some $10^{-4}$ anti-proton for every proton in the cosmic rays that penetrate the upper atmosphere after a $10^8$ years erratic journey since their production by supernova explosions in our galactic disc. This however does not constitute an evidence for a $10^{-4}$ anti-supernova fraction. Indeed, the interstellar dust is dense enough to play the role of fixed target intercepting a small fraction of the primary cosmic proton flux and producing the observed secondary anti-protons. This fraction can be cross-checked against the amount of gamma rays produced by the same collisions. The observed anti-proton fraction $\bar{p}/p \approx 10^{-4}$ is thus only an upper bound on the natural anti-matter fraction. Incidentally, a tighter bound of this type can be obtained by considering heavier nuclei like
deuterium, for which \( \frac{^2D}{p} \approx 1/60 \), while we expect \( \frac{^3D}{p} \approx 10^{-8} \). We can thus assume that in our galaxy, like on earth, there is a total domination of matter over anti-matter, which disappeared by annihilating with neighboring matter.

To quantify the chances of having anything surviving these annihilations, we need a dimensionless number summarizing the baryon asymmetry of the universe. One could take the ratio of the net baryon number density \( n_B \equiv (n_p - n_\bar{p}) \) divided by \( (n_p + n_\bar{p}) \) as is often done in particle physics. However, it is more useful to divide by the entropy density, and define \( Y_B = \frac{n_B}{s} \) which is an invariant if 1) baryon number is conserved and 2) the expansion of the universe is slow and adiabatic enough to avoid irreversible entropy creation in a comoving volume. The baryon to photon ratio \( \eta_B = \frac{n_B}{s} \approx 7Y_B \) is only equivalent at low temperatures\(^4\) \( T < m_e \) when photons and neutrinos dominate the entropy density, electrons and all other particles being non-relativistic.

It is difficult to evaluate \( \eta_B \) by direct observation, because baryons tend to clump together gravitationally, while photons don’t. To get indirect handles on \( \eta_B \), we thus need to identify physical processes sensitive to the baryon density \( n_B \) averaged on cosmological distances, or occurring before the formation of structures like galaxies. A classical example is Big Bang Nucleosynthesis, \( i.e. \) the making of light nuclei in the primeval plasma before star formation. Increasing the baryon density reduces the entropy price to pay for keeping nucleons together in

\(^4\)Explicit values of \( \eta_B \) found in the literature always refer to this present day limit, even when extracted from physics at much earlier times where the relation to the constant \( Y_B \) is different.
a nucleus (instead of letting them fill space) and thus increases the $^4\text{He}/p$ ratio as apparent in the left figure. However, the primordial $^4\text{He}$ abundance is not very sensitive, and is further easily contaminated by Helium later formed in stars. More interesting is the Deuterium, whose much lower binding energy per nucleon allows to understand both its very low production in stars, and its decreasing primordial abundance with increasing $\eta_B$. The ratio of these primordial abundances to hydrogen are extracted from interstellar clouds absorption lines into the emission of $z = 0.1 \to 3.5$ distant quasars. There is some tension between the Helium and Deuterium preferred values, the latter being more trusted and giving baryon to photon ratios around

$$\eta_B \equiv \frac{n_B}{n_\gamma} \approx 5.6 \times 10^{-10} \cong \eta_{10} \times 10^{-10}.$$  

Another cosmological observable that feels the baryon density is the Cosmic Microwave Background radiation, whose temperature fluctuations reflect the way preexisting density perturbations start oscillating when they enter the horizon at the recombination temperature $T_{\text{rec}} \approx 0.1\text{eV}$ where atoms get formed. The RMS spherical harmonics coefficients $C_l$ of the temperature distribution display so called “acoustic peaks” at peculiar values of the inverse angular scale $l$ (see the WMAP report at this conference). Since $m_{p+} \gg m_{e^-}$, the effect of protons self-gravity is not neutral, and tends to enhance the compression peaks ($1^{st}$, $3^{rd}$, ...) and suppress the expansion peaks ($2^{nd}$, ...). Increasing the baryon (and electron) density also decreases the sound velocity in the plasma, which separates the peaks from each other. These effects are illustrated in the upper right figure where all parameters but the baryonic density $\eta_{10} = 274 \Omega_B h^2$ are fixed at some typical values.

Another handle on $\eta_B$ is the baryon fraction deduced from the X-ray emission of large clusters. Analyses of the temperature profiles cannot be understood in terms of baryons alone, and have long been an argument for dark matter on scales larger than the galactic flat rotation curves. On such large scales, it can be argued that the gravitational accretion of baryons and dark matter are similar. Knowing then the cosmological baryon to dark matter fraction, the baryon density is accessible through an estimate of the dark matter density, for instance by combining the CMB and Supernova Ia data.

These three determinations are combined in the lower-right figure. The concordance which emerges is a fascinating confirmation of the adiabatic invariance of $Y_B$: all the way from the nucleosynthesis temperature where we find $\eta_{10}(T_{\text{BBN}} \approx 1\text{MeV}) = 5.6 \pm 0.5$ from Deuterium abundance alone, to essentially today where the X clusters baryon fraction give $\eta_{10}(T_{\text{SN, Ia}} \approx 1\text{MeV}) = 5.1 \pm 1.6$, passing through the recombination where the CMB fluctuations give $\eta_{10}(T_{\text{rec}} \approx 0.1\text{eV}) = 6.0 \pm 0.6$, we find reasonable agreement over 9 orders of magnitude in temperature or comoving volume size. We also show the latest WMAP results reported at this conference $\eta_{10} = 6.1^{+3}_{-2}$ as a hatched band. We should however mention that the quoted errors are given for all other parameters fixed at their global $\chi^2$-minimizing values; minimizing $\chi^2$ for each $\eta_{10}$ would result in wider error bars because of correlations, in particular with $\Omega_{DM}$.

After reviewing some quantitative facts about the baryon asymmetry, let us see why we fancy and care for a dynamical mechanism accounting for this asymmetry. We could of course do without one by simply taking $Y_B \approx 0.8 \times 10^{-10}$ as an initial condition, or rather, a final one. But in this difference lies the whole problem. Indeed, even if we have no direct cosmological signal dating from before nucleosynthesis, nothing prevents us from extrapolating the history of the universe back to temperatures above $200\text{MeV}$, where a relativistic quark-gluon plasma is expected. Now the recipe for setting up the required initial condition in this plasma is the following: count $10 \, 000 \, 000 \, 000$ antiquarks; add to these $10 \, 000 \, 000 \, 014$ quarks, and start over until the 14 extra quarks pile up to the desired chunk of the universe you want to build (e.g. $10^{70}$ times to build our single galaxy). Clearly, even the first step requires a fine tool to achieve the needed $10^{-9}$ relative precision: this is a typical fine tuning problem, equivalent to the fact that every quark today is the lucky winner of a billion to one fatal lottery.
Since the far observations mentioned above are in fact only sensitive to the absolute value \(|\eta_B|\), a way to ease this initial fine tuning problem would be to imagine a baryon separating process that would leave patches having slightly more baryons, and others having slightly more anti-baryons, making up a globally symmetric universe on the whole. However, such a baryon separating process would have to operate before the epoch where baryons start annihilating, namely \(T_{\text{sep}} > 20\text{MeV}\), which means a small causal horizon \(H^{-1}(T_{\text{sep}}) < H^{-1}(20\text{MeV})\) and thus a small baryonic number inside it: \(B_{\text{causal}} < Y_B s H^{-3}|_{20\text{MeV}} \approx 10^{-10}(M_{\text{Pl}}/20\text{MeV})^3 \approx 10^{52} \approx M_{\text{earth}}/m_p\). This maximal “matter island” size is way too small, given that NASA missions to other planets survived, or given the arguments about galactic cosmic rays presented above. Actually, a lower bound on the size of our matter island in a globally symmetric universe can be obtained from the \(\approx 70\text{GeV} \gamma\) rays produced by annihilations at the island boundaries. This matter island minimal size turns out to be of the order of the present causal horizon or visible universe.

This explains why we need baryogenesis, a global mechanism operating everywhere before \(T = 20\text{MeV}\), which can dynamically transform a symmetric initial condition \(Y_B = 0\) into the observed \(Y_B \approx 0.8 \times 10^{-10} \neq 0\), and thus explain “why there is something rather than nothing” after \(p - \bar{p}\) annihilation. The necessary conditions[11] that any such mechanism has to satisfy, have been spelled out as early as 1967 by A. Sakharov.

**S.C.1: Departure from equilibrium** is necessary since at thermal equilibrium,

\[
n_p = \int d^3k (e^{-\sqrt{k^2 + m_p^2}/T + 1})^{-1} = n_p
\]

because of the \(CPT\) relation \(m_p = m_\bar{p}\). Some breakdown of chemical equilibrium is also necessary since otherwise, microreversibility requires the rates for any process and its inverse to be equal.

**S.C.2: C and CP violations** are also necessary, as baryon number is odd under these symmetries. The simplest way to see it is above \(T_{\text{QCD}} \approx 200\text{ MeV}\) where baryon number is carried by quarks:

\[
n_B = \frac{1}{3} (n_{q_L} - n_{\bar{q}_L}) + n_{q_R} - n_{\bar{q}_R}) \Rightarrow \begin{cases} \text{CP} : & q_L \leftrightarrow \bar{q}_L; & B \leftrightarrow -B \\ \text{C} : & q_L \leftrightarrow q_R; & B \leftrightarrow -B \end{cases}
\]

\(C\) violation, like parity, is maximal in the Standard Model and is thus no problem. \(CP\) violation is however also essential, and is responsible for deciding whether to make baryons or anti-baryons out of the CP symmetric initial condition \(Y_B = 0\).

**S.C.3: B violation** is obviously needed to change the baryon number in a comoving volume.

Reviewing these conditions makes it obvious that baryogenesis needs some particle physics inputs from the micro-world, unlike gamma ray bursts, supernova explosions, ultra-high energy cosmic rays or other astroparticle physics puzzles which might ultimately find macroscopic resolutions.

In particular, the 3rd baryon violation condition makes a strong appeal to particle physics beyond the standard model. After the conception of Grand Unified Theories, baryon number violation had a natural niche at energies \(\approx 10^{15}\text{GeV}\), which was then the natural scale for baryogenesis. However, it was soon realized[13] that, even in the Standard Model, an \(SU(2)_L\)

---

[11] It is worth noting that at that time, \(B\) conservation had not yet been ruined by Grand Unification concepts, and that Sakharov immediately sought connections with the proton lifetime, and with the recently discovered \(CP\) violation in \(K_0 - \bar{K}_0\) mixing.
generalization of the triangle anomaly responsible for the process $\pi_0 \to \gamma \gamma$ could violate the conservation of any current of the form:

$$\partial_\mu J_L^\mu = \sum_{i \in \text{doublets}} \partial_\mu [\bar{\psi}_L e_i \gamma^\mu \psi_L] = (\sum_i e_i) \frac{g_W^2}{16\pi^2} F_{\mu\nu}^W F^{\mu\nu}_\nu e_i \gamma^\mu \frac{1-\gamma_5}{2} - \frac{i}{g_W} W_\rho \gamma^\mu \frac{1-\gamma_5}{2}$$

where $e_i$ is any charge assigned to the doublet $i$. In the presence of instantons, i.e. non-zero $W$-field solutions tunneling between different topological vacua and having an integer instanton number $N = \frac{g^3}{2\pi^2} \int d^4xF\tilde{F}$, the charge $Q_L$ associated with this current changes by $\Delta Q_L \approx \Delta [\int d^4x J_L^0] = 2N \sum_i e_i$. So for instance, if we consider the left baryonic charge $Q_L = B_L$ (which corresponds to the choice $e \equiv 0$ except $e_{u_L} = e_{d_L} = \frac{1}{3}$) we find a violation proportional to the number of generations $\Delta B_L = n_{gen}N$: $\beta$ does exist in the SM! Notice that at the same time, the left lepton number $Q_L = L_L$ ($e \equiv 0$ except $\nu_{\mu_L} = \nu_{e_L} = 1$) changes by the same amount $\Delta L_L = n_{gen}N = \Delta B_L$, so that $B - L$ is left intact (as should be the case for a gaugeable symmetry). The rate of these tunneling anomalous processes $\Gamma_{\text{tunnel}} \propto e^{-c N/g_W^2}$ is low enough to preserve the proton stability. At high energies or temperatures \cite{14}, instead of tunneling under the potential barrier that must be present in the complicated field space, there exists a possibility to classically roll over the barrier. In the broken electro-weak phase when $v = \langle H \rangle \neq 0$, this classical rate is governed by a saddle point solution called the sphaleron, $\Gamma_{\text{class.over}}(T) \propto e^{-E_{\text{sphaleron}}/T} \approx e^{-10M_W/T}$ which is again very low. In the unbroken phase $v = 0$ however, it was found \cite{15} to be much faster $\Gamma_{\text{class.over}}(T) \propto \alpha_W^5 T^4$. This important result meant that once above the ElectroWeak Phase Transition $T_{\text{EWPT}} \approx 100\text{GeV}$, S.C.3’s need for baryon number violation beyond the SM can be evaded. The (nearly) known physics at the EWPT being the last chance for an operative baryogenesis, this opened the way to a minimal bottom-up strategy: starting from the well-established standard model, and adding as little ingredients as necessary to obtain a successful baryogenesis.

In the SM, S.C.2 is in principle satisfied by the nearly maximal CKM phase $\delta_{\text{CKM}}$, but only manifest itself in delicate quantum interferences. The required coherence between quarks is likely to be destroyed in a hot plasma undergoing strong interaction collisions, which can GIM suppress the final asymmetry by as much as 15 orders of magnitude \cite{16}. In practice, CKM $\overline{\text{CP}}$ is insufficient. A second failure of SM baryogenesis concerns S.C.1. Indeed, the expansion rate of the universe around $T_{\text{EWPT}} \approx 100\text{GeV}$ is too slow for any relevant SM process to get out of equilibrium. A first order phase transition, where bubbles of the broken EW symmetry phase expand in the middle of an unbroken phase plasma, can efficiently amplify non-equilibrium effects around the critical temperature. However, in the pure SM, this requires a scalar field mass much lower \cite{17} than the present experimental lower bound $m_h > 114\text{GeV}$. The baryon asymmetry $Y_B \approx 10^{-10}$, which seemed too small from an initial conditions point of view, now seems too large for SM baryogenesis.

In the Minimal Supersymmetric extension of the SM, which doesn’t need baryogenesis arguments to be worth considering, both of these problems can in principle find answers. Indeed, the plethoraous scalar fields that come with Supersymmetry can modify the effective potential of the theory and reinforce the strength of the EW phase transition (S.C.1), especially \cite{18} for light $t_R$. Moreover, the MSSM offers new $\overline{\text{CP}}$ violating phases (S.C.2) less prone to GIM suppression than the CKM one. Nevertheless, the naturalness of these ideas has suffered from the rise of the experimental lower bound on $m_h$, to a point that many consider now extremely contrived.

2 Thermal Leptogenesis

At the same time, the fancy for neutrino oscillations and masses has solidified into a more and more established fact, and the inclusion of neutrino masses in the SM changes the shape of
our bottom up program. The reason neutrino masses may have anything to do with baryogenesis have been recognized by Fukugita and Yanagida immediately after the discovery of unsuppressed anomalous processes in the unbroken EW phase. Indeed, just as elastic scatterings (which change particle positions but not numbers) will tend to uniformize the particle density in a box, anomalous processes (which we saw change $B_L$ and $L_L$ but not $B_L - L_L$) will tend to redistribute an asymmetry carried solely by leptons ($L_L = -1, B_L = 0$) into a more evenly shared asymmetry ($L_L = -2/3, B_L = 1/3$). Since the lepton number is today mostly carried by furtive $T \approx 2^4K$ neutrinos, the only observable effect of this redistribution is the generation of a baryon asymmetry. We thus see that the problem of producing a correct baryon asymmetry is solved if we find a mechanism to produce a lepton asymmetry $Y_{L_{LL}} \approx -3 \times 10^{-10}$ anytime before $T_{EWPT}$, i.e. if we find a leptogenesis mechanism. Translating Sakharov Conditions to leptogenesis, it is clear that neutrino masses (and especially Majorana ones) offer a new way to satisfy S.C.3. Before detailing how the others conditions can be met, let us explain why we feel leptogenesis can fit in the minimal bottom up approach we outlined.

As reviewed at length at this conference, neutrino oscillations point to non vanishing neutrino masses. Barring the introduction of a scalar $SU(2)$ triplet which alters the $m_Z/m_W$ ratio, Lorentz invariance then requires the introduction of $SU(2)$ singlet right-handed neutrino fields $N$ which carry no gauge charge, and can thus enjoy Majorana masses a priori disconnected from the EW scale. The most general mass terms for leptons then read:

$$L_{mass} = \bar{\ell}H_v \cdot \frac{1}{\theta} \text{diag}(m_{e,\mu,\tau})\ell_R + \frac{1}{2} \bar{N}^c \cdot \text{diag}(M_{1,2,3})N + \bar{\ell}H_v \cdot \frac{1}{\theta} \text{diag}(m_{R,1,2,3})U_R N \equiv Y_{\ell}$$

where we have decomposed the 18 parameters of the complex Yukawa couplings $Y_{\ell}$ between neutrino flavor $l = e, \mu, \tau$ and Majorana mass eigenstate $N_{i=1,2,3}$ into:

- 3 Dirac mass eigenvalues $m_{1,2,3}$;
- one CKM-like matrix of $SU(3)/U(1)$, $V_{CKM}(\theta_L, \delta_L) = V_{23}(\theta_{L23})V_{13}(\theta_{L13}, \delta_L)V_{12}(\theta_{L12})$, containing 3 angles and 1 phase$^d$;
- one $SU(3)$ matrix $U_R = \text{diag}(e^{i\phi_R}, e^{i\psi_R}, e^{i\psi_R})V(\theta_R, \delta_R).\text{diag}(e^{i\phi_R})$ containing 3 angles and 5 phases (both $\psi^R_i$ and $\phi^R_i$ separately adding up to 0);
- 3 phases $\text{diag}(e^{i\phi_{CKM}})$ which could multiply $V_{CKM}$ from the right and have been reabsorbed by a common rephasing of $L$ and $l_R$.

At this stage, the simplest and most natural way of accounting for the extreme smallness of neutrino masses is to leave Dirac masses $m_D$ around the EW scale where they belong, or at least close to other Dirac e.g. up-quarks masses, and use the fact that $M_i$’s are not a priori related with any scale to raise their value, which reduces the lightest neutrino masses by the see-saw mechanism$^{23,24,25}$. Indeed, for energies below $M_i$’s, the decoupled $N_i$ fields can be integrated out, leaving an effective mass Lagrangian for the fields $\bar{\ell} = (\bar{\ell}_L, \bar{\nu})$:

$$L_{mass} \approx \bar{\ell}_L \text{diag}(m_{e,\mu,\tau})\ell_R + \frac{1}{2} \bar{\nu}^c \mathcal{M} \nu; \quad \mathcal{M} = U_{MNS} \text{diag}(m_{1,2,3})U^T_{MNS} \approx \nu^2 Y.M^{-1}.Y^T$$

$^c$Notice however that without Majorana masses, a clever use of the $L_L$ breaking by Dirac masses can also do$^{21,22}$ even if more contrived.

$^d$These parameters in the lepton sector are a priori unrelated to those in the quark sector, unless imposing a quark-lepton symmetry, natural for instance in $SO(10)$ GUTs.
where the light neutrinos mass matrix $\mathcal{M}$ (as promised inversely proportional to heavy Majorana one $M$) is a symmetric complex matrix which can again be decomposed into 3 eigenvalues, and a mixing matrix $U_{MNS} = V(\theta_{atm}, \theta_{Chooz}, \theta_{sun}, \delta) \text{diag}(e^{i\phi_L})$ containing 3 angles directly measured by oscillations, and only 3 phases (comparing with the decomposition of $U_R$, the difference is that the would-be $\psi_L$ phases can be reabsorbed into a common rephasing of $L$ and $L_R$, like $\phi_{CKM}$). It is worth noting that if we plug $Y$ from equ. \ref{eq:Y} into the see-saw formula \ref{eq:seesaw}, we see that the relation between light neutrino eigenmasses $m_{1,2,3}$ and heavy ones $M_{1,2,3}$ actually involves the product $U_{eff} = V_{CKM} U_{MNS}$ which is closely related to $U_R$ in the sense that if $U_{eff}$ is real or diagonal, so is $U_R$. At this moment, increasing the forbidden Majorana mass of the unavoidable right-handed neutrinos thus seems theoretically the most economical way to account for light neutrino masses. However, as advocated for a long time by Yanagida, these heavy Majorana neutrinos may play an interesting role in early cosmology and offer a natural framework for leptogenesis.

Indeed, S.C.3 is obviously met by the coexistence of decay modes with opposite lepton numbers, while S.C.2 is satisfied if their rates differ because of interferences between the tree and one loop amplitudes:

\[
\begin{align*}
N_i \rightarrow l^- H^+ & : N_i \xrightarrow{Y_{i\ell}} Y_{i\ell}^* \quad l^- \xrightarrow{Y_{ij}} Y_{ij}^* \quad \sum_{\nu,j} Y_{\nu i} Y_{\nu j}^* \quad \sum_{\nu,j} (Y_{\nu i} Y_{\nu j}^* + \nu \leftrightarrow j)Y_{ij}^* \\
N_i \rightarrow l^+ H^- & : Y_{ij}^* \quad \sum_{\nu,j} \Gamma(N_i \rightarrow l^+ H^-) \approx \frac{3}{16\pi} \text{Im}(A_{ij}^2) \frac{M_i}{M_j}
\end{align*}
\]

Each decaying $N_i$ thus generates a leptonic CP asymmetry

\[
\delta_i = \frac{\sum_{l} \Gamma(N_i \rightarrow l + H) - \Gamma(N_i \rightarrow \bar{l} + \bar{H})}{\sum_{l} \Gamma(N_i \rightarrow l + H) + \Gamma(N_i \rightarrow \bar{l} + \bar{H})} \approx \frac{3}{16\pi} \text{Im}(A_{ij}^2) \frac{M_i}{M_j}
\]

with $A_{ij} = (Y^\dagger Y)_{ij} = U_R^\dagger \text{diag}(m_{P,2,3}^2) U_R$ being the relevant combination of Yukawa couplings, whose diagonal terms contain the lifetimes $\Gamma_i \propto A_{ii} M_i$ while off-diagonal terms carry CP violation if $U_R$ (and thus $U_{eff} = V_{CKM} U_{MNS}$) is complex. Notice that for $M_i \approx M_j$, the asymmetry can be enhanced by the resonant self-energy contribution $\propto 1/(M_j - M_i)$ until $\Delta M \approx \Gamma$.

Turning finally to S.C.1, the decay asymmetry $\delta_i$ was here computed in vacuum, but in a hot plasma, the decay products can recombine and wash out the asymmetry if the (inverse) decay rate is much larger than the expansion rate $H$ (S.C.1 fails in local equilibrium) or equivalently if the dimensionless ratio

\[
K_i = \frac{\Gamma(N_i \rightarrow M_i)/H(T = M_i)}{\Gamma(T = M_i)} \approx \frac{1}{(1.66 \times 10^{-8} \text{GeV})} \times A_{ii} M_{pl}/M_i
\]

is much larger than 1. The lepton asymmetry $Y_i$ originating from species $N_i$ in thermal equilibrium is then diluted by a factor $d \propto 1/K$, giving $Y_i = \frac{1}{g^*} d(K_i, M_i) \delta_i$. This lepton asymmetry finally drives anomalous processes to produce a left baryon asymmetry $Y_B \approx -Y_i/3$. Assuming as above that $M_i$’s are hierarchical, the contribution from the lightest $M_1$ often dominates in which case the final result to be confronted with observations takes the simple form:

\[
Y_{B10} = 10^{10} Y_B \approx \frac{10^{10}}{16\pi g^*} \frac{d(K_1, M_1)}{M_{pl}} \frac{M_1}{A_{11}} \sum_{j=2,3} \text{Im}(A_{ij}^2) \frac{M_j}{M_j} \approx 0.8
\]

The dilution factor $d$ is a priori smaller than 1, but a more specific determination requires the numerical solution of the Boltzmann equations describing how various particles depart from
equilibrium in response to the universe expansion. The result depends mostly on the ratio $K_i$, which is often rewritten as a dimensionful seesaw-like mass $\tilde{m}_i = \frac{v^2 A_{ii}}{M_i} = K_i \tilde{m}^*$, where the critical value corresponding to $K = 1$ is $\tilde{m}^* = \sqrt{\frac{512 g^* \pi^3}{90}} \frac{v^2}{M_{pl}} = 1.08 \times 10^{-3}$eV in the SM.

The result is shown in figure 2 for a fixed CP asymmetry $\delta_1 = 10^{-6}$, so that $\eta_B = 10^{-8} d$. On the right, we see the large $K$ dilution effect from inverse decays explained above $d(K > 1, M_i < 10^{15}) \propto 1/K \approx 1/\tilde{m}$. For small $K$, the curves split in two: the lower one $d(K < 1) \propto K$ reflects the difficulty of creating an equilibrium population of right handed neutrino if their only couplings $Y_{li}$ are too small, and the upper curve assumes this population is initially present for some reason. In the first, more natural case, $d$ reaches a maximum of $d_{max} \approx 0.2$ around $K \approx 1$. The dependence on $M_i$ comes from $2 \rightarrow 2$ scattering effects, whose relative importance increase at large $M_i$ for fixed $m_\nu, \tilde{m}_1$.

It would be tempting to try and make a direct connection between the baryon asymmetry $Y_{B10}$ generated by this mechanism and the neutrino oscillations which make it so natural. It is however impossible if the Yukawa couplings $Y$’s are totally free. Indeed, for any set of values $(Y, M)$ transformed by the seesaw into acceptable light neutrino masses and oscillations, but for which leptogenesis produces a wrong asymmetry, say $Y_{B10} \ll 1$, a simple rescaling increasing both Yukawas $Y \rightarrow Y' = Y/\sqrt{Y_{B10}}$ and right handed masses $M \rightarrow M' = M/Y_{B10}$ in such a way as to preserve the see-saw and thus both light neutrinos and $\tilde{m}_i \propto K_i$, would restore a correct asymmetry $Y'_{B10} \approx 1$. Similarly, we saw that the CP phase felt by leptogenesis resides in $U_R$ which, through the see-saw, corresponds to a phase in $U_{eff} = U_{CKM}U_{MNS}$. Clearly, the result depends on the assumption for $V_{CKM}$ in the lepton sector, and short of this assumption, there is no relation possible with the phases in $U_{MNS}$ which affect oscillations or neutrinoless double beta decay.

An interesting upper bound can however be derived without assuming anything other than hierarchical right handed masses $M_i$. Using a parameterization of the Yukawas

$$Y^\dagger = v^{-1} \text{diag}(\sqrt{M}) R \text{diag}(\sqrt{m^B}) U^\dagger_{MNS}$$

with a complex orthogonal matrix $R$, one can write

$$\delta_1 = \frac{-3 M_i}{8\pi v^2} \frac{1}{A_{11}} \text{Im}(Y^\dagger M Y^*)_{11} = \frac{-3 M_i}{8\pi v^2} \sum_j m_j^2 \text{Im}(R_{1j}^2)$$

whose maximization over $R$ gives

$$|\delta_1| \leq \frac{3 M_i}{8\pi v^2} (m_3 - m_1)$$

since the dilution factor $d$ cannot exceed 0.2, the requested $Y_B$ translates into a lower bound on $M_i > \frac{0.06eV}{m_3 - m_1} \times 10^9$GeV. This is turn puts a lower bound on the reheating temperature after inflation $T_{reh} > 10^8$--$10^{10}$GeV. In SUSY models, the overproduction of unstable gravitinos potentially

\*\*In the approximation where we neglect the $M_i$dependence of $d$ above; otherwise, the rescaling factor might take a more complicated form.
destroying nucleosynthesis require $T_{\text{reh,SUSY}} < 10^{9-12}\text{GeV}$, which seems uncomfortably close to the previous number. A possible loophole is to relax the hierarchy requirement under which the bound (3) was obtained, and let $M_2$ get close enough to $M_1$ to benefit from the self-energy diagram enhancement.\[33\]

To illustrate the power of this bound (3), let us consider a simple $SO(10)$-inspired example\[34\] where we extend $b-\tau$ unification and fix Yukawas $Y$ by $m_u^D = m_u/3$. By the see-saw, $M_i$ and $U_R$ are then determined from the light neutrino mass matrix $M$, up to the lightest neutrino mass $m_1$, the MNS element $|U_{e3}|$ and 5 phases which are taken as free parameters. Right-handed masses turn out hierarchical (like the assumed Dirac masses), and $M_1$ can hardly exceed $10^8\text{GeV}$. The maximal asymmetry is nearly 2 orders of magnitude too small, as should follow from (3). Yet this maximal result is obtained for a special value of $|U_{e3}| \approx 0.16$ which approximately cancels CKM Cabbibo mixing and allows for the largest $M_1$ (as seen from the extreme case $U_{\text{eff}} = 1$ on the left figure). All results on figure are for the now disfavored “vacuum” solar neutrino oscillations ($\Delta m_{\odot}^2 = 4.6 \times 10^{-10}\text{eV}^2$), the LMA solution giving similar but yet smaller $M_1$ and asymmetries.\[34\]

## 3 Conclusion

In this short review, we hope to have convinced that the baryon asymmetry of the universe is rather well established and constitutes an important particle physics question worth the considerable amount of work it has attracted. Despite being a single number out of a non-repeatable experiment, it can be seen as one of the only evidences for an incompleteness of the Standard Model. As to the direction this incompleteness points at, thermal leptogenesis relying on the heavy right-handed neutrinos required by the see-saw mechanism, constitutes in our mind a very suggestive minimal, predictive and not excluded solution. Obviously, this opinion is shared by the large number of authors who recently contributed to this field, many of whom may justly feel misrepresented in this review by lack of space. It would be extremely nice to cross-check the existence of heavy right-handed neutrinos by some independent measurement. Unfortunately, apart from lepton flavor violating effects which might show up in certain cases discussed at this conference, this will not be easy. Meanwhile, thanks to the upper bound on CP violation in right-handed neutrino decays, imposing successful leptogenesis takes an non-trivial slice in the see-saw parameter space which usually cuts more than simply one dimension. This might shed useful light in the quest for some order in the present mass anarchy. Finally, the CP violation needed for leptogenesis cannot simply at present be related to the one measurable in future long baselines neutrino experiments, and even less to the one already measured in the
quark sector, unless some relations (like the $SO(10)$-inspired ones discussed above) are imposed: if we have all the phenomenology needed to account for CP violation, a real theory is still badly lacking.

1. M. S. Turner & J. N. Fry, *Baryogenesis without the initial presence of superheavy bosons*, Phys. Rev. **D24** (1981) 3341–3344.
2. D. Maurin *et al.*, *Galactic cosmic ray nuclei as a tool for astroparticle physics*, astro-ph/0212111.
3. WIZARD Coll., M. Boezio *et al.*, *The cosmic ray proton and helium spectra between 0.2-GeV and 200-GeV*, Astrophys. J. **518** (1999) 457–472.
4. M. Ambriola *et al.*, *High-energy deuteron measurement with the CAPRICE98 experiment*, Nucl. Phys. Proc. Suppl. **113** (2002) 88–94.
5. P. Chardonnet, J. Orloff & P. Salati, *The production of anti-matter in our galaxy*, Phys. Lett. **B409** (1997) 313–320, astro-ph/9705110.
6. R. P. Duperray, K. V. Protasov & A. Y. Voronin, *Antideuteron production in proton and proton nucleus collisions*, Eur. Phys. J. **A16** (2003) 27–34, nucl-th/0209078.
7. Particle Data Group Coll., K. Hagiwara *et al.*, *Review of particle physics*, Phys. Rev. **D66** (2002) 010001.
8. S. Burles, K. M. Nollett & M. S. Turner, *Big-Bang Nucleosynthesis Predictions for Precision Cosmology*, Astrophys. J. **552** (2001) L1–L6, astro-ph/0010171.
9. M. Kamionkowski & A. Kosowsky, *The cosmic microwave background and particle physics*, Ann. Rev. Nucl. Part. Sci. **49** (1999) 77–123, astro-ph/9904108.
10. G. Steigman, *The baryon density through the (cosmological) ages*, Fortsch. Phys. **50** (2002) 562–568, astro-ph/0202187.
11. A. G. Cohen, A. De Rujula & S. L. Glashow, *A matter-antimatter universe?*, Astrophys. J. **495** (1998) 539–549, astro-ph/9707087.
12. A. D. Sakharov, *Violation of CP invariance, C asymmetry, and baryon asymmetry of the universe*, Pisma Zh. Eksp. Teor. Fiz. **5** (1967) 32–35.
13. G. ’t Hooft, *Symmetry breaking through Bell-Jackiw anomalies*, Phys. Rev. Lett. **37** (1976) 8–11.
14. V. A. Rubakov & M. E. Shaposhnikov, *Electroweak baryon number non-conservation in the early universe and in high-energy collisions*, Usp. Fiz. Nauk **166** (1996) 493–537, hep-ph/9603208.
15. V. A. Kuzmin, V. A. Rubakov & M. E. Shaposhnikov, *On the anomalous electroweak baryon number nonconservation in the early universe*, Usp. Fiz. Nauk **166** (1996) 493–537, hep-ph/9603208.
16. M. B. Gavela, P. Hernandez, J. Orloff, O. Pene & C. Quimbay, *Standard model CP violation and baryon asymmetry. Part 2: Finite temperature*, Nucl. Phys. **B430** (1994) 382–426, hep-ph/9406289.
17. A. I. Bochkarev, S. V. Kuzmin & M. E. Shaposhnikov, *Electroweak baryogenesis and the Higgs boson mass problem*, Phys. Lett. **B244** (1990) 275.
18. M. Carena, M. Quiros & C. E. M. Wagner, *Electroweak baryogenesis and Higgs and stop searches at LEP and the Tevatron*, Nucl. Phys. **B524** (1998) 3, hep-ph/9710401.
19. J. M. Cline, *Electroweak phase transition and baryogenesis*, in *COSMO-01*, Rovaniemi, 2001, hep-ph/0201286.
20. M. Fukugita & T. Yanagida, *Baryogenesis without grand unification*, Phys. Lett. **B174** (1986) 45.
21. K. Dick, M. Lindner, M. Ratz & D. Wright, *Leptogenesis with Dirac neutrinos*, Phys. Rev. Lett. **84** (2000) 4039–4042, hep-ph/9907562.
22. H. Murayama & A. Pierce, *Realistic Dirac leptogenesis*, Phys. Rev. Lett. **89** (2002)
27. J. M. Frere, F. S. Ling, M. H. G. Tytgat & V. Van Elewyck, *Leptogenesis with virtual Majorana neutrinos*, Phys. Rev. D60 (1999) 016005, [hep-ph/9901337].

28. W. Buchmuller, P. Di Bari & M. Plumacher, *Cosmic microwave background, matter-antimatter asymmetry and neutrino masses*, Nucl. Phys. B643 (2002) 367–390, [hep-ph/0205349].

29. R. Barbieri, P. Creminelli, A. Strumia & N. Tetradis, *Baryogenesis through leptogenesis*, Nucl. Phys. B575 (2000) 61–77, [hep-ph/9911315].

30. K. Hamaguchi, H. Murayama & T. Yanagida, *Leptogenesis from sneutrino-dominated early universe*, Phys. Rev. D65 (2002) 043512, [hep-ph/0109030].

31. S. Davidson & A. Ibarra, *A lower bound on the right-handed neutrino mass from leptogenesis*, Phys. Lett. B535 (2002) 25–32, [hep-ph/0202239].

32. J. A. Casas & A. Ibarra, *Oscillating neutrinos and mu –¿ e, gamma*, Nucl. Phys. B618 (2001) 171–204, [hep-ph/0103065].

33. J. R. Ellis, M. Raidal & T. Yanagida, *Observable consequences of partially degenerate leptogenesis*, Phys. Lett. B546 (2002) 228–236, [hep-ph/0206300].

34. E. Nezri & J. Orloff, *Neutrino oscillations vs. leptogenesis in SO(10) models*, JHEP 04 (2003) 020, [hep-ph/0004227].