BARYOGENESIS WITH LOW SCALE GRAVITY AND EXTRA DIMENSIONS

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A new scenario of baryogenesis in the context of theories with large extra dimensions is proposed. The baryon number is almost conserved at zero temperature by means of a localization mechanism recently analyzed by Arkani-Hamed and Schmaltz: leptons and quarks are located at two slightly displaced positions in the extra space, and this naturally suppresses the interactions which "convert" the latter in the former. We show that this is expected to be no longer true when finite temperature effects are taken into account. The whole scenario is first presented in its generality, without referring to the bulk geometry or to the specific mechanism which may generate the baryon asymmetry. As an example, we then focus on a baryogenesis model reminiscent of GUT baryogenesis. The Sakharov out of equilibrium condition is fulfilled by assuming nonthermal production of the bosons that induce baryon number violation.

1 Introduction

Despite the great success of Quantum Field Theory, a consistent scenario where gravity is also included still lacks. The most promising framework that could help in this task is string theory, whose consistency requires additional dimensions beyond the standard $3+1$. This extra space is usually assumed to be compact, with a small compactification radius of order M_p^{-1}. However, it has been observed in ref. [1] that, having no test of gravity below the millimeter scale, we do not really need such a tiny compactification radius, provided the extra dimensions are accessible only to gravitational interactions. The Standard Model degrees of freedom must indeed be localized on a 3 dimensional wall whose thickness does not exceed the scale of lenghts, of order TeV^{-1}, we currently probe in accelerator experiments. This class of models have the main goal of solving (or at least of rephrasing in geometrical terms) the hierarchy problem, since the weakness of gravity turns out to be due to the largeness of the space available to gravitons.

After the original observation of Arkani-Hamed *et al.* [1], other models based on the same idea were proposed. In particular, in the work [2] the metric is assumed to be nonfactorizable, and in this case even with only one small extra dimension it is possible to solve the hierarchy problem.

There are however some difficulties common to these theories, due to the presence of a low energy cut-off. In particular, both proton stability and baryogenesis may be problematic in models with very low fundamental masses.

For what concerns proton stability in Grand Unified Theories, the standard way to achieve it is to increase the mass of the additional bosons up to about $10^{15} - 10^{16}$ GeV . In the framework of theories with extra-dimensions, an interesting mechanism has been suggested in ref. [3] (see also [4,5,6] for alternative suggestions). In this paper, a dynamical mechanism for the localization of fermions on a wall[7] is adopted: leptons and quarks are however localized at two slightly displaced positions in the extra space, and this naturally suppresses the interactions between the latter and the former.

However, the observed baryon asymmetry requires baryon number (B) violating interactions to have been effective in the first stages of the evolution of the Universe. In this paper we thus wonder how this last requirement can be satisfied in a theory which adopts the idea of[3], to ensure proton stability *now* and baryon production *in the past.* Our proposal is that thermal corrections, which are naturally relevant at early times, may modify the localization of quarks and leptons so to weaken the mechanism that suppresses the B violating interactions. [a]

[a] There exist other proposals for baryogenesis in these theories [6,8,9]: in the work [6], after considering several bounds on baryogenesis with large extra dimensions, a mechanism based on nonrenormalizable operators is proposed; in ref. [8] baryon number is violated by "evaporation" of brane bubbles that carry a net baryonic charge into the bulk, and the

We will indeed see that, although it is not possible to give a complete analysis of the behavior of the system at finite temperature, it is natural to expect that baryon number nonconserving processes have been enhanced in the very eary Universe. In this framework, we will then consider a simple specific example reminiscent of GUT baryogenesis. After taking into account the bounds that apply to such an example (in particular the ones related to the Sakharov out of equilibrium condition), we will estimate the baryon asymmetry obtained in this scheme.

2 Localization

A simple mechanism for localizing fermions on a wall has been recently revisited in ref. [3].

In this paper the idea is illustrated in the easiest case where only one extra dimension is added to the usual four. The main ingredient that is needed is a scalar field ϕ which couples to the fermionic field ψ through the full five dimensional Yukawa interaction $g\,\phi\,\bar{\psi}\,\psi$ and whose expectation value $\langle\phi\rangle$ varies along the extra dimension, but it is constant on our four-dimensional world. [b]

It is possible to show [7] that in this case the fermion field is localized where its total mass $m = m_0 + g\,\langle\phi\rangle$ (m_0 is the bare fermion mass in the five dimensional theory) vanishes, i.e. on a wall with three spatial dimensions characterized by a particular position x_5 in the transverse direction.

For definiteness, we consider the theory described by the lagrangian

$$\begin{aligned} \mathcal{L}_{\phi\psi} &= \bar{\psi}\left(i\,\partial\!\!\!/_5 + \frac{1}{\widetilde{M}_0^{1/2}}\,\phi\,(y) + m_0\right)\psi \\ \mathcal{L}_{\phi} &= \frac{1}{2}\partial_\mu\phi\,\partial^\mu\phi - \left(-\mu_0^2\,\phi^2 + \lambda_0\,\phi^4\right)\ , \end{aligned} \tag{1}$$

where $y \equiv x_5$ is the fifth coordinate, the fields and the parameters have the following mass dimensions

$$[\phi] = 3/2\ ,\ [\psi] = 2\ ,\ [m_0] = [\mu_0] = \left[\widetilde{M}_0\right] = 1\ ,\ [\lambda_0] = -1\ , \tag{2}$$

and where the suffix 0 denotes the value of the parameters at zero temperature.

As we said, the localization position of the fermions depends on the vacuum configuration of the field ϕ. We will assume for $\phi\,(y)$ the kink solution

$$\phi = \frac{\mu_0}{\sqrt{2}\,\lambda_0}\ \tanh\left(\mu_0\,y\right)\ , \tag{3}$$

and we will approximate it with a straight line interpolating between the two vacua (see figure 1)

$$\begin{cases} \phi\,(y) \simeq \frac{\mu_0^2}{\sqrt{2}\,\lambda_0}\,y\ , & |y| < \frac{1}{\mu_0} \\ \phi\,(y) \simeq \pm\frac{\mu_0^2}{\sqrt{2}\,\lambda_0}\ , & |y| > \frac{1}{\mu_0}\ . \end{cases} \tag{4}$$

Thus we see that the localization can occur only if

$$m_0 < \frac{\mu_0}{\sqrt{2\,\lambda_0\,\widetilde{M}_0}}\ , \tag{5}$$

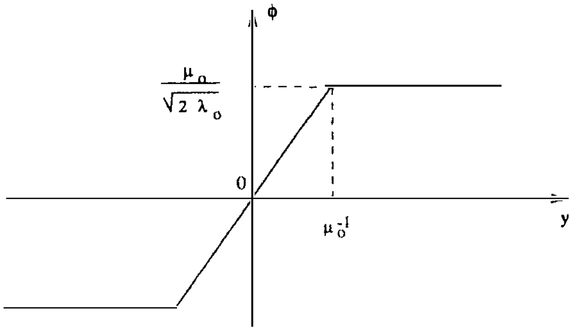
because otherwise the total fermion mass $m_{\rm tot} = \widetilde{M}_0^{-1/2}\,\phi\,(y) + m_0$ never vanishes.

It is possible to show [7] that from the four dimensional point of view, a left handed chiral massless fermionic field results from the localization mechanism, if the above configuration (3) is assumed for the scalar ϕ. The right handed part decouples from the system, being delocalized along the whole

matter–antimatter asymmetry can be due to a primordial collision of our brane with another one, that carried away the missing antimatter; in ref. [9] baryogenesis is obtained via leptogenesis, the latter being due to the existence of sterile neutrinos in the bulk.

[b] In this way the VEV $\langle\phi\rangle$ breaks the full translational invariance, as it is needed to have a preferred direction orthogonal to the wall.

Figure 1:



fifth dimension. This is a positive feature of this scenario, since the Standard Model fermion content has to be limited to chiral fields. The right handed fields can also be localized if a kink-antikink solution is assumed for the scalar ϕ. As a result, the left fields are still localized on the kink, while the right ones are confined to the antikink. If the kink and the antikink are sufficiently far apart, the left handed and right handed fermions however do not interact and again the model reproducing our four dimensional world is built by fermions of a defined chirality. The fermion content of the full dimensional theory is in this case doubled with respect to the usual one, and observers on one of the two walls will refer to the other as to a "mirror world". However most of the physics in one brane is not affected by the presence of the mirror one, and in the most of the present work we will focus on a single wall as if only the kink (3) configuration was present.

As it is shown in ref. [3], this mechanism of localization of fermions on a domain wall can be applied to guarantee proton stability. If indeed one chooses different five dimensional bare masses for different fermion fields, the latter are localized at different positions in the fifth direction. As a consequence, their wave functions do only partially overlap, and increasing the difference between the five dimensional bare masses of two fermions has the effect of suppressing their mutual interactions. In particular, to ensure baryon number conservation it is necessary to suppress interactions between leptons and baryons. We will thus give leptons and baryons different "masses", respectively

$$(m_0)_l = 0 \ , \ (m_0)_b = m_0 \ ,$$

which correspond to the localizations[c]

$$y_l = 0 \ , \ y_b = \frac{m_0 \sqrt{2\,\lambda_0\,\widetilde{M}_0}}{\mu_0^2} < \frac{1}{\mu_0} \ .$$

The shape of the fermion wave functions along the fifth dimension can be cast in an explicit and simple form if we consider the limit $y_b \ll 1/\mu_0$, in which the effect of the plateau for $y > 1/\mu_0$ can be neglected:

$$f_i\,(y) = \left(\frac{\mu_0^2}{\sqrt{2\,\lambda_0\,\widetilde{M}_0 \pi}} \right)^{1/4} \exp\left\{ -\frac{\mu_0^2\,(y-y_i)^2}{2\sqrt{2\,\lambda_0\,\widetilde{M}_0}} \right\} \quad (i = l, b) \tag{6}$$

We expect the Standard Model to be embedded in some theory which, in general, will contain some additional bosons X whose interactions violate baryon number conservation. If it is the case, the four fermion interaction $qq \longleftrightarrow ql$ can be effectively described by

$$\int d^4x\,dy\,\frac{q\,q\,q\,l}{\Lambda\,m_X^2} \ , \tag{7}$$

[c]The last inequality in the next expression comes from (5). We assume quarks of different generations to be located in the same y position in order to avoid dangerous FCNC mediated by the Kaluza-Klein modes of the gluons [10].

where m_X is the mass of the intermediate boson X and Λ is a parameter of mass dimension one related to the five-dimensional coupling of the X-particle to quarks and leptons.

After integration over y, this four fermion operator is thus suppressed in the four dimensional effective theory by a factor

$$\begin{aligned}
\delta &= \frac{1}{\Lambda m_X^2} \int dy \frac{\mu_0^2}{\pi \sqrt{2\lambda_0 \widetilde{M}_0}} \exp\left\{ -\frac{\mu_0^2/2}{\sqrt{2\lambda_0 \widetilde{M}_0}} \left[y^2 + 3\,(y-y_b)^2 \right] \right\} = \\
&= \frac{\mu_0}{\Lambda m_X^2 \sqrt{2\pi} \left(2\lambda_0 \widetilde{M}_0\right)^{1/4}} \exp\left\{ -\frac{3\left(2\lambda_0 \widetilde{M}_0\right)^{1/2}}{8} \frac{m_0^2}{\mu_0^2} \right\} . \qquad (8)
\end{aligned}$$

Current proton stability requires $\delta \lesssim (10^{16}\,\mathrm{GeV})^{-2}$, that is

$$\frac{m_0}{\mu_0} \gtrsim \frac{\sqrt{200 - 6\,\mathrm{Log}_{10}\left(\frac{\Lambda m_X^2}{\mu_0}/\mathrm{GeV}^2\right)}}{\left(2\lambda_0 \widetilde{M}_0\right)^{1/4}} . \qquad (9)$$

The numerator in the last equation is quite insensitive to the mass scales of the model, and – due to the logarithmic mild dependence – can be safely assumed to be of order 10. For definiteness, we will thus fix it at the value of 10 in the rest of our work.

Conditions (5) and (9) give altogether

$$\begin{cases} 2\lambda_0 \widetilde{M}_0 \lesssim 10^{-4} \\ \frac{m_0}{\mu_0} \gtrsim 10^2 \ . \end{cases} \qquad (10)$$

The last limit in eqs. (10) is stronger than the one given in ref. [3] where proton stability is obtained if the ratio of the massive scales of the model is of order 10. However, in ref. [3] the field ϕ simply scales linearly as a function of y, while we expect that whenever a specific model is assumed, conditions analogous to our (5) and (10) should be imposed. Anyway, even when the limits (10) are taken into account, we see that this mechanism allows to achieve proton stability without invoking any strong fine-tuning.

3 Thermal correction to the coefficients

Once the localization mechanism is incorporated in a low energy effective theory – as the system (1) may be considered –, one can legitimately ask if thermal effects could play any significant role. In the present work we are mainly interested in any possible change in the argument of the exponential in eq. (8), that will be the most relevant for the purpose of baryogenesis. For this reason, we introduce the dimensionless quantity

$$a(T) = \frac{m(T)^2}{\mu(T)^2} \sqrt{2\lambda(T) \widetilde{M}(T)} \ . \qquad (11)$$

From eqs. (9) and (10), we can set $a(0) \gtrsim 100$ at zero temperature. Thermal effects will modify this value. There are however some obstacles that one meets in evaluating the finite temperature result. Apart from the technical difficulties arising from the fact that the scalar background is not constant, the main problem is that nonperturbative effects may play a very relevant role at high temperature. As it is customary in theories with extra dimensions, the model (1) is nonrenormalizable and one expects that there is a cut-off (generally related to the fundamental scale of gravity) above which it stops holding. Our considerations will thus be valid only for low temperature effects, and may only be assumed as a rough indication for what can happen at higher temperature.

Being aware of these problems, by looking at the dominant finite-temperature one-loop effects, we estimate the first corrections to the relevant parameters to be

$$\begin{cases} \lambda\,(T) = \lambda_0 + c_\lambda \,\frac{T}{\widetilde{M}_0^2} \\ \widetilde{M}\,(T) = \widetilde{M}_0 + c_{\widetilde{M}}\, T \\ m\,(T) = m_0 + c_m \,\frac{T^2}{\widetilde{M}_0} \\ \mu^2\,(T) = \mu_0^2 + c_\mu \,\frac{T^3}{\widetilde{M}_0} \end{cases} , \tag{12}$$

where the c's are dimensionless coefficients whose values are related to the exact particle content of the theory.

In writing the above equations, the first of conditions (10) has also been taken into account. For example, both a scalar and a fermion loop contribute to the thermal correction to the parameter λ_0. While the contribution from the former is of order $\lambda_0^2\, T$, the one of the latter is of order $T/\widetilde{M}_0^2$ and thus dominates.

Substituting eqs. (12) into eq. (11), we get, in the limit of low temperature,

$$a\,(T) \simeq a\,(0) \cdot \left[1 + \frac{T}{\widetilde{M}_0}\left(\frac{c_\lambda}{2\,\lambda_0\,\widetilde{M}_0} + \frac{c_{\widetilde{M}}}{2} + \frac{2\,c_m\,T}{m_0} - \frac{c_\mu\,T^2}{\mu_0^2}\right)\right] . \tag{13}$$

From the smallness of the quantity $\lambda_0\,\widetilde{M}_0$ [see cond. (10)] we can safely assume (apart from high hierarchy between the c's coefficients that we do not expect to hold) that the dominant contribution in the above expression comes from the term proportional to c_λ.

We thus simply have

$$a\,(T) \simeq a\,(0)\left(1 + c_\lambda\,\frac{T}{2\,\lambda_0\,\widetilde{M}_0^2}\right) . \tag{14}$$

We notice that the parameter c_λ, being related to the thermal corrections to the ϕ^4 coefficient due to a fermion loop, is expected to be *negative*[11]: the first thermal effect is to decrease the value of the parameter $a(T)$, making hence the baryon number violating reactions more efficient at finite rather than at zero temperature.

4 Baryogenesis

We saw in the previous section that thermal effects may increase the rate of baryon number violating interactions of our system. This is very welcome, since a theory which never violates baryon number cannot lead to baryogenesis and thus cannot reproduce the observed Universe. Anyhow baryon number violation is only one of the ingredients for baryogenesis, and the aim of this section is to investigate how the above mechanism can be embedded in a more general context.

A simple scheme which may be adopted is baryogenesis through the decay of massive bosons X. [d] This scheme closely resembles GUT baryogenesis, but there are some important peculiarities due to the different scales of energy involved. In GUT baryogenesis the massive boson X, coupled to matter by the interaction $g\,X\,\psi\,\bar{\psi}$, has the decay rate

$$\Gamma \simeq \alpha\, m_X \quad , \quad \alpha = \frac{g^2}{4\,\pi} . \tag{15}$$

An important condition is that the X boson decays when the temperature of the Universe is below its mass (out of equilibrium decay), in order to avoid thermal regeneration. From the standard equation for the expansion of the Universe, $H \simeq g_*^{1/2}\, T^2\, M_p^{-1}$ (where g_* is the number of relativistic degrees of freedom at the temperature T), this condition rewrites

$$m_X \gtrsim g_*^{-1/2}\,\alpha\, M_p . \tag{16}$$

[d] We may think of these bosons as the intermediate particles which mediate the four fermion interaction described by the term (7).

If X is a Higgs particle, α can be as low as 10^{-6}. Even in this case however the X boson must be very massive. In principle this may be problematic in the theories with extra dimensions we are interested in, whose main feature is a very low fundamental scale.

One natural possibility to overcome this problem is to create the X particles non thermally and to require the temperature of the Universe to be always smaller than their mass m_X. In this way, one kinematically forbids regeneration of the X particles after their decay. In addition, although interactions among these bosons can bring them to thermal equilibrium, chemical equilibrium cannot be achieved.

Nonthermal creation of matter has raised a considerable interest in the last years. In particular, it has been shown that this production can be very efficient during the period of coherent oscillations of the inflaton field after inflation [13,14,15]. The efficiency of this mechanism has also been exploited in the work [16] to revive GUT baryogenesis in the context of standard four dimensional theories. Here, we will not go into the details of the processes that could have lead to the production of the X bosons. Rather, we will simply assume that, after inflation, their number density is n_X. To simplify our computations, we will also suppose that their energy density dominates over the thermal bath produced by the perturbative decay of the inflaton field.

Just for definiteness, let us consider a very simple model where two species of X bosons can decay into quarks and leptons, according to the four dimensional effective interactions

$$g\,X\,\bar{q}\,\bar{q}\quad,\quad g\,e^{-a/4}\,X\,l\,q\;, \tag{17}$$

where (remember the suppression given by the different localization of quarks and leptons) the quantity a is defined in eq. (11). Again for definiteness we will consider the minimal model where no extra fermionic degrees of freedom are added to the ones present in the Standard Model. Moreover we will assume $B-L$ to be conserved, even though the extension to a more general scheme can be easily performed.

The decay of the X bosons will reheat the Universe to a temperature that can be evaluated to be

$$T_{\rm rh}\simeq\left(\frac{30}{\pi^2}\,\frac{m_X\,n_X}{g_*}\right)^{1/4}\;. \tag{18}$$

Some bounds apply to $T_{\rm rh}$:

- since we do not want the X particles to be thermally regenerated after their decay, we require $T_{\rm rh}\lesssim m_X$, that can be rewritten as an upper bound on n_X

$$n_X\lesssim 30\left(\frac{g_*}{100}\right)m_X^3\;; \tag{19}$$

- it is necessary to forbid the B violating four fermion interaction (7) to erase the B asymmetry that has been just created by the decay of the X bosons. We thus require the interaction (7) to be out equilibrium at temperatures lower than $T_{\rm rh}$. From eq. (8) we see that we can parametrize the four fermion interaction with a coupling $g^2\,e^{-3\,a/8}/m_X^2$. Hence, the out of equilibrium condition reads

$$g^4\,e^{-3\,a/4}\lesssim g_*\,\frac{m_X}{M_p}\left(\frac{m_X}{T_{\rm rh}}\right)^3\;; \tag{20}$$

- for the same reason, we require the sphalerons to be out of equilibrium after the baryon asymmetry has been produced. This requirement is necessary only if one chooses the theory to be $B-L$ invariant, while it does not hold for $B-L$ violating schemes. We can approximately consider the sphalerons to be in thermal equilibrium at temperatures above the electroweak scale. Thus, if $B-L$ is a conserved quantity, we will require the reheat temperature to be smaller than about 100 GeV.

We can now estimate the baryon asymmetry obtained in this particular example. If the energy density of the Universe is dominated by the X bosons before they decay, one has

$$\eta_B \simeq 0.1 \ (N_X \, T_{\rm rh}/m_X) \ \langle r - \bar{r} \rangle \ , \tag{21}$$

where N_X is the number of degrees of freedom associated to the X particles and $\langle r - \bar{r} \rangle$ is the difference between the rates of the decays $X \to q\, l$ and $\bar{X} \to \bar{q}\, \bar{l}$.

We denote with X_1 and X_2 the two species of bosons whose interactions (17) lead to baryon number violation, and parametrize by ϵ the strength of CP-violation in these interactions. Considering that e^{-2a} is always much smaller than one, we get [18]

$$\langle r - \bar{r} \rangle \sim 3\, g^2\, e^{-a/2}\, \epsilon \ {\rm Im\, I}_{SS}\, (M_{X_1}/M_{X_2}) \ , \tag{22}$$

where the function ${\rm Im\,I}_{SS}(\rho) = [\rho^2 \ {\rm Log}(1 + 1/\rho^2) - 1]\,/\,(16\,\pi)$ can be estimated to be of order $10^{-3} - 10^{-2}$. It is also reasonable to assume $\epsilon \sim 10^{-2} - 1$.

Assuming N_X to be of order 10, we get the final estimate for the baryon asymmetry:

$$\eta_B \simeq \left(10^{-5} - 10^{-2}\right) g^2\, \frac{T_{\rm rh}}{m_X}\, e^{-a(T_{\rm rh})/2} \ . \tag{23}$$

From the requirement $T_{\rm rh} \lesssim m_X$ we get an upper limit on the baryon asymmetry

$$\eta_B \lesssim \left(10^{-5} - 10^{-2}\right) g^2\, e^{-a(T_{\rm rh})/2} \ . \tag{24}$$

We get a different limit on η_B from the bound (20): assuming $m_X \sim$ TeV and $g_* \sim 100$ indeed one obtains

$$\eta_B \lesssim \left(10^{-6} - 10^{-10}\right) g^{2/3}\, e^{-a(T_{\rm rh})/4} \ . \tag{25}$$

Since the observed amount of baryon asymmetry is of order 10^{-10}, even in the case of maximum efficiency of the process (that is, assuming maximal CP violation and $g \sim 1$), we have that both bounds (24) and (25) imply that $a\,(T_{\rm rh})$ has to be smaller than about 40.

Unfortunately, the temperature at which the condition $a\,(T) \lesssim 40$ occurs cannot be evaluated by means of the expansion of eq. (14), that have been obtained under the assumption $|a\,(T) - a\,(0)| \ll a\,(0)$. On the other hand, it is remarkable that our mechanism may work with a ratio $a\,(T_{\rm rh})\,/a\,(0)$ of order one. We thus expect that a successful baryogenesis may be realized for a range of the parameters of our theory which - although not evaluable through a perturbative analysis - should be quite wide and reasonable.

In scenarios with large extra dimensions and low scale gravity, the maximal temperature reached by the Universe after inflation is strongly bounded from above in order to avoid overproducing Kaluza-Klein graviton modes, which may eventually contradict cosmological observations [19]. For instance, in models with two large extra dimensions the reheating temperature has to be less than about 10 MeV. This value would be exceedingly low for our scenario since η_B is proportional to the ratio $T_{\rm rh}/m_X$, and hence the observed amount of baryons would be reproduced at the price of an unnaturally small value of $a\,(T_{\rm rh})$. However, other schemes with extra dimensions exist where the bounds on $T_{\rm rh}$ are less severe. For example, in the proposals [2,20] the mass of the first graviton KK mode is expected to be of order TeV. The reheating temperature can thus safely be taken to be of order $10 - 100\,$GeV.

5 Conclusions

We have proposed a scheme concerning the issue of baryogenesis in theories with large extra-dimensions. Since the observed proton stability requires to a very high degree of accuracy baryon conservation at zero temperature, this task may be problematic within the above theories, which have very low fundamental scales.

Our proposal relies on the localization mechanism for fermions discussed in ref. [3]. While in this work the present proton stability is due to a different localization (in the transverse direction) of leptons and quarks, we believe that thermal corrections may activate early baryon violating interactions.

We first provided a general discussion of the above scheme, without referring to any particular mechanism of baryogenesis. We found indeed that the first thermal corrections are in the direction of increasing the rate of baryon violations.

We then considered a very specific example, where the matter–antimatter asymmetry is achieved through the decay of a (relatively) heavy boson in a $B-L$ conserving context. In this situation the Sakharov out of equilibrium condition can be obtained in the simplest way by considering nonthermal production of the bosons responsible for $B+L$ violation.

Several bounds apply to the whole mechanism. The most general ones concern the localization procedure (we have found that the limits given in ref. [3] become more stringent once the thickness of the wall is considered). In addition, there are some other constraints which hold in the particular scheme of baryogenesis we adopted. The temperature of the heat bath right after the production of the baryon asymmetry cannot be too high, to avoid thermal regeneration of the bosons that induced baryogenesis. Moreover, this temperature has not to exceed the electroweak scale, in order not to activate the sphaleron transitions that would erase the $B+L$ asymmetry produced at higher energy. Of course, this last bound can be easily overcome by considering some $B-L$ nonconserving process.

We have found that the observed baryon asymmetry can be accomplished quite naturally in our example, and we believe that, because of its generality, our scheme of baryogenesis could be applied to a more general context as well.

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References

1. N. Arkani-Hamed, S. Dimopoulos and G. Dvali, Phys. Lett. **B429** (1998) 263 [hep-ph/9803315].
2. L. Randall and R. Sundrum, Phys. Rev. Lett. **83** (1999) 3370 [hep-ph/9905221].
3. N. Arkani-Hamed and M. Schmaltz, Phys. Rev. **D61** (2000) 033005 [hep-ph/9903417].
4. I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G. Dvali, Phys. Lett. **B436** (1998) 257 [hep-ph/9804398].
5. G. Shiu and S. H. Tye, Phys. Rev. **D58** (1998) 106007 [hep-th/9805157].
6. K. Benakli and S. Davidson, Phys. Rev. **D60** (1999) 025004 [hep-ph/9810280].
7. V. A. Rubakov and M. E. Shaposhnikov, Phys. Lett. **B125** (1983) 136.
8. G. Dvali and G. Gabadadze, Phys. Lett. **B460** (1999) 47 [hep-ph/9904221].
9. A. Pilaftsis, Phys. Rev. **D60** (1999) 105023 [hep-ph/9906265].
10. A. Delgado, A. Pomarol and M. Quiros, JHEP **0001** (2000) 030 [hep-ph/9911252].
11. K. R. Dienes, E. Dudas, T. Gherghetta and A. Riotto, Nucl. Phys. **B543** (1999) 387 [hep-ph/9809406].
12. P. Kanti, I. I. Kogan, K. A. Olive and M. Pospelov, hep-ph/9912266.
13. J. H. Traschen and R. H. Brandenberger, Phys. Rev. **D42** (1990) 2491.
14. L. Kofman, A. Linde and A. A. Starobinsky, Phys. Rev. Lett. **73** (1994) 3195 [hep-th/9405187].
15. Y. Shtanov, J. Traschen and R. Brandenberger, Phys. Rev. **D51** (1995) 5438 [hep-ph/9407247].
16. E. W. Kolb, A. Linde and A. Riotto, Phys. Rev. Lett. **77** (1996) 4290 [hep-ph/9606260].
17. D. J. Chung, E. W. Kolb and A. Riotto, Phys. Rev. **D60** (1999) 063504 [hep-ph/9809453].
18. D. V. Nanopoulos and S. Weinberg, Phys. Rev. **D20** (1979) 2484.
19. N. Arkani-Hamed, S. Dimopoulos and G. Dvali, Phys. Rev. **D59** (1999) 086004 [hep-ph/9807344].
20. N. Kaloper, J. March-Russell, G. D. Starkman and M. Trodden, hep-ph/0002001.