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The small observed baryon asymmetry from a large lepton asymmetry

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Abstract: Primordial Big-Bang Nucleosynthesis (BBN) tightly constrains the existence of any additional relativistic degrees of freedom at that epoch. However a large asymmetry in electron neutrino number shifts the chemical equilibrium between the neutron and proton at neutron freeze-out and allows such additional particle species. Moreover, the BBN itself may also prefer such an asymmetry to reconcile predicted element abundances and observations. However, such a large asymmetry appears to be in conflict with the observed small baryon asymmetry if they are in sphaleron mediated equilibrium. In this paper we point out the surprising fact that in the Standard Model, if the asymmetries in the electron number and the muon number are equal (and opposite) and of the size preferred to improve the agreement between BBN theory with observations, a baryon asymmetry of the Universe of the correct magnitude and sign is automatically generated within a factor of two. This small remaining discrepancy is naturally remedied in the supersymmetric Standard Model.

Keywords: Cosmology of Theories beyond the SM, Baryogenesis, Physics of the Early Universe.

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Contents

1. Introduction

Primordial Big Bang Nucleosynthesis (BBN) is without doubt one of the biggest successes of early universe cosmology. Not only does it provide a stringent test of the Big Bang model, predicting the light element abundences as a function of only a single parameter, $\eta = n_b/n_\gamma$, the cosmological baryon to photon ratio, but it also supplies important constraints on particle physics, the most well-known example being the determination of the number of light neutrino species. Given the consistency between the primordial abundance of light elements (inferred from observation extrapolated back to the primordial values) and theoretical calculations, BBN does not leave much room for extra particles which otherwise could have existed in the early universe.

Many extensions of physics beyond the Standard Model (SM), however, introduce additional relativistic degrees of freedom at the epoch of BBN. A small selection of such new light degrees of freedom include: one or more sterile neutrinos which might be required by the neutrino oscillation data; a light gravitino as a consequence of a low fundamental scale of supersymmetry breaking as in the gauge-mediated scenarios; a hadronic axion in the hot dark matter window, and many other examples. If such new light degrees of freedom exist, the expansion rate at the BBN epoch is faster, resulting in an earlier freeze-out of neutrons and hence a larger number of them, therefore overproducing ^{4}He. Taking the BBN constraint seriously, it is then necessary to modify standard BBN, and the simplest and most elegant possibility is a large lepton asymmetry,[1] a possibility which is not ruled out by current observational limits [1, 2, 3]. Specifically, a large positive asymmetry in the electron number implies an excess in the number of electron neutrinos over that of electron anti-neutrinos, thereby shifting the chemical equilibrium between protons and neutrons towards

[1] One other possibility discussed in the literature is that of a late-decaying ν_τ.

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protons. This results in a smaller number density of neutrons after the freeze-out and hence in a reduced ^{4}He abundance. This effect can therefore compensate the effect of the larger expansion rate due to the additional particle species.

Moreover in recent years, with the advent of new and refined data on the relative abundances of the light elements, there may be appearing a slight but significant discrepancy between the data and the theoretical predictions. In particular, the recent low measurements [4] of the primordial Deuterium abundance is in a conflict with the low ^{4}He abundance as reported in [5]. One possible resolution is that the ^{4}He abundance is higher than that in [5], as claimed by a re-analysis of more-or-less the same data set in [19]. If the conflict persists, on the other hand, then some modification of the standard BBN scenario seems to be required independent of the conjectured existence of new light degrees of freedom. The most promising such modification is again the assumption of a positive chemical potential for electron neutrinos which reduces the final ^{4}He abundance closer to the reported value. It is noteworthy that the preferred sign of the electron asymmetry is the same for both purposes: to compensate the effect of additional particle species and to bring the BBN prediction closer to observations. Of course, given the uncertainties in the data, it is not clear if this is really required by primordial nucleosynthesis. It is, however, useful to explore such modifications of the standard Big-Bang scenario to see if they are either disfavored by other data, or serve some further, unexpected, purpose.

On the other hand, there is an apparent contradiction of an assumption of large lepton asymmetry with the very small observed baryon asymmetry. This arises from the presence of sphaleron mediated transitions at temperatures of the weak scale and above which tend to quickly equilibriate the lepton and baryon asymmetries, resulting in far too large a baryon asymmetry today. There are three logical possibilities for how a large ν_e lepton asymmetry can be compatible with the small baryon asymmetry: (1) Sphalerons were never in equilibrium, (2) The lepton asymmetry is generated after the electro-weak phase transition but before BBN, and, (3) The *total* lepton asymmetry across all three generations is zero.

In this letter we focus on the third possibility — in particular the case where $L_e = -L_\mu \neq 0$ and $L_\tau = 0$ — and show, in section 3, that it has a very pretty and unexpected consequence — the natural generation within the *Standard Model* of a small baryon asymmetry of the observed size, at least within a factor of two, and with the correct sign! This numerical coincidence is quite remarkable, especially given the simplicity and naturalness of the baryon asymmetry generation mechanism. The fundamental reason for the small baryon asymmetry in this case, $L_e = -L_\mu \neq 0$, is quite simple; it is just a consequence of the small muon Yukawa coupling. As we show in section 4, if one goes to the minimal supersymmetric standard model (MSSM) then even the factor of two discrepancy between the predicted and observed baryon asymmetry disappears for large $\tan(\beta)$. (Section 2 contains a more extensive discussion of the reasons for considering a Lepton asymmetry, together with its possible size and sign.)

2. BBN with large lepton number

We will now argue in detail that it is useful to explore the possibility that there may be a slight modification of standard BBN, and that such modifications are certainly not disallowed and are possibly even favored by the light element abundances.

Many particle physics models beyond the SM introduce additional particle species which could be relativistic and thermal at the BBN epoch. Probably the most discussed such example is a sterile neutrino (or many of them, especially in the context of neutrinos from large extra dimensions [6]). If one takes all existent hints for neutrino oscillations seriously, namely the atmospheric neutrino oscillations, solar neutrino deficit and the results from the LSND experiment, the data cannot be accommodated by neutrino oscillations between the three known species: ν_e, ν_μ, ν_τ. The reason is simple. The three hints for oscillations listed above require different values of the mass-squared differences Δm^2, and with three neutrinos only, the sum of Δm^2 should vanish. The only known way to explain the data fully by neutrino oscillations is by introducing an additional "sterile neutrino" ν_s, thereby allowing yet another mass-squared difference to account for three oscillation modes. However, neutrino oscillations should have occurred in the early universe as well, thus producing sterile neutrino states. In order not to overproduce ^{4}He due to the additional sterile neutrino energy density, the quoted bounds are [7, 8]

$$\begin{aligned} \Delta m^2 \sin^4 2\theta &\lesssim 5 \times 10^{-6}\,\mathrm{eV}^2\,, \qquad \nu = \nu_e\,, \\ \Delta m^2 \sin^4 2\theta &\lesssim 3 \times 10^{-6}\,\mathrm{eV}^2\,, \qquad \nu = \nu_{\mu,\tau}\,. \end{aligned} \tag{2.1}$$

These constraints, taken literally, imply that sterile neutrinos cannot be responsible for atmospheric neutrino oscillations or the large angle MSW solution to the solar neutrino problem. The existence of a sterile neutrino exceeding the above bounds would increase the effective number of neutrinos at BBN by one: $\Delta N_\nu = 1$.

In supersymmetric theories, a light gravitino $\tilde{G}$ may be present at the BBN epoch as well. According to the estimate in ref. [9], the gravitinos remain thermal down to the BBN epoch if

$$m_{3/2} \lesssim 10^{-13}\,\mathrm{GeV}\left(\frac{m_{\tilde{l}}}{100\,\mathrm{GeV}}\right), \tag{2.2}$$

due to the process $l^+l^- \to \tilde{G}\tilde{G}$. This roughly corresponds to a primordial supersymmetry breaking scale below a TeV. Such a low scale is not expected in the conventional hidden sector models or gauge mediation, but can occur in models where the supersymmetric standard model is directly involved in the mechanism of dynamical supersymmetry breaking (see, e.g., the model in ref. [10]). Because the produced gravitino states are dominantly helicity $\pm 1/2$ (the would-be Nambu-Goldstino state), they increase the effective number of neutrinos by $\Delta N_\nu = 1$.

Invisible axions are another candidate particle that could be present at the BBN epoch. Despite strong constraints from astrophysics, a hadronic (KSVZ) axion in the mass range 3–20 eV is allowed as long as its coupling to the photon is accidentally suppressed [11]. This is an interesting window for a Hot Dark Matter component of the universe which some recent analyzes of large scale structure prefer [12] (however, for conflicting views, see [13]). The axion in this mass window would contribute to the energy density as an equivalent of $\Delta N_\nu = 0.4$–0.5 [14] and is marginal from the BBN point of view.

Yet another example of an exotic particle which might be in thermal contact during BBN is represented by the majoron, the Goldstone boson associated to the spontaneous breakdown of lepton number. Majorons stay in thermal equilibrium as long as τ-neutrinos, and provide a contribution to ΔN_ν of about 0.6 [15].

Given these important constraints from BBN on particle physics models, it is important to ask how rigid the constraint actually is. In this regard it is interesting to note that the BBN itself may require some modifications.

Specifically, if one takes the low Deuterium measurement [4] and the reported statistical average of the ^{4}He abundance extrapolated to zero metalicity [16], they cannot be reconciled with detailed BBN calculations by choosing an appropriate value of η, the baryon to photon ratio. Of course, it is not yet established that these measurements are reliable. For instance, one should take seriously the conflicting measurement of the Deuterium abundance based on the same technique which returns a high value [17], even though it has been challenged on the basis of a possible overlap with a foreground cloud and less systematic checks than the low abundance observation. (It is interesting to note that by including turbulence effects in the extraction of the D/H ratio [18], all the data is cosistent with a low value of D/He $\simeq 3.5 - 5.2 \times 10^{-5}$.) The "best" determination of the ^{4}He abundance has also been challenged by a re-analysis of the more-or-less the same data set [19]. Nevertheless there is motivation for considering modifications to BBN which can reconcile the "best" determinations of element abundances. Most certainly, such a modification is allowed by current data. (For a recent review see ref. [20].)

It is noteworthy that *both* the presence of additional relativistic degrees of freedom and the apparent inconsistency between the D and ^{4}He abundances prefer a mechanism to reduce the effective number of neutrinos N_ν. Two such possibilities have been proposed in the literature:

1. A late-decaying ν_τ with a mass of $m_{\nu_\tau} \sim 10\,\mathrm{MeV}$ and a lifetime of $\tau \sim 10^{-2}$–$1\,\mathrm{sec}$ [21, 22].

2. A large chemical potential for ν_e [23].

The former proposal is interesting from the collider physics point of view because it is testable in the forthcoming B-factory experiments [25].

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In this letter we focus on the second possibility. Here the idea is that the presence of a large chemical potential for ν_e makes the ν_e number density larger than the thermal number density without chemical potential, which in turn changes the chemical equilibrium of the reaction $\nu_e n \leftrightarrow e^- p$ etc. The presence of a positive chemical potential for ν_e shifts the equilibrium towards the right-hand side, which reduces the neutron number density at the freeze-out. Therefore the ^{4}He abundance is reduced for a given value of η. Since the D abundance [4] prefers a relatively large value of η, which prefers a large ^{4}He abundance, the reduced prediction for the ^{4}He abundance would allow additional relativistic degrees of freedom present at the BBN epoch or reconciles the apparent conflict between the observations and the calculations.[2]

The electron-neutrino chemical potential affects the neutron-to-proton ratio at the freeze-out as

$$\left(\frac{n}{p}\right)_{\xi_{\nu_e}\neq 0} = \left(\frac{n}{p}\right)_{\xi_{\nu_e}=0} e^{-\xi_{\nu_e}}\,, \tag{2.3}$$

where $\xi_{\nu_e} = \mu_{\nu_e}/T$ at the freeze-out temperature. The effect of the extra degrees of freedom on ^{4}He abundance is given by an analytic fit [28]:

$$\begin{aligned} Y_P = \Big[&\left(a_0 + a_1\,\zeta + a_2\,\zeta^2 + a_3\,\zeta^3 + a_4\,\zeta^4\right) + \\ &+\ \Delta N_\nu \left(c_0 + c_1\,\zeta + c_2\,\zeta^2 + c_3\,\zeta^3 + c_4\,\zeta^4\right)\Big] e^{-\xi_{\nu_e}}\,, \end{aligned} \tag{2.4}$$

for $\tau_n = 885.4$ sec, $\zeta = \log_{10}(\eta/10^{-10})$, and

$$\begin{aligned} a_0 &= 0.22292\,, & a_1 &= 0.05547\,, & a_2 &= -0.05639\,, \\ & & a_3 &= 0.04587\,, & a_4 &= -0.001501\,, \\ c_0 &= 0.01276\,, & c_1 &= 0.00409\,, & c_2 &= -0.00703\,, \\ & & c_3 &= 0.00571\,, & c_4 &= -0.00186\,. \end{aligned}$$

The low D measurement requires $\eta \simeq 5\times 10^{-10}$ and hence $Y_P \simeq (0.2495+0.0137(N_\nu - 3))e^{-\xi_{\nu_e}}$ which is beyond the quoted $Y_P = 0.234 \pm 0.002 \pm 0.005$ [16] (see, however, a conflicting number $Y_P = 0.244 \pm 0.002 \pm 0.005$ [19]). This would require $\xi_{\nu_e} \sim 0.064$. This approximate discussion also tells us that an additional degrees of freedom with $\Delta N_\nu = 1$ can be compensated by $\xi_{\nu_e} = 0.053$.

The size of the chemical potential favored to reconcile the observations and the BBN calculations of the light element abundances were studied by intensive numerical analysis in ref. [23] (see also [24]). The result is $\xi_{\nu_e} = (4.3 \pm 4.0) \times 10^{-2}$ at 95% CL,

[2]In the case of neutrino oscillations to a sterile neutrino, the interplay between the neutrino oscillations and thermalization can be quite complicated [26]. However, a large primordial lepton asymmetry which exists from the pre-BBN era does persist [8] and can allow the sterile neutrinos. This differs from the situations discussed in [26] where the lepton asymmetry was assumed to vanish primordially (i.e. before the BBN era).

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quite close to the rough estimate given above. From this the electron-number per photon ratio is given by

$$\frac{n_{\nu_e} - n_{\bar{\nu}_e}}{n_\gamma} = \frac{\pi^3}{12\zeta(3)} \left(\frac{T_{\nu_e}}{T_\gamma}\right)^3 \left(\frac{\xi_{\nu_e}}{\pi}\right) + \mathcal{O}(\xi^3)\,. \tag{2.5}$$

Since $T_{\nu_e} = T_\gamma$ in the relevant temperature regime, and the total entropy density is $s = \frac{43}{4}\frac{2\pi^2}{45}T^3$ (from photons, electrons, positrons and three neutrinos), we find the "preferred" electron-number to entropy ratio L_e to be

$$L_e^{\rm NUC} = \frac{15}{43}\frac{\xi_{\nu_e}}{\pi^2} = (1.52 \pm 1.41) \times 10^{-3}\,. \tag{2.6}$$

For the purpose of allowing an extra relativistic degree of freedom at the epoch of BBN, we would also require an additional contribution the electron-number to entropy ratio of this same magnitude and sign. Thus we take

$$L_e^\star \sim 2L_e^{\rm NUC} = (3.04 \pm 2.82) \times 10^{-3} \tag{2.7}$$

as the favored value of the lepton asymmetry both by compensating an additional relativistic degree of freedom at the BBN epoch and by reconciling the discrepancy between the theory and observation in the BBN itself.

3. Small baryon number from large lepton number

The most uncomfortable aspect of a large chemical potential for ν_e is the consistency with the small observed baryon asymmetry. An almost universal theoretical prejudice is that the baryon asymmetry is a consequence of non-trivial dynamics in the Early Universe, with the three Sakharov conditions being met: (1) the existence of a baryon-number violating interaction, (2) departure from thermal equilibrium, and (3) CP-violation. If there were also a chemical potential for ν_e, or in other words, an asymmetry in the electron number, it should also be a consequence of similar dynamics in the Early Universe. It then appears unnatural that the lepton asymmetry is many orders of magnitude larger than the baryon asymmetry if they are generated by similar mechanisms.

The uncomfortableness mentioned above becomes a conflict in the view of the following consideration. Given the difficulty in generating a large enough baryon asymmetry purely from the electroweak phase transition, the much larger preferred size of the lepton asymmetry from the BBN, eq. (2.6), is highly unlikely to be a consequence of physics at or below the electroweak scale. However, above the electroweak phase transition, neither baryon- or lepton-number is conserved, but only $B - L$ because of sphaleron mediated transitions and the electroweak B and L anomalies [29, 30, 31, 32]. Furthermore, the chemical equilibrium induced by sphaleron transitions enforces the baryon- and lepton-asymmetries to be of the same orders of magnitude.

There are three logical possibilities to avoid this conflict:

1. The large lepton asymmetry is generated below the electroweak scale.
2. The sphaleron transition was never in equilibrium below the temperature at which the lepton asymmetry was generated.
3. The total lepton asymmetry vanishes, while the individual lepton-flavor asymmetries do not.

We already argued that the first possibility is unlikely, even though it is logically possible. The second possibility arises if the large lepton number asymmetry causes a Bose condensate of electroweak-doublet scalar fields [33, 34, 35, 36]. In the Standard Model the preferred value of the lepton asymmetry from nucleosynthesis considerations is below the critical value [37] at which the Higgs doublet acquires a large expectation value and thus at temperatures above the electroweak scale the sphaleron transition is still in equilibrium. The same is true in the case of the MSSM as recently shown in ref. [38]. Note, in particular, that if the squark and slepton masses are heavier than the electroweak phase transition temperature of 100–200 GeV, they are irrelevant to this discussion and the situation is the same as in the SM and hence the sphaleron transitions are active. Moreover, even if one manages to keep sphaleron transitions out of equilibrium, it still does not resolve the question why the lepton asymmetry is so much larger than the baryon asymmetry. From these considerations, we find the third possibility to be the most interesting one, which has not been discussed in the literature so far.

The baryon and the lepton asymmetries are determined by the $B-L$ asymmetry via sphaleron-induced chemical equilibrium. For the Standard Model [32, 39]:

$$B = \frac{8N_G + 4N_H}{22N_G + 13N_H}(B-L)\,, \tag{3.1}$$

where $N_G = 3$ is the number of generations and N_H is the number of Higgs doublets (1 in the SM). In the presence of the supersymmetric particles, the formula is slightly modified [40]. Therefore, if the total lepton asymmetry vanishes, the total baryon number also vanishes. This way, one can obtain a vanishing baryon asymmetry even in the presence of individual flavor-dependent lepton asymmetries.

The above formula is usually assumed to hold above the electroweak phase transition temperature, while it requires modification after the phase transition because of finite mass effects. However, even above the phase transition temperature, the effects of thermal masses need to be considered. Such effects are small and usually ignored, but they cannot be ignored in the presence of the large individual lepton numbers of interest in this letter.

The final resulting baryon asymmetry depends on when the sphaleron transition freezes out, which in turn depends on whether the electroweak phase transition is

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strongly first-order or not [41]. Given the experimental lower bound on the Higgs mass of about 95 GeV together with the results of current large-scale numerical lattice simulations [42] and analytic arguments [43], the phase transition in the Standard Model is certainly not a strongly first order transition, while in the case of the MSSM a weakly first-order transition or smoother is favored over much of the parameter space. In the case that the phase transition is second order, or if the sphalerons are still active after a first order phase transition (i.e. a weakly first-order transition with $\langle\phi(T)\rangle/T \leq 1$, being $\langle\phi\rangle$ the vacuum expectation value of the Higgs field), there are two contributions to the resulting baryon asymmetry. These flavor-dependent effects both arise from the interaction of electrons and muons with the Higgs boson via their Yukawa couplings. (The two effects correspond to the interactions with condensed and real Higgs bosons respectively.) The total flavor-dependent effect was estimated in ref. [32], and in the case of vanishing total lepton asymmetry $L_e + L_\mu = 0$, we find

$$B = A\frac{6}{13\pi^2}\frac{\overline{m}_\mu^2(T)}{T^2}L_e^\star\,, \tag{3.2}$$

where $A \simeq 1$ [44] and

$$\frac{\overline{m}_\mu^2(T)}{T^2} = \frac{1}{6}f_\mu^2 + \frac{1}{3}f_\mu^2\left(\frac{v(T)}{T}\right)^2 \leq \frac{1}{2}f_\mu^2 = 1.8\times 10^{-7}\,. \tag{3.3}$$

The resulting baryon-to-photon ratio in this case is

$$\eta = (1.8 \pm 1.68)\times 10^{-10}\,. \tag{3.4}$$

This should be compared to the preferred value from BBN, e.g. [23] $\eta = (4.0^{+1.5}_{-0.9})\times 10^{-10}$. Thus we find agreement with the required value at the upper edge of the 95% CL region![3]

Notice that, if the electroweak phase transition is strongly first order with $\langle\phi(T)\rangle/T$ larger than unity after the transition, the sphaleron processes are frozen-out and absent after the transition. In this case the chemical equilibrium before the transition determines the baryon asymmetry. The only flavor-dependent effects before the transition are the Yukawa interactions of electrons and muons with the uncondensed Higgs boson. In the case of vanishing total lepton asymmetry $L_e + L_\mu = 0$, we now find

$$\frac{\overline{m}_\mu^2(T)}{T^2} = \frac{1}{6}f_\mu^2 = \frac{\pi\alpha_W}{3}\frac{m_\mu^2(0)}{m_W^2} = 6.0\times 10^{-8}\,. \tag{3.5}$$

This translates into baryon and lepton to entropy ratios of

$$B = \frac{6}{13\pi^2}(6.0\times 10^{-8})L_e^\star = (8.6\pm 8.0)\times 10^{-12}\,. \tag{3.6}$$

This corresponds to a current baryon-to-photon ratio of $\eta = (6.0 \pm 5.6)\times 10^{-11}$, which is off by more than a factor of three.

[3] If one instead use only ^{7}Li and ^{4}He abundances, the preferred region could be as low as $\eta = (1.65\text{–}2.06)\times 10^{-10}$ at 68% CL [45].

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4. Model building

We have seen that the preferred value of L_e from eq. (2.6), $L_e^\star = 2L_e^{\rm NUC}$, together with the relation $L_e = -L_\mu$ gives the correct order of magnitude and sign for the baryon asymmetry. We find this a remarkable coincidence.

Suppose however one takes the preferred value of the lepton asymmetry to be $L_e^{\rm NUC}$, i.e. let us not allow any room for extra degrees of freedom during nucleosynthesis. Then from eq. (3.2) the baryon asymmetry turns out to be correct except for a factor of two or so. A natural question then is if there are corrections that can fix this factor-of-two discrepancy so that the generation of the observed small baryon asymmetry from the magnitude of the lepton asymmetry currently preferred from the BBN is a realistic possibility. We find that there are many ways to achieve this. Another natural question is if there is an appropriate leptogenesis mechanism which can create a large lepton asymmetry with $L_e = -L_\mu$ in a simple way.

The simplest possibility to enhance the baryon asymmetry is to consider the MSSM where all sleptons and squarks are heavier than the electroweak phase transition temperature while the entire Higgs sector, h^0, H^0, A^0 and $H^\pm$ is light. In the approximation where one ignores their masses, the lepton doublets interact only with the H_d doublet with the Yukawa coupling $f_l/\cos\beta$. Here $\tan\beta \equiv \langle H_u\rangle/\langle H_d\rangle$ is the vacuum angle. In this limit, the only change from the case of the SM is to replace the Yukawa couplings f_l by $f_l/\cos\beta$, which enhances the plasma mass effects. The net result is an enhanced baryon asymmetry which brings the predicted value into the required range for a moderate value of $\tan\beta$. If the masses of the Higgs bosons cannot be neglected, the enhancement effect is reduced. But it is clear that a realistic value of the baryon asymmetry can be easily achieved.

There are many other possible enhancement mechanisms of the baryon asymmetry. For instance, light higgsinos and sleptons also contribute to the plasma mass of the lepton doublets. These effects are enhanced by $1/\cos^2\beta$ but Boltzmann-suppressed by their masses $\sim e^{-m/T}$. For suitable values of $\tan\beta$ and the slepton and higgsino masses estimates indicate that the required factor of two is generated. (A detailed quantitative analysis involves the generalization of the formulae in [40] to include the individual lepton asymmetries.) It is therefore clear that there are quite simple extensions of the SM which fairly naturally provide the required factor of two.

We now turn to the question of whether it is possible to generate a large lepton asymmetry with $L_e = -L_\mu$ in a natural and elegant fashion. Such leptogenesis with $L_e = -L_\mu$ can be achieved naturally by utilizing the Affleck-Dine mechanism [46].[4]

[4] It was discussed recently also in [47] that one can generate a large lepton asymmetry by the Affleck-Dine mechanism. The author however required an even larger asymmetry than what we discuss to keep the electromagnetism as well as sphalerons out of equilibrium to solve the monopole problem and avoid the overproduction of baryon asymmetry [33, 34, 35, 36, 38]. Therefore the aim

This requires the operator [48]

$$\int d^4\theta (m_{3/2}\theta^2)(m_{3/2}\bar{\theta}^2)\frac{L_e^* L_\mu H_u^* H_u}{M_X^2}\,, \tag{4.1}$$

where the supersymmetry-breaking spurions are inserted. This operator preserves the total lepton number, while breaking L_e and L_μ individually. The energy scale of this operator, M_X, can be, for example, the (reduced) Planck scale $M_* = 2\times 10^{18}\,\text{GeV}$. The D-flat direction $|L_e|^2+|L_\mu|^2 = |H_u|^2$ is lifted by this operator and the field acquires a large "angular momentum" in the internal space. This corresponds to the generation of individual lepton numbers satisfying $L_e = -L_\mu$. This leads to an estimate of the lepton number,

$$L_e = -L_\mu \simeq \frac{\phi_0^4 T_{RH}}{m_{3/2}M_X^2 M_*^2}\,, \tag{4.2}$$

where T_{RH} is the reheating temperature of primordial inflation, $m_{3/2}$ is the typical mass of the sleptons and Higgs bosons, and ϕ_0 is the initial amplitude of the slepton expectation values. Even taking account of the constraint from gravitino overproduction $T_{RH} \lesssim 10^9\,\text{GeV}$[5] and $m_{3/2}\sim 1\,\text{TeV}$, the initial value of the amplitude can be relatively small $\phi_0 \gtrsim (10^{-3}L_e M_X M_*)^{1/2}$. Taking $M_X \sim M_*$, eq. (4.2) shows that $\phi_0 \sim 10^{15}\,\text{GeV}$ is sufficient to generate the large lepton asymmetry that we require. Note that the detailed mechanism for generating a large initial amplitude, ϕ_0, is model-dependent; it could be a negative mass-squared during the inflationary epoch [50, 51] or quantum effects [52].[6]

One final concern is if this scenario is consistent with the reported atmospheric neutrino oscillation: If the generated asymmetry L_μ is converted partially to an asymmetry in L_τ, it could then generate too large a baryon asymmetry because of the large tau Yukawa coupling f_τ. This fortunately does not happen. By the time of the electroweak phase transition, the probability for neutrino oscillation is suppressed by $\sin^2(\Delta m^2 t/4E_\nu)$, where $t\sim M_*/m_W^2$ and $E_\nu \sim m_W$. Substituting the relevant Δm^2 into this expression then shows that the oscillation to L_τ is negligible.

of the paper is orthogonal to ours.

[5]This bound is obtained considering the thermal production of gravitinos. However, it has been recently pointed out that this mechanism of productio is overcome by the non-thermal generation of gravitinos [49].

[6]Since the total lepton number is preferably conserved within our scenario, the neutrino masses should be Dirac rather than Majorana. The atmospheric neutrino oscillation prefers a small Yukawa coupling of order $h_\nu \lesssim 10^{-12}$. Even though this Yukawa coupling lifts our flat direction, a negative mass squared of, for instance, $-H_{inf}^2$ during inflation, generates an initial amplitude of $\phi_0 \sim H_{inf}/h_\nu$ which is well beyond what we need given the typical value of the Hubble constant during inflation $H_{inf}\sim 10^{11}$–$10^{13}\,\text{GeV}$. Such a small Yukawa coupling could be a natural consequence of a flavor symmetry [53].

5. Conclusion

Over the years, many mechanisms for the generation of the tiny observed baryon asymmetry have been proposed and we have very little idea which if any is the correct one. Furthermore, many of the proposed baryogenesis mechanisms are not able to predict the resulting baryon asymmetry to better than an order of magnitude (sometimes many). On the other hand, so far there is no observational evidence excluding the possibility that the lepton asymmetry in the Universe is almost as large as the present entropy density. On the contrary, the current measurements of the light element abundances may prefer such an asymmetry to reconcile BBN theory with observations. In this paper, we have made a simple observation which seems quite surprising to us: If the asymmetries in electron number and muon number are equal and opposite and of the size indicated by nucleosynthesis considerations, a baryon asymmetry of the observed size is naturally generated within the Standard Model itself due to the small but non-zero muon Yukawa coupling. This might just be a coincidence, but it is quite an intriguing one!

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