Baryogenesis in models with large extra dimensions

ANUPAM MAZUMDAR

The Abdus Salam International Centre for Theoretical Physics,
I-34100, Italy

and

ROUZBEH ALLAHVERDI 1, KARI ENQVIST 2, AND ABDAL PÉREZ-LORENZANA 3

1 Physik Department, TU Muenchen, James Frank Strasse,
D-85748, Garching, Germany.
2 Department of Physics and Helsinki Institute of Physics, P. O. Box 9, FIN-00014,
University of Helsinki, Finland.
3 The Abdus Salam International Centre for Theoretical Physics,
I-34100, Trieste, Italy.
4 Departamento de Física, Centro de Investigación y de Estudios Avanzados del I.P.N.
Apdo. Post. 14-740, 07000, México, D.F., México.

We describe how difficult it is to realise baryogenesis in models where the
fundamental scale in nature is as low as TeV. The problem becomes even
more challenging if we assume that there are only two extra compact spatial
dimensions, because thermal history of such a Universe is strongly constrained
by various cosmological and astrophysical bounds which translate the maximal
temperature of the Universe which must not exceed \( \sim O(10) \) MeV. This simply
reiterates that the observed baryon asymmetry must be synthesised just above
the nucleosynthesis scale. In this talk we address how to construct a simple
model which can overcome this challenge.

PRESENTED AT

COSMO-01
Rovaniemi, Finland,
August 29 – September 4, 2001
1 Introduction

Recently it has been proposed that large extra spatial dimensions can explain the apparent weakness of the electroweak scale compared to gravity in $3 + 1$ dimensions. In such a scenario a four dimensional world is assumed to be a flat hypersurface, called a ‘brane’, which is embedded in a higher dimensional space-time usually known as the bulk. The hierarchy problem is then resolved by assuming that TeV scale can be the fundamental scale in higher dimensions $[1, 2]$. This however requires the compactified radii of the extra dimensions are large. Such a large volume can substantiate the hierarchy in energy scales. The volume suppression $V_d$, the effective four dimensional Planck mass $M_p$, and, the fundamental scale in $4 + d$ dimensions $M_*$ are all related to each other by a simple relation

$$M_p^2 = M_*^{2+d}V_d.$$  

(1)

This automatically sets the present common size of all the extra dimensions at $b_0$. For two extra dimensions, and, $M_* = 1$ TeV, the required size is of order 0.2 mm right on the current experimental limit for the search of deviation in Newton’s gravity $[3]$. Recent astrophysical and cosmological bounds suggest $M_*$ should be larger than 500 TeV for two extra dimensions, and around 30 TeV for three extra dimensions $[4]$. Naturally, such a model has an important impact on collider experiments $[5]$, and on cosmology (see for instance Ref. $[6]$ and references therein). In this talk we address one of the most important question of cosmology and particle physics; how to realise the observed baryon asymmetry $\sim 10^{-10}$ in nature?

The generation of baryon asymmetry requires three well-known conditions; $C$ and $CP$ violation, $B$ and/or $L$ violation, and out of equilibrium decay or scattering processes $[7]$. It is quite likely that the departure from thermal equilibrium is possible at early times when the expansion rate of the Universe is large. Certainly acquiring such a condition becomes more difficult at late times, especially if thermal equilibrium is reached at a temperature below the electroweak scale. Note that the size of the compact extra dimensions must be stabilised and the corresponding mass scale for two extra dimensions is of order $\sim 10^{-2}$ eV. Above this scale the Kaluza-Klein (KK) modes of gaviton can be excited very easily, and as we reach higher in temperature we excite them more and more. We must note that the bulk is also very large which can dilute their number density, and you might suspect that we would never be able to feel their presence. However, this is not correct because these KK modes can decay into light fermions via Planck suppressed couplings. The readers must note that the KK modes once produced can never reach thermal equilibrium and their number density is frozen. Therefore, the KK modes when decay their decay products can have an energy scale proportional to their mass, which could be related to the largest temperature of the thermal bath produced in the early Universe.

One could envisage that the largest cross section of a heavy KK mode is possible at a centre of mass frame; $\gamma + \gamma \rightarrow G$, where $\gamma$ denotes the relativistic species and $G$ signifies
the KK mode. The cross section goes as \( \sigma_{\gamma+\gamma \rightarrow G} \sim (TR)^d/M_p^{2d} \), where \( R \) is the effective size of the extra dimensions. Once the KK modes are produced their evolution can be traced and a simple overclosure of their density; \( n_G/n_\gamma \leq 1 \) limits the initial thermal bath of the Universe which we denote here by \( T_r \), usually known as the largest temperature during the radiation dominated era.

\[
T_r \leq T_c \sim \left( \frac{M_{\ast}^{d+2}}{M_p} \right)^{1/(1+d)}.
\] (2)

where \( T_c \) is the normalcy temperature which states that the Universe better thermalises below this bound. For some preferred values \( M_{\ast} \sim \mathcal{O}(10) \text{TeV} \) for \( d = 2 \), we obtain \( T_c \leq \mathcal{O}(10 - 100) \text{MeV} \). This is an extremely strong constraint on thermal history of the Universe. It reiterates very strongly that the radiation dominated Universe simply cannot prevail beyond this temperature. Above the normalcy temperature our Universe is inevitably dominated by the KK degrees of freedom. Therefore, any physical phenomena such as first order phase transition, out-of-equilibrium decay of heavy particles, if at all taking place beyond \( T_c \) must take the KK degrees of freedom into account. Note that for larger number of compact extra spatial dimensions the bound on normalcy temperature is relaxed quite a lot. Therefore, it is the two extra dimensions scenario which poses the most challenging problems for theorists and also to the experimentalists who are trying their luck in seeking any departure from the Newtons gravity below mm. scale \([3]\).

Apart from the above mentioned problems there exists another major problem with a fast proton decay. The low fundamental scale induces proton decay via dimension 6 baryon number violating operator in the theories beyond the Standard Model (SM). With such a low fundamental scale the usual coupling suppression is not sufficient. We will come back to this issue when we discuss the baryon number violating lepto-quark interactions which is also required for baryon asymmetry \([4]\).

In nut-shell we require an inflationary model which automatically provides very low reheat temperature. Such a possibility is not remote, especially if we assume that the inflaton sector is a SM gauge singlet which resides in the bulk. After compactification the effective four dimensional theory provides a Planck suppressed couplings of the inflaton sector to the SM fields very naturally. Keeping inflaton field in the bulk also aids stabilising the extra dimensions, because at initial times the natural size of all the dimensions is inevitably the string scale; \((\text{TeV})^{-1}\). Therefore, not only our three dimensions must expand, but also the extra dimensions until they are trapped when their radii reach of order mm. for instance in the case of two extra dimensions. The stabilisation process is a dynamical one because the radion field which determines the size of the extra dimensions obtains a dynamical mass through its couplings to the trace of the energy momentum tensor of the bulk fields, see Ref. \([8]\). Therefore, the radion mass during inflation gets a correction during inflation of order \( \sim H \); the Hubble parameter, which makes the radion field heavy and thus its dynamics is frozen and the radion rolls down to its true minimum
very fast in effectively one Hubble time \[^{[8]}\].

The inflaton sector living in the bulk eventually decays after the end of inflation, and since the couplings are Planck mass suppressed, the decay rate of the inflaton into the Higgses, for instance, given as \[^{[4], [8], [9]}\]

\[
\Gamma_{\phi \rightarrow HHH} \sim \frac{g^2 M^3_*}{32\pi M_p^2},
\]

where \(g\) is the coupling constant, and \(M_*\) is the fundamental scale. While deriving the decay rate in Eq. (3), we have implicitly assumed that the mass of the inflaton is roughly the same order as the fundamental scale \(\sim M\), otherwise the decay rate of the inflaton would be even more suppressed for smaller mass scale of the inflaton. The estimated reheat temperature of the Universe is then given by

\[
T_r \sim 0.1 \sqrt{\Gamma M_p} \sim 1(10)\text{MeV},
\]

just right above the temperature required for a successful Big Bang nucleosynthesis. Note that this result is independent of the number of extra dimensions. This is an another important lesson; once we assume that the inflaton resides in the bulk a low reheat temperature is a prediction of the model and then it becomes even more interesting to address the issue of baryogenesis.

At such a low temperature leptogenesis is certainly not possible because the SM sphaleron transitions which preserves \(B - L\) are not in equilibrium below 100 GeV. Any prior lepton asymmetry can not be processed into a net baryon asymmetry. Besides this there is a catch in the leptogenesis scenario. A singlet right handed Majorana neutrino can naturally couple to the SM lepton doublet, and, the Higgs field in a following way: \(h \bar{L} H N\). This leads to a potentially large Dirac mass term unless the Yukawa coupling \(h \sim 10^{-12}\), or, so. Moreover, now the see-saw mechanism fails to work, since, the largest Majorana mass we may expect can never be larger than the fundamental scale. Therefore, given a neutrino mass \(\sim h^2 \langle H \rangle^2 / M_* \sim h^2 \cdot \mathcal{O}(1)\) GeV, we still have to fine tune \(h^2 \leq 10^{-10}\), in order to obtain the right order of magnitude for the neutrino mass. Therefore, the right handed neutrinos if they at all exist are more likely to be residing in the bulk rather than on the brane. Due to the volume suppression, the bulk-brane coupling naturally provides a small coupling \[^{[10]}\]. In any case the decay rate of the right handed neutrino to the SM fields is suppressed by the smallness of \(h\) that gives rise to a decay rate which is similar to Eq. (3). This makes extremely difficult to realise baryogenesis, because eventually when the right handed neutrino decays into the SM fields, the background temperature is of order of the reheat temperature \(\sim \mathcal{O}(1 - 10)\) MeV, and, at this temperature the sphaleron transition is not at all in equilibrium. The sphaleron transition rate is exponentially suppressed. So, a seemingly suitable lepton number might not even get converted to the baryons to produce the desired baryon asymmetry in the
Universe. Indeed, a larger reheating temperature, at least $O(1-100)$ GeV is required for making this scenario viable, which is certainly not bad if the number of extra dimensions are more, because this would relax the normalcy temperature. However, the assumption of inflaton being a bulk field might have to be judged carefully, because being a bulk field the inflaton interaction to the SM fields is always volume suppressed which will automatically ensure a low reheat temperature as we have stressed earlier.

On the other hand in some cases it is possible that the maximum temperature of the Universe is larger than the final reheat temperature of the Universe. This may happen if the decay products of the oscillating inflaton are mainly relativistic species. If the decay products thermalise when the inflaton is still oscillating and dominating the energy density of the Universe, then it changes the usual scaling relationship $T \propto a^{-1}$ between the temperature and the scale factor. The temperature reaches its maximum when $a/a_1 \sim 1.48$, where $a$ denotes the scale factor of the Universe and the subscript 1 denotes the era when inflation comes to an end. In our case the inflationary scale is determined by $H_I \sim M_*$. After reaching the maximum temperature, it decreases as $T \sim 1.3(g_*(T_m)/g_*(T))^{1/4}T_m a^{-3/8}$, where $T_m$ denotes the maximum temperature For $M_* \sim O(10)$ TeV, the maximum temperature could reach $T_m \sim O(10^5)$ GeV [11].

However, the above mentioned situation might not arise in our case because the inflaton coupling to the SM particles is extremely weakly coupled and the comparative decay rates of the inflaton to the Higgses and the SM fermions are equally favourable [9, 14, 6]. Therefore, if there were no initial dominance of relativistic species other than the inflaton, then it is very likely that the thermalisation is happening very late and the final reheat temperature is the maximum temperature one could achieve. However, in an extreme case even if we assume that there exists some initial thermal bath which could allow sphaleron transitions to occur, then one might imagine that sphalerons could reprocess a pre-existing charge asymmetry into baryon asymmetry [12], which might be reflected in an excess of $e_L$ over anti-$e_R$ created during inflaton oscillations. This mechanism requires again $(B + L)$-violating processes are out of equilibrium before $e_R$ comes into chemical equilibrium, such that the created baryon asymmetry could be preserved. Nevertheless, it is important to notice that still decaying inflaton field certainly injects more entropy to the thermal bath, provided the inflaton dominantly decays into the relativistic degrees of freedom. So, an initially large baryon asymmetry has to be created in order to obtain the right amount of asymmetry just before nucleosynthesis. One can easily estimate the amount of dilution that the last stages of reheating era will produce. The entropy dilution factor is given by [14, 6]:

$$\gamma^{-1} = \left( \frac{s(T_i)}{s(T_c)} \right) = \left( \frac{g_*(T_i)}{g_*(T_c)} \right) \left( \frac{T_i}{T_c} \right)^3 \left( \frac{a(T_i)}{a(T_c)} \right)^3,$$  

(5)

where $s$ is the entropy and $T_c$ denotes the electroweak temperature $\sim 100$ GeV. For a low reheating temperature as $T_i \sim 1$ MeV, the above expression gives rise to $\gamma^{-1} \geq 10^{25}$. While
calculating the ratio between the scale factors, we have used $T \propto a^{-3/8}$ and $g_*(T_r) \approx g_*(T_t)$. Therefore, including the entropy dilution factor we concludes that the initial $n_b/s$ has to be extremely large $\geq 10^{15}$ in order to produce the required baryon asymmetry at the time of nucleosynthesis, which is $n_b/s \sim 10^{-10}$. Such a large baryon asymmetry is an extraordinary requirement on any natural model of baryogenesis, which is almost impossible to achieve in any case.

There are couple of important lessons to be learned from the above analysis. First of all the large production of entropy during the last stages of reheating can in principle wash away any baryon asymmetry produced before electroweak scale. The second point is that it is extremely unlikely that leptogenesis will also work because one needs to inject enough lepton asymmetry in the Universe before the sphaleron transitions are in equilibrium. The only simple choice left is to produce directly baryon asymmetry, however, just before the end of reheating. The sole mechanism which seems to be doing well under these circumstances is the Affleck-Dine baryogenesis [13], which we shall discuss in the following section.

2 Affleck-Dine baryogenesis

A scalar condensate which carries non-zero baryonic, or/and leptonic charge survives during inflation and decays into SM fermions to provide a net baryon asymmetry. In our case the AD field; $\chi$, is a singlet carrying some global charge which is required to be broken dynamically in order to provide a small asymmetry in the current density. This asymmetry can be transformed into a baryonic asymmetry by a baryon violating interactions which we discuss later on. In order to break this $U(1)_\chi$ charge we require a source term which naturally violates $CP$ for a charged $\chi$ field, and during the non-trivial helical evolution of the $\chi$ field generates a net asymmetry in $\chi$ over $\bar{\chi}$ [14]. This new mechanism has been recently discussed in the context of supersymmetric inflationary model [15]. The $\chi$ field obtains a dynamical mass through its coupling to the inflaton sector which takes place after the end of inflation.

We remind the readers that the inflaton energy density must govern the evolution of the Universe and the decay products of the inflaton is also responsible for reheating the Universe. This happens once the inflaton decays before $\chi$ decays into SM quarks and leptons. This decay of $\chi$ via baryon violating interaction generates a baryon asymmetry in the Universe which is given by [14]

$$\frac{n_b}{s} \approx \frac{n_b}{n_{\chi}} \frac{T_r}{m_{\chi}} \frac{\rho_{\chi}}{\rho_I}.$$  \hspace{1cm} (6)

The final entropy released by the inflaton decay is given by $s \approx \rho_I/T_r$. The ratio $n_b/n_{\chi}$ depends on the total phase accumulated by the AD field during its helical motion in the
background of an oscillating inflaton field, which can at most be $\approx O(1)$. If we assume that the AD field is a brane-field, then the energy density stored in it can at most be: $ho_\chi \approx m_\chi^2 M_p^2$. On the other hand the energy density stored in the (bulk) inflaton field is quite large $\rho_I \approx M_*^2 M_p^2$, because the projected inflaton energy density on the brane has a Planck enhancement while its couplings to the SM particles has a Planck suppression. This simply means that the amplitude of the inflaton field is large initially, which otherwise would not have been possible if the inflaton were a brane field. Therefore, the final ratio

$$n_b/s \approx \left( \frac{T_r}{M_p} \right) \left( \frac{m_\chi}{M_p} \right) \approx 10^{-34} \left( \frac{m_\chi}{1\text{GeV}} \right) \ll 10^{-10},$$

(7)

for $T_r \sim O(1-10)$ MeV. The conclusion of the above analysis is again disappointing, as it suggests that the AD baryogenesis also leads to a small $n_b/s$. One way to boost this ratio is to assume that the AD field resides in the bulk. In that case one naturally enhances the ratio $\rho_\chi/\rho_I$, however, keeping in mind that it is still less than one in order not to spoil the successes of inflation. The projected AD energy density on the brane has now a Planck enhancement; $\rho_\chi \approx m_\chi^2 M_p^2$. This reiterates that the initial vacuum expectation value (vev) of the AD field is quite large, which were not possible in the earlier discussion where the AD field was a brane field. Therefore, the **maximum** baryon to entropy ratio

$$n_b/s \approx \left( \frac{T_r}{M_*} \right) \left( \frac{m_\chi}{M_*} \right) \sim 10^{-10} \left( \frac{m_\chi}{1\text{GeV}} \right),$$

(8)

where we have evaluated the right hand side for $T_r \sim 10$ MeV and $M \sim 10$ TeV. Although, the mass of the AD field requires some fine tuning, up to the $CP$ phase, the above ratio can reach the observed baryon to entropy ratio quite comfortably. Notice, that the actual predicted value also depends on the initial conditions on $\chi$ that may render $m_\chi$ more freedom. Say for instance, if the initial vev of $\chi_0 \sim M_{GUT}$, we get the right $n_b/s$ provided $m_\chi \sim M$.

We have noticed earlier that due to the violation of $U(1)_\chi$ charge, the dynamics of the AD field generates an excess of $\chi$ over $\bar{\chi}$ fields. This asymmetry is transferred into baryon asymmetry by a baryon violating interaction, such as $\kappa \chi Q\bar{Q}L/M^2 M_p$, however, keeping $B-L$ conserved. We also assume that $\chi$ interactions to SM fields conserve $U(1)_\chi$ symmetry, thus, the quarks and leptons must carry a non zero global $\chi$ charge while the Higgs field does not. This avoids $\chi$ decaying into Higgses, which otherwise will reduce the baryonic abundance and make the above interaction the main channel for its decay. While discussing the decay rate of $\chi$ field one has to take into account all possible decay channels which can be of the order of thousands due to family and color freedom. On the other hand, we assume that the inflaton is decaying mainly into Higgses. Final result is then given by

$$\Gamma_\chi \approx \left( \frac{\kappa}{g} \right)^2 \left( \frac{m_\chi}{M_*} \right)^7 \Gamma_\phi,$$

(9)
By taking $\kappa/g \sim \mathcal{O}(1)$ we can insure that $\chi$ will decay along with the inflaton, provided that its mass is very close to the fundamental scale. This will certainly demand some level of fine tuning in the parameters. We would like to mention that this is perhaps the simplest scenario one can think of for generating baryon asymmetry right before nucleosynthesis takes place. It is worth mentioning that in our model the AD field will not mediate proton decay by dimension six operators as $QQQL$, as long as $\chi$ does not develop any vacuum expectation value. Notice, other processes mediating proton decay, such as instanton effects might still occur. While there is no known solution for such a potential problem yet, our mechanism is at least not adding any new source to proton decay. In the same spirit one may check those operators which induce $n - \bar{n}$ oscillations. Again, effective $\Delta B = 2$ operators of dimensions 9; $UDDUD$, and 11; $(QQH)^2$, can not be induced by integrating out $\chi$.

2.1 The model

In this section we briefly discuss whether we can provide flesh to the above discussed scenario. Without going into much details we describe the AD potential which must have to come in conjunction with the inflaton sector. The AD potential can be given by [6]

$$V_{AD}(\phi, N, \chi_1, \chi_2) = \kappa_1^2 \left(\frac{M_s}{M_p}\right)^2 N^2(\chi_1^2 + \chi_2^2) + \frac{\kappa_2^2}{4} \left(\frac{M_s}{M_p}\right)^2 (\chi_1^2 + \chi_2^2)^2$$

$$+ \kappa_3^2 \left(\frac{M_s}{M_p}\right)^2 \phi N (\chi_1^2 - \chi_2^2),$$

(10)

where $\kappa_1, \kappa_2, \kappa_3$ are constants of order one, and $\chi_1$ and $\chi_2$ are the real and imaginary components of the complex field AD field $\chi$. For the stability of the potential $\kappa_1 \sim \kappa_2 \geq \kappa_3$. Here we have two other dynamical fields which we have denoted by $\phi$ and $N$ which are coming from the inflaton sector. Note that the last term in the above equation which is responsible for breaking global $U(1)$. This could however be possible if both $\phi, N$ have non zero average vev. It might be possible to provide a scenario where initially the last term is vanishing where there is no apparent violation of a global charge. This can be realised in a particular inflationary model where the field is trapped in its own local minimum $N = 0$ during inflation, while responsible for generating a large vev which could drive inflation. The role of other field $\phi$ is to end inflation via instability through its coupling to the $N$ field which we have not written down. This is a perfect example of hybrid inflationary model where after the end of inflation both $\phi, N$ begin oscillations about their true minimum. Therefore, in our picture the first and last term is vanishing during inflation while the AD field is massless and its evolution is determined by the quartic term in the above equation. However, after the end of inflation when both $\phi, N$ begin oscillations.

*In this paper we do not discuss the details of the inflationary model which can be found in Ref. [6].
oscillates all the terms including the last term in Eq. (10) turns on. This is an interesting example of breaking $U(1)_\chi$ symmetry dynamically towards the end of inflation.

Note that in Eq. (10), the effective couplings are Planck suppressed, this is due to the fact that the AD field is a bulk field. If we know the post-inflationary behaviour of $\phi$ and $N$ fields, we shall be able to estimate the total asymmetry in $\chi$ distribution due to the source term which is responsible for $CP$ violation, given by the last term in the above equation. For a charged scalar field this is equivalent to $C$ violation also. The $B$ violation arises via the decay of $\chi$, because the decay products have $\Delta B \neq 0$ as discussed in the earlier section, and we have a non-trivial helical oscillations in $\chi$ which accumulates net $CP$ phase which is transformed into asymmetric $\chi$. The net $\chi$ asymmetry; $n_\chi$ can be calculated by evaluating the Boltzmann equation:

$$n_\chi = \frac{i}{2} (\dot{\chi}^* \chi - \chi^* \dot{\chi}) .$$

(11)

With the help of equations of motion for $\chi_1, \chi_2$ for the potential Eq. (10) we can rewrite the above expression as

$$\dot{n}_\chi + 3Hn_\chi = 4\kappa_3^2 \left( \frac{M_s}{M_p} \right)^2 \langle N(t)\phi(t) \rangle \chi_1(t)\chi_2(t) .$$

(12)

The right hand side of the above equation is a source term which generates a net $\chi$ asymmetry through a non-trivial motion of $\chi_1$ and $\chi_2$ fields. We integrate Eq. (12) from $t_0$ which corresponds to the end of inflation up to a finite time interval $[6]$

$$n_\chi a^3 = 4\kappa_3^2 \left( \frac{M_s}{M_p} \right)^2 \int_{t_0}^{t} \langle N(t')\phi(t') \rangle a^3(t')\chi_1(t')\chi_2(t')dt' .$$

(13)

The upper limit of integration signifies the end of reheating. We assume that the oscillations continue until the fields decay completely. Before we perform the integration, we notice that the integrand decreases in time. This can be seen as follows ; first of all notice that the approximate number of oscillations are quite large before the fields decay $\sim (M_p/M_s)^2$. This allows us to average $\phi$ and $N$ oscillations. The typical time dependence of oscillating fields follow $\langle \phi N \rangle \sim \Phi^2(t) \sim 1/t^2 [1]$. Similarly, the $\chi$ oscillations provide $\langle \chi_1\chi_2 \rangle \sim 1/t^2$ [2] another suppression in time. While taking care of the expansion, where the scale factor behaves like $a(t) \sim t^{2/3}$ during the inflaton oscillations, the overall behavior of the integrand follows $\sim 1/t^2$. This suggests that the maximum contribution

$^1$ In the hybrid inflationary model $N(t) \sim A_0(1 + \Phi(t)/3\cos(m_{\phi}t))$ develops a vev after the end of inflation, however $\phi(t) \sim A_0\Phi(t)/(3\sqrt{2})\cos(m_{\phi}t)$ does not, $A_0$ is the initial amplitude which is of order $\sim M_p$. Such a large amplitude is precisely due to the fact that the inflaton sector is in the bulk, and upon compactification the field can obtain a large vev with energy density $\rho_1 \sim M_s^3 M_p^2$.

$^2$ We must note that $\chi$ field must not develop any vev in order to avoid proton decay via lepto-quark interaction. Therefore $\chi_1(t) \sim \chi_2(t) \sim (|\chi(0)|/t) \cos(m_\chi t)$.
to $\chi$ asymmetry comes only at the initial times when $t_0 \sim 1/H_0$, where $H_0 \sim M_*$ in our case. The right hand side of the above equation turns out to be

$$n_\chi \approx \frac{2}{27} \kappa_3^2 \frac{M_*^2 |\chi(0)|^2}{H_0}.$$  \hfill (14)

We have assumed that the total $CP$ phase, which is given by two factors: an initial phase determined arbitrarily during inflation and the final dynamical phase which is accumulated during the oscillations, is of order $\sim \mathcal{O}(1)$. We are also assuming that the coupling $\kappa_3 \sim \mathcal{O}(1)$.

The final ratio of $\chi$ number density produced and the entropy is given by

$$\frac{n_\chi}{s} \approx \mathcal{O}(1) \left( \frac{|\chi(0)|}{M_p} \right)^2 \left( \frac{T_r}{M_*} \right)$$  \hfill (15)

where we have used the fact that $s \propto a^3$ in our case. This is the final expression for $\chi$ asymmetry produced during the helical oscillations of $\chi$ and this has to be compared with the observed baryon asymmetry $\sim n_B/s \sim 10^{-10}$. Note that the final expression depends on the amplitude of $\chi_1 \sim \chi_2 \sim \chi(0)$ at the time when $U(1)_\chi$ symmetry is broken, or, in our case when $\phi$ and $N$ start oscillating around their minimum. For a particular example; $T_r \sim 10$ MeV, and $M_* \sim 100$ TeV, we require $\chi(0) \sim 10^{16}$ GeV in order to produce a right magnitude of baryon asymmetry. As we have discussed earlier the baryon asymmetry is injected into thermal bath along with the inflaton decay products. It is essential that the thermalisation takes place after AD field has decayed in order not to wash away the total baryon asymmetry.

## 3 Conclusions

In summary, we have presented a natural mechanism of baryogenesis in the context of low quantum gravity scale with large compact extra dimensions. Our mechanism is generic and it is independent of the fundamental scale and the number of compact extra dimensions. This mechanism does not rely on any extra assumption other than invoking a fundamental scalar field that lives in the $4+d$ dimensional space time. The baryogenesis scheme can work at any temperature lower than the electroweak scale because the mechanism does not depend on sphaleron transition and relies on baryon violating interactions, and therefore does not depend on leptogenesis. The presence of non-renormalizable couplings inevitably reheats the Universe with a temperature close to nucleosynthesis, which is really independent of the number of extra dimensions. While performing our calculation we have implicitly assumed that we had a successful model of inflation which could also provide the right amount of density perturbations. This seems to be still alluding. For a recent work we refer the readers Ref. [16]. The main conclusion of Ref. [14] is that
it is difficult to realise a hybrid model which could generate an amplitude for the density perturbations observed by COBE at a scale larger than the size of the horizon with a fundamental scale as low as $\sim 100 - 500\text{TeV}$. It is necessary to push the fundamental scale up by $2 - 3$ orders of magnitude. This certainly relaxes the stringent bound on normalcy temperature for two extra dimensions. Therefore, opening up the chances of other mechanisms of baryogenesis to succeed. In any case the scenario which we have provided here is quite generic and works independent of any preferred fundamental scale and number of extra spatial dimensions.

References

[1] N.Arkani-Hamed, S. Dimopoulos, and G. Dvali, Phys. Lett. B 429, 263 (1998); Phys. Rev. D 59, 086004 (1999); I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos, and G. Dvali, Phys. Lett. B 436, 257 (1998).

[2] For some early ideas see also: I. Antoniadis, Phys. Lett. B246, 377 (1990); I. Antoniadis, K. Benakli and M. Quirós, Phys. Lett. B331, 313 (1994); K. Benakli, Phys. Rev. D60, 104002 (1999); Phys. Lett. B 447, 51 (1999).

[3] C. D. Hoyle et al, Phys. Rev. Lett. 86 1418 (2001).

[4] S. Hannestad, G. G. Raffelt, hep-ph/0110067.

[5] For experimental bounds see for instance: T. G. Rizzo, Phys. Rev. D 59, 115010 (1999); G. F. Giudice, R. Rattazzi and J. D. Wells, Nucl. Phys. B 544, 3 (1999); E. A. Mirabelli, M. Perelstein and Michael E. Peskin, Phys. Rev. Lett. 82, 2236 (1999); J. L. Hewett, Phys. Rev. Lett. 82, 4765 (1999); V. Barger, T. Han, C. Kao and R. J. Zhang, Phys. Lett. B 461, 34 (1999).

[6] R. Allahverdi, K. Enqvist, A. Mazumdar, and A. Perez-Lorenzana, Nucl. Phys. B 618 277 (2001).

[7] A. D. Sakharov, Zh. Eksp. Teor. Fiz. Pis’ma 5, 32 (1967); JETP Lett. B 91, 24 (1967).

[8] A. Mazumdar, and A. Perez-Lorenzana, Phys. Lett. B 508 340 (2001).

[9] R.N. Mohapatra, A. Pérez-Lorenzana and C.A. de S. Pires, Phys. Rev. D 62, 105030 (1999).

[10] For a review see for instance: A. Pérez-Lorenzana, hep-ph/0008333.

[11] R. Kolb, and M. Turner, The Early Universe, Addison-Wesley, (1994).
[12] B.A. Campbell, et al., Phys. Lett. B 297, 118 (1992).

[13] I. Affleck and M. Dine, Nucl. Phys. B 249, 361 (1985).

[14] A. Mazumdar, and A. Perez-Lorenzana; hep-ph/0103215.

[15] Z. Berezhiani, A. Mazumdar, and A. Perez-Lorenzana, Phys. Lett. B 518 282 (2001).

[16] A. M. Green, and A. Mazumdar, hep-ph/0201209.