

## Grand-Unified-Theory Baryogenesis after Preheating

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At the end of inflation the Universe is frozen in a near zero-entropy state with energy density in a coherent scalar field and must be “defrosted” to produce the observed entropy and baryon number. We propose that the baryon asymmetry is generated by the decay of supermassive grand-unified-theory (GUT) bosons produced nonthermally in a preheating phase after inflation. We show that baryogenesis is possible for an inflaton mass of order  $10^{13}$  GeV and a GUT Higgs boson mass of order  $10^{14}$  GeV, thus solving many drawbacks facing GUT baryogenesis in the old reheating scenario. [S0031-9007(96)01670-5]

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In models of slow-roll inflation [1,2], the Universe is dominated by the potential energy density of a scalar field known as the *inflaton*. Inflation ends when the kinetic energy density of the inflaton becomes larger than its potential energy density. At this point the Universe might be said to be frozen: Any initial entropy in the Universe was inflated away, and the only energy was in cold, coherent motions of the inflaton field. Somehow this frozen state must be transformed to a high-entropy hot universe by transferring energy from the inflaton field to radiation. This process is usually called reheating, which may well be a misnomer since there is no guarantee that the Universe was hot before inflation. Since we are confident that the Universe was frozen at the end of inflation, perhaps “defrosting” is a better description of the process of converting inflaton coherent energy into entropy.

In the old reheating (defrosting) scenario [3], the inflaton field  $\phi$  is assumed to oscillate coherently about the minimum of the inflaton potential until the age of the Universe is equal to the lifetime of the inflaton. Then the inflaton decays, and the decay products thermalize to a temperature  $T_F \approx 10^{-1} \sqrt{\Gamma_\phi M_P}$ , where  $\Gamma_\phi$  is the inflaton decay width, and  $M_P \sim 10^{19}$  GeV is the Planck mass. In the simple chaotic inflation model we study, the potential is assumed to be  $V(\phi) = M_\phi^2 \phi^2/2$ , with  $M_\phi \sim 10^{13}$  GeV in order to reproduce the observed temperature anisotropies in the microwave background. If we write  $\Gamma_\phi = \alpha_\phi M_\phi$ , then  $T_F \approx 10^{15} \sqrt{\alpha_\phi}$  GeV [4].

In supergravity-inspired scenarios, gravitinos have a mass of the order a TeV, and a decay lifetime on the order of  $10^5$  s. If gravitinos are overproduced after inflation and decay after the epoch of nucleosynthesis, they would modify the successful predictions of big-bang nucleosynthesis. This can be avoided if the temperature  $T_F$  is smaller than about  $10^{11}$  GeV (or even less, depending on the gravitino mass) [5], which implies  $\alpha_\phi \lesssim 10^{-8}$ .

In addition to entropy, the baryon asymmetry must be created after inflation. There are serious obstacles facing any attempt to generate a baryon asymmetry in an inflationary universe through the decay of baryon number ( $B$ ) violating bosons of grand unified theories (GUTs) [6]. The first problem is that  $B$  violation through sphaleron transitions is expected to be fast at high temperatures, and would erase any preexisting baryon asymmetry produced at the GUT scale [7] unless there is a nonvanishing value of  $B - L$ . But a natural way to overcome this problem is to adopt a GUT-like SO(10), where an asymmetry in  $B - L$  may be generated.

A more serious problem is the low value of  $T_F$  in the old scenario. Since the unification scale is expected to be of order  $10^{16}$  GeV,  $B$  violating gauge and Higgs bosons (referred to generically as “ $X$ ” bosons) probably have masses greater than  $M_\phi$ , and it would be kinematically impossible to produce them directly in  $\phi$  decay.

[Gauge bosons have masses comparable to the unification scale, while  $B$  violating Higgs bosons may have a mass a few orders of magnitude less. For example, in SU(5) there are  $B$  violating “Higgs” bosons in the five-dimensional representation that may have a mass as small as  $10^{14}$  GeV. In fact, these Higgs bosons are more likely than gauge bosons to produce a baryon asymmetry since it is easier to arrange the requisite CP violation in the Higgs decay [8–10]. Furthermore, if  $T_F$  is less than  $10^{11}$  GeV,  $X$  bosons will be exponentially rare in the thermal background after inflation.]

However, it has been recently realized [11,12] that reheating may differ significantly from the above simple picture. In the first stage of reheating, which was called “preheating” [11], nonlinear quantum effects may lead to an extremely effective dissipational dynamics and explosive particle production even when single particle decay is kinematically forbidden. Particles can be produced in a regime of a broad parametric resonance, and it is possible that a significant fraction of the energy stored in the form

of coherent inflaton oscillations at the end of inflation is released after only a dozen oscillation periods.

This Letter demonstrates that preheating may play an extremely important role for the GUT generation of the baryon asymmetry, as first suggested in [11,13] (see also [14,15]). Indeed, we will show that the baryon asymmetry can be produced efficiently just after the preheating era, thus solving many of the problems that GUT baryogenesis had to face in the old picture of reheating.

There are several different ways to resurrect GUT baryogenesis. The simplest way is to take into account that, if particles produced at preheating can rapidly decay, then the reheating temperature may be very large, which may lead to the standard thermal production of superheavy  $X$  particles. However, if the products of parametric resonance are capable of an instantaneous decay and thermalization, then the parametric resonance never happens. Out of all possible ways of development of parametric resonance, nature chooses only those which do not lead to an instantaneous thermalization. In general, it does not preclude sufficiently high reheating temperature and subsequent baryogenesis, which may appear if the bosons produced at preheating decay and thermalize fast, but not fast enough to destroy the resonance. However, by assuming a thermal mechanism for  $X$  boson production, one is losing the advantage of the nonequilibrium nature of preheating, which may allow for a direct nonthermal creation of  $X$  bosons.

Indeed, preheating occurs because the interaction terms of the type  $\lambda_\phi \phi^2 |X|^2$  give the oscillating contribution  $\lambda \phi^2(t)$  to the mass squared of bosons interacting with the inflaton field. This leads to a broad parametric resonance in an expanding universe for  $\lambda_\phi \bar{\phi}^2 > M_\phi^2$ , where  $\bar{\phi}$  is the amplitude of the oscillating inflaton field [11]. The first stage of reheating does not extract all the initial energy of the inflaton field. As the amplitude of the oscillations of the inflaton field decreases, one leaves the resonance regime, and particle production ceases [11].

A crucial observation for baryogenesis is that even particles with mass larger than the inflaton mass may be produced during preheating. While previous studies of preheating concentrated on the creation of light particles, the results can be easily generalized to supermassive particles as well.

Following [11], during preheating, quantum fluctuations of the  $X$  field with momentum  $k$  obey the Mathieu equation:  $X_k'' + [A(k) - 2q \cos 2z]X_k = 0$ , where  $q = \lambda_\phi \phi^2 / 4M_\phi^2$ ,  $A(k) = (k^2 + M_X^2) / M_\phi^2 + 2q$ , and primes denote differentiation with respect to  $z = M_\phi t$ . Particle production in the broad resonance regime occurs above the line  $A = 2q$ . The width of the instability strip scales as  $q^{1/2}$  for large  $q$ , independent of the  $X$  mass. The condition for broad resonance,  $A - 2q \lesssim q^{1/2}$  [11], becomes  $(k^2 + M_X^2) / M_\phi^2 \lesssim \lambda_\phi^{1/2} \phi / 2M_\phi$ , which yields  $E_X^2 = k^2 + M_X^2 \lesssim \lambda_\phi^{1/2} \phi M_\phi / 2$ .

Therefore the typical energy of  $X$  bosons produced in preheating is  $E_X^2 \sim \lambda_\phi^{1/2} \bar{\phi} M_\phi$  [11]. At the end of

the broad parametric resonance, this equation somewhat changes because of the backreaction of produced particles. The resulting estimate for the amplitude of perturbations and for the typical energy of particles at the end of the broad resonance regime for  $M_\phi \sim 10^{-6} M_P$  is  $\langle X^2 \rangle^{1/2} \sim 10^{-1} \lambda_\phi^{-1/4} \sqrt{M_\phi M_P} \sim \lambda_\phi^{-1/4} 10^{15}$  GeV,  $E_X \sim 10^{-1} \lambda_\phi^{1/4} \sqrt{M_\phi M_P} \sim \lambda_\phi^{1/4} 10^{15}$  GeV [11,13].  $X$  bosons can be produced by the broad parametric resonance for  $E_X > M_X$ , i.e., for  $M_X < \lambda_\phi^{1/4} 10^{15}$  GeV. For  $\lambda_\phi \sim 1$ , one would have copious production of particles as heavy as  $10^{15}$  GeV, i.e., 100 times greater than the inflaton mass. In what follows we will consider the model with  $M_X = 10^{14}$  GeV. Such particles can be produced by parametric resonance for  $\lambda_\phi \gtrsim 10^{-3} - 10^{-4}$  [16]. The only problem here is that, for  $\lambda_\phi \gtrsim 10^{-6}$ , radiative corrections to the effective potential of the inflaton field may modify its shape at  $\phi \sim M_P$ . However, this problem does not appear if the flatness of the inflaton potential is protected by supersymmetry.

Thus we assume the first step in reheating is to convert a fraction  $\delta$  of the inflaton energy density into a background of baryon-number violating  $X$  bosons. They can be produced even if the reheating temperature to be established at the subsequent stages of reheating is much smaller than  $M_X$ . Here we see a significant departure from the old scenario. In the old picture, production of  $X$  bosons was kinematically forbidden if  $M_\phi < M_X$ , while, in the new scenario, it is possible as the result of coherent effects. The particles are produced out of equilibrium, thus satisfying one of the basic requirements to produce the baryon asymmetry [17].

The parametric resonance is efficient only if the  $X$  lifetime is greater than the typical time during which the number of  $X$  bosons grows  $e$  times. During the stage of broad parametric resonance this condition typically implies that the lifetime of the  $X$  is greater than about  $10M_\phi^{-1}$ . Assuming the width for  $X$  decay is  $\Gamma_X = \alpha_X M_X$ , this requires  $\alpha_X \lesssim 10^{-2}$ . This is certainly true if  $X$  decay into top quarks is kinematically forbidden. In the beginning of reheating, this condition is satisfied, e.g., if fermions acquire mass greater than  $M_X/2$  due to interaction with the inflaton field. At the end of reheating, the top quark mass receives a large nonthermal correction by means of the interaction with the  $X$  bosons [18],  $m_t \sim h_t \langle X^2 \rangle^{1/2} \sim h_t \lambda_\phi^{-1/4} 10^{15}$  GeV which is typically much greater than  $M_X$ . Also, one can always envisage the situation in which the  $X$  boson generating the baryon asymmetry does not belong to the same representation of the GUT group which gives mass to the third generation. Therefore, from now on we will assume that the  $X$  bosons may decay only to light fermions and that they decay well after the end of explosive particle production, resulting in a reheating temperature much smaller than  $M_X$ .

A self-interaction term in the Lagrangian of the type  $\lambda_X |X|^4$  also provides a nonthermal mass to the  $X$  boson of the order of  $(\lambda_X \langle X^2 \rangle)^{1/2}$ , which we assume to be

smaller than the bare mass  $M_X$ , i.e.,  $\lambda_X \lesssim 10^{-2} \lambda_\phi^{1/2}$ . However, this condition may be somewhat relaxed since the parametric resonance may occur even if the effective mass  $M_X$  grows in its process, because the same happens to the effective mass of the inflaton [11]. Self-interactions do not terminate the resonance effect since most particles remain inside the resonance shell; furthermore, creation of quanta different from  $X$ , e.g., gauge bosons, are suppressed by kinematical reasons if the nonthermal plasma mass of the final states is larger than the initial energy of the  $X$  particles. This happens if  $\lambda_\phi^{1/2} \lesssim g$  [18], where we denote by  $g$  the generic coupling constant between the final states and  $X$ .

The next step in reheating is the decay of the  $X$  bosons. We assume that the  $X$  decay products rapidly thermalize. It is only after this point that it is possible to speak of the temperature of the Universe.

The remaining energy in the inflaton is extracted in the final stage of the reheating process. After the parametric resonance period ends and  $X$  particle production shuts off, the inflaton performs small oscillations around the minimum of the effective potential, and the Universe soon becomes matter dominated. A slow process of particle production continues until the Hubble time becomes comparable to the inflaton decay time, and the inflaton decays. This part of the picture is similar to the old reheating scenario. Note that the above estimate of  $T_F$  did not depend upon the initial energy stored in the inflaton; it only assumed that the energy of coherent oscillations dominated the energy density.

As outlined above, we will consider a three part reheating process, with initial conditions corresponding to the frozen Universe at the end of inflation. The first stage is explosive particle production, where a fraction  $\delta$  of the energy density at the end of preheating is transferred to  $X$  bosons, with  $1 - \delta$  of the initial energy remaining in  $\phi$  coherent oscillation energy. We assume that this stage occurs within a few Hubble times of the end of inflation. The second stage is the  $X$  decay and subsequent thermalization of the decay products. We assume that decay of an  $X - \bar{X}$  pair produces a net baryon number  $\epsilon$ , as well as entropy. Reheating is brought to a close in the third phase when the remaining energy density in  $\phi$  oscillations is transferred to radiation.

The description simplifies if we assume zero initial kinetic energy of the  $X$ 's. This is a good approximation, since, for small  $\lambda_\phi$ , particles are produced with nonrelativistic velocities. We also assume that there are fast interactions that thermalize the massless decay products of the  $X$ . Then in a co-moving volume  $a^3$ , the total number of  $X$  bosons,  $N_X = n_X a^3$ , the total baryon number,  $N_B = n_B a^3$ , and the dimensionless radiation energy,  $R = \rho_R a^4$ , evolve according to

$$\begin{aligned} \dot{N}_X &= -\Gamma_X(N_X - N_X^{EQ}), & \dot{R} &= -aM_X \dot{N}_X, \\ \dot{N}_B &= -\epsilon \dot{N}_X - \Gamma_X N_B (N_X^{EQ}/N_0). \end{aligned} \quad (1)$$

$N_X^{EQ}$  is the total number of  $X$ 's in thermal equilibrium at temperature  $T \propto R^{1/4}$ , and  $N_0$  is the equilibrium number of a massless degree of freedom in a co-moving volume.

Figure 1 shows the results of an integration of Eqs. (1) in a toy model with  $M_\phi = 10^{13}$  GeV,  $M_X = 10^{14}$  GeV,  $\Gamma_X = 5 \times 10^{-6} M_X$ ,  $\Gamma_\phi = 5 \times 10^{-10} M_\phi$ , and two degrees of freedom ( $b$  and  $\bar{b}$ ). Initial conditions were chosen at  $a = a_I$  to be  $\rho_X = \rho_\phi \sim 10^{-4} M_\phi^2 M_P^2$ , and  $R = N_B = 0$ . The  $\rho_X = \rho_\phi$  assumption corresponds to  $\delta = 1/2$ . Since the number of  $X$  bosons produced is proportional to  $\delta$ , the final asymmetry is proportional to  $\delta$ . A more quantitative understanding of particle production in preheating is clearly required. However, we note that  $B/\epsilon \sim 10^{-9}$  can be obtained for  $\delta$  as small as  $10^{-6}$ .

The details of our scenario can be altered by many factors. For example, when the density of  $X$  particles decreases in the expanding Universe, the effective mass of the top quark also decreases, which may open the possibility of  $X$  decay to top quarks. Therefore, at some moment, the decay rate of the  $X$  bosons may suddenly increase. This will change some of our numerical results [18]. However, we believe that our simple model demonstrates the general behavior that might be expected in more realistic/complicated models. The baryon number  $B = n_B/s$  rapidly rises. However,  $B$  decreases as entropy is created and  $X$  inverse reactions damp the baryon asymmetry. After most of the energy is extracted from the initial  $X$  background, the baryon number is further damped as entropy is created during the decay of energy in the  $\phi$  background. In the model illustrated in Fig. 1, the final value of  $B/\epsilon$  is  $5 \times 10^{-4}$ .

We have numerically integrated the equation governing the number density of gravitinos  $n_{3/2}$  [5]. The result for  $G_{3/2} = n_{3/2}/s$  is shown in Fig. 1. Notice that, even though gravitinos are copiously produced at early stages by scatterings of the decay products of the  $X$ ,  $G_{3/2}$  decreases as entropy is created during the subsequent decay of energy in the  $\phi$  background. A similar behavior has

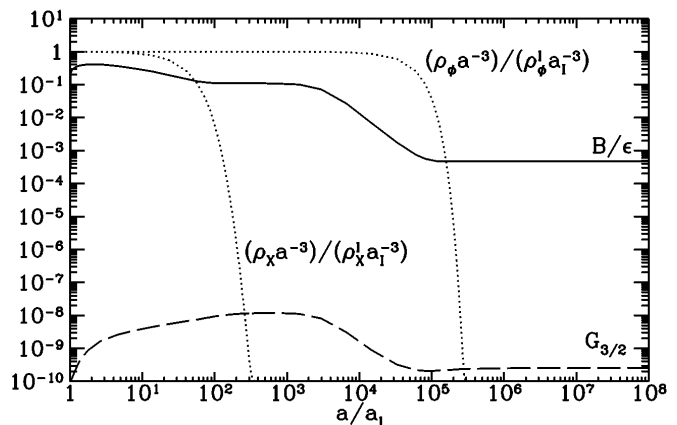


FIG. 1. The evolution of the baryon number, the  $X$  number density, the energy density in  $\phi$  oscillations, and the gravitino-entropy ratio as a function of the scale factor  $a$ .

been found in [19]. Successful nucleosynthesis requires  $G_{3/2} \lesssim 10^{-10}$  which translates into an upper bound on the inflaton decay rate,  $\alpha_\phi \lesssim 10^{-10}$ .

As  $X$  particles decay long after the end of the stage of preheating and their energy density is considerably diminished by the expansion of the Universe, the maximum of the thermalization temperature of their decay products is considerably smaller than the unification scale  $10^{16}$  GeV. This means that GUT symmetry is not restored when  $X$  decay products thermalize. If not the case, the subsequent decrease of the temperature of the thermal bath would be accompanied by a GUT symmetry breaking phase transition and the generation of dangerous topological defects. For the same reason, we require that GUT symmetry is not restored at the early stages of preheating, when nonthermal effects are dominant [13,15,20]. Let  $\Phi$  be the field responsible for GUT symmetry breaking with a potential of the form  $V(\Phi) = -\mu^2\Phi^2 + \lambda_\Phi\Phi^4$ ,  $\mu \sim 10^{16}$  GeV. The  $X$  boson may couple to the  $\Phi$  by an interaction of the type  $\lambda|\Phi|^2|X|^2$ . This interaction induces a mass squared for the  $\Phi$  field of order of  $\lambda\langle X^2 \rangle \sim 10^{-2}\lambda\lambda_\Phi^{-1/2}M_P M_\phi$  [13]. This term is smaller than  $\mu^2$  and does not lead to symmetry restoration for  $\lambda \lesssim 10^2\lambda_\Phi^{1/2}$  [18]. This condition is not difficult to satisfy. Therefore parametric resonance does not lead to GUT phase transitions and to the primordial monopole problem in our scenario.

In conclusion, we have shown that the present baryon asymmetry may be produced after inflation in the decay of nonthermal GUT bosons produced in preheating. Our scenario solves many of the serious shortcomings of GUT baryogenesis in the old theory of reheating, where it was kinematically impossible to produce superheavy particles after inflation. The out-of-equilibrium condition is naturally attained when superheavy quanta are produced in the regime of broad parametric resonance after the stage of inflation and considerably differs from the out-of-equilibrium condition in the GUT thermal scenario [6], where superheavy bosons decouple from the thermal bath when relativistic if  $K = (\Gamma_X/H)_{T=M_X} \ll 1$  and then decay, producing the baryon asymmetry. Gravitinos are subsequently diluted by the entropy released during the late decay of the inflaton field, and their abundance can be easily accommodated to be in agreement with the successful predictions of nucleosynthesis.

Our scenario is based on several assumptions about the structure of the theory and on relations between various coupling constants. For the parameters used to generate the results of Fig. 1, baryon number generation was relatively efficient:  $B/\epsilon \sim 5 \times 10^{-4}$ . Within uncertainties of model parameters, the value of  $\epsilon$ , etc., the present  $B \sim 10^{-10}$  may arise from GUT baryogenesis after preheating. Of course, additional work is needed to implement the ideas discussed above in the context of a more realistic model. However, we feel very encouraged that

recent progress in the theory of reheating has removed many obstacles which precluded successful GUT baryogenesis in inflationary cosmology. We will present more details in a subsequent publication [18].

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- [1] A. D. Linde, Phys. Lett. **108B**, 389 (1982); A. Albrecht and P. J. Steinhardt, Phys. Rev. Lett. **48**, 1220 (1982).
  - [2] A. D. Linde, Phys. Lett. **129B**, 177 (1983).
  - [3] A. D. Dolgov and A. D. Linde, Phys. Lett. **116B**, 329 (1982); L. F. Abbott, E. Fahri, and M. Wise, Phys. Lett. **117B**, 29 (1982).
  - [4] See, for instance, E. W. Kolb and M. S. Turner, *The Early Universe* (Addison-Wesley, Reading, MA, 1990).
  - [5] J. Ellis, J. E. Kim, and D. V. Nanopoulos, Phys. Lett. **145B**, 181 (1984). For recent work, see M. Kawasaki and T. Moroi, Prog. Theor. Phys. **93**, 879 (1995).
  - [6] For a review, see E. W. Kolb and M. S. Turner, Annu. Rev. Nucl. Part. Sci. **23**, 645 (1983).
  - [7] V. A. Kuzmin, V. A. Rubakov, and M. E. Shaposhnikov, Phys. Lett. **155B**, 36 (1985).
  - [8] E. W. Kolb and S. Wolfram, Nucl. Phys. **B172**, 224 (1980).
  - [9] J. Fry, K. A. Olive, and M. S. Turner, Phys. Rev. D **22**, 2953 (1980).
  - [10] D. V. Nanopoulos and S. Weinberg, Phys. Rev. D **20**, 2484 (1979); J. A. Harvey, E. W. Kolb, D. B. Reiss, and S. Wolfram, Nucl. Phys. **B201**, 16 (1982); T. Yanagida and M. Yoshimura, Nucl. Phys. **B168**, 534 (1980).
  - [11] L. A. Kofman, A. D. Linde, and A. A. Starobinsky, Phys. Rev. Lett. **73**, 3195 (1994); L. A. Kofman, A. D. Linde, and A. A. Starobinsky (to be published).
  - [12] Y. Shtanov, J. Traschen, and R. Brandenberger, Phys. Rev. D **51**, 5438 (1995); S. Yu. Khlebnikov and I. I. Tkachev, Phys. Rev. Lett. **77**, 219 (1996).
  - [13] L. A. Kofman, A. D. Linde, and A. A. Starobinsky, Phys. Rev. Lett. **76**, 1011 (1996).
  - [14] M. Yoshimura, Report No. TU-96-500, hep-ph/9605246 (to be published).
  - [15] I. I. Tkachev, Phys. Lett. B **376**, 35 (1996).
  - [16] This estimate based on analytical results of [11,13] was recently confirmed by a numerical investigation performed by S. Yu. Khlebnikov and I. I. Tkachev, Report No. hep-ph/9608458 (to be published).
  - [17] A. D. Sakharov, JETP Lett. **5**, 24 (1967).
  - [18] E. W. Kolb, A. Linde, and A. Riotto (to be published).
  - [19] H. Fujisaki, K. Kumekawa, M. Yamaguchi, and M. Yoshimura, Phys. Rev. D **54**, 2494 (1996).
  - [20] A. Riotto and I. I. Tkachev, Report No. FERMILAB-PUB-96-088-A, hep-ph/9604444 (to be published).