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More Flipped $SU(5) \times U(1)$ Baryosynthesis

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ABSTRACT

We supplement a previous discussion of baryosynthesis in flipped $SU(5) \times U(1)$ GUTs by including 1) the large incoherent field energy density which is likely when $SU(5)$ is broken, and 2) the possibility of additional Higgs triplet fields suggested by four-dimensional string model-building. We consider strong (weak) reheating scenarios in which the Universe is (is not) $SU(5)$ symmetric after inflation. We find an adequate baryon asymmetry subsequent to strong reheating, whatever the number of Higgs triplets (although beware of possible difficulties with quasi-stable relic particles), whereas weak reheating requires at least two Higgs triplets.

Flipped supersymmetric $SU(5) \times U(1)$ is a very appealing GUT.¹⁾ It has a natural and economical mechanism for splitting doublet Higgses from their triplet partners, a naturally long-lived proton ($\tau_p \gtrsim 10^{33} \text{ y}$), and a seesaw mechanism yielding light left-handed neutrinos. Moreover, it is the only known GUT that can be obtained by manifold compactification of the ten-dimensional heterotic string,²⁾ or from formulations of string models directly in four dimensions.³⁾ Flipped $SU(5) \times U(1)$ also has cosmological advantages over other string models, in that strong coupling effects⁴⁾ enable $SU(5)$ to be broken as the Universe cools and expands without an a priori excess of entropy.⁵⁾ To confirm the acceptability of the resulting baryon-to-entropy ratio, one must investigate baryosynthesis in the model and determine the magnitude of the baryon asymmetry that it can produce. This subject was studied in a previous paper,⁶⁾ in which two scenarios were distinguished. In one ("strong reheating"), it was assumed that after inflation the Universe reheated sufficiently to restore $SU(5)$, whereas "weak reheating" means that $SU(5)$ is always broken subsequent to inflation. In both cases, the out-of-equilibrium decays of a massive $SU(3)_C \times SU(2)_L \times U(1)_Y$ singlet ϕ_m were identified as the dominant source of a baryon asymmetry in the minimal $SU(5) \times U(1)$ model.

In this paper we re-examine baryosynthesis in flipped $SU(5) \times U(1)$, introducing two new elements. One is the fact that in the case of strong reheating there is likely to be a large incoherent field energy density when $SU(5)$ is subsequently broken. This produces less entropy than was estimated previously⁶⁾ on the basis of the coherent field oscillations dominant in weak reheating. The other new element in this paper is the observation that realistic $SU(5) \times U(1)$ models extracted³⁾ from fermionic formulations of strings directly in four dimensions contain a non-minimal set of Higgs fields. As in

the cases of minimal SU(5) and its supersymmetric extension,⁷⁾ these additional Higgses provide new mechanisms for baryosynthesis which are likely to be more efficient than those considered previously⁶⁾ in the minimal SU(5)xU(1) model. We find that strong reheating, in which SU(5)xU(1) symmetry is restored subsequent to inflation, may be a viable scenario with or without additional Higgses, modulo other difficulties with the overproduction of quasi-stable relic particles. Weak reheating, in which SU(5)xU(1) symmetry is not restored, is viable if there are additional Higgses.*

The minimal SU(5)xU(1) model^{10,1)} contains three generations of matter fields with the transformation properties

$$F_i = (\underline{10}, 1/2), \bar{F}_i = (\bar{\underline{5}}, -3/2), l_i^c = (\underline{1}, 5/2) \quad (1)$$

the following Higgs representations

$$H = (\underline{10}, 1/2), \bar{H} = (\bar{\underline{10}}, -1/2), h = (\underline{5}, -1), \bar{h} = (\bar{\underline{5}}, 1) \quad (2)$$

and four SU(5)xU(1) singlet fields $\phi_m = (\underline{1}, 0)$. The superpotential of this minimal model is

$$\begin{aligned} W = & \lambda_1^{ij} F_i F_j h + \lambda_2^{ij} F_i \bar{F}_j \bar{h} + \lambda_3^{ij} \bar{F}_i l_j^c h \\ & + \lambda_4 H H h + \lambda_5 \bar{H} \bar{H} \bar{h} \\ & + \lambda_6^{im} F_i \bar{H} \phi_m + \lambda_7^{mh} h \bar{h} \phi_m + \lambda_8^{mnp} \phi_m \phi_n \phi_p \end{aligned} \quad (3)$$

* We do not discuss the Affleck-Dine⁸⁾ mechanism, which exploits coherent squark and slepton energy, because we only expect incoherent field energy density in the case of strong reheating, whilst temperatures large enough to avoid excess entropy generation in the case of weak reheating would in fact destroy coherent squark and slepton fields.⁹⁾

and the GUT symmetry is broken by a large v.e.v. for a combination Φ of H and $\bar{H}$ along a D- and F-flat direction of the supersymmetric potential. This flatness is broken by soft supersymmetry-breaking mass terms of magnitude $\tilde{m} \sim 10^2$ to 10^3 GeV, which also enable the $h, \bar{h}$ and ϕ_m to acquire small v.e.v.'s of order $\tilde{m}$. The λ_1 couplings in (3) generate masses for charge $-1/3$ quarks, the λ_2 couplings for charge $+2/3$ quarks, and the λ_3 couplings for charged leptons. The λ_4 and λ_5 couplings provide the natural doublet-triplet splitting which is a major virtue of this flipped $SU(5) \times U(1)$ model. The λ_6 coupling provides the seesaw mechanism for neutrino masses and λ_7 provides the desired $h\bar{h}$ mixing. (For more details, see Ref. 1.)

This flipped $SU(5) \times U(1)$ GUT is unique in avoiding⁴⁾ the following cosmological pitfall⁵⁾ for models with D- and F-flat potentials. Since any such flat direction is flat to $O(\tilde{m}^2)$, one might expect finite temperature effects to keep the model stuck with zero v.e.v.'s until the temperature $T \lesssim \tilde{m}$. When released, the model would have large coherent field oscillations which would dominate the energy density and produce an entropy excess $\Delta \gtrsim 10^6$ unless the final v.e.v. is $\lesssim 10^7$ GeV.⁵⁾ However, in the $SU(5) \times U(1)$ model, strong coupling effects¹¹⁾ destabilize the origin in field space at a temperature $T \sim \Lambda_c$ (the strong coupling scale) and relax this constraint on the v.e.v. sufficiently to allow $\langle \Phi \rangle \sim 10^{15}$ GeV as required.⁴⁾ Such a destabilization mechanism is essential in any scenario involving strong reheating after inflation, which we now study.

Strong Reheating?

In such a strong reheating scenario, the Φ field starts growing once the temperature T falls below Λ_c . We have previously⁶⁾ considered the possibility that coherent oscillations of Φ around its true minimum at $\langle \Phi \rangle = M_G \sim 10^{15}$ GeV

would dominate the mass density of the Universe. However, if $\Lambda_c > \sqrt{m_\Phi M_G} \sim 10^9$ GeV as we expect, and as we found necessary to avoid excess entropy generation from the decay of Φ oscillations,⁶⁾ then we should expect the $O(\Lambda_c^4)$ incoherent field energy of the Φ field to dominate and destroy any coherent component $\sim m_\Phi^2 M_G^2$ that the Φ field might possess. Assuming incoherent Φ fluctuations dominate the energy density, the Universe remains radiation-dominated until the temperature T_Φ of the Φ field falls below m_Φ , and is then matter-dominated until the non-relativistic Φ particles decay at a temperature $T_\Phi = T_D$. This occurs when the Hubble expansion rate $H = \sqrt{8\pi\rho_\Phi/3}/M_P : \rho_\Phi = \zeta(3)/\pi^2 m_\Phi T_\Phi^3$ obeys

$$H = \Gamma_\Phi = \epsilon m_\Phi^3 / M_G^2 : \epsilon = \frac{9\lambda_{1,2,3\text{or}7}^4}{2048\pi^5} \quad (4)$$

The reheating temperature T_D' produced by these decays is

$$T_D' = \left(\frac{3}{8\pi\tilde{N}_{\text{MeV}}} \right)^{1/4} \frac{\epsilon^{1/2} m_\Phi^{3/2} M_P^{1/2}}{M_G} \approx 50 \text{ MeV } \epsilon^{1/2} \left(\frac{m_\Phi}{1 \text{ TeV}} \right)^{3/2} \left(\frac{10^{15} \text{ GeV}}{M_G} \right) \quad (5)$$

where $\tilde{N}_{\text{MeV}} = 43\pi^2/120$. Conventional cosmological nucleosynthesis requires $T_D' \gtrsim 1 \text{ MeV}$, which imposes a non-trivial constraint on the couplings $\lambda_{1,2,3\text{or}7}$. If we take $m_\Phi = 10^3 \text{ GeV}$ and $M_G = 10^{15} \text{ GeV}$, then $\lambda_{1,2,3\text{or}7} \gtrsim 2$ or 3. Since low-energy phenomenology imposes $\lambda_{1,2,3} \lesssim 1$, this would imply that the Φ decays predominantly into the light ϕ_m field. To determine the entropy produced by these decays, we compare T_D' with the radiation temperature in the absence of Φ decays. Recalling that Φ is ultra-weakly-interacting and hence not in equilibrium with the rest of the Universe once $T < \Lambda_c$, and allowing for thresholds between 1 MeV and Λ_c , we obtain

$$\Delta = \frac{\tilde{N}_{\text{MeV}} T_D^3}{\tilde{N}_c T_D^3} = \left(\frac{8\pi\tilde{N}_{\text{MeV}}}{3} \right)^{\frac{1}{4}} \frac{\zeta(3)}{\pi^2 \tilde{N}_c} \frac{M_G}{\epsilon^{\frac{1}{2}} m_{\Phi}^{\frac{1}{2}} M_P^{\frac{1}{2}}} \quad (6)$$

where $\tilde{N}_c = (\pi^2/30)N_c$ describes the number of effective degrees of freedom at $T = \Lambda_c$. Note that Δ is independent of Λ_c apart from possible weak dependence through N_c . A typical estimate for $m_{\Phi} = 10^3$ GeV, $M_G = 10^{15}$ GeV and $\epsilon = 10^{-3}$ is $\Delta \sim 10^3$. This estimate of Δ is valid in any strong reheating scenario if $\Lambda_c > \sqrt{m_{\Phi} M_G}$ so that incoherent Φ oscillations dominate and we avoid excessive entropy generation from the decay of coherent Φ oscillations.

To judge whether this estimate of Δ is acceptable, we must discuss the possibilities for generating a baryon asymmetry in the flipped $SU(5) \times U(1)$ model. Previously,⁶⁾ we have pointed out that the massive combination of ν_H^c and $\bar{\nu}_H^c$ orthogonal to Φ cannot generate any baryons, and that, as in minimal $SU(5)$ models,¹²⁾ the decays of Higgs triplets in a minimal model with only one conjugate pair of $h, \bar{h}$ multiplets generate too small a baryon asymmetry. Accordingly, we focussed attention on decays of the massive ϕ_m fields, which yield a baryon asymmetry

$$B\epsilon' = \frac{\lambda_1 |\lambda_2|^2 \lambda_5^* \lambda_6^* \lambda_7^* \lambda_8}{2^{11} \pi^5 (|\lambda_7|^2 + |\lambda_8|^2)} \left(\frac{m_{\phi}^{10}}{M_G^4 m_{h, \bar{h}}^6} \right) \quad (7)$$

This is maximized by taking $\lambda_6 \sim \lambda_4 \sim \lambda_5$ so that $m_{\phi} \sim m_{h, \bar{h}}$, and $\lambda_7 \sim \lambda_8$ in which case

$$B\epsilon' = \frac{\lambda_1 |\lambda_2|^2 \lambda_{4,5,6}^7}{2^{11} \pi^5} \quad (8)$$

We know that $\lambda_2^c \sim 1$ and $\lambda_1^b \sim 10^{-1}$, so to maximize (8) one must take $\lambda_{4,5,6}$ as large as possible.

This maximum asymmetry (8) is actually only achieved if the ϕ 's are out of equilibrium when they decay, which occurs if $K_\phi = \Gamma_\phi / (\dot{n}_\phi / n_\phi) < 1$, where the decay rate $\Gamma_\phi = (\lambda_7^2 / 8\pi) m_\phi$, and the equilibrium number density n_ϕ is dropping at a rate of order $\dot{m}_\phi / m_\phi$ and/or $-\dot{T}_\phi / T_\phi$. At the strong coupling phase transition when $T \sim \Lambda_c$, the magnitude of m_ϕ and hence Γ_ϕ is set by $\lambda_6 \Phi \sim \Lambda_c$. The value of $\dot{m}_\phi / m_\phi = \dot{\Phi} / \Phi$ depends on the details of the transition from the confined phase to the Higgs phase. In the neighborhood of the strong phase transition it is reasonable to assume that $\dot{\Phi} / \Phi \sim \Lambda_c$. This is much larger than the Hubble expansion rate which $\sim \Lambda_c^2 / M_P$. Therefore, we also expect $\dot{T}_\phi / T_\phi \sim \dot{m}_\phi / m_\phi \sim \dot{\Phi} / \Phi \sim \Lambda_c$. In this case, $K_\phi \sim \lambda_6 \lambda_7^2 / 8\pi \lesssim 1$ as required, and the baryon-to-entropy ratio

$$n_B/s = \frac{\epsilon' B n_\phi}{s_c \Delta} = \frac{\epsilon' B (\zeta(3)/\pi^2)}{4/3 \tilde{N}_c \Delta} \quad (9)$$

Combining equations (6), (8) and (9) we obtain

$$n_B/s = 3/4 \left(\frac{3}{8\pi \tilde{N}_{\text{MeV}}} \right)^{1/2} \epsilon' B \frac{\epsilon^{1/2} m_\phi^{1/2} M_P^{1/2}}{M_G} \quad (10)$$

which may be evaluated as

$$n_B/s = 10^{-13} \lambda_{4,5,6}^7 \lambda_7^2 \left(\frac{m_\phi}{1 \text{ TeV}} \right)^{1/2} \left(\frac{10^{15} \text{ GeV}}{M_G} \right) \quad (11)$$

To get $n_B/s = 10^{-9}$ to 10^{-11} we need $\lambda_{4,5,6,7,8} \approx 2$ or 3 , which is consistent with our previous requirement on λ_7 .

An alternative mechanism for baryosynthesis exists in non-minimal $SU(5) \times U(1)$ models³⁾ with additional $\underline{10}$, $\overline{\underline{10}}$, $\underline{5}$ and $\overline{\underline{5}}$ Higgs fields, which we denote by $H, \bar{H}, h, \bar{h}$. If they have superpotential couplings of the form $\lambda'_4 H \bar{H} h'$ and $\lambda'_5 \bar{H} \bar{H} \bar{h}'$, the natural doublet-triplet splitting mechanism of the $SU(5) \times U(1)$ model will also furnish large masses for the Higgs triplets $(d_{H^c}^c, D')$ and $(\bar{d}_{\bar{H}^c}^c, \bar{D}')$ while leaving the h and $\bar{h}$ doublets light. The additional Yukawa couplings and phases associated with these extra fields permit a more efficient mechanism for generating a baryon asymmetry via $\bar{h}$ decays.⁷⁾ The leading two-body decay diagrams are shown in Fig. 1: (a) gives $q\bar{q}$ and $q\bar{l}$ final states, while the baryon-number-conserving diagrams (b) and (c) may be suppressed if $m_{\phi_m}, m_h \gtrsim m_{\bar{h}}$. The necessary CP violation ϵ' is provided by the interference between Fig. 1a and the loop diagram involving the extra Higgs fields $\bar{H}'$ and $\bar{h}'$ shown in Fig. 2. The requirement that Fig. 2 contains an absorptive part implies $m_{\phi_m} < m_{\bar{h}}$ and hence $\lambda_6 < \lambda_5$. The interference between Figs. 1a and 2 gives

$$\epsilon' B \approx \frac{1}{8\pi} \frac{\lambda_6^* \lambda_6' \lambda_2'}{\lambda_2} \left(\frac{\lambda_5}{\lambda_5'} \right) \quad (12)$$

where the final factor of (λ_5/λ_5') reflects a possible suppression occurring if $m_{\bar{h}'} > m_{\bar{h}}$. In the strong reheating scenario, n_B/s is computed by substituting (12) into (9). Assuming $\lambda_i \sim \lambda_i'$, we find

$$n_B/s \approx 5 \times 10^{-9} \lambda_6^2 \lambda_{1,2,3 \text{ or } 7}^2 \left(\frac{m_{\phi}}{1 \text{ TeV}} \right)^{1/2} \left(\frac{10^{15} \text{ GeV}}{M_G} \right) \quad (13)$$

where we have previously argued that $\lambda_{1,2,3 \text{ or } 7} \gtrsim 2$ or 3 in order that $T_D' \gtrsim 1 \text{ MeV}$. Hence a baryon asymmetry $n_B/s \sim 10^{-9}$ to 10^{-11} can apparently be obtained for reasonable values of λ_6 .

However, two potential difficulties with extra Higgs generations in $SU(5) \times U(1)$ should be noted. The first is that the additional light Higgs doublets in the renormalization group equations can lead to unacceptable values of M_G and/or $\sin^2 \theta_W$.¹³⁾ The second is that there are color-triplet, weak-doublet components $(u, d)_{H'}$ and $(\bar{u}, \bar{d})_{\bar{H}'}$ of the H' and $\bar{H}'$ multiplets that are light and could lead to unacceptably rapid proton decay¹⁴⁾ if some possible Yukawa couplings are not disallowed. Both of these difficulties would be alleviated by a large Dirac mass term of the form $M_{H, H'} \bar{H}'$. Clearly, it would give large masses to all components of the H' and $\bar{H}'$ multiplets and so prevent rapid proton decay via the $(u, d)_{H'}$ and $(\bar{u}, \bar{d})_{\bar{H}'}$. As for the first difficulty, it would provide a seesaw mass matrix for the $d_H^C, \bar{d}_H^C, D'$ and $\bar{D}'$ triplet fields, yielding one relatively light $D' \bar{D}'$ eigenstate with mass $\sim \lambda_4' \lambda_5' M_G^2 / M_{H'}$ which counteracts the effects of the extra light Higgs doublets in the renormalization group equations.^{* 13)} Just such a mass term appears in the recently derived four-dimensional string version³⁾ of the $SU(5) \times U(1)$ GUT. Extra Higgs fields $h, \bar{h}$ appear in the massless sector along with additional massive $H, \bar{H}$ fields with masses $M_H = O(M_P)$. In the presence of one extra pair of massive $H, \bar{H}$ multiplets, the estimate (12) of $\epsilon' B$ is modified to

$$\epsilon' B \approx \frac{1}{8\pi} \frac{\lambda_6^* \lambda_6' \lambda_2'}{\lambda_2} \left(\frac{\lambda_5 \lambda_5' M_G^2}{M_P^2} \right) \quad (12')$$

which is no problem for the strong reheating scenario discussed above.

* Here we must to avoid rapid proton decay through dimension-5 operators,^{14,15)} which could result from the Dirac mass term $M D' \bar{D}'$.

Thus far, strong reheating appears to be a viable scenario in flipped $SU(5) \times U(1)$, but there is a serious phenomenological difficulty which causes us to reconsider weak reheating. This problem is caused by $\tilde{\Phi}$, the supersymmetric partner of the Φ , which is ultra-weakly-interacting after the GUT symmetry is broken and has a lifetime much longer than the age of the Universe if it is the lightest supersymmetric particle (LSP).¹⁶⁾ In a strong reheating scenario, the number density of the $\tilde{\Phi}$ is large, so that the Universe is matter-dominated during cosmological nucleosynthesis and ultimately becomes overdense.¹⁶⁾ This difficulty may be avoided by having some other sparticle be the LSP, e.g., the more familiar photino $\tilde{\gamma}$ or Higgsino $\tilde{H}$. Then the absence of an exact R parity in the $SU(5) \times U(1)$ model would permit the $\tilde{\gamma}/\tilde{H}$ to decay into light particles prior to nucleosynthesis if it were heavy enough.¹⁶⁾ However, this would require a somewhat heavy sparticle spectrum, which might be unnatural.

Weak Reheating?

In view of the above constraint on strong reheating, we now re-examine the feasibility of baryosynthesis in a weak reheating scenario for flipped $SU(5) \times U(1)$. In this case, the gauge symmetry is not restored after inflation and $T_R < m_{\phi, \bar{h}}$, so that the ϕ 's or $\bar{h}$'s responsible for the baryon asymmetry are not produced in thermal equilibrium, but must be produced by inflaton decays $\psi \rightarrow \phi$ or $\bar{h}$, implying that $m_\psi > m_\phi$ or $m_{\bar{h}}$. Hence the resultant value of n_B depends on n_ψ , the number of inflatons present, given by $n_\psi = \rho_\psi / m_\psi : \rho_\psi = \tilde{N}_R T_R^4$. Since the entropy produced in ψ decay is just $s = 4/3 \tilde{N}_R T_R^3$, we have⁶⁾

$$n_B/s = 3/4 \epsilon' B (T_R/m_\psi) \frac{1}{\Delta} \quad (14)$$

where $\epsilon' B$ is given by equation (8) in the case of ϕ decays, or equation (12)

or (12') in the case of $\bar{h}$ decays. If, following inflation, the effective temperature of the Φ field $T_\Phi < \sqrt{m_\Phi M_G}$, then the entropy enhancement factor Δ in (14) results from the decay of coherent Φ oscillations, and is given by⁴⁾

$$\Delta_{\text{coh}} = \frac{\bar{N}_{\text{MeV}} T_D'^3}{\bar{N}_R T_T^3 (R_R/R_D)^3} = \left(\frac{8\pi \bar{N}_{\text{MeV}}}{3} \right)^{\frac{1}{2}} \frac{m_\Phi^{\frac{1}{2}}}{\bar{N}_R \epsilon^{\frac{1}{2}} M_P^{\frac{1}{2}}} \frac{M_G^3}{T_R^3} \quad (15)$$

where T_D' is given by equation (5) and $(R_R/R_D)^3$ is computed by setting $H = \Gamma_\Phi$ with $\rho_\Phi = m_\Phi^2 M_G^2 (R_R/R_D)^3$. Equation (15) is relevant to low reheating temperatures $T_R < \sqrt{m_\Phi M_G} \sim 10^9$ GeV or if the Φ field is not reheated by decays of the inflaton ψ . Since Φ is weakly interacting and out of thermal equilibrium with the other fields, $T_\Phi < \sqrt{m_\Phi M_G} < T_R$ is a consistent possibility. If, on the other hand, $T_\Phi > \sqrt{m_\Phi M_G}$ (which is still consistent with weak reheating if $T_R < \lambda \Lambda_c$, where λ is the Yukawa coupling giving mass to the lightest superheavy particle), then Δ results from incoherent Φ oscillations and is given by (6).

Then from (14) we have

$$n_B/s = 3/4 \epsilon' B \left(\frac{T_R}{m_\psi} \right) \left[\Delta_{\text{coh}} (T_\Phi < \sqrt{m_\Phi M_G}) \text{ or } \Delta_{\text{inc}} (T_\Phi > \sqrt{m_\Phi M_G}) \right]^{-1} \quad (16)$$

where it is understood that $\Delta \geq 1$.

For $\psi \rightarrow \phi$ decay, to obtain $m_\phi < m_\psi$ we need $\lambda_6 < m_\psi/M_G$ in equation (8) for $\epsilon' B$, so that (16) becomes

$$n_B/s \lesssim 1.2 \times 10^{-7} \left(\frac{m_\psi}{M_G} \right)^7 \left(\frac{T_R}{m_\psi} \right) \left[\Delta_{\text{coh}} (T_\Phi < \sqrt{m_\Phi M_G}) \text{ or } \Delta_{\text{inc}} (T_\Phi > \sqrt{m_\Phi M_G}) \right]^{-1} \quad (17)$$

If we assume $T_R/m_\psi \lesssim 1$ as for standard inflationary potentials,¹⁷⁾ then to get $n_B/s \gtrsim 10^{-11}$ we need $m_\psi \gtrsim 1/4 M_G$ ($m_\psi \gtrsim M_G$) if $T_\Phi < \sqrt{m_\Phi M_G}$ ($T_\Phi > \sqrt{m_\Phi M_G}$),

contrary to extant inflationary theory. We conclude that weak reheating is not viable in the minimal $SU(5) \times U(1)$ model.

For $\psi \rightarrow \bar{h}$ decay in a non-minimal model, n_B/s is obtained by substituting $\epsilon'B$ (12) into equation (16). Assuming that $\lambda_1 \sim \lambda_1'$, this yields

$$n_B/s \approx 3 \times 10^{-2} \lambda_6^2 \left[\frac{T_R}{m_\psi} \right] \left[\Delta_{\text{coh}} (T_\Phi < \sqrt{m_\Phi M_G}) \text{ or } \Delta_{\text{inc}} (T_\Phi > \sqrt{m_\Phi M_G}) \right]^{-1} \quad (18)$$

where $\lambda_6 < \lambda_5 < m_\psi/M_G$ is required for CP violation and for $\psi \rightarrow \bar{h}$ decay.

Assuming $T_R/m_\psi < 1$,^{*} we then have

$$n_B/s \lesssim 3 \times 10^{-2} \left[\frac{m_\psi}{M_G} \right]^2 \left[\Delta_{\text{coh}} (T_\Phi < \sqrt{m_\Phi M_G}) \text{ or } \Delta_{\text{inc}} (T_\Phi > \sqrt{m_\Phi M_G}) \right]^{-1} \quad (19)$$

To get $n_B/s \gtrsim 10^{-11}$ we need $m_\psi/M_G \gtrsim 3 \times 10^{-4}$ (5×10^{-4}) for $T_\Phi < \sqrt{m_\Phi M_G}$ ($T_\Phi > \sqrt{m_\Phi M_G}$), where we have used $\tilde{N}_R = 63\pi^2/8$ for the broken $SU(5) \times U(1)$ phase in evaluating Δ . This is not inconsistent with inflationary models, in which $m_\psi \lesssim 10^{-7} M_P$.¹⁷⁾

The situation is less satisfactory if the estimate (12') of $\epsilon'B$ (corresponding to a large Dirac mass term for $H, \bar{H}'$) is used in equation (16), since $\lambda_5 < m_\psi/M_G$ is required for $\psi \rightarrow \bar{h}$ decay to be kinematically allowed. However, even this model may be rescued by postulating more than one extra massive $H, \bar{H}'$ pair, as occurs in one variant of the model in Ref. 3. In this case, the intermediate-mass $(D, \bar{D}')$, $(D'', \bar{D}'')$, etc. mass eigenstates can yield a baryon asymmetry directly, analogous to standard $SU(5)$ with additional Higgses.

* It is clear from a comparison of equations (18) and (19) that one can get even larger values of n_B/s in a weak reheating scenario if one relaxes the inflationary expectation¹⁷⁾ that $T_R/m_\psi \leq 1$. This requires values of $T_R > \hat{\lambda} \Lambda_C$, which is consistent with weak reheating if $T_R \neq T_\Phi$ and $T_\Phi < \sqrt{m_\Phi M_G}$, in which case Φ can be trapped around its local minimum at $\Phi \sim M_G$.

We have considered various scenarios for generating a baryon asymmetry in flipped $SU(5) \times U(1)$ GUTs. In the absence of extra Higgses, we have found that acceptable baryosynthesis is only possible in a strong reheating scenario where the GUT symmetry is restored subsequent to inflation. Excessive entropy is not produced by Φ decays if $\Lambda_c > \sqrt{m_\Phi M_G}$. However, the overproduction of quasi-stable relic particles may be a serious problem for strong reheating. In the presence of extra Higgses, we have found that an acceptable baryon asymmetry may result from the decay of $\bar{h}$ color-triplet Higgses in both weak and strong reheating scenarios, although the latter still suffers from the above problem of excess relic particles.

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Figure Captions

Figure 1. The dominant two-body decay channels for $\bar{h}$ decay.

Figure 2. One-loop diagram contributing to a net baryon asymmetry from $\bar{h}$ decay.

Figure 1

Figure 2