Baryogenesis in models with a low quantum gravity scale

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We make generic remarks about baryogenesis in models where the scale M_s of quantum gravity is much below the Planck scale. These correspond to M-theory vacua with a large volume for the internal space. Baryogenesis is a challenge, particularly for $M_s \lesssim 10^5$ GeV, because there is an upper bound on the reheating temperature of the Universe, and because certain baryon number violating operators must be suppressed. We discuss these constraints for different values of M_s, and illustrate with a toy model the possibility of using horizontal family symmetries to circumvent them. [S0556-2821(99)06410-3]

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I. INTRODUCTION

There are three experimental observations that might be considered as evidence for beyond-the-standard-model physics: neutrino oscillations [1], the baryon asymmetry of the Universe (BAU) [2], and the temperature fluctuations in the microwave background [3]. Any extension of the standard model (SM) must explain, or at least be consistent with, these data.

One of the reasons to attempt to extend the standard model is the possibility of unifying gravity with the other interactions. Present candidates are believed to be vacua of a single fundamental theory: M theory. The formulation of the latter seems to require adding new degrees of freedom. In a regime where a semiclassical description holds, these degrees of freedom manifest themselves as additional spatial dimensions compactified in an internal space. In its present form, M theory makes no prediction about the size of any spatial dimensions. It allows certain vacua with an arbitrary large size for the internal dimensions limited only by experimental data. If the states propagating in these dimensions have couplings with a size comparable to those of standard model gauge interactions, then the nonobservation of effects associated with Kaluza-Klein excitations leads to lower limits on the size of internal radii of the order of $\sim$ TeV [4]. If, in contrast, all the couplings of these Kaluza-Klein excitations are of the strength of gravitational interactions, then the limit is around 1 mm [5].[1] Mechanisms for the stabilization of the radii of the extra dimensions have been discussed in [6].

Allowing the presence of such large internal dimensions has dramatic effects on phenomenological aspects of M theory. Above the scale where the largest dimensions lie, naive dimensional analysis shows that the strength of gravitational interactions increases rapidly with energies. This implies that gravity and the three other known fundamental interactions will have the same strength and might unify at a scale M_s which can be very low, 1 TeV $\lesssim M_s \lesssim 10^{19}$ GeV. At M_s quantum gravity effects become important and new unknown phenomena might arise. Remnants of these phenomena at low energies are various nonrenormalizable effective operators. The size of the latter, if observed, might provide an indication of the existence and range of values of M_s.

This possibility of a low quantum gravity scale was first suggested in [7] with a scale M_s at $\sim 10^{16}$ GeV leading to unification within the minimal supersymmetric standard model of all the interactions. It was later observed that type I strings [8] (motivated also by a field theoretical proposal in [9] and for which model building was studied in [10]), M theory on S^1/Z_2 [11], and possibly heterotic strings [12] allow $M_s \sim$ TeV. This opens the exciting possibility that extra dimensions could be observed at future colliders [13]. Another proposal is to have M_s at an intermediary scale [11] so as to be associated with neutrino masses, observed ultrahigh energy cosmic rays, or the scale of breaking of a Peccei-Quinn symmetry. In this case the standard unification scenario might also be preserved [14].

In addition to the early phenomenological bounds for large internal dimensions discussed above, other limits on M_s have recently been derived [15] from astrophysical and cosmological considerations. The most significant particle physics constraint on M_s that we are aware of comes from atomic parity violation experiments [16], which determine $\sin^2\theta_W$ at low energy. If we assume[2] that the coefficient of the four-fermion vertex $4G_F/\sqrt{2}$ becomes $4G_F/\sqrt{2}+1/M_s^2$, we get $M_s > 4-6$ TeV. The strongest astrophysical bound estimated in [15] is from supernovas, and requires $M_s \gtrsim 30$ TeV in the case of two large compactified dimensions.

The purpose of this paper is to investigate the consequences of these models for baryogenesis. We will restrict our study to the class of models where matter and gauge fields reside on a (3+1)-dimensional wall and interact only through weak interactions of gravitational strength with fields residing in the $[(3+n)+1]$-dimensional ''bulk.'' The thermodynamics for the case with gauge interactions in higher dimensions (bulk) was recently studied in [17]. In the absence of a precise model, we introduce three mass parameters in various stages of our analysis. The first is M_s where

[1]Notice that the scale suppressing the interactions has increased by 15 orders of magnitude and the experimental limits went down by roughly the same amount.

[2]Note that we do not use the common $2\pi/\Lambda^2$ normalization of the new physics contribution to the four-fermion vertex. Had we done so, we would have found $M_s > 10-14$ TeV.

gravity unifies with the other interactions. It corresponds to the string scale in string models or to the 11-dimensional Planck mass in Hořava-Witten [18] compactification of M theory. The second is $m_{pl(4+n)}$ which is the Planck scale in $(3+n)+1$ dimensions. The relation between $m_{pl(4+n)}$ and M_s involves the volume of the dimensions with smaller radii. If the latter are of order M_s^{-1}, then $m_{pl(4+n)} \sim M_s$. Another parameter that we generically denote by Λ appears as a suppression scale for different nonrenormalizable operators. It is related to M_s through model-dependent coupling constants and numerical factors.

In Sec. II we discuss experimental bounds on nonrenormalizable baryon number violating operators, and which operators need to be forbidden for different values of Λ. In Sec. III we make some remarks about inflation, and discuss the upper bound on the reheat temperature of the Universe, $T_{reh} \ll M_s$, which follows from the production of gravitons in the large internal dimensions. Graviton production during the reheating period is dangerous as its decay products can lead to a greater than observed differential photon flux. In Sec. IV we discuss the difficulties of reconciling baryogenesis with the suppression of baryon number violating operators and the upper bound on the reheat temperature. We consider the possibility of generating baryon asymmetry in the out-of-equilibrium decay of a weakly coupled particle. To provide sizable decay channels we suggest using horizontal family symmetries to suppress dangerous nonrenormalizable operators instead of forbidding them through (discrete) gauge symmetries. A toy model for baryogenesis is formulated to illustrate this scenario. Sec. V summarizes our conclusions.

II. BARYON NUMBER VIOLATING OPERATORS

The presence of new physics at low scales could generate dangerous nonrenormalizable operators. These could, for instance, lead to unobserved baryon number violating processes such as proton decay and neutron-antineutron oscillations. In the absence of a precise model, where such operators can be computed, we make the conservative assumption that every operator that is not forbidden by a (possibly discrete) gauge symmetry could be generated with a coefficient of order 1.[3] This means that nonrenormalizable baryon number violating operators of dimension $4+d$ could appear, suppressed by factors of the scale of new physics, M_s. The precise coefficient of a $(4+d)$-dimensional operator will involve M_s, various coupling constants, and numerical factors, which we absorb into a coefficient Λ^{-d}.

A strong constraint on baryon number violating operators is that the proton must have a lifetime $\tau_p \gtrsim 10^{33}$ y [19]. If the quantum gravity scale is low, this means that one must forbid baryon and lepton number violating operators up to some large dimension [5,20]. For instance the operator $(QQQL)/\Lambda$ in supersymmetry (SUSY) generates proton decay at a rate of order [21]

$$\Gamma \sim 10^{-2} \frac{\alpha^2 m_p^5}{\Lambda^2 m_{SUSY}^2}, \tag{1}$$

which implies $\Lambda \gtrsim 10^{26}$ GeV. For nonsupersymmetric models, the operator is dimension 6 and the bound becomes $\Lambda \gtrsim 10^{15}$ GeV.

Another baryon number violating process that presents a significant constraint for low M_s is neutron-antineutron oscillations. This is a $\Delta B = 2$, $\Delta L = 0$ process that is generated by the dimension-9 operator $udsuds$. The ''lifetime'' for neutron-antineutron oscillations, $\tau_{n\bar{n}} > 1.2 \times 10^8$ s [22], is of order

$$\tau_{n\bar{n}} \simeq \frac{\Lambda^5}{5 \times 10^{-6}\ \mathrm{GeV}^6} \tag{2}$$

in the SM, where the denominator is an estimate of the hadronic matrix element [23,24]. This gives $\Lambda \gtrsim 10^5$ GeV.

A list of baryon and lepton number violating operators in the standard model and the minimal supersymmetric standard model (MSSM) is given in Table I with approximate bounds on the scale Λ. One must forbid with some symmetry all operators that are experimentally constrained to have $\Lambda > M_s$.

We follow [21] to calculate the constraints in the table. We take all supersymmetric particle masses and Higgs vacuum expectation values (VEVs) to be 100 GeV, and the hadronic matrix elements for proton decay to be $\sim 10^{-2}$ (with appropriate mass dimensions provided by the proton mass). The table is not particularly illuminating, because the bounds do not simply scale with dimension. Roughly, operators that violate B and L by one unit each are forbidden up to scales $> 10^{10}$ GeV, operators that violate B alone by one or two units are forbidden up to scales of order 10^5 (10^9) GeV in the SM (MSSM), and operators that violate B by three units are allowed at the TeV scale. An example of a symmetry that forbids $\Delta B = 1$ and 2 processes in the MSSM is the discrete anomaly-free Z_3 symmetry of Ibañez and Ross [25] which conserves B mod 3. The lowest baryon number violating operators it allows are combinations like $(QQQL)^3$, $(U^cU^cD^cE^c)^3$, and $(QQQH_1)^3$.

Note that the bounds on the operators in Table I are usually for first generation quarks and leptons. For low quantum gravity scales, some sort of flavor symmetry presumably should be imposed to remove FCNC operators; so one could imagine that there are flavor-dependent symmetries that forbid or suppress the dangerous baryon and/or lepton number violating operators. For instance, if the hierarchy in the Yukawa couplings is due to a spontaneously broken horizontal symmetry [26], the baryons and leptons can be assigned charges under this symmetry that forbid most of the operators in Table I (e.g., by giving all the SM fermions positive charges). We will discuss this possibility in Sec. IV B.

[3]In this work, we apply this assumption to B and L violating operators, but not, for instance, to flavor changing neutral current (FCNC).

TABLE I. B violating operators of dimension >4 for SM and MSSM particle content, in superfield notation. These are only the ''F terms.'' We list the dimension of the operators, the processes they contribute to, and the best bound we are aware of (in GeV), assuming that the coefficient of a dimension-$(d+4)$ operator is Λ^{-d}. The quark field subscripts are generation indices. We do not include operators of the form (allowed lower dimensional operator) $\times$ (forbidden lower dimensional operators), such as $LH_2H_1H_2$ or $U^cD^cD^cLH_1E^c$, because they are forbidden by whatever removes the unwanted lower dimensional operator.

Operator		Process	SUSY dimension	SUSY bound	SM dimension	SM bound
$Q_1Q_1Q_2L$	$\Delta B=\Delta L=1$	$p\to K\nu$	5	10^{26}	6	10^{15}
$U_1^cU_2^cD_1^cE^c$	$\Delta B=\Delta L=1$	$p\to K\nu$	5	10^{22}	6	10^{12}
$Q_1Q_1Q_2H_1$	$\Delta B=1$	$n-\overline{n}$	5	10^{9}	—	—
$U_1^cU_2^cU_3^cE^cE^c$	$\Delta B=1, \Delta L=2$	?	6		—	—
$U_1^cD_1^cD_2^cH_1H_2$	$\Delta B=1$	$n-\overline{n}$	6	10^{5}	—	—
$D_1^cD_2^cD_3^cLH_1$	$\Delta B=-\Delta L=1$	$n\to\nu\pi$	6	10^{13}	7	10^{9}
$U_1^cD_1^cD_2^cLH_2$	$\Delta B=-\Delta L=1$	$n\to\nu K$	6	10^{14}	7	10^{10}
$U_1^cD_1^cD_2^cU_1^cD_1^cD_2^c$	$\Delta B=2$	$n-\overline{n}$	7	10^{5}	9	10^{5}
$U_1^cD_1^cD_2^cLLE^c$	$\Delta B=-\Delta L=1$	$n\to e^+\mu^-\nu$	7	6×10^{7}	9	5×10^{5}
$U_1^cD_1^cD_2^cLQD^c$	$\Delta B=-\Delta L=1$	$n\to e^+\pi$	7	10^{7}	9	4×10^{5}
$U_1^cU_2^cD_1^cH_2LE^c$	$\Delta B=1$	$n-\overline{n}$	7	$\lesssim10^{3}$	—	—
$U_1^cU_2^cD_1^cH_2QD^c$	$\Delta B=1$	$n-\overline{n}$	7	$\lesssim10^{3}$	—	—
$QQQLLH_2$	$\Delta B=1, \Delta L=2$	?	7	?	—	—
$Q_1Q_1Q_2H_1Q_1Q_1Q_2H_1$	$\Delta B=2$	$n-\overline{n}$	9	10^{4}	11	10^{4}

III. INFLATION AND REHEATING

A. Inflation

A period of inflation is the only known way of generating the temperature fluctuations measured on scales up to 100 Mpc in the microwave background. Since inflation dilutes any preexisting asymmetries, the observed baryon asymmetry of the Universe must be generated afterwards. As we will see, there is an upper bound on the reheat temperature in models with low quantum gravity scale; so the phase transition out of inflation is one of the few places where one can find the out-of-equilibrium dynamics required for baryogenesis.

If we take the energy density of the Universe to be at most M_s^4, then for $n\geqslant 2$ large internal dimensions, the Hubble radius is greater than or equal to the radius of the n dimensions. This means that it is consistent to build an inflation model in 3+1 dimensions. However, a second order inflation model at a scale $\ll 10^{15}$ GeV requires a great deal of fine-tuning to get enough e-foldings and density perturbations of order 10^{-5}. The latter can be estimated as

$$\frac{\delta\rho}{\rho}\sim\frac{V^{3/2}}{m_{pl}^3V'}, \tag{3}$$

where V is the potential energy density of the inflaton, $V'=dV/d\phi$, and both of these are evaluated at the point in the potential where the inflaton was sitting 50–60 e-folds before the end of inflation. If $V\sim M_s^4$, then

$$\frac{V'}{M_s^3}\sim 10^5\left(\frac{M_s}{m_{pl}}\right)^3; \tag{4}$$

so the potential must be very flat. If, for instance, one parametrizes $V=V_0-m^2|\phi|^2+\lambda|\phi|^4+\Sigma\,\phi^{n+4}/M_s^n$, with $V_0\sim M_s^4$, then to get enough inflation [27,28] and the right sized density perturbations, one finds $m\sim M_s^2/m_{pl}$. For $M_s\sim$ TeV, one gets $m\sim 10^{-13}$ GeV. Such a light inflaton might have difficulties reheating the Universe to temperatures ~ 1 MeV, and in any case, $V_0\sim m^4 \ll M_s^4$; so our initial assumptions were inconsistent. To avoid this difficulty, one can build two-field or hybrid inflation models [28] where the mass of the inflaton when it decays is not related to the mass term in the potential when it is generating density perturbations. An *ad hoc* potential of the form $V_0-a_6\phi^6/M_s^2+a_{12}\phi^{12}/M_s^8$ also works, for $a_6\sim a_{12}\sim 10^{-2}$ and $M_s\sim 10$ TeV. For the rest of this work, we will assume that the potential is flat enough to inflate for long enough, and that the mass of the inflaton when it decays might be greater than 1 GeV. This is useful for baryogenesis, if we want to generate the asymmetry in the decay of the inflaton.

B. Gravitons production constraints on T_{reh}

The Universe must at some point get out of its inflationary phase and reheat to a plasma of particles. A safe reheat temperature T_{reh} to ensure that primordial nucleosynthesis takes place as usual is $\gtrsim$ 3 MeV [29]. Baryogenesis at such a low energy scale is hard; so a higher T_{reh} would be desirable.

Getting a high T_{reh} is a challenge in low quantum gravity scale models where matter resides on a (3+1)-dimensional ''wall,'' while gravitons and other very weakly interacting particles reside in the $[(3+n)+1]$-dimensional ''bulk.'' The temperature to which the Universe reheats must be low,

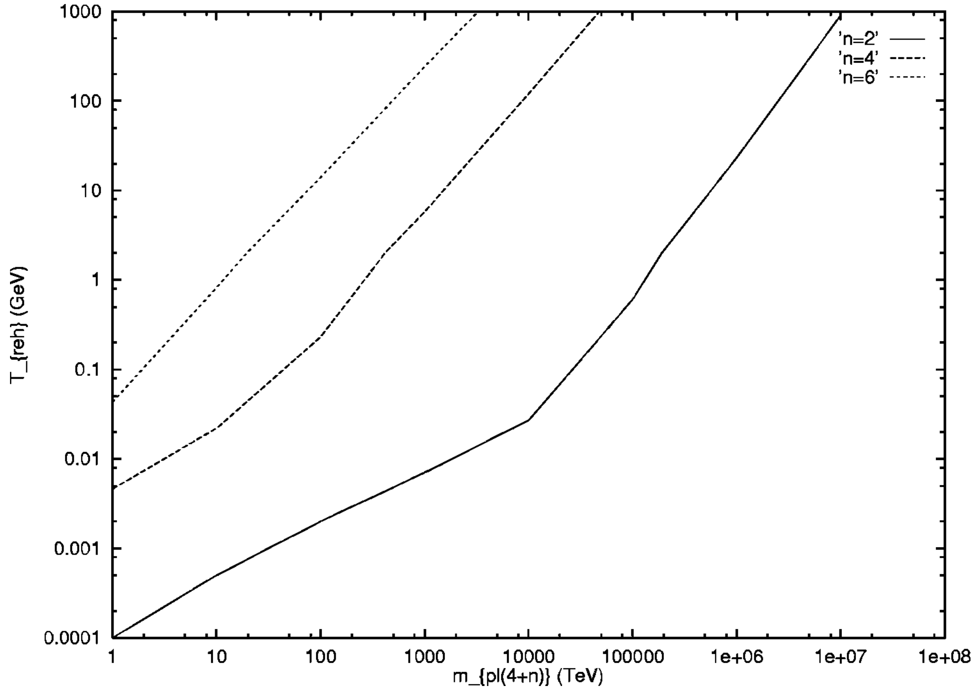

FIG. 1. Maximum allowed reheat temperature T_{reh} as a function of $m_{pl(n+4)}$ for different numbers n of large extra dimensions.

to avoid generating too many ''bulk particles'' (we will generically refer to them as gravitons) in the extra large dimensions. These gravitons can decay into particles in our (3+1)-dimensional boundary. We can set bounds on the number of these decay products from various observations, and therefore set an upper bound on the number of gravitons allowed or, equivalently, an upper bound on the reheat temperature T_{reh}. Below, we estimate this bound as a function of the quantum gravity scale and the number of large extra dimensions.

The behavior of gravitons when $M_s\sim$ TeV was discussed in [15]. Their best bound comes from requiring that photons from graviton decay do not generate a spike in the $E \gg 2.7$ K photon background. For larger M_s, fewer gravitons are produced, and so higher reheat temperatures are allowed. However, as the graviton lifetime becomes shorter, the decay products arrive in our 3+1 dimensions at earlier epochs; so the limit on their number density changes. If the gravitons decay between recombination and today, the photons produced will be in the present photon background. For some period before recombination, photon number changing interactions in the thermal plasma are out of equilibrium; so photons from graviton decay produced at this time would generate a chemical potential for the microwave background. If the gravitons decay before recombination but after nucleosynthesis, they can dissociate light elements. The bound from this is similar to the one from the chemical potential. Gravitons that decay before nucleosynthesis are not a problem. We discuss bounds for all cases below.

One assumption made is that translation invariance in the bulk is broken only at the boundaries. This allows us to speak about momenta and energy of particles residing in the bulk. Such a situation is not generic as the size of other dimensions of the internal space might become larger when going away from our wall towards a hidden one (see, for instance, [11]). It was argued in [15] that gravitons might decay earlier on the hidden wall than on the observable wall, avoiding some of our constraints. We will discuss this situation elsewhere [30].

Consider first the number density n_G of gravitons produced in the bulk. We follow [15] (a similar analysis was done in [31]), and assume that the cross section for particles on the wall to produce gravitons in the n extra large dimensions is of order[4] $\sigma_{\gamma\gamma\to GG}\sim T^n/m_{pl(n+4)}^{n+2}$; so the rate at which gravitons are made is approximately

$$\frac{\partial n_G}{\partial t}-3Hn_G=\sigma n_\gamma \sim \frac{T^{n+6}}{m_{pl(n+4)}^{n+2}}, \tag{5}$$

where H is the Hubble expansion rate $H^2=8\pi\rho/3m_{pl}^2$ and n_γ is the number density of photons. Gravitons made at a temperature T will have momenta in the bulk of order T, and since these momenta do not redshift, the energy of the gravitons remains $\sim T$. The number density of gravitons with energy T at later times (when the photon temperature is T_γ) will therefore be of order

$$n_G(E=T)\simeq \sigma n_\gamma H^{-1}(T)=N\frac{m_{pl}T^{n+1}}{m_{pl(n+4)}^{n+2}}T_\gamma^3, \tag{6}$$

where N is a numerical coefficient which we have not calculated, and m_{pl} is the (3+1)-dimensional Planck mass. We

[4]It is the $(4+n)$-dimensional ''Planck scale'' $m_{pl(n+4)}$ that appears; if we assume that the other internal dimensions have size of the order of M_s^{-1}, then $m_{pl(n+4)}\sim M_s$.

take $N=1$ in Fig. 1. The number and energy of the gravitons increase with T; so the most troublesome ones are those generated at the reheat temperature T_{reh}. We concentrate on these and consider constraints for different values of $m_{pl(n+4)}$.

For the lowest values of M_s, the strongest constraint obtained in [15] on the number density of gravitons is from the decay of gravitons back into photons. We review this bound here. The gravitons of energy E decay to photons of energy $\sim E$ at a rate [32]

$$\Gamma_G = \tau_G^{-1} = D\frac{E^3}{m_{pl}^2}, \tag{7}$$

where D is another unknown numerical factor that we set to 1 in Fig. 1. For $E\sim T_{reh}\lesssim 60D^{-1/3}$ MeV, the lifetime of the graviton τ_G is longer than the age of the Universe, τ_U. The number that will have decayed is therefore of order $n_{G0}\tau_U/\tau_G$. Following [33], one can require that the flux of photons of energy T_{reh} from these decays not exceed the observed differential photon flux $\mathcal{F}$:

$$\frac{n_{G0}}{4\pi}\frac{\tau_U}{\tau_G}\lesssim \mathcal{F}(E) = \frac{\text{MeV}}{E}\ \text{cm}^{-2}\,\text{sr}^{-1}\,\text{sec}^{-1}, \tag{8}$$

where E is the photon energy. This gives

$$\frac{ND}{6\pi}\frac{T_0^3}{H_0 m_{pl}}\left(\frac{T_{reh}}{m_{pl(n+4)}}\right)^{n+2}(T_{reh})^2 < \mathcal{F}(T_{reh}), \tag{9}$$

where T_0 is the microwave background temperature today. This implies

$$(T_{reh})^{n+5} < \frac{7\times 10^{-39}}{ND} m_{pl(n+4)}^{n+2}\ \text{GeV}^3 \qquad (\text{for } T_{reh}<60\ \text{MeV}). \tag{10}$$

For $n=2$ and $T_{reh}\gtrsim 3$ MeV (a safe reheat temperature to ensure that primordial nucleosynthesis takes place as usual [29]), we get $m_{pl(n+4)}>100$ TeV.

For $T_{reh}>60D^{-1/3}$ MeV, the gravitons created at T_{reh} can decay before today. All their energy is therefore in the photon background, but redshifted from when they decayed until now. If this took place after recombination, we can set a bound by requiring that their final products do not exceed the observed photon flux $\mathcal{F}$. The photon temperature T_d when the gravitons decay can be computed from

$$H(T_d)\simeq \frac{2T_{eq}^{1/2}T_d^{3/2}}{m_{pl}} \simeq \Gamma_G \simeq D\frac{T_{reh}^3}{m_{pl}^2}, \tag{11}$$

where $T_{eq}\sim 3$ eV is the photon temperature at matter-radiation equality. This gives

$$T_d \simeq \left(\frac{D}{2}\right)^{2/3}\frac{(T_{reh})^2}{m_{pl}^{2/3}T_{eq}^{1/3}}. \tag{12}$$

The photon flux expected from graviton decay is therefore

$$\frac{n_{G0}}{4\pi}\frac{T_0}{T_d} \simeq \left(\frac{2}{D}\right)^{2/3}\frac{N}{4\pi}T_0^4 m_{pl}^{5/3}T_{eq}^{1/3}\left(\frac{T_{reh}}{m_{pl(n+4)}}\right)^{n+2}(T_{reh})^{-3}\lesssim \mathcal{F}. \tag{13}$$

This gives

$$(T_{reh})^n < 3\times 10^{-33}\frac{m_{pl(n+4)}^{n+2}}{\text{GeV}^2} \quad (60\ \text{MeV}<T_{reh}<2\ \text{GeV}). \tag{14}$$

This applies for $\tau_U>\tau_G>t_{recomb}$, which corresponds to the limit in parentheses (with $D=1$).

Photon number changing interactions of the form $\gamma e \to e\gamma\gamma$ go out of equilibrium at $t_\gamma\sim 10^5$ sec. If the gravitons decay after t_γ, but before recombination, the photons they decay to will induce a chemical potential[5] for the microwave background [34]:

$$\mu \simeq \frac{\rho_G}{\rho_\gamma}. \tag{15}$$

This is in the instantaneous decay approximation, where all the energy of the gravitons is deposited into the photons at $t=\tau_G$. This should be a reasonable approximation for $t_\gamma \ll \tau_G \ll t_{recomb}$ [35]. The present experimental bound is [36] $\mu<3.3\times 10^{-4}$, which implies

$$N\frac{m_{pl}T_{reh}^{n+2}}{m_{pl(n+4)}^{n+2}}\frac{T_\gamma^3}{\rho_\gamma} < 3.3\times 10^{-4} \tag{16}$$

when the gravitons decay. The photon temperature at decay T_d can be determined from $H(T_d)\sim\Gamma_G \simeq DT_{reh}^3/m_{pl}^2$; so one gets

$$m_{pl(n+4)}^{n+2} > 4\times 10^{32}(T_{reh})^{n+1/2}\ \text{GeV}^{3/2} \qquad (2\ \text{GeV}\ll T_{reh}\ll 1\ \text{TeV}). \tag{17}$$

This applies for 10^5 sec$\sim t_\gamma \ll \tau_G \ll t_{recomb}\sim 10^{13}$ sec, or 2 GeV$\ll T_{reh}\ll 1$ TeV.

One of the successes of the big bang model is that it predicts the correct abundances of light elements. ^{4}He, ^{3}He, D, and ^{7}Li are synthesized in the early Universe at temperatures just below 1 MeV, in about the right numbers to agree with present observations [37]. If the gravitons decay after nucleosynthesis, one must check that the their decay products do not destroy or produce too many of these light nuclei. This constraint has been calculated for various particles [38–40]. There are numerical bounds on ρ_G/n_B in [40] for 10^4 sec$<\tau_G<10^7$ sec, which we can simply translate into bounds on M_s as a function of T_{reh}. These turn out to be similar or weaker than Eq. (17).

In Fig. 1 we plot the allowed reheat temperature as a function of the $(4+n)$-dimensional Planck scale $m_{pl(n+4)}$ for different numbers[6] of extra dimensions $n\geqslant 2$. This is a fairly stringent bound; to get a reheat temperature as large as 100 GeV, we need $m_{pl(n+4)}\sim 10^6$ GeV for six extra large dimensions and $m_{pl(n+4)}\sim 10^{10}$ GeV for $n=2$. If the reheat temperature is less than 100 GeV, electroweak baryogenesis

[5]The dimensionless parameter μ is defined as the parameter in the Bose-Einstein distribution function: $1/(e^{E/T+\mu}+1)$.

[6]The case of one extra dimension at 1 mm leads to $T_{reh}\lesssim 10$ MeV, which is easily compatible with primordial nucleosynthesis.

[41] and leptogenesis [42] (generating a lepton asymmetry and then having the ''sphalerons'' reprocess it) are impossible. If $T_{reh} \gg$ TeV, the gravitons generated at T_{reh} will decay before nucleosynthesis and thermalize rapidly; so they are not a problem.

IV. BARYOGENESIS

A. Challenges for baryogenesis models

First let us consider the consequences of the low T_{reh} constraint. For a large choice of M_s and of the number of large internal dimensions, the reheat temperature must be less than ~ 100 GeV; so the electroweak $B+L$ violating processes are not available for baryogenesis. This means that electroweak baryogenesis [41] and leptogenesis [42] are not possible. For larger values of M_s and depending on n, $T_{reh} \gtrsim 100$ GeV is allowed and electroweak baryogenesis is possible. This is attractive because the nonperturbative electroweak $B+L$ violation proceeds through the operator $(qqql)^3$, which does not mediate proton decay because it has $\Delta B = 3$ (as well as being exponentially small at zero temperature).

There has recently been a very interesting suggestion [43] that the BAU could be generated at the QCD phase transition using purely standard model physics (the baryon number and CP violation are spontaneous or nonperturbative). If this model works, then one only needs a reheat temperature of order 1 GeV, which is easier to achieve than 100 GeV, as one can see from Fig. 1. We do not further discuss this mechanism, but it should be kept in mind as a possible way of generating baryon asymmetry in low quantum gravity scale models.

The low T_{reh} creates a generic difficulty. One of the Sakharov [44] conditions for baryogensis is that one needs some out-of-equilibrium dynamics. This can be found at phase transitions or when some interaction is not fast enough to keep up with the expansion of the Universe. However, when the temperature (or energy density) of the Universe is low, the expansion rate is too ($H \sim 10^{-18} T$ at $T \sim 1$ GeV); so interactions have no difficulty keeping up with the expansion. Getting the out-of-equilibrium dynamics anywhere but a phase transition is hard. If the reheat temperature is less than ~ 0.1 GeV, then the only phase transition available appears to be the one out of inflation.

Another difficulty for baryogenesis models is the bounds on baryon number violation discussed in Sec. II. For instance, to avoid fast proton decay through $|\Delta B| = |\Delta L| = 1$ operators and neutron-antineutron oscillations through $\Delta B = 2$ operators, one may assume that B is conserved mod 3. This is problematic for scenarios where the BAU is generated in the out-of-equilibrium decay of a particle X. X must have at least two decay modes with different baryon number in the final state and approximately the same branching ratios [45]. Otherwise, the baryon asymmetry generated will be small.[7] If B is conserved mod 3, then X must decay to final states with $B=1$ and with $B=2$ (or $B=0$ and $B=3$), so that X exchange generates a vertex than conserves B mod 3. But $B=2$ operators are of higher dimension than $B=1$ operators (see Table I); so the branching ratio of X to the $B=2$ final state will be very small. We tried imposing B mod 4, so that X could decay via $\Delta B = 2$ and $\Delta B = -2$ processes, but these operators are of dimension 10 and 12, so that X must have a mass of order 100 GeV to decay before nucleosynthesis.

If the quantum gravity scale is greater than 10^5 GeV in the SM (10^9 GeV in the MSSM), then $\Delta B = 2$ operators do not need to be suppressed or forbidden (see Table I). In this case, B does not need to be conserved, provided that L is; if there are only baryon number violating couplings, and the low energy theory has standard model particle content, the proton cannot decay. This means, for instance, that in SUSY models one can use the interaction $U^c D^c D^c$ to provide the baryon number violation required for baryogenesis. Such a model of low reheat temperature baryogenesis was constructed in [46], where the inflaton decay products include squarks, which then decay via their B violating coupling. They decay before they have time to thermalize or annihilate and so are out of equilibrium and can generate a baryon asymmetry in their decay.

B. Contrived baryogenesis model

Suppose that we are in the ''worst case scenario'' for baryogenesis. This corresponds to the situation with $M_s \lesssim 10^5$ GeV; so symmetries are required to forbid the $n-\bar{n}$ operator $udsuds$ and the fast proton decay vertices. The maximal allowed reheat temperature is much less than 100 GeV; so there is no electroweak $B+L$ violation available. If the motivation for having a low M_s is to solve the hierarchy problem, we can also assume that there is no supersymmetry, since this is also what it is for. This means that Affleck-Dine baryogenesis is not possible. Can the baryon asymmetry be generated in these circumstances?

We first try to construct an out-of-equilibrium decay scenario. For this we need a particle X that decays out of equilibrium to final states with different baryon numbers, with enough CP violation in the decay rates to generate a baryon to photon ratio $\eta \sim 3 \times 10^{-10}$.

Suppose X is the inflaton. This has the advantage that it decays out of equilibrium. Moreover, its width

$$\Gamma_X \sim \frac{T_{reh}^2}{m_{pl}} \tag{18}$$

[7]This is a consequence of CPT: if X decays to a $B=B_1$ final state with a large branching ratio $1-\epsilon$ and a $B=B_2$ state with a small branching ratio ϵ, then one can assign $B=B_1$ to X; so the larger decay is baryon number conserving. By CPT the total decay rates of X and $\bar{X}$ are equal; so the baryon asymmetry created will be proportional to $\epsilon - \bar{\epsilon}$ and therefore very small.

must be small in order to obtain a low reheat temperature.[8] One way to ensure that it has a long lifetime is to make it decay via nonrenormalizable operators. For instance, this can happen via an operator of dimension $4+d$ with coefficient λM_s^{-d}, so that $\Gamma \sim \lambda^2 m_X^{2d+1}/M_s^{2d}$. We would like X to have baryon number violating decays so that it can generate the baryon asymmetry, which also means that X should decay via nonrenormalizable interactions. As it oscillates about its minimum, we suppose that $m_X > 1$ GeV, so that it can produce protons.

Another possibility is that X is a particle generated in the reheating process, with a number density $n_X = \delta n_\gamma$. The annihilation rate for X will be

$$\Gamma_{ann} \sim n_X \sigma_{X\bar{X} \to anything}. \tag{19}$$

If we take $\sigma_{X\bar{X} \to anything} \sim 4\pi\alpha^2/M_s^2$, then requiring that $\Gamma_{ann} < H$ gives

$$4\pi\alpha^2\delta < \frac{M_s^2}{T_{reh} m_{pl}}. \tag{20}$$

If we take M_s to be its minimum value $\gtrsim 3$ TeV and T_{reh} the maximum value possible for $n \le 6$ and $M_s < 10^5$ GeV which is $\lesssim 10$ GeV, then this gives $\alpha^2\delta < 10^{-14}$. This is the condition such that X annihilations will be out of equilibrium at the reheat temperature and thereafter; so all the X's will decay.

We would like to address the possibility of having particles with such small couplings. Consider, for instance, models obtained from type I′ strings after performing T duality on all the internal directions of a type I model. There are two kind of p-branes in these models: three-branes and seven-branes. We assume that the standard model resides on the three-branes with gauge couplings of order 1. The particles that arise from seven-branes have gauge couplings suppressed by the volume of the four-dimensional internal space on which they are wrapped. The corresponding couplings can be arranged to satisfy the above constraints.

To generate the BAU, X needs similar branching ratios to states with different baryon number. As discussed in the previous subsection, this requirement is difficult to implement in models where B is conserved. So instead we consider the possibility that B is not conserved, L is conserved mod 2 (which allows neutrino masses), and there is a horizontal symmetry that suppresses the dangerous $\Delta B = 2$ operators.

We assume that the SM Yukawa couplings are generated by some horizontal U(1) gauge symmetry [26], which is spontaneously broken below M_s. The quarks (q, u^c, d^c) and leptons (l, e^c) carry positive charges under this symmetry, and the charges are higher for the lighter fermions. The Higgs boson that breaks the horizontal U(1) with VEV θ carries negative charge. By choosing the horizontal charges of the fermions Q_f^H with care, one can generate approximately the right structure for the Yukawa matrices, because

[8] We assume that the inflaton is very weakly coupled and so cannot decay by parametric resonance.

TABLE II. Possible charges for the fermions under the horizontal U(1), for three generations. The first generation is u,d,e, and so on. These charges generate approximately the right Yukawa couplings.

Generation	q	u^c	d^c	l	e^c
1	5	5	2	1	5
2	4	2	1	0	4
3	1	0	1	0	1

the interaction $u^c u H \sim m_u \bar{u} u$ appears multiplied by $(\theta/\Lambda)^{Q_{u^c}^H + Q_u^H}$ and $t^c t H \sim m_t \bar{t} t$ appears multiplied by $(\theta/\Lambda)^{Q_{t^c}^H + Q_t^H}$. Such a mechanism is probably required in models with a low M_s to avoid FCNC. It will also suppress the problematic $u^c d^c s^c u^c d^c s^c$ operator: at M_s where θ is zero, it is forbidden by the horizontal symmetry (if all the fermions are positively charged), and once the horizontal symmetry is broken, $u^c d^c s^c u^c d^c s^c$ can appear suppressed by $(\theta/\Lambda)^{2(Q_{u^c}^H + Q_{d^c}^H + Q_{s^c}^H)}$. For $\theta/\Lambda \equiv \epsilon \sim 0.2$ and the charges in Table II, the operator $u^c d^c s^c u^c d^c s^c$ will be multiplied by ϵ^{16}, which is compatible with the experimental limit for $\Lambda \gtrsim$ few TeV. The proton is stable enough provided that L is conserved mod 2.

Suppose that X is a light (~ 10 GeV) gauge singlet scalar with $L=1$. It can decay to SM particles via the dimension-7 operators $Xqqql$ and $Xu^c u^c d^c e^c$. These violate B, respectively, by 1 and -1 units; so a baryon asymmetry could be generated. We suppose that the fermions have the charges under the horizontal U(1) that are listed in Table II. In this case the principle decay rates will be

$$\Gamma_{\bar{p}} \sim \epsilon^{18} \frac{m_X^7}{\Lambda^6}, \quad X \to c^c \ u^c \ b^c \ \tau^c \quad (\bar{D} \ \bar{B} \ \bar{p} \ \tau^+), \tag{21}$$

$$\Gamma_p \sim \epsilon^{18} \frac{m_X^7}{\Lambda^6}, \quad X \to c \ s \ b \ \nu_\tau \quad (D \ B \ K \ p \ \nu_\tau), \tag{22}$$

$$\Gamma_{p2} \sim \epsilon^{20} \frac{m_X^7}{\Lambda^6}, \quad X \to c \ d \ b \ \nu_\tau \quad (D \ B \ p \ \nu_\tau), \tag{23}$$

where $\epsilon = \theta/\Lambda \sim 0.2$. We neglect kinematics, factors of 4π, and so on; so these are very approximate estimates. However, for $\Lambda \sim 3$ TeV and $T_{reh} \sim 3$ MeV, Eq. (18) gives $m_X \sim 25$ GeV. This is heavy enough to decay to B and D mesons, but light enough to (possibly) be produced in the reheating process or to be the inflaton. For larger Λ, we would need a larger m_X.

We have shown that we can construct a scalar particle X that decays before nucleosynthesis, at about the right time to reheat the Universe if it were the inflaton. We now need to consider whether a sufficient baryon asymmetry can be generated in the decays. We assume that $\Gamma_p \gg \Gamma_{p2}$, and so we neglect Γ_{p2} and all other smaller decay modes. The net number of baryons produced per X particle will be

$$\frac{n_b}{n_X} \simeq \frac{\Gamma_p - \Gamma_{\bar{p}} + \bar{\Gamma}_{\bar{p}} - \bar{\Gamma}_p}{\Gamma_p + \Gamma_{\bar{p}}} \equiv \theta_{CP}, \tag{24}$$

where $\bar{\Gamma}$ is the CP conjugate decay. The baryon-to-photon ratio $n_b/n_\gamma \equiv \eta \simeq 3\times 10^{-10}$ [37] will be

$$\eta \simeq \frac{n_X}{n_\gamma} \theta_{CP}. \tag{25}$$

If X is the inflaton, then $n_X/n_\gamma \sim T_{reh}/m_X \sim 10^{-3}$. If X is produced in the reheating process, then $n_X/n_\gamma = \delta$ is a model-dependent parameter. One would not expect to make more than one or two X's in the decay of each inflaton; so in this case $\delta \lesssim 10^{-3}$. This means that we need $\theta_{CP} \gtrsim 10^{-7}$. If we assume that CP violation arises through loop corrections involving new particles at the scale M_s, then $\theta_{CP} \sim (m_X/\Lambda)^2 \sim 10^{-6}$, which is approximately right.

The family symmetry presented here obviously suffers from anomalies. These might be canceled in two different ways. The first is to assume that massive particles in a hidden sector are charged under this U(1), standard model symmetries, and some hidden gauge group. The hidden symmetry might suppress any undesirable nonrenormalizable operator. Another possibility is to appeal to a Green-Schwarz mechanism to cancel the anomaly [26]. If the gauge couplings are all given by the vacuum expectation value of a single modulus (dilaton), then anomaly cancellation implies particular tree level relations between the couplings. For the model at hand, the strong, weak, and hypercharge U(1) couplings are in the ratio 1:1:105/33 at $M_s \sim 1$ TeV instead of the usual relation 1:1:5/3 at 10^{16} GeV. To compare the tree level prediction with experimental measurements we need to know the precise evolution of coupling constants with energy from M_s down to M_Z. Unfortunately for low M_s there is not yet a framework to discuss this running of couplings as these become very sensitive to the spectrum at energies of the order of 1 TeV.[9]

We also imposed L as a spontaneously broken symmetry, to ensure proton longevity; so some additional (heavy) leptons must be included to cancel the anomalies in L [25].

C. Other possibilities

It is clear from the previous section that out-of-equilibrium decay scenarios do not work easily at low scales with SM particle content. Electroweak baryogenesis and leptogenesis will not work in their standard versions if T_{reh} is much below the temperature at which electroweak $B+L$ violation is in equilibrium ~ 100 GeV. However, there are many other baryogenesis mechanisms [2], some of which may work naturally. We will discuss these in a later publication [30]. The most popular mechanism that we have not discussed is the Affleck-Dine mechanism [48], which generates an asymmetry in the cosmological evolution of spartner VEVs. This scenario could be attractive for supersymmetric low string scale models because the reheat temperature can generically be low, and the dimension of the B violating operators is not so relevant. However, the difficulty is that at the end of inflation the spartner VEV should have the same phase over the whole observable Universe.[10] This is not so easy to arrange if the expansion rate H is much smaller than the flat direction's mass $m \gtrsim 100$ GeV, because inflation cannot push the VEV out along the flat direction; so a large VEV that is coherent across the whole Universe could be difficult to generate. It may be possible to resolve this with a small amount of external CP violation. We will pursue this possibility in a subsequent publication [30].

V. CONCLUSION

For traditional models, where the scale M_s of quantum gravity lies far away at energy scales of the order of 10^{19} GeV, the baryon asymmetry can be generated in a plethora of scenarios. In contrast, we found that exhibiting simple scenarios for baryogenesis becomes a challenging problem when $M_s \lesssim 10^5$ GeV. The three Sakharov requirements of baryon number violation, C and CP violation, and out-of-equilibrium dynamics must be satisfied. Baryon number violation is hard to come by because many baryon number violating operators must be forbidden by a symmetry to ensure that they are not generated at M_s. Out-of-equilibrium dynamics is also difficult because there is an upper bound on the reheat temperature of the Universe from requiring that one not overproduce gravitons in the extra large dimensions. We list experimental bounds on baryon number violating operators in Table I, and in Fig. 1 we plot the maximum allowed reheat temperature as a function of $m_{pl(4+n)}$ for different numbers n of large internal dimensions. The T_{reh} bound could possibly be avoided if the bulk fields (gravitons) could decay faster to hidden matter whose energy redshifts.

Standard electroweak baryogenesis and leptogenesis are excluded for low M_s, because the reheat temperature is constrained to be less than 100 GeV. Affleck-Dine baryogenesis is difficult because the Hubble expansion rate is not large enough to drive the flat direction field out to a single VEV with the same phase everywhere.

Out-of-equilibrium decay models are also problematic; the experimental bounds on baryon number violating operators suggest that baryogenesis must proceed through nonrenormalizable operators of very high dimension. An alternative is to suppress baryon number violating operators through a horizontal family symmetry, and ensure that the proton remains stable by conserving L. We implement this idea in a toy model that could generate the correct baryon asymmetry in the decay of a weakly coupled particle (possibly the inflaton).

For larger values of $m_{pl(4+n)}$ we need SUSY to solve the hierarchy problem, in which case the Affleck-Dine mechanism is a possibility. If $m_{pl(4+n)} \gtrsim 10^5$ GeV, baryon number

[9]See [11,47] for a discussion of unification in these models.

[10]The CP violation in the Affleck-Dine scenario is ''spontaneous,'' that is, encoded in the relative phase between the VEV and baryon number violating bumps in the potential.

violation is allowed, provided that L is conserved. For scales $M_s \gtrsim 10^{10}$ GeV the reheat temperature is large and electroweak baryogenesis is possible.

We will return to discuss these issues in a future publication [30].

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[1] For recent data and models on the solar and atmospheric neutrino deficits, see, e.g., *Neutrino 98*, Proceedings of XVIII International Conference on Neutrino Physics and Astrophysics, Takayama, Japan, 1998, edited by Y. Suzuki and Y. Totsuka [Nucl. Phys. B (Proc. Suppl.) (in press)].
[2] For a review, see, e.g., A.D. Dolgov, Phys. Rep. **222**, 309 (1992).
[3] For a review, see, e.g., M. White, D. Scott, and D. Silk, Annu. Rev. Astron. Astrophys. **32**, 319 (1994).
[4] I. Antoniadis, Phys. Lett. B **246**, 377 (1990); I. Antoniadis and K. Benakli, *ibid.* **326**, 69 (1994); see also V.A. Kostelecky and S. Samuel, *ibid.* **270**, 21 (1991).
[5] E. Caceres, V. S. Kaplunovsky, and I. Mandelberg, Nucl. Phys. **B493**, 73 (1997).
[6] I. Antoniadis, C. Muñoz, and M. Quiros, Nucl. Phys. **B397**, 515 (1993); N. Arkani-Hamed, S. Dimopoulos, and J. March-Russell, hep-th/9809124; R. Sundrum, Phys. Rev. D **59**, 085010 (1999).
[7] E. Witten, Nucl. Phys. **B471**, 135 (1996).
[8] J. Lykken, Phys. Rev. D **54**, 3693 (1996); I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos, and G. Dvali, Phys. Lett. B **436**, 257 (1998).
[9] N. Arkani-Hamed, S. Dimopoulos, and G. Dvali, Phys. Lett. B **429**, 263 (1998).
[10] G. Shiu and S.-H.H. Tye, Phys. Rev. D **58**, 106007 (1998).
[11] K. Benakli, hep-ph/9809582.
[12] C. Bachas (unpublished).
[13] I. Antoniadis, K. Benakli, and M. Quiros, Phys. Lett. B **331**, 313 (1994).
[14] C. Bachas, J. High Energy Phys. **11**, 023 (1998).
[15] N. Arkani-Hamed, S. Dimopoulos, and G. Dvali, Phys. Rev. D **59**, 086004 (1999).
[16] C.S. Wood *et al.*, Science **275**, 1759 (1997).
[17] K.R. Dienes, E. Dudas, T. Gherghetta, and A. Riotto, Nucl. Phys. **B543**, 387 (1999).
[18] P. Hořava and E. Witten, Nucl. Phys. **B460**, 506 (1996); **B475**, 94 (1996).
[19] Super-Kamiokande Collaboration, M. Shiozawa *et al.*, Phys. Rev. Lett. **81**, 3319 (1998).
[20] T. Banks and M. Dine, Nucl. Phys. **B479**, 173 (1996); J. Ellis, A. Faraggi, and D.V. Nanopoulos, Phys. Lett. B **419**, 123 (1998).
[21] G.G. Ross, *Grand Unified Theories* (Benjamin-Cummings, Menlo Park, CA, 1984).
[22] Particle Data Group, C. Caso *et al.*, Eur. Phys. J. C **3**, 1 (1998).
[23] F. Zwirner, Phys. Lett. **132B**, 103 (1983).
[24] S.P. Misra and U. Sarkar, Phys. Rev. D **28**, 249 (1983).
[25] L.E. Ibañez and G.G. Ross, Nucl. Phys. **B368**, 3 (1992).
[26] See, e.g., M. Leurer, Y. Nir, and N. Seiberg, Nucl. Phys. **B398**, 319 (1993); L.E. Ibañez and G.G. Ross, Phys. Lett. B **332**, 100 (1994); P. Binetruy and P. Ramond, *ibid.* **350**, 49 (1995).
[27] M.S. Turner, presented at Boulder TASI 92, p. 165, astro-ph/9304012.
[28] For a review, see, e.g., D. Lyth and A. Riotto, Phys. Rep. (to be published), hep-ph/9807278.
[29] P. Delbourgo-Salvador, P. Salati and J. Audouze, Phys. Lett. B **276**, 115 (1992).
[30] S. Abel, K. Benakli, and S. Davidson (in preparation).
[31] S. Abel and S. Sarkar, Phys. Lett. B **342**, 40 (1995).
[32] K. Benakli, J. Ellis, and D. Nanopoulos, Phys. Rev. D **59**, 047301 (1999).
[33] E.R. Kolb and M. Turner, *The Early Universe* (Addison-Wesley, Reading, MA, 1990), Sec. 5.5.
[34] For a review, see, e.g., A.D. Dolgov and Ya.B. Zeldovich, Rev. Mod. Phys. **53**, 1 (1981); P.J.E. Peebles, *Physical Cosmology* (Princeton University Press, Princeton, 1971).
[35] P. Salati, Phys. Lett. **163B**, 236 (1985).
[36] J.C. Mather *et al.*, Astrophys. J. **420**, 439 (1994).
[37] For a review, see, e.g., G. Steigman, Nucl. Phys. B (Proc. Suppl.) **48**, 499 (1996).
[38] D. Lindley, Astrophys. J. **294**, 1 (1985).
[39] J. Ellis, D.V. Nanopoulos, and S. Sarkar, Nucl. Phys. **B259**, 175 (1985).
[40] S. Dimopoulos, R. Esmailzadeh, L.J. Hall, and G.D. Starkman, Nucl. Phys. **B311**, 699 (1998).
[41] For a review, see, e.g., V.A. Rubakov, and M.E. Shaposhnikov, Usp. Fiz. Nauk **166**, 493 (1996) [Phys. Usp. **39**, 461 (1996)].
[42] For a review and references, see, e.g., W. Buchmuller and M. Plumacher, presented at Paris 1997, Phase transitions in cosmology, hep-ph/9711208, pp. 141–160.
[43] R. Brandenberger, I. Halperin, and A. Zhitnitsky, hep-ph/9808471.
[44] A.D. Sakharov, JETP Lett. **5**, 24 (1967).
[45] E. Kolb and S. Wolfram, Nucl. Phys. **B172**, 224 (1980); **B195**, 542(E) (1982).
[46] S. Dimopoulos and L. Hall, Phys. Lett. B **196**, 135 (1987).
[47] K.R. Dienes, E. Dudas, and T. Gherghetta, Phys. Lett. B **436**, 55 (1998); Nucl. Phys. **B537**, 47 (1999).
[48] See, for instance, I. Affleck and M. Dine, Nucl. Phys. **B249**, 361 (1985); J. Ellis, K. Enqvist, D.V. Nanopoulos, and K.A. Olive, Phys. Lett. B **191**, 343 (1987); M. Dine, L. Randall, and S. Thomas, Phys. Rev. Lett. **75**, 398 (1995); Nucl. Phys. **B458**, 291 (1996); M. Gaillard, H. Murayama, and K.A. Olive, Phys. Lett. B **355**, 71 (1995).