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Cloistered Baryogenesis

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Abstract. The cosmological matter-antimatter asymmetry can arise from the baryon number conserving CP asymmetry in two body decays of heavy particles, when the two final states carry equal and opposite baryon number, and one couples directly or indirectly to electroweak sphalerons so that its baryon asymmetry gets partly reprocessed into a lepton asymmetry, while the other remains chemically decoupled from the thermal bath with its baryon content frozen. After sphaleron switchoff the decay of the decoupled particles inject in the thermal plasma an unbalanced baryon asymmetry, giving rise to baryogenesis. We highlight the features of this mechanism in a type-I seesaw model extended by adding a new colored scalar coupled to the heavy Majorana neutrinos. If the colored scalar has an $\mathcal{O}$(TeV) mass, it would leave at the LHC a characteristic signature throughout all layers of the detectors.

Contents

1 Introduction

The Cosmic baryon asymmetry [1, 2] represents an indisputable evidence for physics beyond the standard model (SM), and suggests that in the early Universe new physical degrees of freedom must have been at work.

In the SM baryon (B) and lepton (L) number are violated only by the $B-L$ conserving electroweak (EW) sphalerons. In the early Universe the rates for these processes attain thermal equilibrium at $T \sim 10^{12}$ GeV, and remain in equilibrium until the EW phase transition at around $T \sim 10^2$ GeV. Any $B-L$ asymmetry generated for example by out-of-equilibrium, $B-L$ and CP violating interactions [3], and surviving within this temperature range, will then unavoidably result in a net B asymmetry. This mechanism is at the basis of the standard type-I seesaw [4] leptogenesis [5] as well as of its variants [6]. Among these variants two realizations are particularly intriguing. In the first one, the so-called purely flavored leptogenesis (PFL) [7, 8], the total CP asymmetry in lepton number in the decays of the heavy Majorana neutrinos vanishes exactly. Leptogenesis can still proceed thanks to non-vanishing CP asymmetries in the single lepton flavors, and thanks to the fact that washouts violate total lepton number acting differently along the different flavor directions [7, 9]. A non-vanishing total $B-L$ asymmetry can then result, provided all lepton-flavor-equilibrating processes remain out of equilibrium [10]. Another interesting variant is the so-called Dirac leptogenesis scenario, which can yield successful baryogenesis through leptogenesis even if L remains perturbatively conserved [11, 12]. In Dirac leptogenesis [11] heavy particle decays generate two equal in size and opposite in sign L asymmetries in the left-handed

(LH) lepton doublets and in light right-handed (RH) neutrinos. While lepton doublets participate in EW sphaleron reactions, the RH neutrino singlets do not. The L asymmetry stored in the LH leptons is then partially converted into a B asymmetry through sphalerons interactions. In contrast, as long as the RH neutrinos remain decoupled from the thermal bath, the corresponding L asymmetry remains unchanged. If decoupling holds until temperatures below EW sphaleron freezout then, although globally $B - L = 0$, a non-vanishing B asymmetry results.

Baryogenesis could also proceed via the out-of-equilibrium, C, CP and B violating decays of heavy particles, provided L is conserved in order to guarantee proton stability (see e.g. [13]). It is interesting to see if such a scenario also admits variants similar to the ones mentioned above for leptogenesis, and in particular to verify if a sufficiently low scale, accessible to direct tests, can be reached. In this paper we show that cloistered baryogenesis does represent a viable scenario, although it can only work at a temperature scale above $\sim 10^7$ GeV, thus remaining out of the reach of direct tests.

A baryogenesis scenario similar to PFL, that is a scenario in which the total B-violating CP asymmetry ϵ_B vanishes, is in general not viable. This is because baryon flavors, which get fully distinguished by their respective Yukawa interactions when the temperature drops below $T \sim 10^{11}$ GeV, quickly equilibrate because of intergeneration mixings, driving all baryon flavor asymmetries to zero. Strictly speaking there is in fact a narrow temperature window 10^{13} GeV $\gtrsim T \gtrsim 10^{11}$ GeV when only the Yukawa reactions for the third generation quarks are in thermal equilibrium. A third generation baryon flavor B_3 then does not necessarily equilibrate with the orthogonal flavor combination $B_{3\perp}$, so that in this window a purely flavored baryogenesis scenario is conceivable. However, this appears to us as a bit cumbersome, and we will not consider further this possibility.

The analogous of Dirac leptogenesis is instead a rather interesting possibility. That is, we conceive a baryogenesis scenario in which the decays of a heavy particle do not violate either L or B, but an asymmetry is still generated directly in baryon number. Essentially, two body B conserving decays generate equal in size and opposite in sign B asymmetries in two different sectors. The first one (the "active" sector) is coupled – directly or indirectly – to EW sphalerons. The second one remains (chemically) decoupled from the thermal bath at least until EW sphalerons switchoff, and we will refer to it as the uncommunicated or "cloistered" sector. The initial B asymmetry stored in the active sector gets partially converted into a L asymmetry, so that its net value changes, while the B asymmetry stored in the cloistered sector remains unaffected. After the EW phase transition heavy particles of the cloistered sector decay, injecting their (unbalanced) baryon asymmetry in the thermal bath, giving rise to baryogenesis. Because of the crucial role played by the uncommunicated sector, we will refer to this scenario as *cloistered baryogenesis*. This scenario is in fact similar in many aspects to the so-called WIMPy baryogenesis scenario [14–16] in which, however, the baryon asymmetry is generated from dark matter annihilation instead than from heavy particle decays.

Table 1 resumes, for the sake of illustration, the leptogenesis mechanisms that we have briefly discussed and the corresponding baryogenesis variants.

Lepton sector			Baryon sector		
$\Delta L \neq 0$	Leptogenesis	✔	$\Delta B \neq 0$	Direct baryogenesis	✔
$\epsilon_L = 0$	PFL	✔	$\epsilon_B = 0$	$T \lesssim 10^{11}$ GeV	✘
$\Delta L = 0$	Dirac leptogenesis	✔	$\Delta B = 0$	Cloistered baryogenesis	✔

Table 1. *Different mechanisms for baryogenesis. The left-hand side lists mechanisms in which the matter-antimatter asymmetry is seeded first in the lepton sector, and B is perturbatively conserved. The right-hand side lists the equivalent mechanisms in which the asymmetry is seeded first in the baryon sector. In the second row $\epsilon_L = 0$ and $\epsilon_B = 0$ refer respectively to vanishing total L and B violating CP asymmetries. The first two mechanisms in the baryon sector require perturbative L conservation to ensure proton stability. This is not required for cloistered baryogenesis in the third row. The check-mark indicate the viability of the corresponding scenario.*

In this paper we show that cloistered baryogenesis represents a viable scenario. We implement this mechanism in a simple extension of the type-I seesaw that was recently put forth in ref. [17]. In this setup, the heavy RH neutrinos N couple to the $SU(2)$ singlets up-type quarks u, and to a new colored scalar $\tilde{u}$ which, given that N is a gauge singlet, carries the same gauge quantum numbers than u. In general this scenario is not phenomenologically tenable because both B and L are violated and the nucleon is unstable. However, this can be solved by imposing exact baryon number conservation.

The rest of the paper is organized as follows. In section 2 we derive a lower bound on the scale of cloistered baryogenesis. In section 3 we describe the model and we discuss its phenomenological consistency. In section 4 we discuss baryogenesis within this setup, and derive the chemical equilibrium conditions and the Boltzmann equations for baryogenesis. In section 5 we highlight the role played in cloistered baryogenesis by hypercharge, and finally in section 6 we present our conclusions.

2 The temperature scale for cloistered baryogenesis

In this section we show that assuming a non degenerate RH neutrino spectrum, we can derive a lower bound on the temperature that allows for successful cloistered baryogenesis. This bound follows from the interconnections between the CP asymmetry and the requirement that the cloistered sector will remain uncommunicated with the active sector. While we will be interested in the case in which a Majorana RH neutrino decays in a SM u-type quark and in the complex conjugate of a new scalar $\tilde{u}$ of equal baryon charge, the argument can be presented in a more general form.

Let us consider a generic $U(1)_B$ invariant interaction between two self conjugate particles $X_i = X_i^c$ $(i = 1, 2)$ and other two fields Y and Z carrying opposite $U(1)_B$ charges

$$\mathcal{L} = \sum_{i=1,2} g_i \, X_i \, Y \, Z \; + \text{H.c.} \,, \tag{2.1}$$

with g_1 and g_2 two relatively complex parameters $\mathrm{Arg}(g_1^* g_2) \neq 0$. In general, the X_i can be Majorana fermions, with Y and Z a pair of complex scalar and fermion (as in standard leptogenesis) or alternatively X_i could be real scalars and Y, Z a pair of fermions or a pair of complex scalars (as in soft leptogenesis [18, 19]). In the first two cases g_i are dimensionless couplings, while in the last case they have mass dimension one. Let us now assume the mass ordering $M_{X_2} > M_{X_1} > M_Y + M_Z$ so that the decays $X_1 \to Y\,Z$, $\bar{Y}\,\bar{Z}$ can occur. In general this decay is CP violating, which implies a nonvanishing CP asymmetry in the number of Y and Z particles and antiparticles. This means that B asymmetries in the particle species Y and Z that are equal in size and opposite in sign are generated.

Let us now assume that Y has in-equilibrium chemical reactions with other particles in the thermal bath, while Z does not, and let us further assume that X_1 decays occur before EW sphaleron freezout. The B asymmetry carried by the Y's (ΔB_Y) gets distributed between all SM particles, and because of the partial conversion in a L asymmetry through sphaleron interactions, its overall value is changed $\Delta B_{SM} \neq \Delta B_Y$. In contrast, the B asymmetry carried by the Z's (ΔB_Z) will not change, so that eventually a net total asymmetry given by $\Delta B = \Delta B_{SM} + \Delta B_Z$ arises. After EW sphalerons freezout the Z's decay into SM particles (via B conserving decay modes) and this gives rise to baryogenesis.

The CP asymmetry between the number, say, of Y baryons and $\bar{Y}$ anti-baryons produced in X_1 decays is defined as

$$\epsilon_{X_1} = \frac{\gamma\left(X_1 \to Y Z\right) - \gamma\left(X_1 \to \bar{Y}\bar{Z}\right)}{\gamma^{\mathrm{tot}}} \,, \tag{2.2}$$

where the γ's are thermally averaged decay rates (γ^{tot} is the thermally averaged total decay width). ϵ_{X_1} can be computed from the interference between tree-level and one-loop vertex and wave-function diagrams. For the decays of Majorana fermions into a fermion/scalar pair we have, assuming $M_{X_1} \ll M_{X_2}$ [20]:

$$\epsilon_{X_1}^{(fs)} \simeq -\frac{|g_2|^2}{8\pi}\,\frac{M_{X_1}}{M_{X_2}}\sin\phi \,, \tag{2.3}$$

with $\phi = \mathrm{Arg}\left[(g_1^*\, g_2)^2\right]$. For the two other cases of scalar X particles decaying into fermion pairs or scalar pairs, we have respectively:

$$\epsilon_{X_1}^{(ff')} \simeq -\frac{|g_2|^2}{8\pi}\,\frac{M_{X_1}^2}{M_{X_2}^2}\sin\phi \,, \tag{2.4}$$

$$\epsilon_{X_1}^{(ss')} \simeq -\frac{1}{8\pi}\,\frac{|g_2|^2}{M_{X_2}^2}\sin\phi \,. \tag{2.5}$$

In order to maximize the CP asymmetries we set $\sin\phi \sim 1$. We see that in all three cases the asymmetries increase with the value of g_2. This coupling, however, cannot become arbitrarily large because X_2 mediated $Y Z \leftrightarrow \bar{Y}\bar{Z}$ scatterings would enforce equilibrium for the Y and Z chemical potentials $\mu_Y + \mu_Z = 0$ rendering cloistered baryogenesis ineffective. For the three cases at hand, the $2 \leftrightarrow 2$ scattering rates read:

$$\gamma^{(fs)}(YZ \leftrightarrow \bar{Y}\bar{Z}) \simeq \frac{1}{\pi^3}\frac{T^3}{M_{X_2}^2}|g_2|^4 \to \frac{64}{\pi} M_{X_1}\left(\epsilon_{X_1}^{(fs)}\right)^2 , \tag{2.6}$$

$$\gamma^{(ff')}(YZ \leftrightarrow \bar{Y}\bar{Z}) \simeq \frac{1}{\pi^3}\frac{T^5}{M_{X_2}^4}|g_2|^4 \to \frac{64}{\pi} M_{X_1}\left(\epsilon_{X_1}^{(ff')}\right)^2 , \tag{2.7}$$

$$\gamma^{(ss')}(YZ \leftrightarrow \bar{Y}\bar{Z}) \simeq \frac{1}{\pi^3}\frac{T}{M_{X_2}^4}|g_2|^4 \to \frac{64}{\pi} M_{X_1}\left(\epsilon_{X_1}^{(ss')}\right)^2 . \tag{2.8}$$

where the limiting expressions hold for $T \to M_{X_1}$. We see that in all three cases the equilibrating scattering rates are proportional to the square of the respective (maximum) CP asymmetries. Requiring that around $T \sim M_{X_1}$ these scatterings are out of equilibrium, that is $\gamma(YZ \leftrightarrow \bar{Y}\bar{Z}) \lesssim H(M_{X_1})$ where $H(M_{X_1}) \sim 17 M_{X_1}^2/M_{\rm Planck}$ parameterizes the Universe expansion rate, yields

$$M_{X_1} \gtrsim 10^{19} \times \epsilon_{X_1}^2 \ \mathrm{GeV}\,. \tag{2.9}$$

Given that values of the CP asymmetry smaller than $\epsilon_{X_1} \sim 10^{-6}$ could hardly explain the observed baryon asymmetry, eq. (2.9) implies $M_{X_1} \gtrsim 10^7$ GeV, which constitutes a necessary condition for successful cloistered baryogenesis. [1]

3 General considerations

In the type-I seesaw model, the SM fermion sector is extended by introducing heavy Majorana neutrinos. We assume three of them, and we denote by N the RH components $N = N_R$ while $N^c = N_L^c$ will denote the LH components. Besides a Majorana bilinear (mass) term $\bar{N}^c N$ one can also construct a set of new fermion bilinears by coupling the Majorana neutrinos with the SM fermions as $\bar{\chi}N$, where χ denotes any left-handed SM field: ℓ, Q, e^c, u^c or d^c (the SM RH fields are denoted as $\chi^c = \ell^c$, e, Q^c, d, u). The only bilinear that can be coupled in a gauge and Lorentz invariant way without introducing additional new fields is $\bar{\ell}N$ because it can be coupled to the Higgs field $\tilde{H} = i\sigma_2 H^*$ giving rise to a $SU(2) \times U(1)$ invariant. The seesaw Lagrangian, which contains precisely this term, reads:

$$-\mathcal{L}_{\rm Seesaw} = \overline{\boldsymbol{\ell}}\,\boldsymbol{\lambda}\boldsymbol{N}\tilde{H} + \frac{1}{2}\overline{\boldsymbol{N^c}}\,\boldsymbol{M_N}\boldsymbol{N} + \mathrm{H.c.}\ . \tag{3.1}$$

Henceforth we denote matrices and vectors in boldface, so e.g. $\boldsymbol{N}^T = (N_1, N_2, N_3)$ while $\boldsymbol{\lambda}$ and $\boldsymbol{M_N}$ are 3×3 matrices in flavor space and, without loss of generality, we assume that the seesaw Lagrangian eq. (3.1) is written in the basis in which $\boldsymbol{M}_N$ and the charged leptons Yukawa matrix are both diagonal with real and positive entries.

[1] Similar arguments have been used in [21] to derive numerically a bound on the mass of the lightest RH neutrino in the inert doublet model.

Following ref. [17], by introducing new scalar fields $\tilde{\chi}$ we can construct other invariants involving N and the remaining SM fermions Q, e^c, u^c or d^c. Clearly, since N is a gauge singlet, the gauge quantum numbers of $\tilde{\chi}$ must match the quantum numbers of the corresponding gauge non-singlet fermions. In general, once these new scalars are introduced, new operators beyond those involving the Majorana neutrinos can be constructed by coupling $\tilde{\chi}$ to SM fermions bilinears. The resulting new Lagrangian thus has the general form [17]:

$$-\mathcal{L}_{\tilde{\chi}} = \overline{\boldsymbol{\chi}}\,\boldsymbol{\eta}\,\boldsymbol{N}\,\tilde{\chi} + \sum_{\boldsymbol{\chi}^c\,\boldsymbol{\chi}'} \overline{\boldsymbol{\chi}^c}\,\boldsymbol{y}\,\boldsymbol{\chi}'\,\tilde{\chi} + \text{H.c.}\ , \tag{3.2}$$

where $\boldsymbol{\eta}$ and $\boldsymbol{y}$ are two 3×3 matrices of Yukawa couplings. In the first term it is left understood that different types of scalars have different couplings $\boldsymbol{\eta} = \boldsymbol{\eta}_{\tilde{\chi}}$, while in the second term there are different couplings also for different SM fermion bilinears $\boldsymbol{y} = \boldsymbol{y}_{\tilde{\chi}}^{\chi^c\chi'}$.

Among the various possibilities, those involving the new scalar fields $\tilde{\ell}$, $\tilde{e}$ or $\tilde{Q}$ (one at the time) allow for consistent baryon number assignments for which the $\boldsymbol{\eta}$ and $\boldsymbol{y}$ couplings conserve $U(1)_B$ [17]. In contrast, the inclusion of either $\tilde{u}$ or $\tilde{d}$ yields B and L breaking operators and thus, in their general form, these possibilities yield fast nucleon decay. Nevertheless, as we will discuss below, it is still possible to construct viable models by imposing global $U(1)_B$ conservation as an additional symmetry at the Lagrangian level.

3.1 Adding a $SU(2)$ singlet up-type colored scalar

Let us now study a scenario in which a scalar field $\tilde{u}$ with the same gauge quantum numbers than the RH up-type quarks is added to the SM plus the seesaw. The relevant new Lagrangian terms are:

$$-\mathcal{L}_{\tilde{u}} = \overline{\boldsymbol{u}^c}\,\boldsymbol{\eta}\,\boldsymbol{N}\,\tilde{u}^* + \overline{\boldsymbol{d}^c}\,\boldsymbol{y}\,\boldsymbol{d}\,\tilde{u} + \text{H.c.}\ . \tag{3.3}$$

By assigning conventionally $L=0$ to the RH neutrinos, eq. (3.3) conserves lepton number. However, the two terms in (3.3) cannot be made both $U(1)_B$ invariant by any choice of the baryon charge for $\tilde{u}$, since the first term requires $B(\tilde{u}) = +1/3$, while the second one requires $B(\tilde{u}) = -2/3$. The simultaneous presence of L and B violating terms allows the construction of operators that lead to nucleon ($\mathcal{N} = n, p$) decays. For example, after EW symmetry breaking, the mixing between the heavy sterile and light active neutrinos results in the $\Delta B = \Delta L = 1$ dimension six operator

$$\mathcal{O}_6 = \sqrt{\frac{m_\nu}{M_N}}\frac{\eta\, y}{m_{\tilde{u}}^2}\left(\overline{d^c}\, d\right)\left(\overline{\nu^c}\, u\right)\ , \tag{3.4}$$

where m_ν denotes the light neutrino mass scale, which induces the decay $\mathcal{N} \to \pi\nu$. The nucleon lifetime can be estimated as:

$$\tau_{\mathcal{N}} \sim 10^{32}\left(\frac{10^{-17}}{\eta\, y}\right)^2\left(\frac{M_N}{10^8\,\text{GeV}}\right)\left(\frac{m_{\tilde{u}}}{1\,\text{TeV}}\right)^4\left(\frac{0.1\,\text{eV}}{m_\nu}\right)\ \text{yrs}\ , \tag{3.5}$$

to be compared with the current bounds $\tau_{p\to\pi^+\nu} > 0.25\times 10^{32}$ yrs and $\tau_{n\to\pi^0\nu} > 1.12\times 10^{32}$ yrs [22]. So, if we want to keep the $\tilde{u}$ mass around the TeV (to allow for its possible direct production) we see that even pushing the RH neutrino masses $M_N \gg 10^8$ GeV, the extremely tiny size required for the couplings would render this scenario highly unnatural. Nucleon stability can however be guaranteed if, by imposing global $U(1)_B$ conservation, one of the two terms in $\mathcal{L}_{\tilde{u}}$ is eliminated.[2] In the rest of the paper we assume $B(\tilde{u}) = +1/3$ and thus we drop the second term in eq. (3.3).

4 Cloistered Baryogenesis

The presence of the new interactions in (3.3) open a new channel for RH neutrino decays: $N_i \to u_a\tilde{u}^*$.[3] Despite being B conserving, as long as $\tilde{u}$ remains (chemically) decoupled from the thermal bath this decay can provide a mechanism for baryogenesis. Note that once the lightest RH neutrino mass is fixed to satisfy $M_{N_1} \sim \mathcal{O}(10^7\,\text{GeV})$, standard N_1 leptogenesis can no longer generate a sizable B- L asymmetry because the CP asymmetry is way too small.[4]

For simplicity, let us now assume that the branching ratio for $N_1 \to \ell_\alpha H$ is much smaller than $N_1 \to u_a\tilde{u}^*$ so that to a good approximation the CP asymmetry can be normalized to the sum of the N_1 hadronic decays alone. In short, we assume that while the seesaw Lagrangian still accounts for neutrino masses and mixings, it does not have any role in baryogenesis. Fig. 1 illustrates how baryogenesis can proceed in our scenario. Initially the CP violating out-of-equilibrium N_1 decays produce equal and opposite sign B asymmetries in the up-type quarks (u_a) and in the colored scalars ($\tilde{u}$), that we respectively denote as ΔB_u and $\Delta B_{\tilde{u}}$. In the temperature range when the decays occur ($T \sim 10^7$ GeV), EW sphaleron processes are in thermodynamic equilibrium but, if the $\tilde{u}$'s remain decoupled from the hot plasma until EW sphaleron switchoff, the asymmetry $\Delta B_{\tilde{u}}$ remains unaffected. In contrast, as long as the u Yukawa couplings reactions and/or QCD sphalerons interactions are in thermal equilibrium, ΔB_u gets first transferred to the LH quarks and eventually is partially transformed into a ΔL asymmetry by EW sphalerons. As a result, after EW sphaleron freeze out at $T_{\text{fo}} \approx 80\,\text{GeV} + 0.45\,m_h \sim 135$ GeV [28, 29] (for $m_h = 125$ GeV [30, 31]), a net non-vanishing B asymmetry $\Delta B_{\text{SM}} + \Delta B_{\tilde{u}} \neq 0$ is obtained, although at this stage the total $B - L$ asymmetry still vanishes $\Delta B_{\text{SM}} + \Delta B_{\tilde{u}} - \Delta L = 0$ (final stage in Fig.1).

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[2] Note that one could also ensure nucleon stability by imposing global $U(1)_L$ conservation while allowing for explicit B violation. The resulting setup can allow to generate a B asymmetry through out-of-equilibrium C, CP and B violating decays of N, without the assistance of EW sphalerons. However, in this case one has to give up the possibility of light active Majorana neutrinos.

[3] Whenever necessary we will use Latin indices $i, j, \ldots$ to label RH neutrino generations and $a, b, \ldots$ to denote quark flavors, while lepton flavors will be denoted by Greek indices $\alpha, \beta, \ldots$.

[4] Let us recall that in the temperature regime $T \sim 10^7$ GeV there are no directions in flavor space that remain protected from N_1 washouts, and if N_1 couples sizeably to all flavors (i.e. $|\lambda_{\alpha 1}|^2 v^2/M_{N_1} \geq 10^{-3}$ eV) any pre-existing asymmetry will be erased [23]. In this case our mechanism for baryogenesis could be particularly relevant. Alternative possibilities for baryogenesis with $M_{N_1} \sim \mathcal{O}(10^7\,\text{GeV})$ include scenarios based on N_2 leptogenesis [24], resonant leptogenesis [25], and models with slightly broken L [26, 27].

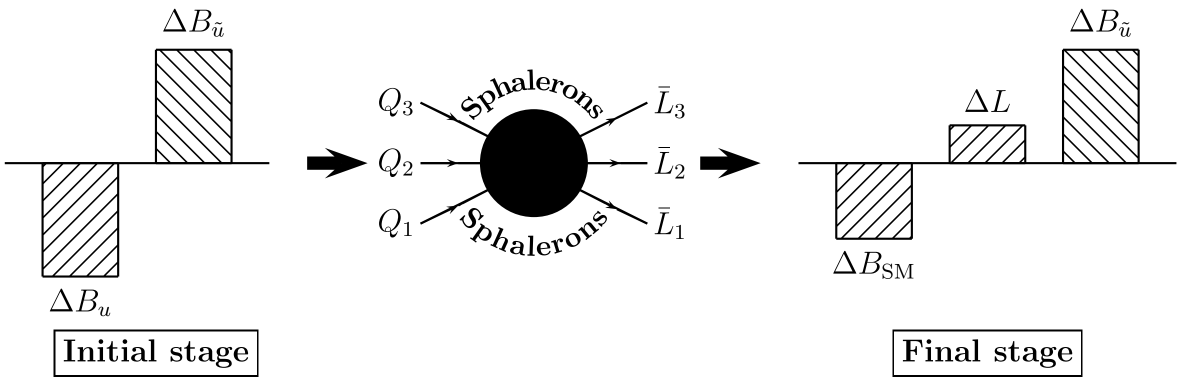

Figure 1. *Sketch of the cloistered baryogenesis mechanism. The equal and opposite sign B asymmetries respectively in u and $\tilde u$ in the initial stage are denoted by $\Delta B_{u,\tilde u}$. At EW sphaleron decoupling the B asymmetry in SM particles $\Delta B_{\rm SM}$ is no longer equal in magnitude to the opposite sign asymmetry $\Delta B_{\tilde u}$ due to the EW sphaleron processes which transfer part of the initial ΔB_u to the lepton sector. A net non-vanishing asymmetry $\Delta B_{\rm SM} + \Delta B_{\tilde u} \neq 0$ then results.*

Now, given that astrophysical arguments rule out the possibility of cosmologically stable heavy colored particles [32], $\tilde u$ must eventually decay. It is a feature automatically embedded in our model that they can do so only after EW symmetry breaking, that is when the baryon asymmetry they release cannot be affected any more by sphaleron interactions. Decays occur because the Dirac entries in the seesaw neutrino mass matrix, which are proportional to the vev of the Higgs field, induce N-ν mixing, and this opens up the decay $\tilde u \to u\nu$. The last step of baryogenesis thus occurs after EW symmetry breaking when $\Delta B_{\tilde u}$ is released in the plasma. This asymmetry remains largely unbalanced by the baryon asymmetry already present (see section 4.2) and in this way a net cosmological baryon asymmetry results.

4.1 The viability of B-conserving baryogenesis

The CP asymmetry in N_1 decays arises from the interference between the tree-level decay and the one-loop vertex and wave function corrections, as shown in Fig. 2. Assuming a hierarchical RH neutrino mass spectrum ($M_{N_i} < M_{N_j}$ for $i < j$), summing over quark flavors and taking into account color factors, the CP asymmetry between the number of $\tilde u$ and $\tilde u^*$ scalars produced in N_1 decays is

$$\epsilon_{N_1}^{\tilde u} \simeq -\frac{1}{4\pi}\frac{1}{(\boldsymbol{\eta}\boldsymbol{\eta}^\dagger)_{11}}\sum_{j\neq 1}\mathbb{I}\mathrm{m}\left[\left(\boldsymbol{\eta}\boldsymbol{\eta}^\dagger\right)^2_{1j}\right]\frac{M_{N_1}}{M_{N_j}}\ . \tag{4.1}$$

In addition to a sufficiently large CP asymmetry, the success of cloistered baryogenesis requires that the following three conditions are satisfied:

(i) The decays $\Gamma_{N_1} \equiv \Gamma(N_1 \to \sum_a u_a \tilde u^*)$ should occur out of equilibrium, and the rate for the N_1 semileptonic decays should satisfy $\Gamma(N_1 \to \sum_\alpha \ell_\alpha H^*) < \Gamma_{N_1}$;

(ii) The scalars $\tilde u$ should remain chemically decoupled from the thermal bath (of course strong interactions will keep them in kinetic equilibrium);

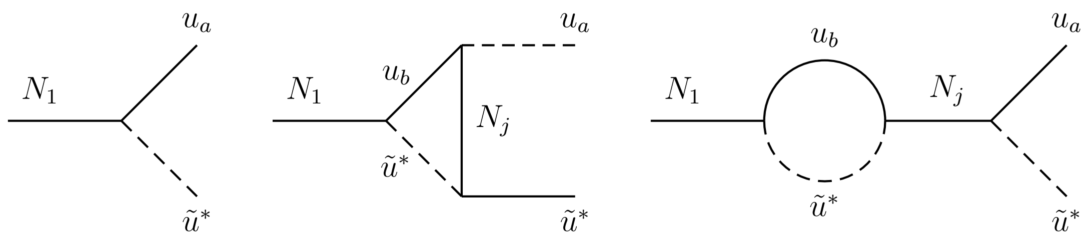

Figure 2. *Tree-level and one-loop vertex and wave function corrections Feynman diagrams responsible for the CP asymmetry in the colored scalar scenario.*

(*iii*) The decays $\tilde{u} \to u\nu$, which eventually fix the final amount (and sign) of the baryon asymmetry, should occur well before the Big Bang Nucleosynthesis (BBN) era.

All together these conditions enforce constraints on the relevant model parameters.

Condition (*i*) is satisfied provided that $\Gamma_{N_1} = \frac{3}{8\pi}|\eta_{a1}|^2 M_{N_1}$ is smaller than the Universe expansion rate at $z \equiv M_{N_1}/T \sim 1$, which implies

$$|\eta_{a1}| \lesssim 1 \times 10^{-5} \left(\frac{M_{N_1}}{10^7\,\text{GeV}}\right)^{1/2} . \tag{4.2}$$

Here and henceforth we normalize the RH neutrino mass to 10^7 GeV, as suggested by the condition for successful baryogenesis discussed in section 2.

Condition (*ii*) implies specific requirements on the rates of the s and t channel scattering process $u_a\tilde{u}^* \leftrightarrow \bar{u}_b\tilde{u}$, on the (inverse) decay rates of the heavier RH neutrinos $u_a\tilde{u}^* \to N_{2,3}$, and on the rates of the three-body decays $\tilde{u} \to \ell_\alpha\, H\, u_a$:

- $N_{2,3}$ **mediated** s **and** t **channel** $2 \leftrightarrow 2$ **scatterings**. As argued in section 3, $2 \leftrightarrow 2$ processes can place tight constraints on baryogenesis. The role played in our specific case by $N_{2,3}$ mediated $u_a\,\tilde{u}^* \leftrightarrow \bar{u}_b\,\tilde{u}$ scatterings can be readily understood from the one-loop diagrams in Fig. 2, since they involve the same couplings as the $2 \leftrightarrow 2$ reactions. Requiring that these reactions are out of equilibrium enforces constraints on the ratio between the Yukawa couplings and the heavier neutrino masses, and in turn this can imply a too large suppression of the CP asymmetries. Considering for definiteness only N_2 and one single channel, the $u_a\,\tilde{u}^* \leftrightarrow \bar{u}_b\,\tilde{u}$ scattering rate is approximately given by

$$\Gamma(u_a\,\tilde{u}^* \leftrightarrow \bar{u}_b\,\tilde{u}) \simeq \frac{1}{\pi^3}\frac{M_{N_1}^3}{M_{N_2}^2}\,|\eta_{a2}|^2\,|\eta_{b2}|^2 \ , \tag{4.3}$$

and demanding that this reaction to be decoupled at $z \sim 1$, implies the following constraint on the Yukawa couplings:

$$|\eta_{a2}|\,|\eta_{b2}| \lesssim 2 \times 10^{-5} \left(\frac{M_{N_2}}{M_{N_1}}\right)\left(\frac{M_{N_1}}{10^7\,\text{GeV}}\right)^{1/2} . \tag{4.4}$$

The analogous limits for N_3 mediated reactions are obtained by substituting $\eta_{a2} \to \eta_{a3}$ and $M_{N_2} \to M_{N_3}$.

- $N_{2,3}$ **inverse decays**: At $T \sim M_{N_1} \ll M_{N_{2,3}}$, $N_{2,3}$ inverse decays are Boltzmann suppressed, but one has to ensure that this suppression is sufficient to avoid depleting the asymmetry from N_1 decays. The thermally averaged inverse decay rate can be approximately written as

$$\gamma(u_a\,\tilde{u}^* \to N_{2,3}) \simeq \Gamma(N_{2,3} \to u_a\,\tilde{u}^*)\,\left(\frac{M_{N_{2,3}}}{M_{N_1}}\right)^{3/2}\,e^{-M_{N_{2,3}}/T}\,. \tag{4.5}$$

In terms of the RH neutrino masses and Yukawa couplings, the condition $\gamma(u_a\,\tilde{u}^* \to N_{2,3}) \lesssim H(z \sim 1)$ translates into

$$|\eta_{a(2,3)}| \lesssim\, 1.5 \times 10^{-5}\,\left(\frac{M_{N_1}}{M_{N_{2,3}}}\right)^{5/4}\,\left(\frac{M_{N_1}}{10^7\,\mathrm{GeV}}\right)^{1/2}\,e^{M_{N_{2,3}}/2M_{N_1}}\,. \tag{4.6}$$

For example, by taking $M_{N_1} = 10^7$ GeV and $M_{N_1}/M_{N_j} = 0.04$, we obtain $|\eta_{aj}| \lesssim 7 \times 10^{-2}$. Because of the exponential factor, as soon as the ratio M_{N_1}/M_{N_j} falls below $\sim 10^{-2}$ this constraint becomes completely irrelevant with respect to the constraints from $2 \leftrightarrow 2$ scatterings eq. (4.4), which are only power suppressed.

- **Three-body decays**: Already above the EW symmetry breaking scale, the colored scalars can decay via the RH-neutrino-mediated three body channel $\tilde{u} \to u_a\,\ell_\alpha H$. If sufficiently fast, this process would spoil the generation of the baryon asymmetry because $\Delta B_{\tilde{u}}$ is re-injected in the thermal bath too early, that is when EW sphalerons are still active. This decay channel, however, involves not only the $\boldsymbol{\eta}$ couplings but also the parameters responsible for neutrino masses and mixings $\boldsymbol{\lambda}$. The corresponding constraint reads

$$|\lambda_{\alpha j}|\,|\eta_{aj}| \lesssim 2 \times 10^{-3}\left(\frac{T_{\mathrm{fo}}}{135\,\mathrm{GeV}}\right)\left(\frac{M_{N_j}}{10^7\,\mathrm{GeV}}\right)\left(\frac{1\,\mathrm{TeV}}{m_{\tilde{u}}}\right)^{3/2}\,. \tag{4.7}$$

For $\mathcal{O}(M_{N_j}) \sim 10^7$ GeV consistency with a neutrino mass scale below, say, a few tenths of eV already requires $|\lambda| \lesssim 10^{-3}$, so that the constraint eq. (4.7) is easily satisfied and basically of no importance.

After EW symmetry breaking the active-RH neutrino mixing induces the decays $\tilde{u} \to u_a\,\nu$ which release the asymmetry $\Delta B_{\tilde{u}}$ in the thermal bath. Condition (*iii*) requires that these decays occur at temperatures safely above the temperature where the n/p ratio freezes out and BBN starts. Note that BBN constraints on hadronically decaying massive particles [33] assume in general that no baryon asymmetry is generated in these decays, and thus involve a different type of effects. In our case the requirement that has to be imposed is that the correct value (and sign) of the baryon-to-photon ratio is established as the initial condition for BBN. This yields the following constraint:

$$|\eta_{aj}| \gtrsim 3 \times 10^{-4}\,\left(\frac{T_{\mathrm{BBN}}}{10\,\mathrm{MeV}}\right)\,\left(\frac{M_{N_j}}{10^7\,\mathrm{GeV}}\right)^{1/2}\,\left(\frac{0.1\,\mathrm{eV}}{m_\nu}\right)^{1/2}\,\left(\frac{1\,\mathrm{TeV}}{m_{\tilde{u}}}\right)^{1/2}\,. \tag{4.8}$$

Note that the out-of-equilibrium condition require the η_{a1} couplings to be smaller than $\sim 10^{-5}$ (see eq. (4.2)). As shown in eq. (4.4), constraints on the $\eta_{a(2,3)}$ couplings are much weaker, implying that $\tilde{u} \to u_a \nu$ decays can occur at a sufficiently early stage thanks to the contributions from $\nu - N_{2,3}$ mixing. For example, fixing the $\eta_{a(2,3)} \sim 0.03$ and $M_{N_{(2,3)}} \sim 10^8$ GeV, one obtains for the colored scalar a mean lifetime $\tau_{\tilde{u}} \sim 10^{-4}$ sec. which ensures that all $\tilde{u}$ will have decayed much before BBN.

4.2 Chemical equilibrium conditions and kinetic equations

A more quantitative analysis of cloistered baryogenesis requires writing down the relevant Boltzmann equations, while taking into account the chemical equilibrium conditions enforced by those reactions that, in the range of temperatures relevant for the production of $\Delta B_{\rm SM}$ and $\Delta B_{\tilde{u}}$, are faster than the Universe expansion rate. In the following we fix this temperature at $T \sim 10^7$ GeV that, as was discussed in section 2, within our scenario is the lowest value compatible with successful baryogenesis.

We start by recalling some well known relations and by introducing notations. The number density asymmetry of bosons and fermions $\Delta n_{b,f} \equiv n_{b,f} - \bar{n}_{b,f}$ is related to the corresponding chemical potentials $\mu_{b,f}$. In the relativistic limit ($m_{b,f} \ll T$) and at first order in $\mu_{b,f}/T \ll 1$ the corresponding relations read:

$$\Delta n_b = \frac{T^3}{3}\left(\frac{\mu_b}{T}\right), \qquad \Delta n_f = \frac{T^3}{6}\left(\frac{\mu_f}{T}\right) . \tag{4.9}$$

Note that above we have defined $\Delta n_{b,f}$ as particle number asymmetries for degree of freedom. Then the number of degrees of freedom $g_{b,f}$ of each particle has to be taken into account when constructing global asymmetries for example in baryon or in lepton number. To remove the effect of the expansion of the Universe it is customary to normalize the particle number densities to the entropy density $s = g_\star\,(2\pi^2/45)\,T^3$ i.e. $Y_{\Delta n} = Y_n - Y_{\bar{n}} \equiv \Delta n/s$. In principle, for each non-self conjugate particle there is one chemical potential. However, the overall number of independent chemical potentials is drastically reduced by the different constraints imposed by the chemical equilibrium conditions and/or conservation laws, and eventually it turns out to be equal to the number of conserved charges. We follow ref. [10] and adopt the notation $X \equiv \mu_X$, where X is either a SM field or $\tilde{u}$ and μ_X its corresponding chemical potential. The constraints on the chemical potentials are:

1) Chemical potentials for the gauge bosons vanish $W = B = g = 0$ and hence all the particles belonging to the same $SU(3)_C$ or $SU(2)_L$ multiplets have the same chemical potential [34].

2) The Yukawa reactions for the second and third generations of SM fermions are in thermodynamic equilibrium. For simplicity we assume equilibrium also for the first generation (numerical differences do not exceed the ten percent level [35]). Also, intergenerational quark mixing ensures $Q_a = Q$, so we get:

$$\ell_\alpha - e_\alpha - H = 0 \qquad (\alpha = e, \mu, \tau) , \tag{4.10}$$

$$Q - u_a + H = 0 \qquad (u_a = u, c, t) , \tag{4.11}$$

$$Q - d_a - H = 0 \qquad (d_a = d, s, b) . \tag{4.12}$$

3) Equilibrium of EW sphaleron interactions yields

$$9Q + \sum_\alpha \ell_\alpha = 0. \tag{4.13}$$

4) In terms of chemical potentials, the condition of cosmological hypercharge neutrality $\sum_X \mathcal{Y}_X g_X \Delta n_X$ (with $\mathcal{Y}_X$ the X-particle hypercharge and g_X its number of degrees of freedom) translates into:

$$3Q + \sum_a (2u_a - d_a) - \sum_\alpha (\ell_\alpha + e_\alpha) + 2H + 4\tilde{u} = 0\ . \tag{4.14}$$

Note that when all quarks Yukawa reactions are assumed to be in thermodynamic equilibrium QCD sphalerons do not impose an independent constraint [35]. All in all, the initial 15 chemical potentials $u_a, d_a, e_\alpha, \ell_\alpha, Q, H, \tilde{u}$, are reduced to 4 by the $9+1+1=11$ conditions implied by 2, 3 and 4, namely, by Eqs. (4.10-4.12), (4.13) and (4.14). As mentioned above, this could have been expected simply from symmetry considerations. In the approximation in which the $N_1 \leftrightarrow u_a \tilde{u}^*$ reactions are completely out of equilibrium, there are four conserved charges corresponding to global $U(1)_{\tilde{u}}$ [5] and to the three global $U(1)_{\Delta_\alpha}$ where $\Delta_\alpha \equiv \Delta B_{\rm SM}/3 - \Delta L_\alpha$. Hence the normalized number density asymmetries of all particle species can be expressed in terms of the asymmetries in the four charges $Y_{\Delta_{\tilde{u}}}$ and Y_{Δ_α}. We obtain:

$$\begin{aligned}
Y_{\Delta\ell_\alpha} &= -\frac{3}{79}Y_{\Delta\tilde{u}} + \frac{16}{711}Y_{\Delta(B_{\rm SM}-L)} - \frac{1}{3}Y_{\Delta_\alpha}\,, & Y_{\Delta u_a} &= -\frac{12}{79}Y_{\Delta\tilde{u}} - \frac{5}{237}Y_{\Delta(B_{\rm SM}-L)}\,,\\
Y_{\Delta e_\alpha} &= \frac{10}{79}Y_{\Delta\tilde{u}} + \frac{52}{711}Y_{\Delta(B_{\rm SM}-L)} - \frac{1}{3}Y_{\Delta_\alpha}\,, & Y_{\Delta d_a} &= \frac{14}{79}Y_{\Delta\tilde{u}} + \frac{19}{237}Y_{\Delta(B_{\rm SM}-L)}\,,\\
Y_{\Delta Q} &= \frac{1}{79}Y_{\Delta\tilde{u}} + \frac{7}{237}Y_{\Delta(B_{\rm SM}-L)}\,, & Y_{\Delta H} &= -\frac{26}{79}Y_{\Delta\tilde{u}} - \frac{8}{79}Y_{\Delta(B_{\rm SM}-L)}\,,
\end{aligned} \tag{4.15}$$

where $Y_{\Delta(B_{\rm SM}-L)} = \sum_\alpha Y_{\Delta_\alpha}$. It is important to note, as could be readily verified from the previous relations, that since the hypercharge condition eq. (4.14) is different from the SM case because of the presence of the contribution from the $\tilde{u}$ scalars, the relation between the amount of baryon asymmetry and $B-L$ asymmetry stored in SM particles is also changed, and reads

$$Y_{\Delta B_{\rm SM}} = \frac{28}{79}Y_{\Delta(B_{\rm SM}-L)} + \frac{12}{79}Y_{\Delta\tilde{u}}\,, \tag{4.16}$$

where the first term on the RH side is the usual SM result, while the second is new.

Now, the dynamical equations for baryogenesis get largely simplified in the approximation in which $N \leftrightarrow \ell_\alpha H$ interactions are neglected, and we will adopt this approximation in the last part of this section. Although at $T \sim 10^7$ GeV the three

[5] Note that the presence of a global $U(1)_{\tilde{u}}$ can be in fact taken as an *operative* definition of having $\tilde{u}$ decoupled from the thermal plasma.

lepton flavors are neatly distinguished by their Yukawa interactions [7, 9], in this approximation all dynamical processes become symmetric under a relabeling of the lepton flavor index α, and this means that the asymmetries Y_{Δ_α} evolve in the same way and must be equal at all times. Thus we can simply set $Y_{\Delta_\alpha} = \frac{1}{3}Y_{\Delta(B_{\rm SM}-L)}$. Another simplification stems from the fact that at this stage the total $B-L$ is a conserved quantity, that is

$$Y_{\Delta(B_{\rm SM}-L)} + Y_{\Delta\tilde{u}} = 0\,, \tag{4.17}$$

where the second term in the RH side is the contribution to total ΔB from the $\tilde{u}$ scalars. Eq. (4.17) implies that to estimate the baryon asymmetry yield of cloistered baryogenesis is sufficient to solve a system of just two Boltzmann equations:

$$\dot{Y}_{N_1} = -(y_{N_1}-1)\,\gamma_{N_1}\ , \tag{4.18}$$

$$\dot{Y}_{\Delta\tilde{u}} = (y_{N_1}-1)\epsilon^{\tilde{u}}_{N_1}\,\gamma_{N_1} + \frac{1}{2}\left(y_{\Delta u} - y_{\Delta\tilde{u}}\right)\gamma_{N_1}\ , \tag{4.19}$$

where γ_{N_1} denotes the thermally averaged N_1 decay rate, the time derivative is defined as $\dot{Y} \equiv sHz\,dY/dz$, the density asymmetries have been normalized as $y_{\Delta\tilde{u}} = Y_{\Delta\tilde{u}}/Y_b^{\rm Eq}$ and $y_{\Delta u} = Y_{\Delta u}/Y_f^{\rm Eq}$ with the respective boson and fermion equilibrium abundances $Y_b^{\rm Eq} = 2Y_f^{\rm Eq} = \frac{15}{4\pi^2 g_*}$, we have dropped the RH up-type quark flavor index by setting, according to eq. (4.15), $u_a = u$, and finally we have neglected on-shell and off-shell contributions from $N_{2,3}$. Note that $y_{\Delta u}$ appearing in the washout term in eq. (4.19) has to be evaluated by means of the first relation in the right side column in eq. (4.15) together with eq. (4.17). This yields

$$y_{\Delta u} = -\frac{62}{237}\,y_{\Delta\tilde{u}}. \tag{4.20}$$

According to eq. (4.16) and eq. (4.17), once the era of N_1 decays is ended, but before the $\tilde{u}$ scalars start decaying (let us say, for definiteness, at temperatures around the EW phase transition), the amount of baryon asymmetry stored in SM particles is

$$Y^{\rm EW}_{\Delta B_{\rm SM}} = -\frac{16}{79}Y^{\rm EW}_{\Delta\tilde{u}}\,, \tag{4.21}$$

that is about 20% of the final value of $Y_{\Delta\tilde{u}}$ but of opposite sign. However, what should be confronted with cosmological measurements is the baryon asymmetry *after* all the $\tilde{u}$ scalars have decayed (say, for definiteness, at temperatures around the BBN era) which is given by:

$$Y^{\rm BBN}_{\Delta B} = Y^{\rm EW}_{\Delta B_{\rm SM}} + Y^{\rm EW}_{\Delta\tilde{u}} = \frac{63}{79}Y^{\rm EW}_{\Delta\tilde{u}}\,. \tag{4.22}$$

Confronting eq. (4.21) and eq. (4.22) shows that the main contribution to the present cosmological baryon asymmetry as well as its sign, are determined by the asymmetry stored in the colored scalars $\tilde{u}$, which remain decoupled from the thermal bath down to temperatures well below the EW phase transition. This asymmetry could in fact be released at temperatures as low as $\mathcal{O}(10~{\rm MeV})$, right before the onset of BBN.

5 The role of hypercharge

The analysis of the previous section indicates that in our baryogenesis model the small amount of perturbative L violation does not play any crucial role. As regards baryon number, apart from sphaleron interactions, at the Lagrangian level it remains conserved at all stages. It is then interesting to ask which is the fundamental charge whose asymmetry is feeding all particle asymmetries, and eventually baryogenesis. As we will now argue, the answer is that this role is played by the asymmetry in the total hypercharge of the SM particles.[6] The following example will help to make this point more clear. Let us assume the following setup:

- The two baryon asymmetries $\Delta B_{\rm SM}$ and $\Delta B_{\tilde{u}}$ are generated in the out-of-equilibrium decays of N_2, with the usual condition $\Gamma(N_2 \to \ell H) \ll \Gamma(N_2 \to u\tilde{u}^*)$.
- The N_1 decay rate $\Gamma(N_1 \to u\tilde{u}^*)$ is instead negligible, while the L violating decays and inverse decays $N_1 \leftrightarrow \ell H, \bar{\ell} H^*$ are in full thermal equilibrium.

This second assumption implies one additional condition, which should be added to the set of chemical potential relations eqs. (4.10)-(4.13). Recalling that the Majorana states N have vanishing chemical potential, this condition reads:

$$\ell_\alpha + H = 0 \qquad (\alpha = e, \mu, \tau)\,. \tag{5.1}$$

Now, from the hypercharge neutrality condition eq. (4.14) we have that the sum of the SM particle number asymmetries weighted by the hypercharge of the corresponding particles, and written in terms of chemical potentials, should add to $-2g_{\tilde{u}}\,\mathcal{Y}_{\tilde{u}}\,\tilde{u} = -4\tilde{u}$ (with $g_{\tilde{u}} = 3$ the color degrees of freedom of $\tilde{u}$, and $\mathcal{Y}_{\tilde{u}} = 2/3$ its hypercharge), that is it should exactly balance the amount of hypercharge asymmetry stored in the cloistered sector. The solution of the set of chemical potential conditions eqs. (4.10)-(4.13) and eq. (5.1) is straightforward: since all SM reactions as well as $N_1 \leftrightarrow \ell H, \bar{\ell} H^*$ conserve exactly hypercharge, the chemical potential of the SM particles must be simply proportional to the particle hypercharges:

$$\mu_\phi = \kappa\,\mathcal{Y}_\phi \qquad (\phi = \ell_\alpha, e_\alpha, Q, u_a, d_a, H)\,. \tag{5.2}$$

The coefficient κ can then be directly evaluated from total hypercharge conservation:

$$\kappa = -\frac{2\,g_{\tilde{u}}\,\mathcal{Y}_{\tilde{u}}}{\sum_\phi\, g_\phi\,\mathcal{Y}_\phi^2}\,\tilde{u}\,. \tag{5.3}$$

Note that, with a slight abuse of notation, within the sum in the denominator $g_H = 2\times 2$ where the first factor is from the Higgs $SU(2)$ degrees of freedom, and the second from boson/fermion statistics $\Delta n_H/\Delta n_f = 2\mu_H/\mu_f$. This allows us to write the chemical potentials of all the SM particles in terms of $\tilde{u}$, which in turn is obtained by integrating the Boltzmann equations (4.18) and (4.19). We thus see that even when L is violated by in-equilibrium reactions and B is perturbatively conserved, still, in order to balance the net amount of hypercharge stored in the cloistered sector, all SM particles carrying hypercharge develop non-vanishing asymmetries.

[6] That such an asymmetry could drive baryogenesis was noted already long ago in ref. [36].

6 Conclusions

We have studied a scenario where the cosmological matter/antimatter asymmetry stems from an asymmetry in baryon number related to heavy particle decays. To ensure nucleon stability, baryon number is imposed as a symmetry at the Lagrangian level; however, baryogenesis can still proceed because a certain amount of baryon asymmetry generated from the B conserving decays of heavy particles is confined into a *cloistered* sector that remains chemically decoupled from the thermal bath until $B+L$ violating sphaleron reactions are switched off. An initial equal amount of baryon asymmetry stored in the SM sector gets instead partially transformed into a lepton asymmetry. When the asymmetry in the cloistered sector is eventually released into the thermal bath (in our model this can naturally occur at temperatures not far above the onset of BBN) the unbalance between the two asymmetries gives rise to baryogenesis.

We have studied some necessary conditions to allow for successful baryogenesis within this scenario. For example we have found that sufficiently large CP asymmetries together with the requirement that the cloistered sector will remain chemically decoupled from the thermal bath, require that the mass of the initial heavy decaying particles must be at least of $\mathcal{O} \sim 10^7$ GeV. While this is about two orders of magnitude lower than the scale required for successful leptogenesis, it remains well above the TeV scale, thus excluding the possibility of direct tests at colliders.

We have implemented cloistered baryogenesis within a specific setup, based on a straightforward extension of the standard seesaw model to which a colored scalar $\tilde{u}$ with the same quantum numbers of the up-type RH quarks is added. We have illustrated in detail the viability of this realization, we have analyzed various constraints showing that they can all be satisfied, we have derived the chemical equilibrium conditions that relate the SM particle asymmetries, and we have written down the kinetic equations whose solution allows to estimate the present amount of cosmological baryon asymmetry. Finally, we have highlighted the fundamental role played in our setup by hypercharge conservation [36].

If the new colored states which are the clue ingredient of cloistered baryogenesis have, as we have assumed, masses of $\mathcal{O}(\text{TeV})$, they would be well within the LHC reach even with moderate luminosity, given that their production rates are governed by α_s. The requirement that they will keep decoupled from the thermal bath implies, as a specific signature, a relatively long lifetime. Thus, they could be produced at the LHC in large numbers, and leave a characteristic signature throughout all layers of the detectors, much alike the long lived colored particles studied in [37]. The experimental observation of colored scalars with these characteristics will clearly not suffice to identify cloistered baryogenesis as the mechanism responsible for the cosmic baryon asymmetry, but it would certainly support this idea.

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