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Cosmology of 3-Generation Calabi-Yau Superstring Models

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Abstract

The cosmology of 3-generation Calabi-Yau (CY) superstring models is discussed. A model independent lower bound $\sim 10^{14}\text{GeV}$ on the largest intermediate scale in models with $SU(3)^3$ gauge group is derived by requiring that the magnetic monopole flux in the Galaxy does not exceed the Parker bound. Using this bound we argue that the baryon asymmetry in 3-generation CY models should necessarily be created at temperatures $\sim 1\text{GeV}$ at most. A late baryogenesis scenario is presented in the context of a model with naturally suppressed proton decay which is generalizable to others with Z_3 matter parity and low $U(1)_{B-L}$ breaking scale.

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The Calabi-Yau (CY) approach [1] to the compactification of the $E_8 \times E_8$ heterotic superstring [2] has certainly many positive features. The four-dimensional models that it gives rise to possess $N=1$ supersymmetry and chiral fermions belonging to 27 or $\bar{27}$ representations of the gauge group E_6 . For multiply connected CY spaces the Wilson group mechanism can be used to further break E_6 down to a proper subgroup already at the compactification scale M_c . The net number of light chiral fermion generations is a topological invariant of the CY manifold. In addition, the compactified theory often has discrete symmetries related in a very well-defined way to the complex structure of the internal space.

Clearly the above positive features are not sufficient to guarantee successful confrontation of CY superstring models with experiment due to the large number of phenomenological constraints. Many of these constraints are, more or less naturally, satisfied by the numerous models based on a single and so far unique acceptable 3-generation CY manifold [3]. The problem, however, which proved the most difficult to take care of in a natural manner is the sufficient suppression of proton decay processes occurring through dimension five ($d=5$) operators. Such processes are suppressed by only one power of a heavy mass scale and lead to an unacceptable proton decay rate unless the average Yukawa couplings involved in the tree-level diagram giving rise to the $d=5$ operator are smaller than $\sim 10^{-5} - 10^{-6}$ for values of the heavy scale $\sim 10^{14}$ GeV.

Some time ago we identified a class of 3-generation superstring models with $SU(3)^3$ gauge group in four dimensions in which, as a result of a $Z_3 \times Z_3$ discrete global symmetry combined with a low $U(1)_{B-L}$ breaking scale, all proton decay processes occurring through tree-level $d=5$ or $d=6$ operators are naturally suppressed [4]. It turns out that the dominant proton decay processes occur through one-loop $d=6$ operators. Such processes are adequately suppressed with only moderate intermediate scales $M_I \sim 10^{10}$ GeV and average values $\sim 10^{-2}$ for the 27^3

superpotential couplings involved in the loops [5].

Our past experience with grand unified models has taught us that besides the phenomenological constraints a model should satisfy several others coming from cosmological considerations which sometimes turn out to be even more stringent. Typical examples are the constraints resulting from the formation of topological defects (magnetic monopoles, cosmic strings or domain walls) accompanying the phase transitions in the early Universe, the constraints resulting from the domination of the energy density of the Universe by stable massive states, or the constraints that nucleosynthesis places on the number of massless degrees of freedom and on the baryon asymmetry n_b/s (the ratio of the net baryon number density over the entropy density) of the Universe. The purpose of the present note is to examine some of the cosmological consequences of 3-generation CY superstring models in general and more specifically of the above mentioned class possessing a natural suppression mechanism of the proton decay rate.

An immediate consequence of the mechanism that generates the intermediate scales M_I in CY superstring models [6] is that the associated phase transition has a rather low critical temperature $T_c \sim 10^{3 \pm 1}$ GeV (M_c is the supersymmetry breaking scale in the visible sector) and generates, for large intermediate scales, an enormous amount of entropy [7]. The entropy increase factor $\Delta \sim f^{-1}(M_c/M_P)^{1/2}(M_I/T_c)^3$ where $M_P \sim 10^{19}$ GeV is the Planck mass and f is determined from the decay width $\Gamma \equiv f^2 M_c^2/M_I^2$ of the standard model singlet acquiring the intermediate vacuum expectation value (vev). Such a low T_c makes it implausible that the phase transition will experience the primordial inflationary phase which is usually invoked in order to solve the homogeneity, isotropy and flatness problems of the standard conventional cosmological model. Still, we have demonstrated that the formation of global topological defects, like domain walls associated with the spontaneous breaking of discrete global symmetries, could be avoided due to an inflationary

mechanism involving the gauge singlet sector of the theory [8]. This mechanism does not apply to the formation of topological defects like gauge cosmic strings or magnetic monopoles associated with the spontaneous breaking of local gauge symmetries.

The most popular flux breaking in the 3-generation CY space leaves unbroken in four dimensions the gauge group $SU(3)_c \times SU(3)_L \times SU(3)_R$. Spontaneous breaking of this symmetry at a scale M_I leads to the appearance of topologically stable 't Hooft - Polyakov monopoles with mass $\sim M_I$ and three Dirac units of magnetic charge. The late formation and subsequent evolution of these magnetic monopoles in a matter dominated Universe has been discussed before [9] assuming a scale $M_I \sim 10^{15}$ GeV. This value of M_I seemed at the time to be dictated by considerations related to the proton lifetime in specific early models [10]. Our discussion above in connection with the proton decay rate shows that there exist today 3-generation CY models in which the scale M_I could be considerably lower than $\sim 10^{15}$ GeV. It therefore makes sense to treat M_I as a free parameter.

Monopole-antimonopole annihilation becomes more effective for larger values of the magnetic charge h or equivalently for smaller values of the $SU(3)_L \times SU(3)_R$ coupling constant g . It turns out that one can put a lower bound on g using the one-loop renormalization group equations (RGEs) for $\alpha_{LR} = g^2/4\pi$ and $\sin^2\theta_w$ independently of the value of M_I and the exact mass spectrum between M_s and M_I . The RGEs are of the form $\sin^2\theta_w = 3/8 - \alpha_{em}(M_w)(16\pi)^{-1}[(110/3)\ln(M_s/M_w) + X - Y]$ and $\alpha_{LR}^{-1} = (3/8)\alpha_{em}^{-1}(M_w) - (16\pi)^{-1}[10\ln(M_s/M_w) + X + Y']$ where X , Y , Y' are non-negative and Y and Y' correspond to contributions from the same states. For a given value of M_s and $\sin^2\theta_w$ (less than $3/8$) we can minimize α_{LR} by dropping Y and Y' and expressing X as a function of $\sin^2\theta_w$. We thus obtain a lower bound $\alpha_{LR} = [(5/3\pi)\ln(M_s/M_w) + \sin^2\theta_w\alpha_{em}^{-1}(M_w)]^{-1}$ which for $\sin^2\theta_w = 0.24$ and $M_s = 10$ TeV is equal to 0.03. Correspondingly the largest

possible value of the monopole magnetic charge is $h = 50$.

For monopole masses many orders of magnitude larger than the critical temperature the dominant annihilation mechanism is capture by emission of radiation. To simplify the discussion we will overestimate the annihilation by identifying the monopole mass with M_I and by assuming the largest possible value of h independently of the value of M_I . The final value of the monopole to photon number density is $\tau_f \sim \tau_{in}\Delta^{-1}[1 + 10^2(10^2\tau_{in})(M_P/M_s)(M_s/M_I)^{21/10}]^{-1}$ where $\tau_{in} \sim 10^{-2}$ is its initial value, Δ^{-1} represents the dilution factor due to the entropy production and the last factor incorporates the effect of annihilations. If $10^2(10^2\tau_{in})(M_P/M_s)(M_s/M_I)^{21/10} \geq 1$ monopole annihilation is effective. The condition $10^{-30} \geq \tau_f$, corresponding roughly to the Parker bound [11] on the monopole flux in the Galaxy, leads to the relation $10^{-18}(M_I/M_s)^{9/10}(M_s/10^3\text{GeV})^{-1/2} \geq f$. Taking into account that $M_I \geq M_s$ we obtain for the decay width $\Gamma = f^2 M_s^2/M_I^2$ the relation $10^{-33}\text{GeV} \geq \Gamma$. Since the lifetime Γ^{-1} is much longer than the cosmic time $\sim 1\text{sec}$ corresponding to the beginning of the process of nucleosynthesis we exclude all values of M_I for which there is an effective annihilation through emission of radiation. For $1 \geq 10^2(10^2\tau_{in})(M_P/M_s)(M_s/M_I)^{21/10}$ we have $\tau_f \sim \tau_{in}\Delta^{-1}$. From $10^{-30} \geq \tau_f$ follows that $10^{-36}(10^2\tau_{in})^{-1}(M_I/M_s)^2(M_s/10^3\text{GeV})^{1/2} \geq f$. Combining the last inequality with the requirement that $\Gamma \geq 10^{-24}\text{GeV}$ (decay before $\sim 1\text{sec}$) we obtain the bound $M_I \geq 10^{14.25}(10^2\tau_{in})^{1/2}(M_s/10^3\text{GeV})^{1/2}$ GeV. Clearly this bound should be regarded as only approximate since the connection between the bound on τ_f and the Parker bound on the monopole flux involves assumptions about the way monopoles behave when matter collapses to form galaxies.

For scales M_I only a few orders of magnitude above $T_c \sim M_s$ one should consider an annihilation mechanism involving diffusion of monopoles through the dense plasma of charged particles toward an antimonopole [12]. Combining refs. [9]

and [12] we obtain $\tau_f \sim 4\pi(2\pi)^{1/2}(2^{1/2}3\alpha_{LR})^{-1/2}h^{-4}(M_I/M_P)(M_I/M_s)^{1/2}\Delta^{-1}$. The condition $10^{-30} \geq \tau_f$, which is very conservative for light (relativistic) monopoles, leads to the condition $0.7 \times 10^{-17}(M_I/M_s)^{3/2}(M_s/10^8 \text{ GeV})^{-1/2} \geq f$ which in turn gives $0.5 \times 10^{-31}(M_I/M_s) \text{ GeV} \geq \Gamma$. Requiring that $\Gamma \geq 10^{-24} \text{ GeV}$ we obtain $M_I/M_s \geq 2 \times 10^7$. However, the present mechanism is only operative for $2^{1/2}\alpha_{LR}B^2h^4/(4\pi)^2 \geq M_I/M_s$, with $B = (3/4\pi^2)\zeta(3)\Sigma_i(h\tilde{q}_i/4\pi)^2$ (the sum running over all relativistic charged particle spin states). This gives us the inequality $10^5 \geq M_I/M_s$, which is inconsistent with the previous one. We conclude that our earlier bound is not altered.

It is important to note that this cosmological bound on M_I does not depend on details of the superpotential and relies only on the assumption that the gauge group in four dimensions is $SU(3)^3$. For the second possible group in the 3-generation manifold, namely flipped $SU(6)$, whose breaking does not give rise to magnetic monopoles the gauge boson mediated proton decay gives an even stronger model independent bound on M_I [13], namely M_I larger than $\sim 10^{16} \text{ GeV}$.

The necessity for very heavy scales in the 3-generation manifold has as a consequence an enormous entropy production during the period between the beginning and the completion of the associated phase transition. It becomes then obvious that unless n_b/s is initially very large ($\gg 1$), the entropy produced will dilute it to a value much below $10^{-10} - 10^{-11}$. Therefore, only scenarios that generate the baryon asymmetry near or after the completion of the intermediate scale phase transition, most probably from the decay products of the standard model singlet responsible for it, should be considered. The natural decay width Γ of this singlet is $\Gamma \sim M_s^3/M_I^2$ and consequently the expected final temperature T_f after the completion of the phase transition is $T_f \sim (\Gamma M_P)^{1/2} \sim (M_s M_P)^{1/2} (M_s/M_I)$. Assuming $M_s \sim 10^8 \text{ GeV}$ and $M_I \geq 10^{14} \text{ GeV}$ we see that $T_f \sim 1 \text{ GeV}$ or less. Such a very low temperature baryogenesis cannot have anything to do with the sphaleron related

strong baryon number violating effects in the standard model [14] and the physics beyond the standard model has to be somehow involved. The only other known source of baryon number violation in such theories are the baryon number violating couplings in the superpotential. All such couplings involve one of the extra (not existing in the standard model) $SU(2)_L$ singlet, color triplets (antitriplets) in the 27 of E_6 which are usually called g (g^c) quarks (antiquarks). In a low temperature baryogenesis scenario it is almost inevitable that the value of n_b/s will be suppressed by powers of M_s/M_g or T_f/M_g (M_g being the g or g^c mass). On the other hand M_g is strongly constrained by the proton decay rate. Viable scenarios have to make sure that g 's are light enough with sufficiently strong baryon number violating couplings and still the proton does not decay very fast.

In an earlier paper [15] we constructed such a 3-generation GY model in which the out-of-equilibrium decay of a relatively heavy lepton with the quantum numbers of an antineutrino produced a late baryon asymmetry. The g quarks involved in the decay diagrams were relatively light and had strong couplings to ordinary light quarks without an unacceptable proton decay rate. This was achieved because the leptons belonging to the same $SU(3)^3$ multiplet as the heavy ν^c , although they did not have any conserved quantum number distinguishing them from ordinary leptons, had only a tiny mixing angle with them. Here we would like to point out that there exists an analogous baryogenesis mechanism in a relatively larger class of models with Z_3 matter parity in which the $U(1)_{B-L}$ breaking scale is very low. In such models the most dangerous proton decay processes occurring through tree-level $d=5$ operators are mediated by the $g - g^c$ quarks which remain invariant under the Z_3 matter parity. Those are the ones which have to be superheavy. There exist, however, others transforming non-trivially under the Z_3 matter parity which do not contribute to tree-level $d=5$ operators for proton decay because in all baryon number violating vertices involving one of them there

is necessarily an additional quark or lepton superfield with a fermionic component heavier than the proton. As a result of this property they only contribute to one-loop $d=5$ operators and the constraints on their masses and couplings are more loose. It is natural to try to generate the baryon asymmetry using diagrams involving only the latter. For a baryogenesis scenario even details are important and for this reason it can only be discussed in connection with a specific model. In the following we will concentrate on such a scenario using one specific model of the class mentioned in the beginning with $SU(3)^3$ gauge symmetry and natural suppression of the most dangerous for proton decay $d=5$ operators.

The 3-generation CY space is defined as the complete intersection 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 Z_3 symmetry group. For the polynomial constraints we choose $P_1 = \sum_{i=0}^2 x_i^3 + ax_0x_1x_2 = 0$, $P_2 = \sum_{i=0}^2 y_i^3 + by_0y_1y_2 = 0$ and $P_3 = x_0y_0 + c_1x_1y_1 + c_2x_2y_2 = 0$ ($|c_1| \sim 1$, $|c_2| \sim 30$, $|a|, |b| \ll 1$), with x_i and y_i ($i = 0, \dots, 3$) the homogeneous coordinates of the two CP^3 's. The freely acting Z_3 symmetry is generated by $g = \text{diag}(1, \alpha^2, \alpha, \alpha, 1, \alpha, \alpha^2, \alpha^2)$ acting on $[x_i, y_i]$ with $\alpha = \exp(2\pi i/3)$. In the limit $c_1 = 1$ and $a = b = 0$ the 3-generation space has a maximal 108 element discrete symmetry [16] generated by $B: x_1 \rightarrow \alpha x_1$ and $y_2 \rightarrow \alpha^2 y_2$, $F: x_{0,1} \rightarrow \alpha^2 x_{1,0}$ and $y_{0,1} \rightarrow \alpha y_{1,0}$, $G: x_3 \rightarrow \alpha x_3$ and $y_3 \rightarrow \alpha y_3$, $S: x_i \rightarrow y_i$ and $y_i \rightarrow x_i$. The unique harmonic $(3,0)$ form ω of the CY space transforms under the R-symmetry G like α^2 . As the value of c_1 departs from 1 the symmetry generated by F^3 is broken. If we further assume $a = b \neq 0$, B and F^2 get broken with the combination F^2B remaining unbroken. Finally $a, b \neq 0$, $a \neq b$ breaks S and leaves the $Z_3 \times Z_3$ symmetry generated by G and F^2B unbroken. The chosen values $|c_1| \sim 1$, $|c_2| \sim 30$, $|a|, |b| \ll 1$, leave exactly unbroken the generators G and F^2B , break strongly F^3 and only slightly F^2 , B , and S .

For the flux breaking we assume the embedding of the freely acting Z_3 into

E_8 defined by the homomorphism $g \mapsto U_g \equiv (I_3, \alpha I_3, \alpha I_3) \in SU(3)_c \times SU(3)_L \times SU(3)_R \equiv H$ which leaves unbroken below M_c the gauge group H . Under H the left-handed lepton (λ), quark (Q) and antiquark (Q^c) fields in a 27 of E_8 transform as follows: $\lambda = (1, \bar{3}, 3) = \begin{pmatrix} H^1 & H^2 & 1 \\ e^c & \nu^c & N \end{pmatrix}$, $Q = (3, 3, 1) = \begin{pmatrix} q \\ g \end{pmatrix}$, $Q^c = (\bar{3}, 1, \bar{3}) = (u^c, d^c, g^c)$. The massless gauge-non-singlet spectrum and its transformation under the maximal discrete symmetry group generators are given in table 1. In addition there could exist many light gauge singlets.

We choose a $\text{vev} \sim 10^{14.5} \text{GeV}$ in the N direction for one linear combination consisting mainly of N_7 with an $\sim 10^{-4.5}$ contribution from N_1 and an $\sim 10^{-6.5}$ contribution from N_3 and for $\bar{N}_1$ in the $\bar{N}$ direction. The $SU(2)_R \times U(1)_{B-L}$ symmetry is broken at low energies by $\langle \bar{\nu}_4^c \rangle \sim 10^4 \text{GeV}$ and $\langle \nu_2^c \rangle \sim 10^6 \text{GeV}$. We also assume the existence of a gauge singlet S transforming like $(\alpha, 1, \alpha)$ under G , F^2B and B respectively with $\langle S \rangle \sim 10^8 \text{GeV}$. Alternatively, one could use a singlet transforming like $(\alpha^2, 1, \alpha)$.

The present model is only a small perturbation around models whose phenomenology has been studied extensively [4,5]. Here we simply state that there is no problem with proton decay if one assumes the values for the Yukawa couplings given explicitly in [5] in terms of the complex parameters c_1 and c_2 . These values correspond to a calculation [17] for the model with an exact 54 element discrete symmetry generated by F^3 , B , G and S . This calculation is exact modulo an overall constant A , which we choose to have a value $A \sim 10^{-1} c_2^{-1}$, but does not take properly into account the normalization of the kinetic terms of the fields. Although the missing factors could change the results we still feel that these values constitute a sufficiently restrictive parametrization of the Yukawa couplings (there are many more calculated couplings than unknown field normalizations). In spite of the fact that the calculation of the effects of the parameters a and b which are slightly away from zero (slightly broken generators B and S) is not

available, these effects are not expected to change the results since departures from the calculated values will be proportional to (some power of) a and b and therefore tiny. All couplings which are forbidden by B and S but are allowed by $F^2 B$ and G will be assumed to be small. It is also important to remember that the combination of the Z_3 symmetry $F^2 B$ with the Z_3 subgroup of $SU(3)_R$ generated by $\text{diag}(\alpha^2, \alpha, 1)$ is left unbroken and plays the role of matter parity. The light quarks (antiquarks) are the q (u^c, d^c)'s which transform trivially under $F^2 B$. The light leptons (charged antileptons) are linear combinations of the l 's (e^c 's) which transform trivially under $F^2 B$ (mostly l_1, l_2, l_3 (e_1^c, e_2^c, e_3^c)).

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 irrelevant. We start at the point where the temperature $T \sim M_s \sim 10^{3.5} \text{GeV}$ 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 $\phi \equiv \frac{1}{2} (N_I + \bar{N}_I)^\dagger$ (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 $\sim M_I$. This is also the time of the superheavy monopole formation. 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. During the same period there is continuous entropy production at a rate $\sim t^{3/4}$ which dilutes any preexisting baryon asymmetry and the monopole to photon number density. The baryon asymmetry will then have to be generated by the decay products of the ϕ particles.

In our scenario the baryon asymmetry is generated from the out-of-equilibrium decays of $\nu_2^c(f)$. [We indicate by (f) and (b) the fermionic and bosonic components

of a superfield.] For the ν_2^c 's to be produced from the decays of the ϕ 's their mass has to be smaller than the mass of $\phi(b)$ $M_\phi \sim M_s$. The mass of ν_2^c is mainly due to the term $10^{-8} M_c^{-1} \nu_2^c \bar{\nu}_4^c N_7 \bar{N}_2$ ($M_c \sim 10^{17.5} \text{GeV}$) which is also responsible for the decay of $\phi(b)$ with decay width $\Gamma_\phi \sim 10^{-1} (10^{-8} M_I / M_c)^2 M_\phi \sim 10^{-10} \text{GeV}$. [Notice that the above term is forbidden by the slightly broken symmetry generator B and therefore is naturally tiny. We also assume that the term $\nu_2^c \bar{\nu}_6^c S$ gives ν_2^c a mass $\sim 10^3 \text{GeV}$ or less.] 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 arrange for this channel to be the dominant decay channel of $\phi(b)$. Provided the ν_2^c 's decay or annihilate sufficiently fast (compared to Γ_ϕ^{-1}) into light (compared to 1GeV) states, the final temperature of the radiation dominated Universe immediately after the decay of the $\phi(b)$ particles will be $T_f \sim 1 \text{GeV}$.

For the generation of the baryon asymmetry we consider the decay channels of $\nu_2^c(f)$ given in fig.1. The relevant superpotential terms are: $A c_1 c_2 (1 + c_1^3)^{-1} \nu_2^c g_s d_3^c$, $A g_s^2 d_1^c u_3^c$, $-A c_2 c_1^{-1} [(1 + c_1)(1 + c_1^3)^{-1} + (c_1^2 + c_2^2)^{-1}] \nu_2^c g_s d_3^c$, $A [1 + c_1^2 (c_1^3 + c_2^3)^{-1}] \nu_2^c g_1 d_1^c$, $-A c_2 (1 + c_1^3)^{-1} g_1^2 d_1^c u_3^c$. The masses M_g of the g 's and g^c 's involved are all $M_g \sim 10^{6.5} \text{GeV}$. Then the decay width of $\nu_2^c(f)$ is $\Gamma_{\nu_2^c(f)} \sim 10^{-2} |c_1|^2 |c_2|^2 |A|^4 M_g^{-2} M_s^3 \sim 10^{-9.5} \text{GeV}$ corresponding to the first decay channel in fig.1. This is certainly large compared with the Hubble constant $\sim \Gamma_\phi \sim 10^{-10} \text{GeV}$ and the $\nu_2^c(f)$ annihilation rate. 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 $-1/3$ to $u_3^c(f)$ and $d_3^c(b)$, $+2/3$ to $d_1^c(b)$ and $d_3^c(b)$, and 0 to $\nu_2^c(f)$. These assignments are determined by the dominant processes by which the asymmetries in the above particles are converted into ordinary quark asymmetries. Squark asymmetries are converted into quark asymmetries through decays to the corresponding quark and a photino or another neutralino. Also $\nu_2^c(f)$ decays fast to an antilepton and

$q_3(f) + u_3^c(f)$ via a light H_1^+ Higgs boson. The $d_1^c(b)$ decays first to a neutralino and a $d_1^c(f)$ which later decays primarily to an anti- $u_3^c(b)$ and an anti- $d_3^c(f)$ via the coupling $A g_2^c d_1^c u_3^c$ and the $g_2^c - d_3^c$ mixing $\sim c_1 c_2 (1 + c_1^2)^{-1} < \nu_2^c > / < N_3 >$. The $d_3^c(b)$ asymmetry is converted first to $d_3^c(f)$ asymmetry and later, through the coupling $M_c^{-2} < N_7 > < \bar{N}_2 > d_3^c q_7 H_2^+$ and the $H_2^+ - l_3$ mixing $\sim < \nu_2^c > / < N_3 >$, is converted to anti- $q_7(b)$ asymmetry which gives rise to anti- $q_7(f)$ asymmetry. Each anti- $q_7(f)$ gives rise to two quarks through the term $-A c_2 (1 + c_1^2)^{-1} q_3 q_7 g_3$ and the $g_3 - q_3$ mixing. The final result is that each $d_3^c(b)$ gives rise to two quarks before the time of nucleosynthesis.

The baryon number produced in each $\nu_2^c(f)$ 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 $\nu_2^c(f)$ and is readily estimated to be $n_b/n_{\nu_2^c(f)} \sim \text{Im}[c_1^{-1}] |c_2|^2 [1 + c_1^2]^{-2} |A|^2 M_g^2 / M_g^2 \sim 10^{-2} M_g^2 / M_g^2$. The number of 'photons' (entropy) produced by each $\nu_2^c(f)$ decay is roughly $s/n_{\nu_2^c(f)} \sim M_*/T_f$. Therefore, the baryon asymmetry is estimated to be $n_b/s \sim 10^{-2} M_* T_f M_g^{-2} \sim 10^{-9.5}$.

The rate of the baryon-number-violating 2-2 scatterings $\tau_{2-2}^{-1} \leq 10^{-2} \alpha_* T_f^5 M_g^{-2} M_*^{-2}$ is smaller than the Hubble constant and the baryon asymmetry is not erased.

We expect a further dilution of both n_b/s and the final monopole to photon number density τ_f by a factor ~ 10 at most due to the somewhat delayed decay of $\phi(f)$.

In spite of the fact that the Z_3 matter parity does not strictly forbid the decay of the 'lightest supersymmetric particle' (LSP), when combined with the Z_3 R-symmetry G and the low $< \nu^c >$ makes its lifetime longer than the present lifetime of the Universe. Therefore the LSP is a dark matter candidate.

To summarize, we presented cosmological considerations which place a lower bound on the magnitude of the dominant intermediate scale in 3-generation CY superstring models and using this bound we argued that there is necessity for a

very low temperature ($\sim 1\text{GeV}$) baryogenesis scenario. We also presented such a scenario in the context of a specific phenomenologically attractive 3-generation CY superstring model which is generalizable to other models with Z_3 matter parity and low $U(1)_{B-L}$ breaking scale. It is, in our opinion, very important that we finally have CY superstring models based on specific 3-generation internal space constructions which are phenomenologically and cosmologically acceptable.

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	B	F	S	G	F ² B	B	F	S	G	F ² B
λ_1	α	α	λ_3	1	1	λ_1	1	λ_1	1	1
λ_2	1	α	λ_4	α	α^2	λ_2	1	$-\lambda_2$	1	1
λ_3	α^2	α^2	λ_1	1	1	λ_3	$-\alpha^2$	λ_4	α	α^2
λ_4	1	α^2	λ_2	α	α	λ_4	$-\alpha^2$	λ_5	α^2	α^2
λ_5	1	-1	λ_6	1	1	λ_5	$-\alpha^2$	λ_6	α^2	α^2
λ_6	1	1	λ_7	1	1	λ_6	$-\alpha$	λ_7	α^2	α
λ_7	1	1	λ_8	α^2	1	λ_7		λ_8		
λ_8	α	1	λ_9	α	α^2	λ_8		λ_9		
λ_9	α^2	1	λ_9	α	α^2	λ_9		λ_9		
Q_1	α	$\alpha^2 Q_2^2$	Q_1^c	α	α^2	Q_1	Q_4^c	Q_1^c	α	α
Q_2	α^2	αQ_1^c	Q_2^c	α	α	Q_2	Q_3^c	Q_2^c	α^2	α^2
Q_3	1	Q_3^c	Q_3^c	1	1	Q_3	Q_2^c	Q_3^c	α^2	α
Q_4	α^2	$\alpha^2 Q_6^c$	Q_4^c	1	1	Q_4	Q_1^c	Q_4^c	α	α^2
Q_5	1	$\alpha^2 Q_7^c$	Q_5^c	α	α	Q_5				
Q_6	α	αQ_4^c	Q_6^c	1	1	Q_6				
Q_7	1	αQ_5^c	Q_7^c	α	α^2	Q_7				
Q_1^c	α^2	αQ_2	Q_1	α	α	Q_1^c	Q_4	Q_1	α^2	α^2
Q_2^c	α	$\alpha^2 Q_1$	Q_2	α	α^2	Q_2^c	Q_3	Q_2	α	α
Q_3^c	1	Q_3	Q_3	1	1	Q_3^c	Q_2	Q_3	α^2	α^2
Q_4^c	α	αQ_6	Q_4	1	1	Q_4^c	Q_1	Q_4	α	α
Q_5^c	1	αQ_7	Q_5	α	α^2	Q_5^c				
Q_6^c	α^2	$\alpha^2 Q_4$	Q_6	1	1	Q_6^c				
Q_7^c	1	$\alpha^2 Q_5$	Q_7	α	α	Q_7^c				

Table 1: Transformation properties of the 'massless' superfields under the symmetries generated by B , F , S , G and $F^2 B$.

Fig. 1. (a) The decay $\nu_2^c(f) \rightarrow \bar{d}_3^c(b) + \bar{u}_3^c(f) + \nu_3^c(f)$; (b) the decay $\nu_2^c(f) \rightarrow \bar{d}_3^c(b) + \bar{d}_5^c(b) + \nu_3^c(f)$.

Fig. 2. Radiative corrections to (a) fig. 1(a) and (b) fig. 1(b).