COSMOLOGICAL BARYON PRODUCTION VERSUS INTERMEDIATE MASS SCALES?

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Received 1 October 1981

We analyze the conditions for a late cosmological matter–antimatter asymmetry production at some intermediate energy scale $\sim 10^6$ GeV. We envisage a new class of models where the compatibility between enough ΔB production and decoupling of subsequent B violating interactions at $T \sim 10^6$ GeV is automatically achieved through the assumption of weak coupling between B nonconserving particles and ordinary matter. Possible consequences for nucleon stability are analyzed. The mechanism could be interesting for dynamical schemes (technicolor, composite models) where B production seems very difficult to be obtained before $T \sim \Lambda_H$. Furthermore a late origin of matter–antimatter asymmetry could be essential in the resolution of some cosmological puzzles, such as the excess of monopole density.

Grand unified theories for the first time provided a dynamical mechanism to explain the cosmological matter–antimatter asymmetry [1]. Starting from a symmetric universe (i.e., baryons and antibaryons in equal amount), baryon (B) violating interactions may cause a small excess of baryons over antibaryons which will survive after the subsequent matter–antimatter annihilation. A necessary condition for such a scenario to take place is that these B violating interactions are slow in comparison with the expansion rate of the universe. The implementation of this requirement in grand unified theories leads to a remarkable insight in the possible energy scale of grand unification: the masses of the superheavy Higgs and gauge bosons, responsible for B production in their decays, must be greater than $\sim 10^{13}$ GeV and $\sim 10^{16}$ GeV, respectively. Indeed, not only bosons with masses $< 10^{13}$ GeV could not give rise to any B asymmetry in their decays, but, even more dramatically, if they mediate B violating reactions, they would wipe out any preexisting B excess [2]. Can one conclude from these results that there exists a sort of no-go theorem for cosmological B production at T much less than the grand unification mass (GUM)? Our paper deals essentially with this question. We argue that it is indeed possible to find conditions in the GU models to allow B generation at intermediate mass scales (IMS), i.e. at scales such that: $M_W <$ IMS $\ll$ GUM. Before plunging into our analysis, we would like to remind the reader that the above question is far from being a mere academical curiosity. There are at least two motivations for such an investigation. First of all, the dynamical theories today on the market (technicolor, composite models, ...) show a common trend to introduce a scale Λ_H (far below the GUM), at which some new kind of force becomes strong, determining the low-energy physics we are testing. In these theories, in particular in the composite models, it is hard to conceive a B production at $T > \Lambda_H$. Thus, it becomes of utmost interest to work out the conditions for B generation at $T \lesssim \Lambda_H$. Secondly, there are some specific cosmological problems pointing in favour of a late creation of B asymmetry: in particular, a prolonged first-order phase transition could provide a strong dilution of the monopole density and offer a hope for some explanation of the flatness and horizon problems. Thus many authors have recently considered the possibility of a supercooling of the universe, with a sudden reheating when the phase transition is accomplished [3]. If this scenario takes place, the issue of cosmological B asymmetry becomes very delicate: the B production at GUM is enormously diluted by this increase of entropy ($n_B/n_\gamma \ll 10^{-9}$), so

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that some new B creation must arise. Now, if after the reheating T is somewhat lower than 10^{13} GeV, the usual mechanism is ruled out. Apart from this consideration, an early B production ($n_B/n_\gamma \sim 10^{-9}$ at $T \sim 10^{13}$ GeV) may always be jeopardized if, for some reason, some heavy or superheavy particle (boson or fermion) could not decay for a long time after becoming non-relativistic, thus dominating the whole energy density of the universe with its rest mass energy [4]. In summary, while it is true that the "classical" cosmology GUT's marriage is not compellingly in favor of a B production at IMS, it is also fair to say that it would be much safer to have a late B asymmetry.

We conceive of two different pictures for a B production at IMS: (i) B are produced through "non-renormalizable" interactions [4,5] or (ii) in the decay of weakly coupled particles.

The first option has been recently explored by Mohapatra and one of the authors [5]: they analyze the B production in the decays of right-handed neutrinos in the context of left–right-symmetric models at IMS. They conclude that the following recipe may ensure B generation at IMS: the processes responsible for matter–antimatter asymmetry must be given by effective *non-renormalizable* interactions scaling with T as T^n, $n \geqslant 5$, whereas particles whose decays violate B via *renormalizable* interactions must be heavy enough so that they are highly non-relativistic by the time ΔB creating processes fall out of equilibrium.

In this class of models one must be very careful about the compatibility of achieving enough ΔB and, at the same time, pushing out of equilibrium all the possible B violating interactions taking place after such a B production.

In the second option, this condition may be automatically met, as we shall illustrate in a specific example. The major characteristics of this second class of models with B production at IMS is that matter–antimatter asymmetry originates from B violating *renormalizable* interactions (two-body decays): the condition of being out of equilibrium is achieved through *weak couplings* of B violating particles to ordinary matter.

We want to illustrate this latter strategy in the following model. Let us introduce the right-handed partner of ν_L (we call it N) in the minimal SU(5) model [6]. Clearly N is a singlet of SU(5) and it couples at

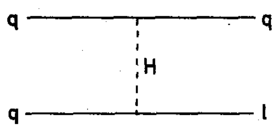

Fig. 1. Proton decay mediated by the H ≡ (3, 1, −1/3) Higgs scalar.

tree level only to the ordinary fermions constituting a **5**-plet of SU(5) through the **5**-plet of Higgs scalars. N possesses B violating decays as it presents two decay channels with different B number in the final state:

$$\mathrm{N} \to \ell_L \varphi \quad \text{(a)}, \qquad \mathrm{N} \to \bar{q} + \mathrm{H} \to \bar{q} + \bar{q} + \bar{q} \quad \text{(b)}, \qquad (1)$$

where $\ell_L \equiv (\nu, e)_L$, $\varphi \equiv (\varphi^+, \varphi^0)$ is the usual Weinberg–Salam Higgs doublet and H is the colored triplet contained in the **5**-plet of Higgs scalars.

We want to investigate the possibility of B production at IMS through N decays. For this reason, we take tentatively $m_{N_1} \sim 10^7$ GeV, where N_1 indicates the N of the first generation and we assume it to be the lightest of N's. We shall later on return to this choice. In (1), H must then be virtual: indeed, H mediates the proton decay process in fig. 1, so that $m_H > m_{N_1}$. However, we cannot get enough B number because of the suppression factor $(m_N/m_H)^6$ and the high power of Yukawa couplings (see ref. [7]). We must look for other Higgs bosons carrying nonvanishing B number and with the possibility of coupling to N. As an analysis of the transformation properties of the couplings N-ordinary fermions under SU(3) ⊗ SU(2) ⊗ U(1) immediately reveals, the only Higgs scalar (Φ) playing such a role carries (3, 1, 2/3) quantum numbers with respect to the above group. The Higgs Φ = (3, 1, 2/3) is contained in the representation **10** of SU(5): therefore, we introduce a 10-plet of Higgs scalars, thus allowing also for a coupling of N to the fermions contained in the **10** of SU(5). For our strategy it is essential that the condition $m_\Phi < m_N$ holds true: in this case, the two-body decay:

$$\mathrm{N} \to \bar{q} + \Phi \qquad (2)$$

is allowed. The Φ's will eventually decay into $\bar{d}\bar{d}$, so that the decay mode (2) carries $B = -1$. As the decay (1b) is extremely suppressed, the decays involved in B production are then (1a) and (2).

The condition of out of equilibrium for N_1 decays at $T \sim m_{N_1}$ gives the relation:

$$\Gamma_{tot} = (m_{N_1}/32\pi)[3(\hat{f}^+\hat{f})_{11} + 2(\hat{h}^+\hat{h})_{11}] < \sqrt{N} m_{N_1}^2/m_*, \quad (3)$$

where the h's are the Yukawa couplings in the process (1a), N represents the suitably weighted number of all particles with masses $< m_{N_1}$ ($\sqrt{N} \sim 10^2$), m_* is the Planck mass and the hat ^ indicates a matrix.

Taking $f \sim h$, (3) provides an upper bound on $(\hat{f}^+\hat{f})_{11}$:

$$(\hat{f}^+\hat{f})_{11} < 10^{-9}. \quad (4)$$

It is known that the Yukawa coupling responsible for the electron mass must be $\sim 10^{-6}$, so that relation (4) seems natural for the couplings of N_1.

Turning now to the B asymmetry calculation [8], the leading contribution is given by the interference in fig. 2 [7]:

$$\Delta B = (3/4\pi)\,\mathrm{Im}(f_{1i} f_{ij}^* f_{jk}^* f_{k1})\, k(m_{N_j}^2/m_{N_1}^2)/(\hat{f}^+\hat{f})_{11}, \quad (5)$$

with

$$k(x) = \sqrt{x}\{1 + (1+x)\ln[x/(1+x)]\}.$$

For simplicity, we have taken $m_\Phi \ll m_{N_1}$ in (5).

In analogy with the Yukawa couplings responsible for the fermion masses, we assume f_{33} to be the largest entry in the matrix $\hat{f}$. Thus, the main term in the sum (5) is:

$$\Delta B \simeq (3/8\pi) f_{33}^2 m_{N_1}/m_{N_3}, \quad (6)$$

provided that $m_{N_3} > m_{N_2}, m_{N_1}$ and $f_{13}^2 \sim (\hat{f}^+\hat{f})_{11}$. In (6), maximal *CP* violation is assumed. The constraint $\Delta B \sim 10^{-9}$ implies:

$$f_{33}^2/m_{N_3} \sim 10^{-15}\ \mathrm{GeV}^{-1}. \quad (7)$$

Taking a hierarchical mass pattern for the N's:

$$m_{N_1} : m_{N_2} : m_{N_3} = m_u : m_c : m_t, \quad (8)$$

which we shall motivate later on, we get $m_{N_3} \sim 10^{11}$ GeV and, thus, from eq. (7), $f_{33} \sim 10^{-2}$.

It is important to notice that our mechanism for B production is based on the assumption of weak couplings of N_1 to fermions, however, this does not imply a too small ΔB, as in the diagram of fig. 2 also the higher generation particles participate in the intermediate states and their couplings have no reason to be particularly small.

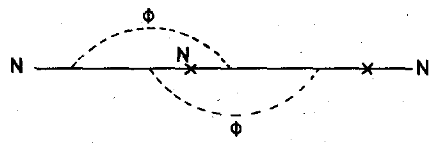

Fig. 2. The interference giving rise to the ΔB amount calculated in eq. (5). The cross denotes a Majorana mass of N.

In order to be sure that the B production in the N_1 decays can survive at later stages, we must ascertain that no B violating interactions are in equilibrium at $T < m_{N_1}$. Notice that at this temperature the channel $\Phi \to N + q$ is energetically forbidden, so that Φ has only one possible decay channel, $\Phi \to \bar{d}\bar{d}$. Thus, there is no risk of B dilution from Φ decays.

In fig. 3 we represent two typical processes with B violating interactions at $T < m_{N_1}$. The reaction rate is roughly $f_{33}^4 T^3/m_{N_3}^2 \sim (\Delta B)^2 T^3/m_{N_1}^2 < (\Delta B)^2 m_{N_1}$. Since $\Delta B \sim 10^{-9}$, clearly all these processes are completely out of equilibrium. Therefore, in our model there is an *automatic compatibility* between the requirements of producing enough ΔB and the decoupling of B violating interactions after B generation.

We turn now our attention to the Higgs Φ. As it mediates B violating interactions, one should worry about nucleon stability since its mass is very low ($m_\Phi < m_N \sim 10^7$ GeV) in comparison with GUM. The analysis of Φ-mediated nucleon decay will provide a lower bound on m_Φ. First of all, let us notice that the coupling matrix of Φ to $d_R d_R$, with d representing a $Q = -1/3$ quark, must be antisymmetric. Let us call g the coupling matrix of Φ to the quark physical states (mass eigenstates) [‡1]. Then the diagram in fig. 4 gives the leading contribution to Φ mediating proton decay. Notice that it is a $\Delta S = 1$ p-decay channel, $p \to K + \nu + X$: this is indeed the only possibility, since

[‡1] In fact, $d_R d_R$ is antisymmetric with respect to color and Lorentz indices, so that it must also be antisymmetric in generation indices to respect Fermi statistics. Clearly, both the coupling matrices of Φ to the quark current and mass eigenstates share the property of being antisymmetric. We shall deal only with the couplings of Φ to the quark physical states.

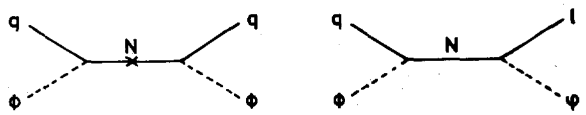

Fig. 3. Examples of B violating interactions taking place at $T < m_{N_1}$.

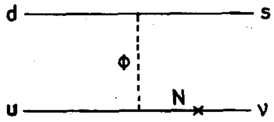

Fig. 4. Proton decay mediated by the $\Phi \equiv (3, 1, 2/3)$ Higgs scalar. Notice that Φ couples to the quarks u and d of the first generation and, thus, these couplings are small $\sim 10^{-5}$. The cross denotes a mixing between N and ν.

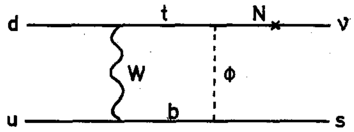

Fig. 5. Proton decay where Φ couples to the quarks of higher generations. The transition from the proton quark constituents to t and b is realized through a weak gauge boson (W) exchange.

g is an antisymmetric matrix and the decay into beauty particles is not energetically allowed. We must impose that the decay rate of the process in fig. 4 should respect the experimental bound on proton stability.

$$g_{12}^2 f_{1i}^2 (m_{\mathrm{D}}^{\nu}/m_{\mathrm{N}_i})^2 m_{\mathrm{P}}^5/m_{\Phi}^4 < 10^{-58}\,\mathrm{GeV}. \tag{9}$$

Assuming $f_{12}^2 \sim g_{12}^2 \sim (\hat{f}^+\hat{f})_{11} \sim 10^{-10}$ [‡2], we obtain:

$$m_{\Phi} > 10^4\text{–}10^5\ \mathrm{GeV}. \tag{10}$$

We emphasize that it is the need of the neutrino mass insertion in fig. 4 which allows for $m_{\Phi} < 10^7$ GeV. The diagram in fig. 1, which shows the analogous process mediated by H, does not have to present the suppression $(m_{\mathrm{D}}^{\nu}/m_{\mathrm{N}})^2 \sim 10^{-20}$, so that m_{H} must be $> 10^{12}$ GeV to guarantee enough p-stability. Since $f_{33} \sim 10^{-2}$ is much larger than the couplings in eq. (9), one could wonder whether it is even more dangerous than fig. 4 to take a process where Φ couples only to fermions of higher generations, whereas the transition from proton quark constituents (u and d) to t and b is realized through weak-interaction processes. This idea is graphically presented in fig. 5. Even taking the coupling of Φ to bs and tN of order f_{33}, the process in fig. 5 can dominate over the process in fig. 4 only if the mixing angles at the weak vertices are larger than 10^{-1}.

‡2 Strictly speaking, eq. (4) furnishes only an upper bound on $(\hat{f}^{\dagger}\hat{f})_{11}$, so that one could wonder whether it is possible to take f_{1i} arbitrarily small. We have a cosmological argument to prevent f_{1i} to be smaller than 10^{-5}. One should remember that at a temperature $T \sim m_{\mathrm{N}_1}/3N$, the N_1 rest mass energy, $\rho_{\mathrm{N}_1} \sim n_{\mathrm{N}_1} m_{\mathrm{N}_1}$, would start dominating the full energy density in the universe if the N_1's have not yet decayed [4]. This in turn would imply a reheating of the universe, when they eventually decay, up to a temperature T_* such that: $n_{\mathrm{N}_1} m_{\mathrm{N}_1} \sim 3NT_*^4/\pi^2$ with a consequent dilution in the B asymmetry by a factor: $d \sim (T_{\mathrm{D}}/T_*)^3$, where T_{D} indicates the decay temperature of N's. In order to avoid a substantial dilution, one has to take f_{1i} not much smaller than 10^{-5}.

Incidentally, we observe that Φ can also mediate neutron–antineutron oscillations (fig. 6): owing to the g antisymmetric properties one needs at least two W-exchanges, so that the process becomes completely negligible. Other effects of Φ, such as the peculiar b decay in fig. 7, are again enormously suppressed with respect to the usual weak-interaction processes.

We would like to close with a remark on the mass scales we have introduced in our model. As no intermediate partial unifying group is present here, symmetry breaking is connected with only two fundamental mass scales, $\mathrm{M_W}$ and GUM [6]:

$$\mathrm{SU(5)} \xrightarrow[\mathrm{GUM}]{} \mathrm{SU(3)_c} \otimes \mathrm{SU(2)} \otimes \mathrm{U(1)}_y \xrightarrow[M_{\mathrm{W}}]{} \mathrm{SU(3)_c} \otimes \mathrm{U(1)_{em}}.$$

For the mass $m_{\Phi} \sim 10^5$ GeV we can always invoke a fine tuning of the parameters in the Higgs potential similarly to what is done to bring m_{φ} down to M_{W} [6]. What appears more awkward is the presence of the N masses at an intermediate scale $\sim 10^7$–10^{12} GeV. As the N's are singlets of SU(5), they get mass directly and one would naturally suppose the dimensional parameter M associated with this mass term of the order GUM: just to use the appropriate name, one would say that the "survival hypothesis" [9] naturally pushes m_{N} to be at GUM. This situation is avoided once our SU(5) model is embedded into a SO(10) scheme [10] breaking to SU(5) at some scale $M_{\mathrm{SO(10)}} > \mathrm{GUM}$. Since the N's are no longer singlets

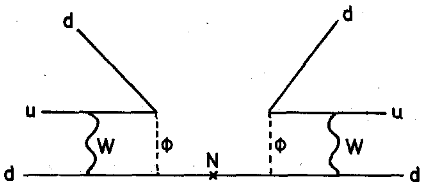

Fig. 6. Highly suppressed neutron–antineutron oscillations mediated by Φ.

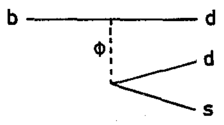

Fig. 7. Exotic b decay mediated by Φ. This process is much slower than the usual b decay mediated by weak gauge bosons.

of SO(10), they cannot be given a direct mass, but two ways are available to provide their mass: either one introduces a **126** Higgs representation of SO(10), so that the N's can get mass from the vacuum expectation value of this Higgs scalar [11], or one can induce only radiatively a mass to N's [12]. This latter option provides a hierarchical pattern in the N sector, with masses of order [12]:

$$m_{N_i} \simeq (\alpha/\pi)^2 (m_{q_i}/M_W) M_{SO(10)}$$

where q_i is the $Q = 2/3$ quark of the ith generation. For $M_{SO(10)} \sim 10^{18}$ GeV, one gets:

$$m_{N_1} \simeq 10^8 \text{ GeV}, \quad m_{N_2} \simeq 10^{11} \text{ GeV},$$

$$m_{N_3} \simeq 10^{12} \text{ GeV}.$$

Our essential $\Phi \equiv (3, 1, 2/3)$ Higgs scalar is not contained in the **10** representation of SO(10), so that one should think at a SO(10) model a la Witten with a **10** and a **120** (or two **120**'s) to provide with their vev's mass to all the fermions, except for the N's which would get a radiatively induced mass.

In conclusion, we have presented a model where an alternative strategy to produce a cosmological B asymmetry at IMS is pursued: baryons are produced in renormalizable interactions (two-body decays) through the decay of particles which are weakly coupled to the ordinary matter. Our observation allows for a safe B production in grand unified models with a supercooling in the early universe and it might shed new light on the difficult issue of B generation in composite models [13].

We wish to acknowledge stimulating discussions and useful comments from R.N. Mohapatra, S. Nussinov, R.D. Peccei and M. Yoshimura. We also thank R.D. Peccei for reading the manuscript.

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