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Baryogenesis in Three-Generation Calabi-Yau Superstring Models

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Abstract

Baryogenesis in the context of 3-generation Calabi-Yau superstring models is discussed and a concrete example is presented in which an acceptable baryon asymmetry is generated through the out-of-equilibrium decay of 'heavy' particles at temperatures $\sim 1\text{GeV}$.

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One of the successes of the hot-big-bang cosmology is a rather accurate prediction of the abundances of the light elements under the assumption that the ratio of the net baryon number density over the entropy density n_b/s , usually called the baryon asymmetry, is of order $10^{-10} - 10^{-11}$ [1]. If one is not satisfied with merely postulating this value as an initial condition, one must investigate the necessary conditions under which such an asymmetry will be generated in a universe that starts off in a matter-antimatter symmetric state. It has been known for a long time that the essential ingredients for baryogenesis include the existence of physical processes that violate baryon number, C and CP, and departure from thermal equilibrium [2]. C and CP were known to be violated in the $K^0 - \bar{K}^0$ system and departures from thermal equilibrium could have occurred during some early stages in the evolution of the universe. However, the development of concrete baryogenesis scenarios awaited the advent of grand unified models which reconciled the apparent stability of the proton with the violation of the symmetry associated with baryon number. Among these scenarios the one which became distinguished as the standard one is the out-of-equilibrium decay of heavy particles [3].

The purpose of the present note is to discuss an analogue to the standard baryogenesis scenario in the context of four-dimensional models arising from the compactification of the ten-dimensional heterotic superstring [4] on a six-dimensional Calabi-Yau (C-Y) space [5]. Scenarios of this type for C-Y superstring-inspired models have been considered earlier [6]. The one described here is a variant of one of the scenarios discussed in ref.[6]. The fundamental difference between the present work and ref.[6] is that in our model the compactification is made on a specific version of the only concrete 3-generation C-Y space known today [7]. As a result of this further specification the spectrum and the discrete or continuous symmetries (including the gauge group) of the four-dimensional model are uniquely determined by the internal space. This enormously constrains the freedom, essen-

tial in the models of ref.[6], to choose fields or symmetries at will. Here we have the reduced freedom to choose the (complex structure of the) internal space and to assume the symmetry breaking pattern of our preference (essentially exploiting our present ignorance of the exact four-dimensional superpotential).

A common feature of all models published so far based on Yau's 3-generation manifold is the necessity for rather high intermediate scales $M_I (\geq 10^{13} GeV)$, mainly in order to adequately suppress the proton decay rate by giving rather high masses to the leptoquarks which mediate proton decay processes. An immediate consequence of the mechanism that generates the intermediate scales in these models [8] is that the associated phase transition has a rather low critical temperature $T_c \sim M_*$, $\sim 10^{3\pm 1} GeV$ (M_* is the supersymmetry breaking scale in the visible sector) and generates an enormous amount of entropy [9]. The entropy increase factor $\Delta \sim (M_*/M_P)^{1/2}(M_I/T_c)^3$ ($M_P \sim 10^{19} GeV$ is the Planck mass). For $M_I \geq 10^{10} GeV$ one can immediately see that, assuming a value of n_b/s before the phase transition not larger than ~ 1 , the entropy produced will dilute it to a value below $10^{-10} - 10^{-11}$. Therefore, only scenarios that generate the baryon asymmetry after the completion of the intermediate scale phase transitions should be considered. For all such scenarios [6] the masses of some leptoquarks have to be rather low ($\leq 10^{10} GeV$) and at the same time these leptoquarks should participate in baryon number violating couplings involving ordinary 'light' quarks which are not very small. These requirements are not met by any of the models published so far based on the only concrete 3-generation C-Y manifold known today. They are satisfied in our model because the relatively light leptoquarks have couplings involving simultaneously ordinary 'light' quarks and ordinary 'light' leptons which are heavily suppressed.

The 3-generation C-Y space is defined as the complete intersection K_0 of three polynomial constraints in $CP^3 \times CP^3$, of bidegree (3,0), (0,3) and (1,1) respectively, divided out by a freely-acting order 3 discrete symmetry group. For the polynomial constraints we choose $P_1 = \sum_{i=0}^3 x_i^3 = 0$, $P_2 = \sum_{i=0}^3 y_i^3 = 0$ and $P_3 = \sum_{i=0}^3 c_i x_i y_i + c_5 x_3 y_2 = 0$, with x_i and y_i ($i = 0, \dots, 3$) the homogeneous coordinates of the two CP^3 's. The complex parameters c_j ($j = 0, \dots, 5$), which are assumed to be different from each other and from zero, define the complex structure and are only subject to the transversality condition $dP_1 \wedge dP_2 \wedge dP_3 \neq 0$. The freely-acting Z_3 symmetry G is generated by $g = \text{diag}(1, \alpha^2, \alpha, \alpha, 1, \alpha, \alpha^2, \alpha^2)$ acting on $[x_i, y_i]$ with $\alpha = \exp(2\pi i/3)$. The 3-generation space $K = K_0/G$ is invariant under the Z_3 symmetry R generated by $r = \text{diag}(1, 1, \alpha, \alpha, 1, 1, \alpha^2, \alpha^2)$ acting on $[x_i, y_i]$.

For the flux breaking we assume the embedding of the freely acting Z_3 into E_6 defined by the homomorphism $g \mapsto U_g \equiv (I_3, \alpha I_3, \alpha I_3) \in SU(3)_c \times SU(3)_L \times SU(3)_R$. At energies below the compactification scale $M_c \sim M_P$ the gauge group is $H \equiv SU(3)_c \times SU(3)_L \times SU(3)_R$. Under H the left-handed lepton (λ), quark (Q) and antiquark (Q^c) fields in a $\mathbf{27}$ of E_6 transform as follows: $\lambda = (1, \bar{3}, 3) = \begin{pmatrix} H^1 & H^2 & l \\ e^c & \nu^c & N \end{pmatrix}$, $Q = (3, 3, 1) = \begin{pmatrix} q \\ g \end{pmatrix}$, $Q^c = (\bar{3}, 1, \bar{3}) = (u^c, d^c, g^c)$. With the transformation of each field under the generator r of the honest Z_3 discrete symmetry R given in parenthesis, there are 9 lepton fields $\lambda_1(\alpha), \lambda_2(\alpha), \lambda_3(\alpha^2), \lambda_4(\alpha^2), \lambda_k(1)$ ($k = 5, \dots, 9$), 6 mirror lepton fields $\bar{\lambda}_1(1), \bar{\lambda}_2(1), \bar{\lambda}_3(\alpha), \bar{\lambda}_4(\alpha), \bar{\lambda}_5(\alpha^2), \bar{\lambda}_6(\alpha^2)$, 7 quark (antiquark) fields $Q_1(\alpha), Q_2(\alpha), Q_3(\alpha), Q_4(\alpha^2), Q_5(\alpha^2), Q_6(\alpha^2), Q_7(1)$ ($Q^c_1(\alpha), Q^c_2(\alpha), Q^c_3(\alpha), Q^c_4(\alpha^2), Q^c_5(\alpha^2), Q^c_6(\alpha^2), Q^c_7(1)$) and 4 mirror quark (antiquark) fields $\bar{Q}_l(1)$ ($\bar{Q}^c_l(1)$) ($l = 1, \dots, 4$). In addition there could exist many light gauge singlets.

We choose a vacuum expectation value $\langle \text{vev} \rangle \sim 10^{4.5} \text{GeV}$ in the N direction for one linear combination consisting mainly of $N_1(\alpha)$ with a small ($\sim 10^{-5.5}$)

admixture of $N_3(\alpha^2)$ and for $\bar{N}_3(\alpha)$ in the $\bar{N}$ direction. In the ν^c direction we choose a $\text{vev} \sim 10^4 \text{GeV}$ for $\nu^c_6(1)$ with an $\sim 10^{-2}$ admixture of $\nu^c_4(\alpha^2)$. For the supersymmetry breaking scale we choose a value $M_s \sim 10^4 \text{GeV}$.

With the above choices it is rather trivial to work out the mass spectrum. The ordinary light quarks and antiquarks are mainly among the ones transforming like α^2 or 1 under r while the ordinary light leptons are among the ones transforming like 1. There are 3 pairs of naturally light (masses $\sim < N_1 >^3 / M_P^2$) electroweak Higgs doublets transforming like 1 under r , which therefore could give masses to the leptons and at most one generation of quarks. The leptoquarks (g, g^c) transforming like α or 1 under r and one linear combination of the ones transforming like α^2 acquire masses $\sim < N_1 > \sim 10^{4.5} \text{GeV}$ while the two remaining linear combinations of the ones transforming like α^2 are relatively light with masses $M_g \sim < N_3 >$ $\sim 10^{-5.5} < N_1 > \sim 10^9 \text{GeV}$. The mixing between 'heavy' and 'light' states with different transformation properties under r is naturally $\sim (< N_1 > / M_P)^2 \sim 10^{-9}$ or $\sim < N_3 > / < N_1 > \sim 10^{-5.5}$. For adequate suppression of the proton decay rate, however, mixings as small as $\sim 10^{-15}$ between 'heavy' leptons transforming like α^2 under r and the ordinary 'light' ones transforming like 1 are needed. The natural value of this mixing is $\sim (< N_1 > / M_P)^2$. It is perfectly plausible that the further suppression required will be provided by the small coefficients of the relevant higher order superpotential couplings of the form $M_P^{-3} \lambda^3 \bar{\lambda}^3$.

The situation concerning proton decay is not very different from most models of this type, i.e. the most dangerous proton decay processes occur through dimension-five operators and are mediated at the tree level by the fermionic components of the leptoquark superfields. They are adequately suppressed if the couplings involved are $\sim 10^{-5}$ assuming that the masses of the exchanged leptoquarks are $\sim 10^{4.5} \text{GeV}$. The new feature here is that, for the baryogenesis scenario that we have in mind, some couplings of the form $Q(\alpha^2)Q^c(\alpha^2)\lambda(\alpha^2)$ and $Q^c(\alpha^2)Q^c(\alpha^2)Q^c(\alpha^2)$

are of order 1. Here $Q(\alpha^2)$ and $Q^c(\alpha^2)$ should be only the ones containing the 'light' leptoquarks $g(\alpha^2)$ and $g^c(\alpha^2)$ respectively with masses $M_g \sim 10^9 \text{GeV}$. These 'light' leptoquarks, when exchanged between the vertices corresponding to the above strong couplings, mediate proton decay processes which are adequately suppressed provided the l and e^c components of $\lambda(\alpha^2)$ are mainly 'heavy' with only $\sim 10^{-15}$ mixing with ordinary 'light' states.

We are now ready to discuss the generation of the baryon asymmetry. For this discussion the details of the very early evolution of the universe are not of much relevance. We start at the point where the temperature $T \sim M_s \sim 10^4 \text{GeV}$ which is of the order of the critical temperature T_c of the intermediate scale phase transition associated with the dominant intermediate scale $M_I \sim 10^{14.5} \text{GeV}$ in the theory. For T just above T_c the energy density of the universe is dominated by the energy density difference $\sim M_s^2 M_I^2$ between the true and the false vacuum. For $T = T_c$ the field $\phi \equiv \frac{1}{2^{1/2}}(N_1 + \bar{N}_3)^\dagger$ (bosonic component) rolls down from the origin (false vacuum) towards the true vacuum at M_I and starts oscillating coherently with frequency $\sim M_s$ and amplitude $\sim M_I$. For cosmic times t between $t_c \sim M_s^{-1} M_I^{-1} M_P$ and Γ_ϕ^{-1} , where Γ_ϕ is the decay width of ϕ , the universe is dominated by the ϕ particles. The completion of the phase transition is accompanied by an enormous entropy production, as we discussed earlier, which dilutes any pre-existing baryon asymmetry. [A more accurate statement [10] is that there is continuous entropy production for $t_c \leq t \leq \Gamma_\phi^{-1}$ at a rate $\sim t^{5/4}$.] The baryon asymmetry will then have to be generated by the decay products of the ϕ particles.

The scenario that we will present in the following generates the baryon asymmetry from the out-of-equilibrium decays of $\nu_3^c(\alpha^2)(f)$ (we indicate by (f) and (b) the fermionic and bosonic components of a superfield). For the ν_3^c 's to be produced from the decays of the ϕ 's after the intermediate scale phase transition their mass $M_{\nu_3^c(t)}$ has to be smaller than the mass of $\phi(b)$ $M_\phi \sim M_s$. The ν_3^c mass is mainly

due to the term $10^{-6} M_P^{-1} \nu_3^c \bar{\nu}_3^c N_1 \bar{N}_3$ which is also responsible for the decay of $\phi(b)$ to $\nu_3^c(f) + \bar{\nu}_3^c(f)$ with decay width $\Gamma_\phi \sim 10^{-1}(10^{-6} M_I/M_P)^2 M_\phi \sim 10^{-18} \text{GeV}$. This width is of the same order of magnitude as the Hubble constant of a radiation dominated universe with temperature $\sim 1 \text{GeV}$. We can easily arrange for this channel to be the dominant decay channel of $\phi(b)$. Provided the ν_3^c 's decay or annihilate sufficiently fast (compared to Γ_ϕ^{-1}) into light (compared to 1GeV) states, $T_f \sim 1 \text{GeV}$ will be the temperature of the radiation dominated universe immediately after the decay of the $\phi(b)$ particles.

For the generation of the baryon asymmetry we consider the decay channels of $\nu_3^c(f)$ given in fig.1. We assume that the coupling constants of the terms $g(\alpha^2)d^c(\alpha^2)\nu_3^c(\alpha^2)$, $g(\alpha^2)d^c(\alpha^2)\nu_4^c(\alpha^2)$ and $g^c(\alpha^2)u^c(\alpha^2)d^c(\alpha^2)$ involving only the light leptoquarks are of order 1. Then the decay width of $\nu_3^c(f)$ is $\Gamma_{\nu_3^c(f)} \sim 10^{-2} M_\sigma^{-2} M_{\nu_3^c(t)}^3 \sim 10^{-8} \text{GeV}$ corresponding to the decay channels in fig.1. This is certainly large compared with the Hubble constant $\sim \Gamma_\phi \sim 10^{-18} \text{GeV}$. Also annihilations of $\nu_3^c(f)$'s are suppressed because of their small number density $\sim 10^{-8} \text{GeV}^3$. CP non-conservation is introduced through the interference of the tree graphs of fig.1 with the one-loop graphs of fig.2. The channels (a) and (b) correspond to different baryon numbers in the final states if we assign baryon number 2/3 to $d^c(b)$ and baryon number 1 to $\nu_4^c(f)$. These assignments are determined by looking at the dominant processes by which the asymmetries in antiquarks and $\nu_4^c(f)$'s are converted into ordinary quark asymmetries. This is done through decays of antiquarks into two quarks (using an effective term $u^c(\alpha^2)d^c(\alpha^2)d^c(\alpha^2)$ obtained from the term $u^c(\alpha^2)d^c(\alpha^2)g^c(\alpha^2)$ by invoking $g^c(\alpha^2) - d^c(\alpha^2)$ mixing) and decays of $\nu_4^c(f)$'s into three quarks (using a diagram analogous to the one in fig.1(a) with a gluino exchanged between the two antiquarks and assuming that the mass of $\nu_4^c(f)$ is smaller than the mass of a (anti)squark or slepton).

The baryon number produced in each $\nu_3^c(f)$ decay is given by the imaginary part

of the interference between the diagrams of fig.1 and fig.2 over the total decay width of $\nu_3^c(f)$ and is readily estimated to be $n_b/n_{\nu_3^c(f)} \sim (10^{-1} - 10^{-2})M_{\nu_3^c(f)}/M_g$. The number of 'photons' (entropy) produced by each $\nu_3^c(f)$ decay is roughly $s/n_{\nu_3^c(f)} \sim M_{\nu_3^c(f)}/T_f$. Therefore, the baryon asymmetry is estimated to be

$$n_b/s \sim (10^{-1} - 10^{-2}) \frac{M_{\nu_3^c(f)}}{M_g} \frac{T_f}{M_{\nu_3^c(f)}} \sim (10^{-1} - 10^{-2}) \frac{T_f}{M_g} \sim 10^{-10} - 10^{-11},$$

which is acceptable.

The rate of the baryon-number-violating 2-2 scatterings

$$\tau_{2-2}^{-1} \leq 10^{-2} \alpha_s^2 \frac{T_f^5}{M_g^2 M_f^2} \sim 10^{-30} GeV$$

is smaller than the Hubble constant which prevents the baryon asymmetry from being erased.

A cosmologically interesting consequence of the large value of the baryon number violating coupling $u^c(\alpha^2)d^c(\alpha^2)g^c(\alpha^2)$ and the $g^c(\alpha^2) - d^c(\alpha^2)$ mixing is the lack of stability of the 'lightest supersymmetric particle' which therefore is not a dark matter candidate in the present model.

The cosmology discussed above makes sense only if the intermediate scale phase transition does not give rise to catastrophic domain walls associated with the spontaneous breaking of the global Z_3 symmetry R at low temperatures $\sim T_c$. Fortunately an inflationary mechanism for avoiding the domain wall problem by invoking the gauge singlet sector of these models has recently been proposed [11].

In conclusion, we presented a 3-generation Calabi-Yau compactification of the heterotic superstring in which there is a possibility of generating a more or less computable baryon asymmetry through a mechanism analogous to the well-understood, in the context of grand unified models, out-of-equilibrium decay of heavy (compared to the temperature) particles. One very interesting aspect of our scenario is the low temperature $\sim 1 GeV$ at which the mechanism is operative.

References

- [1] A.M. Boesgaard and G. Steigman, Ann. Rev. Astron. Astrophys. **23** (1985) 319.
- [2] A.D. Sakharov, Zh. Eksp. Teor. Fiz. Pis'ma **5** (1967) 32;
S. Weinberg, in Lectures on particles and fields, ed. S. Deser, K. Ford (Prentice-Hall, Englewood Cliffs, NJ) p. 682.
- [3] M. Yoshimura, Phys. Rev. Lett. **61** (1978) 281; (E) **62** (1978) 746;
S. Dimopoulos and L. Susskind, Phys. Rev. **D18** (1978) 4500;
A.Yu. Ignatiev, N. Krasnikov, V. Kuzmin and A. Tavkhelidze, Phys. Lett. **76B** (1978) 686;
S. Weinberg, Phys. Rev. Lett. **42** (1979) 850;
D. Toussaint, S. Treiman, F. Wilzek and A. Zee, Phys. Rev. **D19** (1979) 1036;
J. Ellis, M.K. Gaillard and D. Nanopoulos, Phys. Lett. **80B** (1979) 360; (E) **82B** (1979) 464;
S.M. Barr, Phys. Rev. **D19** (1979) 3803;
E. Kolb and M. Turner, Ann. Rev. Nucl. Sci. **33** (1983) 645.
- [4] D.J. Gross, J. Harvey, E. Martinec and R. Rohm, Phys. Rev. Lett. **54** (1985) 502;
M. Green and J.H. Schwartz, Phys. Lett. **149B** (1984) 117.
- [5] P. Candelas, G. Horowitz, A. Strominger and E. Witten, Nucl. Phys. **B258** (1985) 46;
E. Witten, Nucl. Phys. **B258** (1985) 75.
- [6] G. Lazarides, C. Panagiotakopoulos and Q. Shafi, Nucl. Phys. **B307** (1988) 937.