Baryogenesis and the Gravitino Problem in Superstring Models

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(Received 2 December 1985)

Superstring models with an intermediate mass scale possess a novel mechanism for generating the baryon asymmetry in the Universe. It involves the out-of-equilibrium decay of heavy particles at temperatures close to the electroweak scale. The gravitino problem encountered in many supergravity models is automatically resolved.

PACS numbers: 12.10.Gq, 14.80.Ly, 98.80.Cq

A wide variety of phenomenological and cosmological considerations can be effectively employed to impose constraints on candidate superstring models arising from the $E_8 \otimes E_8$ superstring theory.[1] Relatively obvious requirements such as acceptable values of $\sin^2\theta_W$, suppression of rapid proton decay, and sufficiently small neutrino masses already greatly limit the number of viable models.[2] The imposition of additional constraints like the absence of spontaneously broken discrete symmetries and harmful axions rule out several other models.[3] It was pointed out in Ref. 3 that the above constraints together with some additional ones probably cannot be satisfied by the vast majority of the models. Examples of models based on a rank-five[4] subgroup of E_6 which do, however, satisfy the constraints were presented in Ref. 3.

The purpose of this Letter is to consider two important and often related cosmological problems in the context of superstring models. These are (i) the origin of the baryon asymmetry in the Universe, and (ii) the gravitino problem usually encountered in models in which the gravitino mass is on the order of the electroweak scale. We wish to show that in a class of superstring models, examples of which were presented in Ref. (3), these two problems could be resolved in a novel way.

The merits of superstring models which possess an intermediate scale $M_I \sim 10^9$ GeV have been extolled elsewhere.[2,3] It turns out that not only the existence but also the mechanism whereby M_I arises play an essential role in resolving both the baryogenesis issue and the gravitino problem in superstring models.

In order to see how baryogenesis can arise in superstring models, let us recall that the compactification of the $E_8 \otimes E_8$ heterotic theory on suitable Calabi-Yau spaces predicts, besides the "known" quark, lepton, and Higgs superfields, the existence of additional fields which include color-triplet, SU(2)-singlet superfields g_i.[5] For simplicity we restrict our attention to the boson component of g (also denoted by the same symbol with the family index i not explicitly displayed). We would like to show that the out-of-equilibrium decay of massive g bosons at temperatures close to the electroweak scale can produce the desired baryon asymmetry. This presumes, of course, that the other two requirements for successfully generating the baryon asymmetry, the existence of baryon-number-nonconserving couplings of g and the presence of CP nonconservation in these couplings, are fulfilled. We will display models in which all three conditions can be met.

We first present a simplified account of the model-independent features of the baryogenesis scenario. The g boson has a positive mass-squared term $\sim M_s^2 g^* g$ ($M_s \simeq 1$ TeV is of the order of the supersymmetry-breaking scale in the known sector) which presumably arises from supersymmetry breaking in the hidden sector. The important point is that it also acquires a mass through its coupling $\sim g^* g \phi^* \phi$ to a $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$-singlet scalar field ϕ, whose zero-temperature effective potential is of the form[2]

$$V_0(\phi) = -(M_s^2/2)\phi^*\phi + (\lambda/6M_c^2)(\phi^*\phi)^3, \tag{1}$$

Here $M_c \sim 10^{17}$ GeV denotes the compactification scale. The zero-temperature expectation value of ϕ is given by $|\langle\phi\rangle| = M_I \sim (\lambda^{-1} M_s M_c)^{1/2}$. Hence the g supermultiplet, as well as the gauge supermultiplet to which ϕ couples, acquire masses of order M_I.

For nonzero temperatures, we will add to the zero-temperature effective potential $V_0(\phi)$ the one-loop contribution[6]

$$V_T(\phi) = (T^4/2\pi^2)\sum_i (-1)^F \int_0^\infty dx\, x^2 \ln(1-(-1)^F \exp\{-[x^2+M_i^2(\phi)/T^2]^{1/2}\}). \tag{2}$$

The sum in (2) is over all helicity states, $(-1)^F$ is ± 1 for bosonic and Fermionic states, respectively, and $M_i(\phi)$ is the field-dependent mass of the ith state.

The salient features of the effective potential $V(\phi) = V_0(\phi) + V_T(\phi)$ are as follows. For $\phi << T$, there is a temperature-dependent mass term $\sigma T^2 \phi\phi^*$. Hence $V(\phi)$ has a minimum at $\phi = 0$ for $T > T_c = \sigma^{-1/2} M_s$. For $\phi \gtrsim T$, the temperature-dependent mass term is exponentially suppressed, and $V(\phi)$ develops a second minimum at $\phi \simeq M_I$ for $T \lesssim M_I$.

The minimum at $\phi = 0$ is the absolute minimum for $M_I \gtrsim T \gtrsim \mu = (M_s M_I)^{1/2} \sim 10^6$ GeV. This follows because (i) in this temperature range, the radiation energy density ($\propto T^4$) dominates over the zero-temperature vacuum energy density $\sim \mu^4$ in the $\phi = 0$ phase, and (ii) the number of massless degrees of freedom in the $\phi = 0$ phase exceeds that in the $\phi \simeq M_I$ phase.

For $T \lesssim \mu$, the $T = 0$ vacuum energy density dominates and the $\phi \simeq M_I$ phase becomes the absolute minimum of $V(\phi)$. This minimum is separated from the local minimum at $\phi = 0$ by a barrier of height $\sim T^4$ and width $\gtrsim T$. For $T_c < T < \mu$, the phase transition from the false vacuum at $\phi = 0$ to the true one at $\phi \simeq M_I$ could take place, in principle, through barrier penetration. Detailed analytical and numerical studies[7] of such processes with the Coleman-Weinberg potential have revealed that a huge amount of supercooling precedes the phase transition. Since the width of the potential $V(\phi)$ exceeds that in the Coleman-Weinberg case, we can safely assume that the universe remains in the $\phi = 0$ phase for $T > T_c$. Note that in the temperature range $T_c \lesssim T \lesssim \mu$ the vacuum energy density $\sim \mu^4$ dominates over the radiation energy density and the Universe experiences a modest amount of inflation.

When T reaches T_c the minimum at $\phi = 0$ disappears together with the barrier, and ϕ starts to roll toward the minimum at $\phi \simeq M_I$.

The g boson has mass $\sim M_s$ and its number density $n_g = n_{g^*} = n_\gamma \sim T^3$ for $T > T_c$. As ϕ starts the rollover, g acquires an additional mass $\sim |\langle\phi\rangle|$. To compute the mass M_g^* of g when it decays, we compare its lifetime with the time δt needed for $\langle\phi\rangle$ to grow from $\langle\phi\rangle \sim T_c$ to $\langle\phi\rangle \sim M_g^*$. (The reason we are not interested in the time needed for ϕ to grow from zero to T_c is the fact that n_g remains of order n_γ for ϕ in this range.)

The classical evolution of the ϕ field is governed by the equation

$$\ddot{\phi} + 3H\dot{\phi} = -dV/d\phi. \tag{3}$$

For $\phi > T$ the temperature-dependent mass terms in $V(\phi)$ can be safely ignored, and for $\phi < M_I$, the term $\lambda(\phi^*\phi)^3/6M_c^2$ in Eq. (1) can be neglected. Also, the Hubble constant $H \sim \mu^2/M_p \sim 10^{-7}$ GeV $<< M_s$ ($M_p \simeq 1.2\times 10^{19}$ GeV is the Planck mass). Equation (3) then becomes

$$\ddot{\phi} \simeq M_s^2 \phi, \tag{4}$$

which gives, for the time-dependent mass of g,

$$M_g \sim \phi(\delta t) \simeq T_c \exp(M_s \delta t). \tag{5}$$

The probability p for g to decay at time δt is given by

$$p = 1 - \exp(-\int_0^{\delta t} dt/\tau_g). \tag{6}$$

Here $\tau_g \sim f^{-2} M_g^{-1}$ is the lifetime of g, and f denotes the appropriate coupling constant. It follows that the mass of g at decay is given by

$$M_g^* \simeq T_c(1 + f^{-2} M_s/T_c). \tag{7}$$

With $f \sim \frac{1}{5}$ and $M_s/T_c \simeq \frac{1}{3}$, for instance, $M_g^* \sim 10 T_c$.

It follows from the above discussion that the g's are well out of equilibrium when they decay. Thus, at least one of the crucial ingredients for successful baryogenesis[8] is present in superstring models with an intermediate scale.

For consistency we must ensure that the g's do not annihilate before they can decay. The time needed for g-g^* annihilation through the color gauge interactions is $\tau_{an} \sim \alpha_s^{-2} M_g^2/T^3$, where $\alpha_s \simeq 0.1$ denotes the QCD coupling at $T \sim T_c$. Clearly, $\tau_{an} >> \delta t$ for $M_g \gtrsim T_c$ and so there is no g-g^* annihilation.

The field ϕ reaches the bottom of the potential $V(\phi)$ in a time which is estimated from Eq. (5) to be roughly of order $M_s^{-1}\ln(M_I/T_c) \sim 10 M_s^{-1}$. It then oscillates about the minimum with a frequency M_s which is much greater than the expansion rate of the Universe. The coupling of ϕ to other fields in the theory will result in the conversion of the scalar-field energy into radiation. The entropy production dilutes the baryon asymmetry produced in the g decays by a factor Δ which turns out to be at least as large as $\sim 10^6$. This provides an important constraint on model building. One needs to ensure that the baryon asymmetry initially produced is sufficiently large to sustain the subsequent dilution.

We now discuss the baryon asymmetry produced in the models introduced in Ref. 3. They are based on the gauge group $G = SU(3)_c \otimes SU(2)_L \otimes U(1)_L \otimes U(1)_R$ which can arise from the compactification of the $E_8 \otimes E_8$ superstring theory on a suitable

Calabi-Yau space.[5] It was shown in Ref. (3) that models based on G (supplemented by an additional discrete symmetry) possess many desirable properties. In particular they satisfy all of the constraints mentioned at the beginning of this Letter. For our purposes here, the relevant terms in the superpotential are gD_cN, gg_cS_1, and $g_cg_cU_c$.[3] Here g (g_c) denote $SU(3)_c$-triplet (-antitriplet), SU(2)-singlet superfields, while U_c and D_c carry the quantum numbers of up and down antiquarks, respectively. The N and S_1 fields are singlets with respect to $SU(3)_c \otimes SU(2)_L \otimes U(1)_\gamma$, and transform nontrivially only with respect to $U(1)_{\gamma'}$, where γ' denotes the generator orthogonal to γ.

The scalar component of S_1 has a vacuum expectation value of order $M_I \sim (M_s M_c)^{1/2}$ at zero temperature and plays the role of the scalar field ϕ from our earlier discussion. The masses of S_1 and N particles are of order M_s.

Consider the decay of the scalar g. There are two baryon-number-nonconserving channels, $g(b) \to \bar{D}_c(f)\bar{N}(f)$ and $g(b) \to U_c(b)D_c(b)N(b)$, corresponding to Figs. 1(a) and 1(b). Here b and f denote bosonic and fermionic fields, respectively. CP nonconservation can be introduced through an interference of these diagrams with those in Figs. 2(a) and 2(b). It is important to note that these diagrams give rise to CP nonconservation only if more than one generation is involved.

Although we only discussed the scalar bosons g, it should be clear that their fermionic partners provide a comparable contribution to the baryon asymmetry. This also holds for any other relevant superfields in the theory (in our case the g_c superfields).

A rough estimate shows that n_b/s initially cannot be much larger than about 10^{-5}. With a favorable choice of the parameters the decay width of S_1 can be arranged to be of the order of the Hubble constant. This minimizes the dilution of the baryon asymmetry due to entropy production.[9] The reheat temperature T_r is estimated to be $T_r \sim 3\times10^5$ GeV. If we take $T_c \simeq 3$ TeV, say, the dilution factor $\Delta \simeq 10^6$. The baryon asymmetry consequently is estimated to be less than or of order 10^{-11}. In order that the baryon asymmetry does not undergo further dilution it is important that all subsequent phase transitions do not produce any significant entropy. Of course, it may be possible to construct alternative models which produce a much larger initial baryon asymmetry.

Finally, one also may check that the characteristic times for all baryon-number-nonconserving scatterings are much greater than $\delta t \sim M_s^{-1}\ln(M_g/T_c)$. Hence the baryon asymmetry generated cannot be erased.

It should now be evident that the gravitino problem is neatly resolved in the present superstring models. For the convenience of the reader we briefly recall what the problem is. In many supergravity models (and also presumably in some superstring models) the gravitino is not the lightest supersymmetric particle and has mass on the order of the electroweak scale. Cosmological arguments rule out this possibility,[10] unless inflation is invoked to dilute away the primordial gravitinos.[11] However, one also must require that the reheat temperature T_R after inflation be below $\sim 10^8$ GeV to avoid the regeneration of too many gravitinos.[12] Since the particles that produce the baryon asymmetry typically possess masses much larger than T_R, the last requirement often makes it difficult to produce sufficient baryon asymmetry in supersymmetric theories.

The problem is easily evaded in the present models since the baryon asymmetry is generated by the decays of the g's which remain in abundance till temperatures of order T_c. The temperature T_R can therefore be as

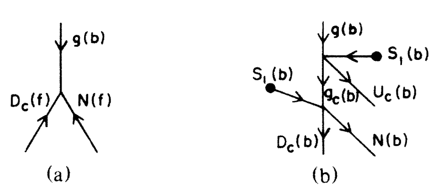

FIG. 1. (a) The decay $g(b) \to \bar{D}_c(f) + \bar{N}(f)$; (b) the decay $g(b) \to U_c(b) + D_c(b) + N(b)$.

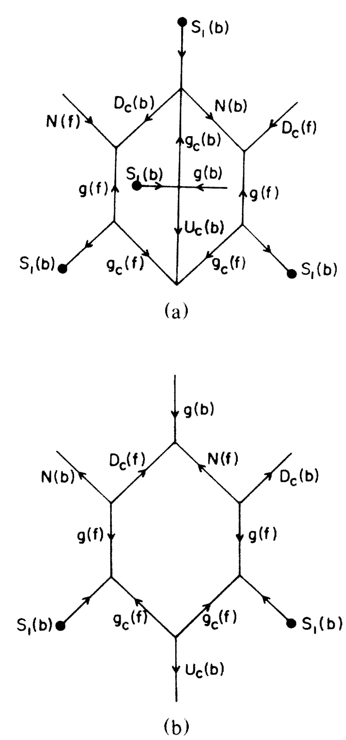

FIG. 2. Radiative corrections to (a) Fig. 1(a) and (b) Fig. 1(b).

low as is necessary to suppress the gravitino number density to acceptable levels.

To conclude, we have shown that superstring models with an intermediate mass scale possess a novel mechanism for generating the observed baryon asymmetry in the Universe. They also neatly evade the gravitino problem. Thus, a gravitino mass on the order of the electroweak scale, which may well turn out to be the case in many superstring models, is perfectly acceptable.

One of us (C.P.) would like to thank Louise Dolan for a useful discussion. One of us (G.L.) thanks the Bartol Research Foundation for hospitality during the completion of this work. This work was supported in part by the U. S. Department of Energy under Contracts No. DE-AC02-81ER40033B.000 and No. DE-FG02-84ER40176.

Note added.—After this paper was submitted for publication and went into limited circulation, K. Yamamoto kindly sent us a copy of his paper "Phase Transitions Associated with Intermediate Symmetry Breaking in Superstring Models," Johns Hopkins University Report No. JHU-HET 8508. The two papers are in agreement as far as the nature of the phase transition is concerned. However, Yamamoto does not consider the novel baryogenesis scenario proposed by us.

[1]D. J. Gross, J. Harvey, E. Martinec, and R. Rohm, Phys. Rev. Lett. **54**, 502 (1985); M. B. Green and J. Schwarz, Phys. Lett. **149B**, 117 (1984).

[2]M. Dine, V. Kaplunovsky, M. Mangano, C. Nappi, and N. Seiberg, Nucl. Phys. **B259**, 549 (1985).

[3]G. Lazarides, C. Panagiotakopoulos, and Q. Shafi, Phys. Rev. Lett. **56**, 432 (1986).

[4]E. Witten, Nucl. Phys. **B258**, 75 (1985).

[5]P. Candelas, G. Horowitz, A. Strominger, and E. Witten, Nucl. Phys. **B258**, 46 (1985).

[6]L. Dolan and R. Jackiw, Phys. Rev. D **9**, 3320 (1974); S. Weinberg, Phys. Rev. D **9**, 3357 (1974).

[7]E. Witten, Nucl. Phys. **B177**, 477 (1981); A. Billoire and K. Tamvakis, Nucl. Phys. **B200**, 329 (1982).

[8]A. D. Sakharov, Zh. Eksp. Teor. Fiz. Pisma **5**, 32 (1967) [JETP Lett. **5**, 24 (1967)]. For a recent review and other useful references see E. Kolb and M. Turner, Annu. Rev. Nucl. Sci. **33**, 645 (1983).

[9]M. Turner, Phys. Rev. D **28**, 1243 (1983).

[10]S. Weinberg, Phys. Rev. Lett. **48**, 1303 (1982); H. Pagels and J. Primack, Phys. Rev. Lett. **48**, 223 (1982).

[11]J. Ellis, A. Linde, and D. Nanopoulos, Phys. Lett. **118B**, 59 (1982).

[12]J. Ellis, D. Nanopoulos, and S. Sarkar, CERN Report No. TH4057, 1984, and references therein (to be published).