Cosmology II: From the Planck Time to BBN

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Abstract

Progress in early Universe cosmology, including strings, extra dimensions, inflation, phase transitions, and baryogenesis, is reviewed.

1 Introduction

In 1984, as astro-particle physics was beginning to thrive, the first Inner Space/Outer Space workshop was organized [1]. The contents of that first volume reflect well on the state of the field at that time. Concerning pre-big bang nucleosynthesis (BBN) cosmology, the major focuses were centered on inflation, magnetic monopoles, Kaluza-Klein cosmology, supersymmetry, supergravity, and quantum gravity. Inflation was still new enough that there was a strong interest in finding working models of inflation using realistic particle physics. Particular emphasis was placed on inflationary models in the context of supergravity. After the demise of old and new inflation in the context of SU(5), the inflaton was created to allow a framework for constructing toy models [2]. After fifteen years, the search for a realistic inflationary model continues.

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Magnetic monopoles were a big topic at ISOS I. The report of a possible discovery of one of these cosmological GUT relics sparked an enormous amount of activity, which in contrast to inflation, has largely died away.

Beginning with SU(5) and the possibility of generating the baryon asymmetry of the Universe, unification has been an integral part of early Universe cosmology. At ISOS I, while there were only a few contributions on supersymmetry/supergravity, there was an active session on Kaluza-Klein unification. All of these avenues have been incorporated (more or less) in string theory and in its parent (mother) theory, M-theory.

In the concluding remarks of ISOS I, we made some predictions regarding the future of the particle physics/cosmology interface. I quote from that paper, “That brings us to the future. First, on the theoretical side, it seems likely that inflation, or at least some offshoots of the inflationary paradigm, will continue to be promising avenues to pursue.” (Given the volumes of publications on inflation in the ensuing years, this is at best an understatement.) “One of the most promising approaches to unification of all forces seems to be through additional spatial dimensions. An area still in its infancy, cosmology with extra dimensions adds yet another puzzling, but perhaps not unrelated, fact to our list: why are all but three of the spatial dimensions so small? Superstring cosmology opens a Pandora’s box of new problems - - whence came geometry, was there an initial singularity, does the Universe after all have a limiting temperature?” It is remarkable that in the last fifteen years, these are precisely the same questions we continue to ask. To be sure, much progress has been made on the technical side, but cosmology at the Planck time remains a holy grail.

While it is not possible in the context of this contribution in memory of Dave Schramm to completely review the state of the cosmology from the Planck time to BBN, I will attempt to touch on some key issues that are of particular interest today.

2 The Planck Time

The Planck time in cosmology is certainly the most challenging epoch to study, as the framework in which it must be described is still lacking. Though there are several different approaches to this problem, I will limit myself to that of string/M-theory.

Much of the work on string cosmology has focussed on the problem of inflation, and I will return to that subject in the next section. The realization that all of the different string theories, together with 11-dimensional supergravity, can be related through dualities has had a major impact on how we view the Planck epoch in cosmology. It has even led to the ultimate question as to what we mean by the Planck epoch. In a standard 4-dimensional
model, the gravitational action can be written as

\[ S = \frac{1}{2\kappa^2} \int d^4x \sqrt{g} (R + \cdots) \]  

(1)

where \( \kappa^2 = 8\pi G_N = 8\pi/M_P^2 \). In the context of a four dimensional theory, there is nothing ambiguous here as the Planck mass, determined from Newton’s constant is simply \( M_P = 1.2 \times 10^{19} \text{ GeV} \). Even in the context of an “old” 10-dimensional string theory, the 10-dimensional gravitational constant, \( \kappa_{10}^2 = 8\pi/M_{10}^8 \) is equivalent to its 4-dimensional counterpart if the size of the 6-dimensional compact space is Planck scale in extent, i.e. \( M_{10} = M_P \). Of course, the gravitational action in string theory must be augmented by the presence of additional fields in the gravitational sector, most notably by the universal coupling of the string dilaton [6]. Restricting our attention to only the dilaton gravity action, we can write

\[ S = \frac{1}{2\kappa_{10}^2} \int d^{10}x \sqrt{g_s}\{e^{-2\phi}(R_s + 4(\nabla\phi)^2 + \cdots)\} \]  

(2)

Upon reduction to 4 dimensions, we recover a relatively simple dilaton gravity system, which for a fixed dilaton (fixed by some potential which is perhaps induced by supersymmetry breaking), is just Einstein gravity in 4 dimensions along with some higher derivative curvature corrections (not shown) and additional moduli fields (also not shown).

As noted above, M-theory is emerging as the single underlying theory capable of unifying all particle interactions [7]. Although our understanding of M-theory is still incomplete, its various low energy limits, and the links between them, are known. These are the consistent string theories and 11D supergravity, related by dualities. A consequence of these developments is that the dilaton can be viewed as another modulus field in 11D supergravity. This could have important consequences for cosmological applications. The troubles with implementing conventional inflationary scenarios in string theory arise because of the dilaton and its couplings to the other modes in the string spectrum (see below). In 11D supergravity, such couplings are absent, and thus some of the obstacles for inflation in dilaton-plagued string theories could perhaps be resolved by way of M-theory.

One of the key points in the application of M theory to phenomenology is the reconciliation of the bottom-up calculation of \( M_{GUT} \sim 10^{16} \text{ GeV} \) with the string unification scale, which is close to the four-dimensional Planck mass scale \( M_4 \sim 10^{19} \text{ GeV} \). This is achieved by postulating a large fifth dimension \( R_5 \gg M_{GUT}^{1/4} \), which is not felt by the gauge interactions, but causes the gravitational interactions to rise with energy much faster than in the conventional four dimensions [8]. If we assume that the compact 6-space (Calabi-Yau?) is of the size determined by the fundamental Planck scale, \( M_{11} \) so that \( V_6 \sim M_{11}^{-6} \), then the fundamental Planck scale, \( M_{11} \) is related to the four dimensional Planck scale by \( M_{11}^9 V_6 R_5 \sim M_4^2 \). Or,

\[ M_{11}^3 R_5 \sim M_4^2 \]  

(3)
For $R_5 \gg M_{11}^{11}$, it is possible to achieve $M_4 \gg M_{11} \sim M_{GUT}$. In this type of scenario, one could expect that inflation should be considered within a five-dimensional framework.

Within this general five-dimensional framework, two favored ranges for the magnitude of $R_5$ can be distinguished. One is relatively close to $M_{GUT}^{11}$: $R_5^{-1} \sim 10^{12}$ to $10^{15}$ GeV, and the other could be as low as $R_5^{-1} \sim 1$ TeV [8]. In this latter case, the large dimension is not necessarily the conventional fifth dimension of $M$ theory. Indeed, in models studied in [8] the large dimension may be related to what is normally considered as one of the six “small” dimensions that is conventionally compactified a là Calabi-Yau. Of course the physics of a very large ($\lesssim 1$ mm) extra spatial dimension has received considerable attention lately [9, 10].

As a starting point therefore, one should consider the 11D supergravity action

$$S = \frac{1}{2\kappa_{11}^2} \int d^{11}x \sqrt{g} \left\{ R + \cdots \right\}$$

(4)

where $R$ is the scalar curvature of the 11D metric (terms involving the 3-form potential and its 4-form field strength are not shown). Reducing to 10D is done easily by assuming that the 11th direction is compact. In this case, we can carry out Kaluza-Klein reduction of (4) to find

$$S = \frac{1}{2\kappa_{11}^2} \int d^{10}x \sqrt{g_{10}} R_{11} \left\{ R_{10} + \cdots \right\}$$

(5)

After a conformal rescaling $g_{10} = R_{11}^{-1} g_s$, and defining the dilaton by $\exp(2\phi/3) = R_{11}$, we find exactly the action given in eq. (4). This is precisely the effective action which describes the low energy limit of the IIA superstring (though to be sure of the identification, one would be required to keep track of the terms shown here as $\cdots$). It is easy to rewrite this action in the ten-dimensional Einstein frame, by a further conformal rescaling $g_s = e^{\phi/2} g_E$. The action (4) can be reduced further to make contact with type IIB and heterotic theories.

If the 6-space is of the fundamental scale, then as described above, the Universe will have passed through a phase where it can be effectively described by a 5-dimensional space time. A simple ansatz for the 5D metric is that of the FRW form

$$ds^2 = -dt^2 + a^2(t)d\bar{x}^2 + c^2(t)d\varphi^2$$

(6)

In this way, the dilaton expectation value $\langle \phi \rangle$ is related to the scale factor $c(t)$. In the specific case of Horava-Witten type compactifications on $S^1/Z_2$, the Universe is described by two 4D branes at the end points of the line segment. This type of compactification is particularly well suited for a reduction of M-theory to heterotic string theory, where the matter and hidden sectors sit on opposing branes. The cosmology of these theories has been studied at length in [10].
3 Inflation

It has been known for quite some time that it is very difficult to incorporate conventional inflationary scenarios based on (de Sitter expansion) into the low-energy limit of string theory [11]. The principal obstacle has been the fact that the low-energy dynamics of the theory contains massless scalar fields with non-minimal couplings to gravity whose coupling constants are precisely given by the conformal symmetry and/or the dualities of string theory. In an expanding universe, these fields roll during the course of the expansion as dictated by their equations of motion, consuming the available energy and hence decreasing the rate of expansion. That is, besides the Minkowski space solution, there are no other solutions to Einstein’s equations (when the Gauss-Bonnet curvature squared terms are kept) of the dilaton-gravity system with constant curvature and a stationary dilaton. de Sitter solutions are possible if the dilaton sits in a potential minimum (with $V \neq 0$) due to supersymmetry breaking effects, but in this case, the theory suffers a graceful exit problem.

Instead, one typically finds solutions where the scale factor of the universe grows as a power of time, with the power determined by the scalar coupling constants. Once the numerical values of these constants, fixed by string theory, are taken into account, it has been found that the resulting power laws are too slow to give an inflationary universe [12]. For example, if we ignore the compact 6-space of eq.(2) and include a cosmological term $\Lambda$ we have a 4D action

$$ S = \frac{1}{2\kappa_4^2} \int d^4x \sqrt{|g_s|} \left\{ e^{-2\phi} \left( R_s + 4(\nabla \phi)^2 + \cdots + \Lambda \right) + \cdots \right\} \quad (7) $$

where now $\kappa_4^{-2} = M_4^2/8\pi$. In the Einstein frame this becomes

$$ S = \frac{1}{2\kappa_4^2} \int d^4x \sqrt{|g_s|} \left\{ (R - 2(\nabla \phi)^2 + \cdots + e^{2\phi} \Lambda) + \cdots \right\} \quad (8) $$

In the Einstein frame, the solution to the equations of motion yield a scale factor which grows linearly in time [13]. However, this expansion turns out not to be physical – the corresponding scale factor in the string frame is constant. To resolve this quandary, one must compare the scale factor in each frame with some physical length scale, such as the Compton wavelength of a massive particle. Had we included a massive scalar field in the action (7), we would find that the scale factor relative to the Compton wavelength $\lambda \sim m^{-1}$ is also constant. However, in the Einstein frame, the mass term would also be dilaton dependent with mass $e^{\phi} m$, implying that the Compton wavelength also grows linearly in time ($\phi \sim -\ln t$) and hence relative to a physical measure of length, the Universe according to (8) is also not expanding [12].

An interesting alternative to the standard inflationary picture in string theory is the pre-big bang scenario [14]. The solutions to the equations of motions in the string frame yield
two distinct branches often labeled (+) and (−) corresponding to solutions which either evolve towards singularities and are singularity-free in the past, or evolve from singularities and are singularity-free in the future. One of the goals of the pre-big bang scenario is the connection of the expanding solutions in the two branches, in such a way that the (+)-branch chronologically precedes the (−)-branch and hence the cosmological singularity would be removed. It still remains to be seen if a coherent and fully consistent description of branch-changing can be found. In fact there are strong arguments showing the precise difficulty of a branch change [15] though some progress has been made to resolve this graceful exit problem [16]. It is interesting to note that in the string frame, both the expanding and contracting metrics are degenerate to a single Einstein frame metric, and that the only difference between the two subclasses of solutions is the sign of the dilaton field. In the context of a full 11D theory, there are several solutions which can be found [17] all degenerate with the known pre-big bang solutions.

There is of course no possibility in this contribution to review all of the work on inflation in the context of string theories and extra dimensions. However, before moving to the subject of baryogenesis, I will mention two possibilities in which extra dimensions aid the implementation of inflation. For other recent work see [18].

In the first [19], one makes use of the higher derivative curvature terms in the action. Among the first utilizations of higher-derivative curvature terms is the Starobinsky model [20], which is based on obtaining a self-consistent solution of Einstein’s equations when they are modified to include one-loop quantum corrections to the stress-energy tensor $T_{\mu\nu}$. In its simplest form, the model is equivalent to a theory of gravity with an $R^2$ correction which can be written as [21]

$$S = \frac{1}{2\kappa^2} \int d^4x \sqrt{g} (R + R^2/6M^2) \quad (9)$$

It is well known that this theory is conformally equivalent to a theory of Einstein gravity plus a scalar field [22]. The potential for the resulting scalar is extremely flat for field values $\phi \gg M_4$ and has a minimum at $\phi = 0$ with $V(\phi = 0) = 0$. For large initial values of $\phi$, one can recognize this as an excellent model for chaotic inflation [23].

In general, quantum corrections to the right-hand side of Einstein’s equation in the absence of matter can be written as [24]

$$\langle T_{\mu\nu}\rangle = \left(\frac{k_2}{2880\pi^2}\right)(R^\rho_{\mu\rho\nu} - \frac{2}{3}RR_{\mu\nu} - \frac{1}{2}g_{\mu\nu}R^\rho\sigma R_{\rho\sigma} + \frac{1}{4}g_{\mu\nu}R^2)$$

$$+ \frac{1}{6}\left(\frac{k_3}{2880\pi^2}\right)(2R_{\mu\nu} - 2g_{\mu\nu}R^\rho_{\rho\nu} - 2RR_{\mu\nu} + \frac{1}{2}g_{\mu\nu}R^2) \quad (10)$$

where $k_2$ and $k_3$ are constants that appear in the process of regularization. $k_2$ is related to the number of light spin states, which can be very large in variants of string theories
based on $M$ theory. On the other hand, the coefficient $k_3$ is independent of the number of light states. This term is equivalent to the variation of the $R^2$ term in the effective action. The theory admits a de Sitter solution which can be found from the 00 component of gravitational equation of motion [25]. Defining $H' = 2880\pi^2/k_2$ and $M^2 = 2880\pi^2/k_3$, and setting the spatial curvature $k = 0$, one finds [26]

$$H^2(H^2 - H'^2) = (H'^2/M^2)(2\dot{H}H + 2H^2\dot{H} - \dot{H}^2)$$  \hspace{1cm} (11)$$

where $H$ is the Hubble parameter. The de Sitter solution corresponds to $H = H'$ and of course $\dot{H} = \ddot{H} = 0$.

In order to avoid the overproduction of gravitons there is a lower limit on the parameters $k_{2,3}$ [27, 26]: $k_2 \gtrsim 10^{10}$ implying the need for billions of spin degrees of freedom to be present. While this seems like an inordinately large number, it is possible to generate very large numbers of degrees of freedom in theories with extra dimensions. Particularly if $N \sim M_{\text{GUT}}R_5 \gtrsim 10^8$ [19]. The bound for $k_3$ is $k_3 \gtrsim 10^9$, corresponding to $M \lessapprox 10^{14}$ GeV.

In general $R^2$ corrections to the action do not appear in 5 (11) dimensions. Therefore the first correction is of order $R^4$. Like the curvature squared correction, the action

$$S = \int d^5x \sqrt{G_5}\left\{\frac{M_5^3}{16\pi}R_5 + \alpha M_5^{-3}R_4^4\right\}$$  \hspace{1cm} (12)$$
can be conformally transformed to a 5D Einstein theory with an additional scalar field ($\chi$) and a potential $U(\chi)$ [28]. Making the KK reduction to 4D, we find in addition to $\chi$, a second scalar field ($\phi$) which is the modulus of the 5th dimension and is related to the dilaton. The resulting potential takes the form [19]

$$U(\phi, \chi) \sim M_4^2 M_5^2 e^{\sqrt{2\kappa_5}\phi} e^{-\frac{2\sqrt{2\kappa_5}}{3}\phi} e^{\frac{2\sqrt{2\kappa_5}}{3}\chi} \left(e^{\frac{4\sqrt{2\kappa_5}}{3}\chi} - 1\right)^{\frac{4}{3}}$$  \hspace{1cm} (13)$$

The dilaton potential here, as in most string descendant models, is problematic unless it is fixed by an additional potential. In the absence of a dilaton potential, this model will not inflate. The remaining potential (of $\chi$) differs from that in the $R^2$ model. As one can see it is no longer flat, but rather takes the shape similar to that of the double well potential. As a result, chaotic inflation [23] with a large initial value of $\chi$ is impossible here. Nevertheless there may remain the possibility to realize inflationary expansion in this model by using the potential energy around the local maximum, $V(\chi_m)$, as in the topological inflation scenario of Linde and Vilenkin [29]. In this scenario, if the scalar field $\chi(x)$ is randomly distributed initially with a large dispersion, some part of the universe will roll to $\chi = 0$, while in other parts it will run away to infinity. Between any two such regions there will appear domain walls, containing a large energy density, $\rho \sim V(\chi_m)$. If the wall is thicker than the Hubble radius of this energy density, there will exist a sufficiently large quasi-homogeneous region, filled with large potential energy, where inflationary expansion naturally sets in.
It can be shown [19] that the potential (13) does satisfy the conditions for topological inflation. In addition, there is no problem with the graceful exit. However, the spectrum of density fluctuations can be calculated and the spectral index is given by

\[ n_s = 1 - \frac{m^2}{H_m^2}, \]

where \( m^2 \) and \( H_m^2 \) are the curvature (of the potential) and Hubble parameter determined at the maximum. In this model, \( n_s = 3/8 \). Furthermore, the magnitude of density fluctuations requires \( H_m \sim M_5 \sim 10^{-14} M_4 \) which are probably unacceptably small.

These problems can be remedied by either the inclusion of an \( R^2 \) correction to the action, or by a large tower of KK states. The resolution of the problem can be traced to the magnitude of the Hubble parameter, which is determined by \( H^2 \sim U(\chi)/M_4^2 \). In the presence of a large number of degrees of freedom, quantum corrections modify this relation along the line of eq. (11) in which case \( H \) is driven to \( H' \) and need not be exceptionally small. It suffices now that \( H, M_5 \sim 10^{-5} M_4 \) with \( R^{-1}_5 \sim O(100) \) GeV. In this case, the spectral index is \( \approx 0.95 \) and the magnitude of fluctuations is acceptably small.

The second possibility [30, 31], also makes use of the tower of KK states. The idea of using multiple fields to drive inflation where the parameters of the theory would normally not lead to sufficient inflation is called assisted inflation [32]. For example, it is well known that scalar fields with exponential potentials of the form \( V(\phi) = e^{-\lambda \phi} \), lead to power law expansion with the cosmological scale factor growing as \( R(t) \sim t^p \) and \( p = 2/\lambda^2 \). Density fluctuations are no longer scale invariant but scale as \( |\delta \rho/\rho(k)|^2 \sim k^{n-1} \) with \( n = 1 - 2/p \). Sufficient inflation along with \( n \approx 1 \), requires \( p \) to be large. In [32], it was shown that a system of scalar fields each with exponential potentials, (even if the individual powers, \( p_i \), are not sufficiently large to generate inflation) has an attractor solution in which the universe power-law expands with a power given by \( p = \sum p_i \).

Assistance, can in fact be generalized to other types of potentials [30]. For example, we can consider a general field theory of multiple, self-interacting scalar fields of the form

\[
-\mathcal{L} = \sum_{i=1}^{N} \left\{ \frac{1}{2} (\partial \phi_i)^2 + \frac{m_i^2}{2} \phi_i^2 + \frac{\lambda_3}{3!} \phi_i^3 + \frac{\lambda_4}{4!} \phi_i^4 \right\} + \sum_{i=1}^{N} \left\{ \frac{\lambda_3}{3!} \phi_i^3 + \frac{\lambda_4}{4!} \phi_i^4 \right\} .
\]

This system consists of \( N \) completely equivalent, decoupled scalar fields. As a result, the Lagrangian can be written as

\[
-\mathcal{L} = N \left\{ \frac{1}{2} (\partial \bar{\phi})^2 + \frac{m^2}{2} \bar{\phi}^2 + \frac{\tilde{\lambda}_3}{3!} \bar{\phi}^3 + \frac{\tilde{\lambda}_4}{4!} \bar{\phi}^4 \right\}
\]

\[
= \frac{1}{2} (\partial \bar{\phi})^2 + \frac{m^2}{2} \bar{\phi}^2 + \frac{\tilde{\lambda}_3}{3!} \bar{\phi}^3 + \frac{\tilde{\lambda}_4}{4!} \bar{\phi}^4 ,
\]

where

\[
\bar{\phi} = \sqrt{N} \phi_1 , \quad \tilde{\lambda}_3 = \frac{\lambda_3}{\sqrt{N}} , \quad \tilde{\lambda}_4 = \frac{\lambda_4}{N} .
\]
The resulting theory describes a single scalar field with the same type of self-interactions compared to the fields in the original theory. However, these self-interactions are considerably weaker since both of the coupling constants now scale with the number of scalar fields \( N \). As a result, as the number of scalar fields that we include in the theory becomes larger, the coupling constants become smaller and the corresponding fine-tuning becomes milder. Thus the same basic idea expounded in \([32, 33]\) carries over very simply to chaotic inflation \([23]\) based on a quartic potential. While \( \hat{\lambda} \) must still be of order \( 10^{-12} \), the fundamental coupling in the theory \( \lambda \) can now be much larger if \( N \) is large. In addition, the problem associated with chaotic inflation \([34]\), large \( (\phi \gg M_p) \) initial field values and the necessary fine tuning of non-renormalizable interactions, is cured \([31]\). Because of the rescaling in eq. \((16)\), in a theory of multiple fields, chaotic inflation will operate, even though none of the fields have expectation values greater than \( M_p \).

In \([30, 31]\), it was suggested that the source of the multiple fields is the KK reduction from a higher dimensional theory to 4D. For example, the reduction of a 5D theory results in \( N \simeq R_5 M_5 \) fields. Though these fields are not decoupled as in the example above, and hence assistance is not guaranteed \([30, 35]\), attractor solutions do exist and assistance is generated primarily due to the rescaling of the fields (from 5D to 4D) which is necessary due to the dimensional reduction. To see this, consider the action

\[
S_\phi = -\int d^5x \sqrt{G_5} \left\{ G_5^{AB} \partial_A \hat{\phi} \partial_B \hat{\phi} + \frac{\hat{\lambda}}{4!M_5} \hat{\phi}^4 \right\}
\]  

(17)

When KK-reduced to 4D, \( \sqrt{G_5} \rightarrow \sqrt{G_4} R_5 \) and canonical 4D scalar fields must be defined by \( \phi = \sqrt{2R_5} \hat{\phi} \) and hence the quartic potential becomes \( \lambda \phi^4 / 4! \) with \( \lambda = \hat{\lambda}/2R_5M_5 = \hat{\lambda}/N \).

4 Baryogenesis

The exact time period for baryogenesis is not known but most certainly it is completed no later than the electroweak phase transition. There are of course many mechanisms for baryogenesis which have been proposed in the literature \([36]\), too many to be comprehensive here. All require baryon number violation, C and CP violation, and a departure from thermal equilibrium \([37]\). The original out-of-equilibrium decay (OOED) scenario \([38]\) is probably still the simplest of all mechanisms. Originally formulated in the context of grand unified theories, OOED involved the decay of a superheavy gauge or Higgs boson with baryon number violating couplings. For example, the \( X \) gauge boson of SU(5) couples to both a \( \bar{u}u \) pair and \( (e^-d) \). The decay rate for \( X \) will be

\[
\Gamma_D \simeq \alpha M_X
\]  

(18)
However decays can only begin occurring when the age of the Universe is longer than the $X$ lifetime $\Gamma_D^{-1}$, i.e., when $\Gamma_D > H$

$$\alpha M_X \gtrsim N(T)^{1/2} T^2 / M_P$$

(19)

The out-of-equilibrium condition is that at $T = M_X, \Gamma_D < H$ or

$$M_X \gtrsim \alpha M_P (N(M_X))^{-1/2} \sim 10^{18} \, \text{GeV}$$

(20)

In this case, we would expect a maximal net baryon asymmetry to be produced, $n_B/s \sim 10^{-2}\epsilon$, where $\epsilon$ is a measure of the CP violation in the decay.

In the context of inflation, OOED requires either strong reheating (or preheating) [39], or that the decaying particles have masses less than the inflaton mass so that they can be produced (necessarily out of equilibrium) by inflaton decays. For example, if the inflaton potential carries only a single scale which is fixed by the magnitude of density fluctuation as measured in the microwave background radiation, then we can write [40]

$$V(\eta) = \mu^4 (1 + O(\eta/M_P)^n + \cdots)$$

(21)

Typically,

$$\frac{\delta \rho}{\rho} \sim O(100) \frac{\mu^2}{M_P^2}$$

(22)

so that

$$\frac{\mu^2}{M_P^2} = \text{few} \times 10^{-8}$$

(23)

In this case a relatively light Higgs is necessary since the inflaton is typically light ($m_\eta \sim \mu^2/M_P \sim O(10^{11}) \, \text{GeV}$) and the baryon number violating Higgs would have to be produced during inflaton decay. Clever model building could allow for such a light Higgs, even in the context of supersymmetry [41]. In this case, the baryon asymmetry is given simply by

$$\frac{n_B}{s} \sim \epsilon \frac{n_\eta}{T_R^3} \sim \epsilon \frac{\rho_\eta}{m_\eta} \sim \epsilon \left( \frac{m_\eta}{M_P} \right)^{1/2} \sim \epsilon \frac{\mu}{M_P} \sim 10^{-4}\epsilon$$

(24)

where $T_R$ is the reheat temperature after inflation, $n_\eta = \rho_\eta/m_\eta \sim \Gamma^2 M_P^2 / m_\eta$, and $\Gamma = m_\eta^3 / M_P^2$ is the inflaton decay rate.

In the context of supersymmetry, there is an extremely natural mechanism for the generation of the baryon asymmetry utilizing flat directions of the scalar potential [42]. One can show that there are many directions in field space such that the scalar potential vanishes identically when SUSY is unbroken. That is,

$$V = |F|^2 + |D|^2 = 0$$

(25)
One such example is

\[ u_3^c = a \quad s_2^c = a \quad -u_1 = v \quad \mu^* = v \quad b_1^c = e^{i\phi}\sqrt{v^2 + a^2} \]  \(26\)

where \(a, v\) are arbitrary complex vacuum expectation values. SUSY breaking lifts this degeneracy so that

\[ V \simeq \tilde{m}^2 \phi^2 \]  \(27\)

where \(\tilde{m}\) is the SUSY breaking scale and \(\phi\) is the direction in field space corresponding to the flat direction. For large initial values of \(\phi\), \(\phi_o \sim M_{\text{gut}}\), a large baryon asymmetry can be generated\[42, 43\]. This requires the presence of baryon number violating operators such as \(O = qqql\) such that \(\langle O \rangle \neq 0\). The decay of these condensates through such an operator can lead to a net baryon asymmetry.

When combined with inflation, it is important to verify that the AD flat directions remain flat. In general, during inflation, supersymmetry is broken. The gravitino mass is related to the vacuum energy and \(m_{3/2}^2 \sim V/M_P^2 \sim H^2\), thus lifting the flat directions and potentially preventing the realization of the AD scenario as argued in \[44\]. To see this, consider a minimal supergravity model whose Kähler potential is defined by

\[ G = zz^* + \phi_i^* \phi^i + \ln |W(z) + W(\phi)|^2 \]  \(28\)

where \(z\) is a Polonyi-like field\[45\] needed to break supergravity, and we denote the scalar components of the usual matter chiral supermultiplets by \(\phi^i\). \(W\) and \(\overline{W}\) are the superpotentials of \(\phi^i\) and \(z\) respectively. In this case, the scalar potential becomes

\[ V = e^{z^* + \phi_i^* \phi^i} \left[ |W_z + z^*(\overline{W} + W)|^2 + |W_{\phi^i} + \phi_i^* (\overline{W} + W)|^2 - 3|(|\overline{W} + W)|^2 \right] \]  \(29\)

Included in the above expression for \(V\), one finds a mass term for the matter fields \(\phi^i\), \(e^G \phi_i^* \phi^i = m_{3/2}^2 \phi_i^* \phi^i\) \[40\]. As it applies to all scalar fields (in the matter sector), all flat directions are lifted by it as well. The above arguments can be generalized to supergravity models with non-minimal Kähler potentials.

There is however a special class of models called no-scale supergravity models, that were first introduced in \[47\] and have the remarkable property that the gravitino mass is undetermined at the tree level despite the fact that supergravity is broken. No-scale supergravity has been used heavily in constructing supergravity models in which all mass scales below the Planck scale are determined radiatively \[48, 49\]. These models emerge naturally in torus \[50\] or, for the untwisted sector, orbifold \[51\] compactifications of the heterotic string.

In no-scale models (or more generally in models which possess a Heisenberg symmetry \[52\]), the Kähler potential becomes

\[ G = f(z + z^* - \phi_i^* \phi^i) + \ln |W(\phi)|^2 \]  \(30\)

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Now, one can write
\[ V = e^{f(n)} \left[ \left( \frac{f''}{f'} - 3 \right) |W|^2 - \frac{1}{f'} |W_i|^2 \right] \]  
(31)

It is important to notice that the cross term $|\phi_i^* W|^2$ has disappeared in the scalar potential. Because of the absence of the cross term, flat directions remain flat even during inflation [53]. The no-scale model corresponds to $f = -3 \ln \eta$, $f'^2 = 3 f''$ and the first term in (31) vanishes. The potential then takes the form
\[ V = \left[ \frac{1}{3} e^{\frac{2}{3}} |W_i|^2 \right], \]  
(32)

which is positive definite. The requirement that the vacuum energy vanishes implies $\langle W_i \rangle = \langle g_a \rangle = 0$ at the minimum. As a consequence $\eta$ is undetermined and so is the gravitino mass $m_{3/2}(\eta)$.

The above argument is only valid at the tree level. An explicit one-loop calculation [54] shows that the effective potential along the flat direction has the form
\[ V_{eff} \sim \frac{g^2}{(4\pi)^2} \langle V \rangle \left( -2 \phi^2 \log \left( \frac{\Lambda^2}{g^2 \phi^2} \right) + \phi^2 \right) + O(\langle V \rangle)^2, \]  
(33)

where $\Lambda$ is the cutoff of the effective supergravity theory, and has a minimum around $\phi \simeq 0.5 \Lambda$. Thus, $\phi_0 \sim M_P$ will be generated and in this case the subsequent sfermion oscillations will dominate the energy density and a baryon asymmetry will result which is independent of inflationary parameters as originally discussed in [42, 43] and will produce $n_B/s \sim O(1)$. Thus we are left with the problem that the baryon asymmetry in no-scale type models is too large [55, 53, 56].

In [53], several possible solutions were presented to dilute the baryon asymmetry. These included 1) Moduli decay, 2) the presence of non-renormalizable interactions, and 3) electroweak effects. Moduli decay in this context, turns out to be insufficient to bring an initial asymmetry of order $n_B/s \sim 1$ down to acceptable levels. However, as a by-product one can show that there is no moduli problem [57] either. In contrast, adding non-renormalizable Planck scale operators of the form $g^{2n-2} / M_P^{2n-6}$ leads to a smaller initial value for $\phi_0$ and hence a smaller value for $n_B/s$. For dimension 6 operators ($n = 4$), a baryon asymmetry of order $n_B/s \sim 10^{-10}$ is produced. Finally, another possible suppression mechanism is to employ the smallness of the fermion masses. The baryon asymmetry is known to be wiped out if the net $B - L$ asymmetry vanishes because of the sphaleron transitions at high temperature. However, Kuzmin, Rubakov and Shaposhnikov [58] pointed out that this erasure can be partially circumvented if the individual $(B - 3L_i)$ asymmetries, where $i = 1, 2, 3$ refers to three generations, do not vanish even when the total asymmetry vanishes. Even though there is still a tendency that the baryon asymmetry is erased by the chemical equilibrium due to the sphaleron transitions, the finite mass of the tau lepton
shifts the chemical equilibrium between $B$ and $L_3$ towards the $B$ side and leaves a finite asymmetry in the end. Their estimate is

$$B = -\frac{4}{13} \sum_i \left( L_i - \frac{1}{3} B \right) \left( 1 + \frac{m_i^2}{\pi^2 T^2} \right)$$

(34)

where the temperature $T \sim T_C \sim 200$ GeV is when the sphaleron transition freezes out (similar to the temperature of the electroweak phase transition) and $m_\tau(T)$ is expected to be somewhat smaller than $m_\tau(0) = 1.777$ GeV. Overall, the sphaleron transition suppresses the baryon asymmetry by a factor of $\sim 10^{-6}$. This suppression factor is sufficient to keep the total baryon asymmetry at a reasonable order of magnitude in many of the cases discussed above.

Finally, it is necessary to mention one other extremely simple mechanism based on the OOED of a heavy Majorana neutrino \[59\]. This mechanism does not require grand unification at all. By simply adding to the Lagrangian a Dirac and Majorana mass term for a new right handed neutrino state,

$$L \ni M \nu^c \nu^c + \lambda H L \nu^c$$

(35)

the out-of-equilibrium decays $\nu^c \rightarrow L + H^*$ and $\nu^c \rightarrow L^* + H$ will generate a non-zero lepton number $L \neq 0$. The out-of-equilibrium condition for these decays translates to $10^{-3} \lambda^2 M_P < M$ and $M$ could be as low as $O(10)$ TeV. (Note that once again in order to have a non-vanishing contribution to the C and CP violation in this process at 1-loop, at least 2 flavors of $\nu^c$ are required. For the generation of masses of all three neutrino flavors, 3 flavors of $\nu^c$ are required.) Sphaleron effects can transfer this lepton asymmetry into a baryon asymmetry since now $B - L \neq 0$. A supersymmetric version of this scenario has also been described \[40, 60\].

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As the Inner Space/Outer Space II workshop, and these proceedings are in memory of Dave Schramm, it is fitting to acknowledge Dave’s role. Dave’s research interests were predominantly concerned with the epoch of Big Bang Nucleosynthesis and the post BBN Universe. His research interests are known to have been extremely diverse covering such areas as chemical evolution, cosmic rays, cosmochronology, dark matter, galaxy formation and mergers, the gamma-ray background, magnetic fields, mass extinctions, neutrinos, (late-time) phase transitions, supernovae, ... In addition, he was a master at synthesizing and finding relationships between these topics. Though pre-BBN cosmology was not Dave’s mainstay, he did of course make important contributions in areas such as baryogenesis and topological defects. But Dave’s most important contribution was the early recognition of the impact of cosmology on particle physics. He can legitimately be considered a founding father of astro-particle physics. On a more personal note, it is difficult to put into words
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